Electronic spin precession and interferometry from spin-orbital entanglement in a double quantum dot

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A double quantum dot inserted in parallel between two metallic leads can entangle the electron spin with the orbital (dot index) degree of freedom. An Aharonov-Bohm orbital phase can be transferred to the spinor wavefunction, providing a geometrical control of the spin precession around a fixed magnetic field. A fully coherent behavior occurs in a mixed orbital/spin Kondo regime. Evidence for the spin precession can be obtained, either using spin-polarized metallic leads or by placing the double dot in one branch of a metallic loop.

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Control of the electron spin is important for the realization of novel nanoelectronic devices for spintronics or for quantum information processing. In the latter case, manipulation of individual spins is necessary. The use of time-dependent gates, put forward some years ago \[\text{[1]},\] has seen considerable progress \[\text{[2]}.\] Yet another possibility is to build single or two-spin operations (gates) into a given device geometry, with static control parameters only. This may allow faster processing speed and facilitate integration into more complex devices. One way of controlling the spin in quantum dots is through energy filtering by applied gate voltages, as proposed for spin entanglement \[\text{[3]},\] teleportation \[\text{[4]},\] and spin filtering \[\text{[5]}].

Spin precession has also been put forward in metallic electronics which have not yet been realized with single electrons. The basic unit is a double quantum dot in parallel, coupled to metallic leads. As previously shown in \[\text{[6]},\] the basic set-up in a loop enclosing an Aharonov-Bohm (AB) flux gives rise to new physics. Indeed, an incoming state \[|\uparrow\rangle + |\downarrow\rangle\] becomes \[|\uparrow\rangle \otimes |\uparrow\rangle + be^{i\varphi}|\downarrow\rangle \otimes |\downarrow\rangle\] at the output where \[\varphi = 2\pi \frac{\phi}{\hbar}\] is the AB phase between the two branches. The intermediate state involves entanglement between spin and orbital degrees of freedom, thus applying an orbital phase causes a rotation of the spinor around the fixed magnetic field. This permits a purely geometrical control of the spin precession angle, by static parameters such as the gate voltages and the AB flux. This set-up indeed belongs to the class of "Stern-Gerlach interferometers," first considered as gedanken experiments \[\text{[7]},\] then realized with neutron interferometry \[\text{[8]}\]. The success of spin precession relies on quantum coherence between the two branches, and no "which-path" information being gained. Thus our proposal based on a nanoelectronic device provides a very sensitive test of decoherence effects.

Let us consider two small quantum dots in parallel, keeping only one orbital level, with charge number states \(0, 1, 2\) in each of them. Under an applied Zeeman field, it has been shown \[\text{[9]}\] that a single dot may filter spins "up" or "down," depending on the applied gate voltage \(V_{g\nu}(\nu = 1, 2)\) via the capacitances \(C_{g\nu}\). Indeed, in a resonant regime, transitions between number states \(0, 1\) involve spins \(\uparrow\) only (states \(0, \uparrow\)) while transitions between number states \(1, 2\) instead involve spins \(\downarrow\) only (states \(\uparrow, \downarrow\)). We assume the two dots to be coupled by a large capacitance \(C_0\) and label the double-dot (DD) states as \((\Psi_1, \Psi_2)\). The gate voltages are chosen such that the lowest-lying states be \((\uparrow, \uparrow)\) and \((0, \downarrow)\) and be degenerate. Single-electron transitions from \(0\) to \(1\) the leads involve higher-energy states such as \((0, \uparrow)\) and \((\uparrow, \downarrow)\). One can achieve a Kondo regime where the resonance between states \((\uparrow, \uparrow)\) and \((0, \downarrow)\) involves spin-up transitions in

![Schematic representation of the proposed setup](image-url)

**FIG. 1:** Left: Schematic representation of the proposed setup: two small quantum dots coupled by a capacitance \(C_0\) and connected to left and right partially polarized reservoirs. Depending on the choice of gate voltages, the upper branch filters spins up and the lower ones down, or vice versa. A magnetic flux \(\phi\) threads the whole device. Right: The polarization axis \(\hat{n}_L\) and \(\hat{n}_R\) make an angle \(\theta/2\) with the \(\hat{e}_z\) axis and \(\vec{B} = B\hat{e}_z\).
dot 1 and spin-down transitions in dot 2. This is an orbital/spin Kondo effect entangling spin and orbital degrees of freedom. It achieves a novel sector of the Kondo physics, completing the already existing pure spin Kondo \cite{10} and pure orbital Kondo \cite{11, 12} ones. Here, both dots are connected to the same left and right lead with tunneling hopping parameters $t_{L\nu}$ and $t_{R\nu}$. In view of a direct detection of the spin precession, we allow for spin-polarization in the leads. The corresponding tunnel junction capacitances are $C_L$ (left), $C_R$ (right). A magnetic field $B$, oriented in the plane of the setup, is applied to each dot. The optimum Kondo regime is reached at the symmetric point where the lowest-lying excited states $(0, \uparrow)$, $(\uparrow, \downarrow)$ involve the same charge excitation energy $E_c = \frac{e^2 C_0}{2C(C+2C_0)}$ where $C = C_L + C_R + C_g$ (we assume $C_g = C_L \approx C_R$). The isolated DD system may be described \cite{13} at low energy by

$$H_{dot} = -\delta ET^z - tT^z - g\mu_B(B_1S^z_1 + B_2S^z_2),$$

where we have defined the orbital pseudospin $T^z = (n_1 - n_2 + 1)/2 = \pm 1/2$ from the charge occupations $n_\nu$. Here $E_c = \frac{e^2 C_0}{2C(C+2C_0)}[C_R(V_{g2} - V_{g1}) - \epsilon]$. The second term in $H_{dot}$ represents a small parasitic tunneling amplitude between the dots \cite{2}. The last term expresses the effective Zeeman splitting. Due to exchange contributions with the leads \cite{14}, the local fields $B_\nu$ in the dots may be different. Notice that a large level spacing $\delta E$ (or equivalently small quantum dots) is necessary to eliminate the triplet states $(0, \downarrow)$. Under the condition $\delta E = \frac{1}{2}g\mu_B(B_1 + B_2)$, the states $(\uparrow, \uparrow)$ and $(0, \downarrow)$ are degenerate. The total spin $S^z = S^z_1 + S^z_2 - \frac{1}{2}$ is entangled with the orbital pseudospin $T^z$, e.g. a spin flip is locked to an orbital pseudo-spin flip. Therefore the Kondo screening of the spin involves spin-up electrons in branch 1 and spin-down electrons in branch 2. Notice that the gate voltage difference compensates for the Zeeman splitting between spins up in dot 1 and down in dot 2, and no splitting of the Kondo zero-bias conductance peak occurs.

The system Hamiltonian is $H = H_{leads} + H_{tun} + H_{dot}$. The leads are described by $H_{leads} = \sum_{k,\gamma,\pm} \varepsilon_{k,\pm} c^\dagger_{k,\gamma,\pm} c_{k,\gamma,\pm}$, where $c^\dagger_{k,\gamma,\pm}$ creates an electron with energy $\varepsilon_{k,\pm}$ in lead $\gamma = L, R$ with spin along $\uparrow$ or $\downarrow$. Spin polarization in the leads would result in a spin asymmetry in the density of states $\rho_{\gamma,\pm}(\omega)$. We further neglect the energy dependence in the density of states and also suppose $\rho_{\gamma,\pm} \approx \rho_{\pm}$. The ratio $\rho = (\rho_+ - \rho_-)/(\rho_+ + \rho_-)$ denotes the degree of spin polarization in the leads, which may have a noncollinear polarization. The tunneling junctions between the leads and the dots are described by $H_{tun} = \sum_{k,\gamma,\nu,\pm,\pm,\pm} (t_{\gamma,\nu,\pm} c^\dagger_{k,\gamma,\nu,\pm} d_{\nu,\pm} + H.c.)$ where $d_{\nu,\pm}$ destroys an electron in dot $\nu = 1, 2$ with spin $s = \pm$.

In addition to the applied field, the polarized electrodes may also generate effective magnetic fields $B_{eff,\nu}$. They depend on $p$, on the dot internal parameters and gate voltages, and on the tunnel-coupling strength $\Gamma_{\nu} = \Gamma_{LV} + \Gamma_{RV}$ where $\Gamma_{\gamma\nu} = \pi \sum_{t_{\gamma\nu}} |t_{k,\gamma\nu}|^2 \rho_{\gamma}$. An explicit mean-field calculation of $B_{eff,\nu}$ is derived in \cite{14} where it was shown that $B_{eff,\nu}$ is zero for a particle-hole symmetric situation. In what follows, we assume that the fields $B_{eff,\nu}$ can be neglected. We suppose that the lead magnetization axis lies in the dot plane $(\vec{e}_x, \vec{e}_y)$, with a relative angle $\theta$ (see Fig. 1). Following \cite{14}, it is convenient to quantize the dot spin along the $B$ axis corresponding to $\vec{e}_z = (\vec{n}_L \times \vec{n}_R)$, the other vector components being $e_x = (\vec{n}_L + \vec{n}_R)/|\vec{n}_L + \vec{n}_R|$, $e_y = (\vec{n}_L - \vec{n}_R)/|\vec{n}_L - \vec{n}_R|$. In this rotated basis, the tunneling Hamiltonian then reads \cite{14}

$$H_{tun} = \sum_{k,\nu} \left( t_{LV} e^{i(-1)^\nu\pi/4} c^\dagger_{k,L,\sigma} d_{\nu,\sigma} + H.c. \right) (2)$$

$$+ (L \rightarrow R, \varphi \rightarrow -\varphi).$$

The $c^\dagger_{k,L,\pm}$ are linear combinations of the $c^\dagger_{k,L,\uparrow/\downarrow}$ and are related by $(c^\dagger_{k,L,R,+}, c^\dagger_{k,L,R,-}) = (c^\dagger_{k,L,R,+}, c^\dagger_{k,L,R,-}) e^{i\delta\theta\sigma+\pi/4}$. In order to determine the effective coupling between the double dot and the leads, we consider virtual excitation to both excited states $(0, \uparrow)$ and $(\uparrow, \uparrow)$ generated by $H_{tun}$. Using a Schrieffer-Wolf (SW) transformation, the Kondo Hamiltonian $H_K$ is obtained:

$$H_K = \frac{1}{2} T^z (J_{LL} e^{-i\varphi/2} \psi^\dagger_L \psi_L + J_{RR} e^{i\varphi/2} \psi^\dagger_R \psi_R)$$

$$+ \frac{1}{2} T^z (J_{LR} \psi^\dagger_L \psi_R + J_{RL} \psi^\dagger_R \psi_L) + H.c.$$

$$+ \frac{1}{2} T^z (J_{LL} \psi^\dagger_L \psi_L + J_{RR} \psi^\dagger_R \psi_R)$$

$$+ \frac{1}{2} T^z (J_{LR} \psi^\dagger_L \psi_R - J_{RL} \psi^\dagger_R \psi_L)$$

$$+ \frac{1}{2} T^z \left( e^{-i\varphi/2} J_{LR} \psi^\dagger_L \psi_R - e^{i\varphi/2} J_{LR} \psi^\dagger_R \psi_L + H.c. \right),$$

where $\psi^\dagger_L/R, \sigma = \sum_k c^\dagger_{k,L/R,\sigma}$ and $T^z = d^\dagger_{1L} d_{1R}$ flips both the spin and the orbital pseudo-spin. We have introduced several Kondo couplings $J_{LL/RR} \approx \frac{t_{LL/RR}}{\mu_B \hbar}$ and $J_{LR} \approx t_{LR} \sqrt{t_{LL} t_{RR}}$ with $\beta, \gamma = L, R$. As usual the $k$-dependence of the Kondo couplings is neglected. The $J_{LL/RR}$ Kondo couplings are in general spin-dependent, due to intrinsic asymmetries in the two branches, except when $t_{\gamma1} \sim t_{\gamma2} \sim t_{\gamma}$. However, due to the entanglement of orbital and spin degrees of freedom, a geometrical asymmetry is easily compensated with an orbital field, i.e. with a fine-tuning of the dot gate voltages $V_{g1}$ and $V_{g2}$. If not stated, $t_{\gamma1} \sim t_{\gamma2} \sim t_{\gamma}$ is therefore assumed in the following. Other cotunneling terms involving higher energy states like $(\uparrow, \uparrow)$, turn out to be irrelevant under renormalization group (RG) in the low-energy limit and thus do not spoil the spin filtering. In the unitary limit, the spin $S$ or equivalently the pseudo-spin $T$ is completely screened and an entangled spin/orbital singlet is formed with the left and right electrodes. We can get rid of the phase in the Hamiltonian in eq. (3) by defining the new
rotated basis \( (\psi_{L/R}^\dagger \psi_{L/R}^\dagger) = (\psi_{L/R}^\dagger \psi_{L/R}^\dagger) R^2(\pm \varphi/2) \). In this spin-rotated basis, the Kondo Hamiltonian takes the simpler form: \( H_K = \sum_{\beta, \gamma, s, s', \pm} J_{\beta \gamma} \psi_{\beta s}^\dagger \tilde{T}_{ss'}^\dagger \psi_{\gamma s'} \). The AB phase \( \varphi \) has disappeared from the Kondo Hamiltonian and has been swapped onto the spin direction in the source and lead. In this basis the angles \( \theta \) and \( \varphi \) therefore play a similar role \[12\].

For weak polarization, the Kondo temperature is well approximated by \( T_K \approx \Delta \exp\left(-\frac{1}{\rho R L + V_{DD}}\right) \) with \( \Delta \approx \delta \epsilon \) the dot level spacing \[10\]. At \( T < T_K \), the system reaches the unitary limit \[10\] and the \( T = 0 \) transmission amplitude \( \mathcal{T}_\sigma \) reads:

\[
\mathcal{T}_\sigma(\epsilon_F) = 2i \frac{t_L t_R}{t_L^2 + t_R^2} e^{i\delta_\sigma(\epsilon_F)} \sin(\delta_\sigma(\epsilon_F)),
\]

where \( \delta_\sigma(\epsilon_F) = \delta(\epsilon_F) = \delta \sim \pi/2 \) is the Kondo phase shift and \( \epsilon_F \) the Fermi energy. Assume one prepares an incoming state \(|\eta\rangle_{in}\) in a superposition of \( \uparrow \) and \( \downarrow \) states, \(|\eta\rangle_{in} = a|\uparrow\rangle_L + b|\downarrow\rangle_L\), the outgoing state reads

\[
|\eta\rangle_{out} = 2i \frac{t_L t_R}{t_L^2 + t_R^2} e^{i(\delta - \varphi/2)} \left(a|\uparrow\rangle_R + e^{i\varphi} b|\downarrow\rangle_R\right).
\]

One sees that by manipulating the AB flux, one can perform a coherent precession of the spin of the incoming wave function along the magnetic field axis. This effect has a signature in the low-temperature conductance through the DD:

\[
G_1 = G_0 \left(1 - \frac{2p^2}{1 + p^2} \sin^2 \left(\frac{\theta - \varphi}{2}\right)\right),
\]

where \( G_0 = \frac{\epsilon^2}{2h} \frac{4e^2}{(1 + p^2)} \) is the maximum conductance obtained for parallel lead polarization. In Eq. \[6\], the conductance reaches its maximum value for \( \varphi = \theta \), showing that a small AB flux compensates the mismatch of non-collinear polarizations.

The above results were derived assuming \( T = 0 \) and neglecting environmental fluctuations inherent to any mesoscopic setup. While the DD operates close to the unitary regime \( T < T_K \), inelastic processes induced by a finite \( T \) are rather small and their main effects are to reduce the amplitude of the outgoing wave vector and give a dephasing time \( \tau_\varphi^{-1} \sim T^2 \) \[17\]. Another source of decoherence is brought by the circuit electromagnetic fluctuations, that couple to tunneling events to and from each of the dots. We have used an equivalent circuit representation of our Stern-Gerlach interferometer \( \text{(SGI)} \) following Ref. \[18\]. By modeling the environment by an external impedance \( Z \), we have shown using a RG analysis that only a large impedance of order \( R_K \) is able to destroy the global coherence of the electronic SGI. This will be detailed elsewhere \[19\].

The above proposal of detecting the precession directly by spin filtering the source and drain electrodes is similar to polarized neutron interferometry experiments \[9\]. Still, it might be difficult to control the spin filtering in the leads, as high Zeeman fields are necessary for the orbital/spin Kondo effect. These fields can be estimated as \( B \sim \text{few} \ T \) for GaAs-based dots, and one order of magnitude less for InAs dots owing to the much larger g-factors.

Alternatively, there exists another way of detecting the precession. Suppose one inserts the DD in the branch (II) of a larger loop including a flux \( \Phi \) \[20\] (Figure 2). The whole loop is supposed to be phase-coherent. The transmission through the upper branch (I), as well as the scattering matrices at the loop extremities do not involve any spin dependence. AB interference in the large loop amounts to adding a spin-conserving amplitude in branch (I) and a spin-precessing amplitude in branch (II).

Consider first for simplicity an open AB interferometer \( \text{(ABI)} \) realizing the equivalent of Young’s double-slit experiments (see Ref. \[21\] for an experimental realization) and assume perfect transmission probability in both branches (this implies \( t_L = t_R \) and \( \delta_\sigma = \pi/2 \)). The conductance then reads

\[
G = \frac{e^2}{h} \left(1 + \cos(\varphi/2) \sin(2\pi e \Phi/h)\right).
\]

Two different periodicities appear with the orbital field \( B_o \) responsible for the fluxes. First, a fast oscillation, of period \( \frac{2\pi}{e S} \), is mainly due to the flux in the section \( S \) of the large loop. Second, a slow one comes from the flux in the small DD loop of area \( s \). Its period is \( \varphi = 4\pi \) instead of \( 2\pi \), owing to the spinor nature of the wavefunction. This results in slower (beating) oscillations, with period \( \frac{2\pi}{e S} \). Notice that the visibility of the fast AB oscillations is minimal for a rotation of \( \varphi = 2\pi \), for which the spinor changes its sign. This very striking consequence of quantum mechanics was proposed \[22\] and verified \[23\] with neutron interferometry. Let us for example compare this result to the situation in which the dots are not capacitively coupled and are tuned independently in the Kondo regime. The conductance then reads

\[
G = \frac{e^2}{2h} (1 + \cos^2(\varphi/2) + 2 \sin(2\pi e \Phi/h) \cos(\varphi/2)).
\]

The conductance is also 4\(\pi\)-periodic in \( \varphi \) but here the 4\(\pi\) period can be traced back to the spin-independent interference between the upper arm of the large ABI and both branches in the lower arm. Nevertheless, this expression of the conductance is clearly different from the one in Eq. \[7\] and precession can in principle be detected.
Another striking consequence of precession in branch II is the spin polarization in the output, even when the incoming electrons are not polarized. Indeed one finds that \( \langle S_z \rangle = -\cos(\pi e(2\Phi)/h) \sin(\varphi/2) \) while \( \langle S^2 \rangle = 0 \) without spin precession. Maximum polarization comes from a destructive interference occurring for one spin direction only. Testing the latter prediction requires spin filtering only in the output and should be easier than the previous test based on Fig. 1.

For a closed interferometer, one may also try to compare the conductance in the large ABI obtained when the spin is precessing (with a SGI in the lower arm) to a reference case where no such precession is present (like for two independent quantum dots in the Kondo regime). We have described both forks of the large ABI by 3 × 3 S-matrices as in Ref. [24]. Nevertheless, one does not observe clear signatures through the conductance. Furthermore, the latter strongly depends on our choice of parameters entering the S-matrices and we did not find any universal feature able to unambiguously distinguish between the two situations. It is therefore preferable to use an open large ABI.

Let us briefly discuss the feasibility of this proposal. Concerning the first experiment, one possibility is to use ferromagnetic semiconductors (Ga,Mn)As as lead electrodes coupled to InAs quantum dots [25]. Another promising route is to use carbon nanotube (CN) quantum dots where Kondo effects have been shown [26]. They can be coupled to ferromagnetic electrodes and large magnetoresistance effects have been observed recently [27]. Concerning the second experiment, two dots in parallel can be fabricated and inserted in an AB-loop [28], and a strong mutual capacitive coupling could be achieved (the residual tunneling amplitude only needs to be smaller than the Zeeman energy [21]). This does not require any lead polarization if one searches for beating effects in the AB interferences. This latter experimental proposal is achievable and may be easier than the first one.

In summary, we have shown how spin interferometry can be performed using orbital/spin entanglement in a double dot. The unitary transmission obtained in the Kondo regime is accompanied by spin precession, leading to a novel periodicity in an AB experiment. The authors acknowledge useful discussions with G. Zarand, P. Nozières, C. Balseiro and T. Kontos. This work was partially supported by the contract P.Nano “QuSpins” of Agence Nationale de la Recherche.

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