Magnetic field switching in parallel quantum dots

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Abstract – We show that the Coulomb blockade in parallel dots pierced by magnetic flux Φ completely blocks the resonant current for any value of Φ except for integer multiples of the flux quantum Φ₀. This non-analytic (switching) dependence of the current on Φ arises only when the dot states that carry the current are of the same energy. The time needed to reach the steady state, however, diverges when Φ → nΦ₀.

The system of two quantum dots coupled in parallel to two reservoirs has attracted a great deal of attention as a realization of a mesoscopic Aharonov-Bohm interferometer [1–3]. Indeed such a system pierced by an external magnetic field (fig. 1) is an interference device whose transmission can be tuned by varying the magnetic field. In the absence of the interdot electron-electron interaction, the interference effects in the resonant current through this system are quite transparent. This is not the case, however, for interacting electrons [4,5].

Consider for instance a strong interdot electron-electron repulsion — a Coulomb blockade. While the two dots may be occupied simultaneously in the non-interacting model, the Coulomb blockade prevents this. At first sight one might not expect that this repulsion could dramatically modify the resonant current’s dependence on the magnetic field. We find, however, that the resonant current is completely blocked for any value of the magnetic flux except for integer multiples of the flux quantum Φ₀ = ħ/e. This striking effect has so far been overlooked in the literature, despite the current-switch in tunnel-coupled quantum dots, induced by coherent radiation of two microwaves, was proposed earlier [6].

In this letter we show how such a switching effect can be seen in the exact analytical solution for certain symmetric choices of the quantum dots’ parameters. This solution can be found in the infinite bias limit, which we justify by comparing the finite- and infinite-bias cases for non-interacting electrons. We then present an explanation of the switching effect as a self-trapping phenomenon in non-linear systems. This explanation allows us to formulate general conditions for this phenomenon and clearly displays its generic nature.

Consider a double dot (DD) connected in parallel to two reservoirs, as shown in fig. 1. For simplicity, we consider spinless electrons. We also assume that each of the dots contains only one level, E₁ and E₂, respectively. In the presence of a magnetic field, the system can be described by the following tunneling Hamiltonian:

$$H = H₀ + H_T + \sum_{\mu=1,2} E_\mu d_\mu d_\mu + U d_1 d_1 d_2 d_2.$$  (1)
Here $H_0 = \sum_k |E_{kL}a_{kL}^{\dagger}a_{kL} + E_{kR}a_{kR}^{\dagger}a_{kR}|$ and $H_T$ describes the reservoirs and their coupling to the dots,
\[H_T = \sum_{\mu,k} (t_{\mu L}a_{\mu L}^{\dagger}a_{k L} + t_{\mu R}a_{\mu R}^{\dagger}a_{k R}) + \text{H.c.},\] (2)
where $\mu = 1, 2$ and $a_{\mu k}^{\dagger}$ are the creation operators for the electrons in the reservoirs while $d_{\mu}^{\dagger}$ are the creation operators for the DD. The last term in eq. (1) describes the interdot repulsion. We assume that there is no direct transmission between the dots and that the couplings of the dots to the leads, $t_{\mu L(R)}$, are independent of energy. In the absence of a magnetic field all couplings are real. In the presence of a magnetic flux $\Phi$, however, the tunneling amplitudes between the dots and the reservoirs are in general complex. We write $t_{\mu L(R)} = t_{\mu L(R)}e^{i\phi_{\mu L(R)}}$, where $t_{\mu L(R)}$ is the coupling without the magnetic field. The phases around the closed circle are constrained to satisfy $\phi_{L} + \phi_{R} - \phi_{L} - \phi_{R} = \phi$, where $\phi \equiv 2\pi\Phi/\Phi_0$.

Let the initial state of the system correspond to filling the left and right reservoirs at zero temperature with electrons up to the Fermi energies $\mu_L$ and $\mu_R$, respectively (fig. 1). Consider first non-interacting electrons, $U = 0$. In this case the problem can be solved exactly for any values of the bias voltage, $\mu_L - \mu_R$, and of the couplings to the leads [7]. Indeed, the total wave function for the non-interacting electrons, $|\Psi(t)\rangle = \exp(-iHt)|\Psi(0)\rangle$, can be written at all times as a product of single-electron wave functions, $|\Psi(t)\rangle = \prod_{\mu} |\psi_{\mu}(t)\rangle$. Here $|\psi_{\mu}(t)\rangle$ describes a single electron initially occupying one of the levels $E_{\mu L} \leq \mu_L$, or $E_{\mu R} \leq \mu_R$ in the left or right lead. It can be written as

$$|\psi_{\mu}(t)\rangle = \left[ \sum_{k,\alpha = L, R} b_{\mu k}^{\alpha}(t)a_{\mu k}^{\dagger} + \sum_{\mu = 1, 2} b_{\mu \alpha}(t)d_{\mu}^{\dagger} \right]|0\rangle, \quad (3)$$

where $b_{\mu k}^{L(R)}(t)$ and $b_{\mu \alpha}(t)$ are the probability amplitudes for finding the electron at the level $E_{\mu L(R)}$ in the left (right) reservoir, or at the level $E_{\mu}$ inside the DD system. The total probability of finding the electron in the right lead is therefore $P_{\mu}(t) = \sum_{k \mu}|b_{\mu k}^{R}(t)|^2$.

Consider $\mu_L > \mu_R$. Then the total average charge $Q(t)$ accumulated in the right lead by time $t$ is a sum of $P_{\mu}(t)$ over all electrons with initial energy within the potential bias, $\mu < E_{\mu L(R)} < \mu_L$. The average current is $I(t) = \dot{Q}(t)$. (We adopt units where the electron charge $e = 1$.) Let us take the continuum limit $\sum_k \rightarrow \int \rho dE_k$, where $\rho$ is the density of states in the corresponding lead. Then the current can be written as

$$I(t) = \int_{\mu_L}^{\mu_R} \rho_L I_{\mu L}(t) dE_{k L}, \quad (4)$$

where $I_{\mu L}(t) = \partial_t \int_{-\infty}^{\infty} |b_{\mu k}^{L}(t)|^2 dE_k$ is a single electron current.

Substituting eq. (3) into the Schrödinger equation $i|\dot{\psi}_{\mu}(t)\rangle = H|\psi_{\mu}(t)\rangle$ we obtain the following equations for the amplitudes $b(t)$:

$$\dot{b}_{\mu k}^{L} = E_{\mu L}b_{\mu k}^{L} + \sum_{\mu = 1, 2} t_{\mu L}b_{\mu \alpha}^{R}, \quad (5a)$$

$$\dot{b}_{\mu \alpha} = E_{\mu \alpha}b_{\mu \alpha} + \sum_{k} (t_{\mu L}b_{\mu k}^{L} + t_{\mu R}b_{\mu k}^{R}), \quad (5b)$$

$$\dot{b}_{\mu k}^{R} = E_{\mu R}b_{\mu k}^{R} + \sum_{\mu = 1, 2} t_{\mu R}b_{\mu \alpha}. \quad (5c)$$

Replacing sum over the lead states by integrals, these equations can be solved analytically. Taking for simplicity $t_{\mu L(R)} = t_{\mu L(R)}e^{i\phi_{\mu L(R)}}$ we obtain for a single electron stationary current ($t \rightarrow \infty$) the following result:

$$I_{\mu} = 4\Gamma L/|\rho_L|^2 \left[ |f_1|^2 + |f_2|^2 + 2\Re(e^{-i\phi}f_1f_2) \right]/D^2, \quad (6)$$

where $\Gamma L(R) = 2\pi|\rho_L(R)|/|\langle L(R)|d_L(R)\rangle|^2$ and

$$f_1,2 = (E_{k L} - E_{1,2}) \pm \frac{i}{2} (\Gamma - \Gamma_\phi),$$

$$D = (2E_{k L} - E_1 - E_2 + i\Gamma^2 - \varepsilon^2 + |\Gamma_\phi|^2).$$

Here $\varepsilon = E_1 - E_2$, $\Gamma = \Gamma_L + \Gamma_R$, and $\Gamma_\phi = \Gamma_L + \Gamma_R e^{i\phi}$.

In the case of large bias, $|\mu_L - E_{1,2}| \gg \Gamma$, the integration limits in eq. (4) can be extended to infinity. As a result, we find for the total current $I = I(\phi)$,

$$I(\phi) = \frac{e^2}{4\Gamma L} \left[ \Gamma L \sin^2 \phi \frac{e^2}{\Gamma} + 4\Gamma L \sin^2 \phi \frac{e^2}{\Gamma} \right], \quad (7)$$

where $I(0) = 2\Gamma L \Gamma R/\Gamma$ is the resonant current for non-interacting electrons in the absence of the magnetic field. The $\phi$-dependence in eq. (7) is an example of the Aharonov-Bohm effect; we illustrate this for finite and infinite bias in fig. 2a. We find the infinite bias limit eq. (7) is a very good approximation for a finite bias whenever the level is well inside the bias window, $|\mu_L - E_{1,2}| \gg \Gamma$.

The entire treatment can in fact be simplified in the large bias limit by transforming eqs. (5) into Bloch-type master equations for the reduced density matrix of the DD system, $\sigma_{\mu \nu}(t)$. Here $j^\prime = \{0, 1, 2, 3\}$ label the DD states in order: the empty DD, the first dot occupied, the second dot occupied, and both dots occupied. This density matrix is related to the amplitudes $b_{\mu \nu}(t)$ via $[8]$

$$\sigma_{\mu \nu}(t) + \delta_{\mu \nu} \sigma_{33}(t) = \int_{-\infty}^{\infty} b_{\mu \nu}(t)b_{\mu \nu}^{*}(t)\rho_L dE_{k L}, \quad (8)$$

for $\mu = 1, 2$ and $\sigma_{00} = 1 - \sigma_{11} - \sigma_{22} - \sigma_{33}$. Then multiplying eqs. (5) by $b_{\mu \nu}^{*}(t)$ and integrating over $E_{k L}$ gives the
where the total current is given by eq. (10) with $\sigma_{33} = 0$. Equations (11) can in fact be rigorously derived by employing the non-equilibrium Greens function in the weak coupling limit [4,9], or directly from the many-body Schrödinger equation in the large bias limit [10] by assuming that $\max(\Gamma, T)/|U_{L,R} - E_{1,2}| \ll 1$, where $T$ is the temperature of the reservoirs [11,12]. Under this assumption they are valid to all orders in the tunneling couplings. (In fact, the master equations approach has been already shown to be very useful for description of electron transport in coupled dots [6,13,14].)

Solving eqs. (11) in the steady-state limit we obtain for the total current

$$I(\phi) = I_C \frac{e^2}{\varepsilon^2 + I_C \left(2\Gamma_R \sin^2 \frac{\phi}{2} - \varepsilon \sin \phi \right)}.$$

(12)

where $I_C = 2\Gamma_L\Gamma_R/(2\Gamma_L + \Gamma_R)$ is the total current (with the Coulomb blockade) in the absence of the magnetic field. Comparing eq. (12) with eq. (7) for $\varepsilon \neq 0$ we find that both currents display the Aharonov-Bohm oscillations, fig. 2. Nevertheless, eq. (12) shows an asymmetric behavior with respect to the magnetic flux $\phi$ (the dashed curve in fig. 2b). It looks contradictory to the Onsager relation that locks the current peaks at $\phi = 2\pi n$ in any two-terminal linear transport [15]. Yet, this relation is not applicable in our case, corresponding to interacting system under finite bias voltage [4].

More strikingly is the current behavior in the interacting case for $\varepsilon \to 0$. While the resonant current for the non-interacting electrons keeps oscillating with the magnetic field, fig. 2a, it becomes non-analytic in $\phi$ in the case of Coulomb blockade. Indeed, one finds from eq. (12) that $I = I_C$ for $\phi = 2\pi n$, where $n = \Phi/\Phi_0$ is an integer, but $I = 0$ for any other value of $\Phi$ (fig. 2b).

Such an unexpected “switching” behavior of the electron current in the magnetic field can be explained in the following way. Let us disentangle the Coulomb blockade and quantum interference effect, which interplay in a non-trivial way in the electron current. This can be done by defining new basis states of the DD, $d_{1,2}^\dagger \to d_{1,2}^\dagger \to d_{1,2}^\dagger \to d_{1,2}^\dagger$, chosen so that they do not interfere in the electron current. For instance, if the state $d_{1,2}^\dagger \to d_{1,2}^\dagger$ is not coupled to the right reservoir, i.e., $t_{2R} \to t_{2R} = 0$, then the current would flow only through the state $d_{1,2}^\dagger \to d_{1,2}^\dagger$. Obviously, no interference between these two states would appear in the total current.

Such a basis can be found for $\varepsilon = E_1 - E_2 = 0$ when the DD Hamiltonian, $H = d_{1,2}^\dagger \to d_{1,2}^\dagger$, is invariant under $SU(2)$ transformations. Then the unitary transformation

$$\begin{pmatrix} d_1 \\ d_2 \end{pmatrix} = \frac{1}{N} \begin{pmatrix} t_{1R} & t_{2R} \\ -\bar{t}_{2R} & t_{1R} \end{pmatrix} \begin{pmatrix} d_1 \\ d_2 \end{pmatrix}.$$

(13)

with $N = (t_{1R}^2 + \bar{t}_{2R}^2)^{1/2}$, results in $\bar{t}_{2R} = 0$.

Consider now the coupling of the state $d_{1,2}^\dagger \to d_{1,2}^\dagger$ to the left lead. One obtains from eq. (2) for $t_{2L} \equiv t_{2L} \equiv 0$:

$$\bar{t}_{2L}(\phi) = -e^{i(\phi_{2L} - \phi_{1L})}(t_{1L}^* t_{2R} - t_{2L}^* t_{1R})/N.$$

(14)
It follows from this expression that $\tilde{t}_{2L} = 0$ for $\phi = 2n\pi$ provided that $t_{1L}/t_{2L} = t_{1R}/t_{2R}$, or for $\phi = (2n + 1)\pi$ if $t_{1L}/t_{2L} = -t_{1R}/t_{2R}$ [16]. Then the state $\tilde{d}_{1}^1|0\rangle$ decouples from both leads. This would not affect the resonant current for non-interacting electrons ($U = 0$), since the state $\tilde{d}_{1}^2|0\rangle$ is already decoupled from the right lead.

In the case of Coulomb blockade, however, the coupling to the left lead becomes the point of crucial importance. Indeed, any discrete state coupled to an infinite reservoir is going to be totally occupied, no matter how weak the coupling [7]. Then the state $\tilde{d}_{1}^1|0\rangle$, carrying the current, will be blocked by the Coulomb interdot repulsion. As a result, the total current **vanishes**. However, if the state $\tilde{d}_{2}^1|0\rangle$ is decoupled from both leads, it remains unoccupied, so that the current can flow through the state $\tilde{d}_{2}^2|0\rangle$. As shown above, this takes place precisely for $t_{1L}/t_{2L} = \pm t_{1R}/t_{2R}$. If this condition is not fulfilled, the current is always zero, even for $\phi = 2m\pi$.

It is clear from our explanation that the switching phenomenon disappears for $\varepsilon \neq 0$, as can be seen from eq. (12). Indeed, in this case the DD Hamiltonian $\sum_{\nu} E_{\nu} d_{\nu}^\dagger d_{\nu}$ is not invariant under $SU(2)$ transformations. Therefore any unitary transformation that eliminates one link between the DD and the reservoirs, like eq. (13), would generate direct coupling ($\tilde{t}_{12} = \varepsilon t_{1R} t_{2R} / (|t_{1R}|^2 + |t_{2R}|^2)$) between the states $\tilde{d}_{1,2}^1|0\rangle$ (see footnote 1). As a result, the current is not interrupted by the Coulomb blockade, so that its behavior for $\varepsilon \neq 0$ would be similar to that in the non-interacting case, fig. 2.

The switching effect of the magnetic field on the electric current takes place only in the stationary regime, $t \to \infty$, where one of the states $\tilde{d}_{1,2}^1|0\rangle$ is fully occupied. However, it takes a certain transition time, $t_{tr}$, during which the total current is not zero. One estimates from eq. (11) that $t_{tr} \sim 1/\tilde{t}_{2L}$, where $\tilde{t}_{2L} = 2\pi p_L |\tilde{R}_{2L}|^2$. Using eq. (14) we find that $t_{tr} \sim \Gamma^{-1}(\Phi_0/\Phi)^2$ for $\Phi \ll \Phi_0$. Hence $t_{tr}$ becomes much longer than the usual relaxation time $\Gamma^{-1}$ for a very small magnetic field (or in general when $\Phi \to n\Phi_0$). This is illustrated in fig. 3, which shows the transient current $I(t)$ as a function of $t$ for different values of $\phi$. One finds from this figure that the current always increases for small $t$. However, it eventually disappears for $t \gg t_{tr}$. This explains the non-analyticity of $I(\phi)$ at $\phi = 2m\pi$. Indeed, at any finite time $t$ there is no discontinuity at $\phi = 2m\pi$. It appears only in the limit of $t \to \infty$ because $t_{tr} \to \infty$ for $\Phi \to 0$.

As we demonstrated above, the switching effect becomes very transparent in a particular basis of the DD states. Still, it is very surprising how such a basis emerges dynamically? Indeed, an electron from the left lead can enter the DD system in any of $SU(2)$ equivalent superpositions of its states. Therefore there exists a probability for each electron to enter the DD in the superposition that eliminates one of the links with the right lead. When it happens, the electron would be trapped in this state. Even if the probability of this event for one electron is very small, the total number of electrons passing the DD goes to infinity for $t \to \infty$. Therefore the trapping event is always realized for large enough time. In the presence of Coulomb blockade this would lead to the switching effect, as explained above. (A similar effect of the Coulomb blockade, leading to divergency in the shot-noise power has been discussed in different publications [14,17–19]. However, this phenomenon is not related to vanishing of current, but rather to electron bunching leading to a system’s bistability [18,19].)

Our interpretation of the switching effect allows us to determine the necessary conditions for its realization in real experiment. First, we need the total occupation of one of the states $\tilde{d}_{1,2}^1|0\rangle$, decoupled with the right lead. This takes place only if the energy level is far from the corresponding Fermi energy, $|\mu_L - E_1| \gg \Gamma$. For instance, if $E_1 = \mu_L$, the occupation probability reaches only $1/2$ at $t \to \infty$. As the bias increases, however, other levels of the dots can enter into the bias window. In this case the transport would proceed through several levels. If the dots are identical, the same “rotation” (13), applied to each pair of levels with the same energy, would result in a simultaneous decoupling of the corresponding “rotated” states from the right lead. Then the switching effect is expected to take place in this case as well, even if the interlevel spacing is very small.

In real system, the surrounding environment will cause dephasing between the two dots due to fluctuations of the dots parameters, like energy levels, tunneling couplings, etc. Since a particular origin of these fluctuation is irrelevant for evaluation of the corresponding dephasing rate [20], one can model the environment by isolated fluctuators interacting with the system and vibrating its energy levels [19]. Using such a model for the fluctuating environment one finds for the stationary current (for $\varepsilon = 0$ and $\Gamma_L = \Gamma_R = \Gamma/2$) [19],

$$I = \frac{I_c}{1 + I_c \tau_d \sin^2(\phi/2)},$$

where $I_c = \Gamma/3$ and $\tau_d$ the dephasing time. Therefore in order to observe the switching effect one requires $I_c \tau_d \gg 1/\sin^2(\phi/2)$. For small $\phi$ it is equivalent to $\tau_d \gg t_{tr}$. 

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1No such coupling is produced by the interdot repulsion term in eq. (1), since it is always invariant under the transformation (13): $U d_1^\dagger d_1 d_2^\dagger d_2 \to U d_1^\dagger d_2^\dagger d_2$. 

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In general, for a weak coupling to the environment, the dephasing rate $\gamma_d = 1/\tau_d$ can be evaluated as $\gamma_d \sim (\delta \epsilon)^2 S(0)$, where $\delta \epsilon$ is an average fluctuation of dots levels and $S(\omega)$ is the corresponding spectral density [20,21]. In the case of the thermal environment $S(\omega) \propto T$, where $T$ is temperature [21]. Therefore by decreasing the temperature one can make the corresponding decoherence rate arbitrary small in order to reach the switching effect even for small values of the magnetic flux.

As we explained above the experimental realization of the switching effect would require fulfillment of different conditions which should be met. The most essential of them are large bias and interdot Coulomb repulsion, which should exceed the bias.

With respect to the bias, it would be hard to realize the large bias voltage (on scale of $\Gamma$) with one level inside the bias window. However, in the case of two identical dots, the switching effect is expected even if many levels are inside the bias. Therefore the large bias condition should not create essential experimental problem.

A realization of the large interdot repulsion condition would represent a more complicated problem. Indeed, it implies that two dots are very close. This condition is still not met in present experiments [1]. However, it can be met by decreasing the dots size. An another way to achieve proximity of two parallel dots is to use different materials. For instance the quantum dots in graphene system can be a very promising set-up for an investigation of the switching effect [22].

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