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Kondo effect in Aharonov-Casher interferometer

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We consider a model describing a spin field-effect transistor based on a quantum nano-wire with tunable spin-orbit interaction embedded between two ferromagnetic leads with anti-collinear magnetization. We investigate a regime of strong interplay between resonance Kondo scattering and interference associated with Aharonov-Casher effect. Using Keldysh technique at weak coupling regime we calculate perturbatively the charge current. It is predicted that the effects of spin-orbit interaction result in a non-vanishing current for any spin polarization of the leads including the case of fully polarized anti-collinear contacts. We analyze the influence of the Aharonov-Casher phase and degree of spin polarization in the leads onto a Kondo temperature.

I. INTRODUCTION

Kondo effect is known to play a very important role for a charge transport through nano-structures facilitating the maximal conductance of a nano-device at zero bias [1]. Having a spin nature, Kondo effect is associated with a resonance scattering accompanied by the spin flip through the multiple cotunneling processes in Coulomb blockaded nano-devices [2]. Kondo effect in GaAs based semiconductor nano-structures (quantum dots, quantum point contacts, quantum wires etc) attracted enormous attention of both experimental and theoretical communities during last two decades [2-5]. Recently, semiconductor quantum wires fabricated on InAs and InSb heterostructures started to be widely used in the new quantum technological devices [6],[7]. One of the most important property of these materials is related to effects of a strong spin-orbit interaction (SOI) which does not conserve spin in the resonance scattering processes, see, e.g., [8]. The high tunability of the interplay between SOI and resonance Kondo effect and its influence on the charge and spin transport through nano-structures paves a way for practical applications of these materials in spintronics devices. It is known that contrast to effects of external magnetic field, the effect of SOI on electrons scattering is preserving a time-reversal symmetry. While magnetic field is destructive for the Kondo effect due to suppression of the spin-flip processes, the influence of SOI on the resonance scattering is more delicate.

One of the most remarkable manifestation of SOI in quantum devices (e.g. Datta-Das spin field-effect transistor) is associated with an accumulation of a spin-dependent phase difference in the electron (spinor) wave function [9],[10]. This phase accumulation being controlled by an external electric field applied to a nano-

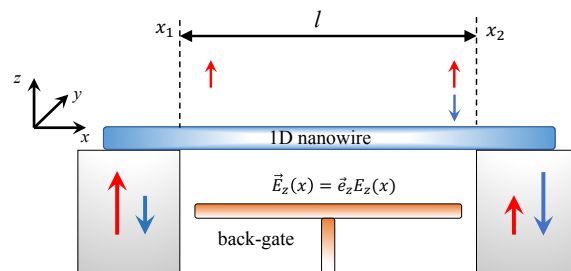


FIG. 1. Scheme of a nano-device. 1D nanowire of length L is placed between two massive magnetically polarized electrodes. Polarization is chosen to be collinear antiparallel (AP). Spin-dependent density of states in the leads is defined through $\nu_{L\uparrow} = \nu_{R\downarrow} = \nu(1+p)$ and $\nu_{L\downarrow} = \nu_{R\uparrow} = \nu(1-p)$. Back-gate electrode situated near the nanowire creates an electric field in z -direction inducing a spin-orbit interaction (SOI) in the nanowire. Short nanowire is treated as a quantum dot (QD) (see the main text for a discussion). We assume that QD is in a strong Coulomb blockade regime. Odd Coulomb valleys provide an access to Kondo physics. SOI leads to an accumulation of Aharonov-Casher phase in the electron wave function which is equivalent semiclassically to an electron spin precession.

device is known as Aharonov-Casher effect [9]. The electric field manipulation of the Aharonov-Casher interference provides a big advantage compared to external magnetic field control. In particular, no magnetization currents are generated both in nano-device and in the leads and no extra decoherence associated with an extra heating appears. Additional degree of control associated with use of ferromagnetic leads allows to open (enhance) and close (suppress) a charge current through the nano-structure [10],[11] similar to effects of a spin valve [12],[13],[14].

In this paper we present an example when strong Coulomb blockade and established quantum coherence of electrons are simultaneously present and controlled

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in a quantum nano-wire. We consider Kondo tunnelling of electrons through the nano-wire in the presence of strong SOI in it. We show that quantum interference originating from the Aharonov-Casher phase accumulated during tunneling process affects qualitatively many-body Kondo transmission and results in strong renormalization of Kondo temperature and significant enhancement of the charge current.

The paper is organized as follows: In Sec.II we introduce a model Hamiltonian of a spin-orbit active one-dimensional nanowire placed between spin-polarized electrodes and derive effective Kondo model. In Sec.III we analyse the charge current through nanowire calculated in the lowest order of perturbation theory. In Sec.IV we obtain the contribution to the charge current in second order of perturbation theory and evaluate a Kondo temperature as a function of Aharonov-Casher phase and degree of spin polarization in the leads. In Sec.V we analyze Kondo temperature in particular limit of fully polarized anti-collinear contacts.

II. MODEL

We investigate a Datta-Das spin field-effect transistor [10] with spin-orbit active weak-link in the Kondo regime. We consider a one-dimensional (1D) nanowire embedded between two magnetically polarized electrodes in anti-parallel configuration (AP). Back gate is controlling a spin-orbit interaction (SOI) in the nanowire, see Fig.1. One dimensional nanowire can be treated as a quantum dot (QD) in a regime when temperature and bias voltage are smaller compared to a mean level spacing in the QD $\delta\varepsilon \sim \hbar v_F L^{-1}$. We assume the odd number of electrons in the QD to access the Kondo regime. The model is described by the Hamiltonian: $H = H_0 + H_{tun}$, where

$$H_0 = \sum_{k,\alpha,\sigma} (\varepsilon_k - \mu_\alpha^\sigma) c_{k\alpha\sigma}^\dagger c_{k\alpha\sigma} + \sum_\lambda (\varepsilon_0 d_\lambda^\dagger d_\lambda + U_C \hat{n}_\lambda \hat{n}_{\bar{\lambda}}) \quad (1)$$

characterizes 1D nanowire and the magnetically polarized left (right) electrodes with chemical potentials $\mu_{L(R)}^\sigma$. Here ε_0 stands for the energy of first half-filled level of the dot counted from the Fermi level of the leads, U_C is the charging energy in the nanowire. The annihilation (creation) operators of the conduction electrons are denoted by $c_{k\alpha\sigma}$ ($c_{k\alpha\sigma}^\dagger$), where $\alpha = L, R$. The electron states in the leads are characterized by spin quantum number $\sigma = (\uparrow, \downarrow)$. The two-fold degenerate quantum level in the dot represented by linear superposition of the states with $\sigma = \uparrow$ and $\sigma = \downarrow$ is described by the pseudospin quantum number with two eigenvalues $\lambda, \bar{\lambda}$. We use notations d_λ (d_λ^\dagger) for the electrons in the QD, $\hat{n}_\lambda = d_\lambda^\dagger d_\lambda$, see Appendix A. To describe partial spin polarization of the electrodes we introduce spin-dependent density of states at Fermi energy: $\nu_{L\uparrow} = \nu_{R\downarrow} = \nu(1+p)$ and $\nu_{L\downarrow} = \nu_{R\uparrow} = \nu(1-p)$, where parameter p defines a degree of polarization, see Fig.1. If magnetizations in the leads

are collinear and oriented anti-parallel, the net magnetic field produced by leads at the position of the nanowire is zero. Therefore, the net magnetic field does not lift a two-fold degeneracy of the pseudospin state in the QD. If the orientation of magnetizations is parallel, the net magnetic field at the position of the nanowire is non-zero resulting in the time-reversal symmetry breaking effects.

While the spin is a good quantum number in the leads, it can not be used for the characterization of the state in the nano-wire due to presence of SOI. Thus, the tunnel matrix element computed using wave functions of electrons in the leads and in QD (see, e.g., [15]) is characterized by two indices σ (spin) and λ (pseudospin). The most general form of the tunneling Hamiltonian is given by

$$H_{tun} = \sum_{k,\alpha;\sigma\lambda} \left(V_{k\alpha}^{\sigma\lambda} c_{k\alpha\sigma}^\dagger d_\lambda + h.c \right). \quad (2)$$

We solve 1D Schrödinger equation for an electron in the nanowire in the presence of SOI, see details in Appendix A, and express the tunnel matrix amplitudes in terms of the SOI parameters. The two component electron wave function $\vec{\psi}_\lambda$ at different points are connected through the operator \hat{U} in such way $\vec{\psi}_\lambda(x_2) = \hat{U}(x_2, x_1) \vec{\psi}_\lambda(x_1)$, where

$$\hat{U} = \exp \left[\frac{i\hat{\sigma}^y \vartheta(x_1, x_2)}{2} \right], \quad \vartheta = \frac{2\alpha p_F l}{\hbar v_F}, \quad (3)$$

characterizes accumulation of Aharonov-Casher phase, $l = |x_2 - x_1|$, see Fig.1. In Eq.(3) $\alpha \alpha E_z$ is the SOI coupling constant, E_z is the electric field in z direction produced by the back gate electrode, $\hat{\sigma}^y$ is the y -Pauli matrix. Using Eq.(3) and assuming that the tunneling occurs at the points x_1, x_2 , we parametrize the tunnel matrix elements as follows:

$$V_{k\alpha}^{\tau\tau'} = V_{k\alpha} (\delta_{\tau\tau'} \cos(\vartheta/4) \mp i\hat{\sigma}_{\tau\tau'}^y \sin(\vartheta/4)). \quad (4)$$

Here $-/+$ stands for the L/R lead correspondingly. Effect of the SOI on the tunneling processes is characterized by the parameter ϑ . The SOI vanishes for $\vartheta=0$ and reaches its maximal value at $\vartheta=\pi$, when both tunneling processes, diagonal and off-diagonal in spin (pseudospin) indices, contribute equally, see Eq.(4). We assume full symmetry in tunneling junction $V_{kL} = V_{kR} = V_{tun}$.

The mapping of the Anderson impurity model Eq.(1) onto a Kondo-like model is done using standard Schrieffer-Wolff transformation [16], see Appendix B. We assume single occupied two-fold degenerate level in the QD and consider the energy level width to be smaller compared to the charging energy $\Gamma = 2\pi\nu |V_{tun}|^2 \ll U_C$.

Effective Hamiltonian $H_{eff} = H_{dir} + H_{ex}$ contains H_{dir} describing a direct (potential) electron scattering be-

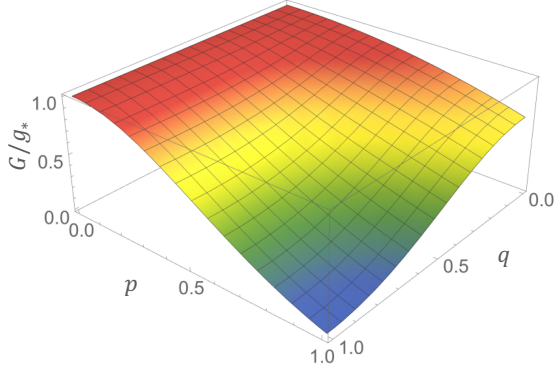


FIG. 2. Differential conductance in the units of $g_* = 3e^2\pi^2(\nu J)^2/(4\pi\hbar)$ as a function of both polarization in the leads p and effects associated with the accumulation of the Aharonov-Casher phase ϑ parametrized by $q = \cos(\vartheta/2)$ for the case of the bias voltage $eV/T \gg 1$.

tween the leads and

$$\begin{aligned}
H_{ex} = & \sum_{\alpha} J [(s_{\alpha\alpha}^z S^z + s_{\alpha\alpha}^x S^x) \cos(\vartheta/2) + s_{\alpha\alpha}^y S^y] \\
& + \sum_{\alpha \neq \alpha'} J [s_{\alpha\alpha'}^z S^z + s_{\alpha\alpha'}^x S^x + s_{\alpha\alpha'}^y S^y \cos(\vartheta/2)] \\
& - J \sin(\vartheta/2) [(s_{LL}^z - s_{RR}^z) S^x - (s_{LL}^x - s_{RR}^x) S^z] \\
& - \frac{J}{2} \sin(\vartheta/2) j_{LR} S^y
\end{aligned} \quad (5)$$

constituting the effective exchange interaction between pseudospin-1/2 in the QD, \vec{S} , and spin $\vec{s}_{\alpha\alpha'} = \sum_{kk'} (1/2) c_{k\alpha\sigma}^\dagger \hat{\sigma}_{\sigma\sigma'} c_{k'\alpha'\sigma'}$ of the conduction electrons in the L/R leads and charge transfer $j_{LR} = i \sum_{kk'} (c_{kL\uparrow}^\dagger c_{k'R\uparrow} + c_{kL\downarrow}^\dagger c_{k'R\downarrow} - h.c.)$. We used the following notations for the exchange interaction constant in Eq. (5):

$$J = 2U_C \frac{|V_{tun}|^2}{|\varepsilon_0|(U_C - |\varepsilon_0|)}. \quad (6)$$

We concentrate below on the case of electron-hole symmetry, $\varepsilon_0 \rightarrow (-U_C/2)$, and ignore irrelevant processes of potential scattering [17].

The influence of SOI effects onto the Kondo scattering has several facets. First, SOI is responsible for the different types of the spin anisotropies in the terms diagonal and off-diagonal in the lead indices (first and second lines in Eq.(5)). Second, SOI produces an additional coupling between the pseudospin in QD and spin density of the conduction electrons aka Dzyaloshinsky-Moriya (DM) interaction ($\propto \vec{e}_y \cdot [\vec{s}_{\alpha\alpha} \times \vec{S}]$) [18]. Third, SOI mediates the interaction between the pseudospin in the QD and the charge transfer.

III. COTUNNELING CURRENT

Assuming a high temperature (compared to some emerging energy scale to be defined below) regime we

calculate the current through the nano-wire perturbatively in $\nu J \ll 1$. The first non-vanishing contribution to the charge current is $\propto (\nu J)^2$. The cotunneling current given by Eq.(7) can be straightforwardly derived either through equation of motion method or using the nonequilibrium Keldysh Green function technique (we adopt below the units $k_B=1$):

$$\begin{aligned}
I^{(2)} = & \frac{e\pi^2}{4\pi\hbar} (\nu J)^2 eV \left\{ (1-p^2)(2-q^2) \right. \\
& \left. + (1+q^2)(1+p^2) - 8pq \coth\left(\frac{eV}{2T}\right) \langle S^z \rangle \right\}, \quad (7)
\end{aligned}$$

where we use shorthand notations $q = \cos(\vartheta/2)$ for parametrization of the accumulated Aharonov-Casher phase ($q=1$ for the case of the absence of SOI, $\vartheta = 0$, and $q = 0$ is when the SOI is maximal, $\vartheta = \pi$). In Eq.(7) $\langle S^z \rangle$ denotes an out-of-equilibrium QD (nanowire) magnetization [19],[20]. The QD magnetization that appears because of applied bias voltage in the presence of a finite polarization p is non-vanishing even without external magnetic field, see [21]. The temperature T in the Eq. (7) stands for the temperature in the contacts which are assumed to be in the equilibrium. The expression for QD magnetization $\langle S^z \rangle$ is obtained from the steady-state solution of the quantum Langevin equation of motion [21],[22] for the QD spin-1/2 in the lowest order of perturbation theory in νJ :

$$\langle S^z \rangle = \frac{pq(eV/T)}{2(1-q^2p^2) + \varphi(\frac{eV}{T})(p^2+q^2)}, \quad (8)$$

here $\varphi(x) = x \coth(x/2)$. Non-equilibrium QD magnetization described by Eq.(8) is limited by $\langle S^z \rangle = \pm pq/(p^2+q^2)$ achieved at large bias-voltage $eV \gg T$.

The appearance of the non-equilibrium QD magnetization $\langle S^z \rangle$ in the Eq. (7) influences the shape of the peak in the differential conductance, $G^{(2)} = dI^{(2)}/dV|_{V \rightarrow 0}$, see [23],[21]. At large bias voltages $eV \geq T$ the effects of saturation of the QD magnetization result in suppression of the charge current. At low bias voltages $eV \ll T$ the contribution to the current proportional to the QD magnetization is vanishing and the current reaches its maximum value. The height of the conductance peak depends on the QD magnetization slope $\partial \langle S^z \rangle / \partial(eV) = pq/[2T(1+p^2+q^2-p^2q^2)]$.

The p - and q - dependence of the differential conductance given by the Eq.(7) is illustrated in Fig.2 at $eV \gg T$. The leading contribution to the differential conductance calculated in the lowest non-vanishing order of the perturbation theory for the case of non-magnetic leads ($p=0$) in the absence of the SOI ($q=1$) saturates at $G^{(2)} = g_*$, where $g_* = (e^2/\pi\hbar)(3/4)(\pi\nu J)^2$. The conductance peak at zero bias voltage for the case of partial polarization of the leads ($p \neq 1$) is $G^{(2)} = g_*(1-p^2)$. The linear response (voltage independent) part of the differential conductance at large bias voltage ($eV \gg T$) is asymptotically given by $G^{(2)} = g_*(3-4p^2+p^4)/[3(1+p^2)]$, see [23]. Spin-dependent tunneling induced by SOI at $q < 1$ enhances the charge transport through QD at

any collinear AP of the reservoirs ($p \neq 0$). The zero-bias conductance in the case of fully polarized leads $p=1$ is given by $G^{(2)}|_{V \rightarrow 0} = g_*(2/3)(1 - q^2)$, while the linear response conductance at large bias voltages is $G^{(2)}|_{V \rightarrow \infty} = g_*(2/3)(1 - q^2)^2/(1 + q^2)$. The effect of SOI on the charge current is maximal at $q = 0$. The polarization dependence of the differential conductance in this case is given by $G^{(2)} = g_*(1 - p^2/3)$.

IV. KONDO CONTRIBUTION TO THE CHARGE CURRENT

The next non-vanishing contribution to the charge current $\propto (\nu J)^3$ depends on the spin-flip processes and therefore described by the Kondo physics. We apply the non-equilibrium Keldysh Green's function technique and Abrikosov's *pseudofermion* representation [24], see details in [20] to proceed with the calculations. The current $I^{(3)} = I_K^{(3)} + I_{an}^{(3)}$ consists of two parts: i) $I_K^{(3)}$ is originating from the anisotropic Kondo model (first two lines in Eq.(5))

$$\frac{I_K^{(3)}}{g_* 4\nu J} = \frac{1}{3} \left\{ (1 - p^2) [(1 + q^2)V - 2pq\mathcal{S}^z] + q^2(1 + 3p^2)V - pq(1 + q^2)(3 + p^2)\mathcal{S}^z \right\} \log\left(\frac{D}{T^*}\right), \quad (9)$$

and ii) $I_{an}^{(3)}$ is accounting for both Dzyaloshinsky-Moriya and charge transfer processes in Eq.(5)

$$\frac{I_{an}^{(3)}}{g_* 4\nu J} = \frac{(1 - q^2)}{3} (1 - p^2) \{2V - 3pq\mathcal{S}^z\} \log\left(\frac{D}{T^*}\right) \quad (10)$$

Here D is the bandwidth of the leads. We use the shorthand notations $\mathcal{S}^z = V \coth(eV/2T) \langle S^z \rangle$ and $T^* = \max[|eV|, T]$. The third order in $(\nu J)^3$ correction to the charge current logarithmically grows with the decreasing of both the temperature and the applied bias voltage, revealing a *Kondo anomaly*. The validity of perturbation theory approximation (weak coupling regime) of Eqs.(9),(10) determines the energy scale T_K , the Kondo temperature. Perturbation theory breaks down at $T \lesssim T_K$. The effective coupling constants in this (strong coupling) regime flow towards the strong coupling fixed point.

The dependence of the Kondo temperature on the parameters p and q is in general determined by the solution of the system of coupled Renormalization Group (RG) equations [25],[26]. Without losing a generality we parametrize T_K by the function $f(p, q)$ [27],[28]

$$T_K = D \exp\left(-\frac{f(p, q)}{2\nu J}\right). \quad (11)$$

The form of $f(p, q)$ is known for several limiting cases [29-31]. In particular, $f(p, 1) = 1$ for all $|p| \leq 1$ [29],[30]. Besides, in the case of non-magnetic leads ($p = 0$) we get $f(0, q) = 1$. The form of $f(p, q)$ on a line $p=1$ for $r=1-q \ll 1$ is found perturbatively from the condition

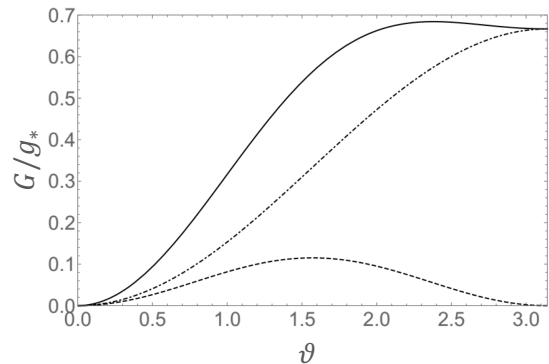


FIG. 3. Dependence of zero-bias differential conductance (in units of $g_* = (3/4)(e^2/\pi\hbar)(\pi\nu J)^2$) on the Aharonov-Casher phase for the case of full AP polarization of the leads $p = 1$. Dot-dashed line denotes differential conductance $G^{(2)}$ obtained from Eq.(7), dashed line is the third order $\propto (\nu J)^3$ perturbative correction to the conductance $G^{(3)}$ determined by Eqs.(9),(10). Solid line represents the sum of $G^{(2)}$ and $G^{(3)}$. Figure is plotted for the following values of model parameters: $D = 1$ eV, $T = 1$ K and $\nu J = 0.1$.

of breaking down perturbation theory for the differential conductance: $f(1, r) = 1 + r$ (here $r = \vartheta^2/4$). With the same logic we found $f(p \ll 1, 0) = 1 + 2p^2/3$. Perturbative analysis leads to a conclusion that $f(p, q) \geq 1$. Moreover, the absolute minimum of $f(p, q)$ is reached at symmetry lines (points), where original Hamiltonian (5) is mapped onto isotropic Kondo model. We conclude that Kondo temperature in the Datta-Das transistor with AP polarized electrodes becomes function of the Aharonov-Casher phase and degree of spin polarization in the leads.

The most striking effect of the influence of the Aharonov-Casher interference onto Kondo scattering is manifested in the case of full AP polarization, $p=1$. In particular, the non-vanishing charge current is controlled by the Aharonov-Casher phase.

The expression for the differential conductance $G = G^{(2)} + G^{(3)}$ derived in the limit of small bias voltage $eV \ll T$ is given by:

$$G = \frac{e^2}{2\pi\hbar} (\pi\nu J)^2 (1 - q^2) \left(1 + q^2 4\nu J \log \frac{D}{T}\right). \quad (12)$$

The conductance dependence on the Aharonov-Casher phase ϑ is shown in Fig.3. Charge current at $p=1$ and $q=1$ is blocked by the Pauli principle. However, absence of the charge current is not in a contradiction with the presence of resonant Kondo scattering. Kondo effect (and T_K) are fully determined by multiple pseudospin flip processes on a QD and a single (L or R) contact [29]. As it is seen from Eq.(12), the logarithmic corrections to the conductance are positive for $J > 0$ when $q \neq 1$ and increase with decreasing of T . While $G^{(2)}$ is monotonously increased with ϑ , the behavior of $G^{(3)}$ is non-monotonous due to interplay between Kondo effect and Aharonov-Casher interferometer. As a result,

the maximal current is reached at some critical value $q_{cr} = \sqrt{[1 - (4\nu J)^{-1} \log(T/D)]/2}$ which depends on T and initial parameters of the model.

V. SOI INFLUENCE ON KONDO TEMPERATURE

The Hamiltonian Eq. (5) casts a simple form for $p=1$

$$\tilde{H}_{ex} = J \{ (s^z S^z + s^y S^y) \cos(\vartheta/2) + s^x S^x \} - \frac{1}{2} J \sin(\vartheta/2) \sum_{kk'} c_{k\gamma}^\dagger c_{k'\gamma'} S^x. \quad (13)$$

The Hamiltonian (13) is derived from Eq. (5) by retaining the operators $c_{k\gamma} (c_{k\gamma}^\dagger)$ with $\gamma=1$ for (L, \uparrow) , $\gamma=-1$ for (R, \downarrow) and $\vec{s} = \sum_{kk'} (1/2) c_{k\gamma}^\dagger \hat{\sigma}_{\gamma\gamma'} c_{k'\gamma'}$. We omit all other terms in Eq.(5) with zero expectation values. Eq. (13) describes an *anisotropic Kondo-like model* with an additional term, which couples the spin in the QD with the charge density in the leads. The last term in (13) can be viewed as an extra potential scattering and therefore disregarded for the particle-hole symmetric limit.

As it is known both from the RG and exact Bethe ansatz solution of the Kondo model the maximal Kondo temperature is achieved in isotropic case [27]. The Kondo temperature in anisotropic case is defined through the Bethe ansatz solution [27] by an equation equivalent to (11) with the dimensionless function $f(1, q)$ dependent on the anisotropy parameter $q = \cos(\vartheta/2)$:

$$f(1, q) = -\frac{1}{\sqrt{1-q^2}} \log \left(\frac{1 - \sqrt{1-q^2}}{q} \right). \quad (14)$$

The anisotropy controlled by SOI suppresses Kondo temperature. Asymptotic behavior of Eq.(14) in the limit of weak ($\vartheta \rightarrow 0$) SOI is given by $f(1, \vartheta) - 1 \approx \vartheta^2/12$. Similar behavior is also obtained from the Renormalization Group treatment under condition $\nu J \ll 1$. While the cases of small anisotropy can be accessed perturbatively, see discussion after Eq.(11), or by Bethe ansatz solution ($p=1$), see Eq.(14), the solution for function $f(p, q)$ for arbitrary values of its arguments $-1 < (p, q) < 1$ remains an interesting and unsolved problem [32].

VI. SUMMARY AND OUTLOOK

The interplay between resonance Kondo scattering in the quantum wire, effects of SOI in the tunnel barriers and partial spin polarization in the leads provides high level of controllability for the charge transport through the nano-device. In particular, the fine-tuning of the Kondo temperature is achieved by control of three independent tunable parameters of the system. First, the Aharonov-Casher phase is tuned by the electric field applied to the area of the nano-wire. Second, the degree

of spin-polarization in the leads is manipulated by the spin-valve [12-14]. Third, the local out-of-equilibrium QD magnetization of the nano-device is controlled by the source-drain voltage. The central result of the paper is a prediction of a finite charge current through the nano-wire even at full AP polarization of the leads in the presence of non-zero spin-orbit interaction. Besides, perturbative (weak coupling) calculations demonstrate pronounced (logarithmic) effects of enhancement of the current by SOI at any given partial polarization of the leads. Competition between the resonance scattering resulting in maximal Kondo temperature in the absence of SOI at $\vartheta=0$ and quantum interference due to Aharonov-Casher effect maximal at $\vartheta=\pi$ allows to find an optimal strength of SOI at $\vartheta_{cr}(q_{cr})$ under condition of maximizing the electric current.

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Appendix A: Spin-orbit interaction and spin-dependent tunnel matrix elements

In this Appendix we calculate the Aharonov-Casher phase ϑ , see Eq.(3). We start from a Schrödinger equation for the electron in 1D nanowire in the presence of external homogeneous electric field E_z produced by the back-gate and directed perpendicular to nanowire:

$$-\frac{\hbar^2}{2m} \frac{\partial^2 \vec{\psi}}{\partial x^2} - \alpha \hat{\sigma}^y i \hbar \frac{\partial \vec{\psi}}{\partial x} = (E_F - \varepsilon) \vec{\psi}, \quad (A.1)$$

where $\vec{\psi}$ is a two component electron wave function (spinor), $\alpha \propto E_z$ is a spin-orbit interaction coupling constant. Since $\varepsilon \leq \hbar v_F L^{-1} \ll E_F$, we can present the electron's wave function in terms of right and left moving parts:

$$\vec{\psi} = e^{ip_F x / \hbar} \vec{\psi}_+(x) + e^{-ip_F x / \hbar} \vec{\psi}_-(x), \quad (A.2)$$

where $p_F = \sqrt{2m(E_F - \varepsilon)}$ is a Fermi momentum. Substituting the wave function Eq.(A.2) into Eq.(A.1) and neglecting the second derivative we get:

$$-iv_F \hbar \frac{\partial \vec{\psi}_\pm}{\partial x} = \alpha p_F \hat{\sigma}^y \vec{\psi}_\pm. \quad (A.3)$$

From this equation one can see that spinor function $\vec{\psi}_\pm$ satisfy relation

$$\vec{\psi}_\pm(x) = \hat{U}(x) \vec{\psi}_\pm(0) \quad , \quad \hat{U}(x) = e^{i\hat{\sigma}^y \vartheta(x)/2}, \quad (A.4)$$

where $\vartheta(x) = 2\alpha xp_F/(\hbar v_F)$ is an Aharonov-Casher phase.

Despite the fact that in the presence of SOI electronic spin is unsuitable quantum number for the classification of the electronic states, the energy levels continue to be double degenerate. If $\vec{\psi}_\lambda$ is the eigenstate with energy ε , then $\vec{\psi}_{\bar{\lambda}} = i\hat{\sigma}^y \vec{\psi}_\lambda^*$ is also eigenstate with the same energy. Here we used notations $\vec{\psi}_{\lambda(\bar{\lambda})}$ for the normalized electron wave function in the nanowire. As this takes

$$\vec{\psi}_\lambda(x_1) = \hat{U}(x_1, 0) \begin{pmatrix} 1 \\ 0 \end{pmatrix} = \begin{pmatrix} \cos(\vartheta/4) \\ \sin(\vartheta/4) \end{pmatrix}, \quad \vec{\psi}_\lambda(x_2) = \hat{U}(x_2, 0) \begin{pmatrix} 1 \\ 0 \end{pmatrix} = \begin{pmatrix} \cos(\vartheta/4) \\ -\sin(\vartheta/4) \end{pmatrix}, \quad (\text{A.5})$$

$$\vec{\psi}_{\bar{\lambda}}(x_1) = \hat{U}(x_1, 0) \begin{pmatrix} 0 \\ 1 \end{pmatrix} = \begin{pmatrix} -\sin(\vartheta/4) \\ \cos(\vartheta/4) \end{pmatrix}, \quad \vec{\psi}_{\bar{\lambda}}(x_2) = \hat{U}(x_2, 0) \begin{pmatrix} 0 \\ 1 \end{pmatrix} = \begin{pmatrix} \sin(\vartheta/4) \\ \cos(\vartheta/4) \end{pmatrix}. \quad (\text{A.6})$$

Amplitude of electron tunneling from the left lead to

$$V_{kL}^{\uparrow\lambda} = \sum_k V_{kL} \langle (\uparrow, 0) | \vec{\psi}_\lambda(x_1) \rangle, \quad V_{kL}^{\downarrow\lambda} = \sum_k V_{kL} \langle (0, \downarrow) | \vec{\psi}_\lambda(x_1) \rangle, \quad (\text{A.7})$$

$$V_{kL}^{\uparrow\bar{\lambda}} = \sum_k V_{kL} \langle (\uparrow, 0) | \vec{\psi}_{\bar{\lambda}}(x_1) \rangle, \quad V_{kL}^{\downarrow\bar{\lambda}} = \sum_k V_{kL} \langle (0, \downarrow) | \vec{\psi}_{\bar{\lambda}}(x_1) \rangle, \quad (\text{A.8})$$

where V_{kL} is the transition amplitude. The tunnel matrix element for tunneling processes from the right lead can be defined in similar way by using wave functions $\vec{\psi}_{\lambda(\bar{\lambda})}(x_2)$.

Appendix B: Schrieffer-Wolff transformation

In this section we derive effective Kondo Hamiltonian for the general case of spin-dependent tunnel matrix elements. The mapping of the Anderson-like impurity model Eqs.(1-2) onto a Kondo-like model is done using a standard Schrieffer-Wolff transformation:

$$H_{eff} = e^S H e^{-S} \equiv H + [S, H] + \frac{1}{2}[S, [S, H]] + \dots (\text{B.1})$$

The first step is to eliminate the first order in tunneling amplitude terms using the following condition:

$$[S, H_0] = -H_{tun}, \quad (\text{B.2})$$

place one can classified the states by the spin structure of the wave functions at fixed point, for example, at $x = 0$ middle point of the nanowire between left and right electrodes. Assuming that eigenstates $\vec{\psi}_{\lambda(\bar{\lambda})}(0)$ in the middle of nanowire correspond to the state with spin up and spin down, we define the value of the wave function in the point x_1, x_2 where tunneling into the leads occurs ($x_{1(2)} = \pm l/2$):

the nanowire in state $\lambda, \bar{\lambda}$ can be found as follows

As a result, the effective Hamiltonian is transformed to

$$H_{eff} = H_0 + \frac{1}{2}[S, H_{tun}] + \dots (\text{B.3})$$

We choose operator S in the following form:

$$S = \left[\sum (A_{\alpha\sigma\lambda} + B_{\alpha\sigma\lambda} \hat{n}_{\bar{\lambda}}) c_{k\alpha\sigma}^\dagger d_\lambda - h.c. \right]. \quad (\text{B.4})$$

Using condition given by Eq.(B.2) one can determine the constants $A_{\alpha\sigma\lambda}, B_{\alpha\sigma\lambda}$. After straightforward calculations using Eqs.(B.3-B.4) we obtain the effective Hamiltonian $H_{eff} = H_{dir} + H_{ex}$. The first term responsible for the electron potential scattering between leads is given by

$$H_{dir} = \sum \frac{K_{\alpha\alpha'}}{4} \left[\frac{n_{\alpha\alpha'}}{2} \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* + V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* + V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\uparrow})^* \right) \right. \\ + s_{\alpha\alpha'}^z \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* - V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* - V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\uparrow})^* \right) \\ + s_{\alpha\alpha'}^x \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* + V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\downarrow\uparrow})^* + V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* \right) \\ \left. + i s_{\alpha\alpha'}^y \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* - V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\downarrow\uparrow})^* - V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* \right) \right]. \quad (\text{B.5})$$

The matrix elements used in Eq.(B.5) are given by

$$K_{\alpha\alpha'} = \frac{1}{\varepsilon_{k\alpha} - \varepsilon_0} + \frac{1}{\varepsilon_{k'\alpha'} - \varepsilon_0} - \frac{1}{U_C + \varepsilon_0 - \varepsilon_{k\alpha}} - \frac{1}{U_C + \varepsilon_0 - \varepsilon_{k'\alpha'}}, \quad (\text{B.6})$$

The next step is to find the Hamiltonian responsible for the exchange processes:

$$H_{ex} = H_1 + H_2 + H_3, \quad (\text{B.7})$$

where the Hamiltonian H_1

$$\begin{aligned} H_1 = \sum \frac{J_{\alpha\alpha'}}{2} & \left[s_{\alpha\alpha'}^z S^z \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* + V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* - V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* - V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\uparrow})^* \right) \right. \\ & + s_{\alpha\alpha'}^x S^x \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* + V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* + V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\uparrow})^* \right) \\ & \left. + s_{\alpha\alpha'}^y S^y \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* + V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* - V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* - V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\uparrow})^* \right) \right], \quad (\text{B.8}) \end{aligned}$$

is equivalent to an anisotropic Kondo model. The Hamiltonian H_2

$$\begin{aligned} H_2 = \sum \frac{J_{\alpha\alpha'}}{4} & \left[n_{\alpha\alpha'} S^z \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* - V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* + V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* - V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\uparrow})^* \right) \right. \\ & + n_{\alpha\alpha'} S^x \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* + V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* \right) \\ & \left. + i n_{\alpha\alpha'} S^y \left(-V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* + V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* - V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* \right) \right] \quad (\text{B.9}) \end{aligned}$$

describes the coupling between the spin-1/2 on the dot and the "charge" density in the leads $n_{\alpha\alpha'}$.

The Hamiltonian H_3

$$\begin{aligned} H_3 = \sum \frac{J_{\alpha\alpha'}}{2} & \left[i s_{\alpha\alpha'}^y S^x \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* - V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* - V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\uparrow})^* \right) \right. \\ & + i s_{\alpha\alpha'}^x S^y \left(-V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* + V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* - V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\uparrow})^* \right) \\ & + s_{\alpha\alpha'}^x S^z \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* - V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* - V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* + V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* \right) \\ & + s_{\alpha\alpha'}^z S^x \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* - V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* - V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* \right) \\ & + i s_{\alpha\alpha'}^y S^z \left(V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* - V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* + V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* - V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* \right) \\ & \left. + i s_{\alpha\alpha'}^z S^y \left(-V_{k\alpha}^{\uparrow\uparrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\uparrow\downarrow} (V_{k'\alpha'}^{\uparrow\uparrow})^* - V_{k\alpha}^{\downarrow\downarrow} (V_{k'\alpha'}^{\uparrow\downarrow})^* + V_{k\alpha}^{\downarrow\uparrow} (V_{k'\alpha'}^{\downarrow\downarrow})^* \right) \right], \quad (\text{B.10}) \end{aligned}$$

accounts for the Dzyaloshinsky-Moriya-like interaction.

Exchange coupling constant reads as follows

$$J_{\alpha\alpha'} = \frac{1}{\varepsilon_{k\alpha} - \varepsilon_0} + \frac{1}{U_C + \varepsilon_0 - \varepsilon_{k\alpha}} + \frac{1}{\varepsilon_{k'\alpha'} - \varepsilon_0} + \frac{1}{U_C + \varepsilon_0 - \varepsilon_{k'\alpha'}}. \quad (\text{B.11})$$

Substituting parametrization Eq.(4) into Eqs.(B.6)-(B.7) and considering exchange coupling constants of

conduction electrons at Fermi energy, $\varepsilon_{k\alpha} = \varepsilon_k - \mu_\alpha^\sigma \approx 0$, we can obtain Hamiltonian Eq.(5).

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