Kondo Force in Shuttling Devices: Dynamical Probe for a Kondo Cloud

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We consider the electromechanical properties of a single-electronic device consisting of a movable quantum dot attached to a vibrating cantilever, forming a tunnel contact with a nonmovable source electrode. We show that the resonance Kondo tunneling of electrons amplifies exponentially the strength of nanoelectromechanical (NEM) coupling in such a device and make the latter insensitive to mesoscopic fluctuations of electronic levels in a nanodot. It is also shown that the study of a Kondo-NEM phenomenon provides additional (as compared with standard conductance measurements in a nonmechanical device) information on retardation effects in the formation of a many-particle cloud accompanying the Kondo tunneling. A possibility for superhigh tunability of mechanical dissipation as well as supersensitive detection of mechanical displacement is demonstrated.

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Recent progress in the fabrication of nanoelectromechanical systems (NEMS) based on suspended carbon nanotubes [1] as well as on suspended Si [2] and SiN [3] nanowires vibrating at radio frequencies (rf) resulted in a rapidly growing amount of theoretical work [4–8] addressing the issues of the interplay between spin or charge transport and nanomechanics [9]. The observation of Coulomb blockade [2] in NEMS opened a possibility to consider the influence of strongly correlated and resonance effects on a behavior of nano-oscillators.

Usually, NEM regime implies strong coupling between the electronic and mechanical degrees of freedom. The coupling is provided by two main mechanisms. Motion of the movable dot (shuttle) between two metallic banks results in time-dependent tunneling amplitudes. On the other hand, the electron charge transport between the banks in the presence of magnetic field results in the appearance of a Lorentz force acting on the shuttle, which should also be taken into account [4,5]. The aggregate dynamics of a shuttle is that of a periodic oscillator with decrement or increment and an electron tunneling (cotunneling) parametrically dependent on this slow classical motion. In some sense the problem may be treated as a tunneling through an anharmonic vibronic system. In many cases, e.g., in shuttling devices including bending carbon nanotubes [5], the vibronic language completely describes the physical situation.

NEM coupling, like other nanometer length scale phenomena, is strongly affected by mesoscopic fluctuations. Spatial quantization of electronic motion in a quantum dot makes electromechanical transduction a sample sensitive phenomenon [2]. An exception to this rule is electromechanical coupling due to the many-body Kondo tunneling. Indeed, in this case the charge transfer is controlled by the singularity of the tunneling density of states at the energy pinned to the Fermi level of the injector and thus protected against mesoscopic fluctuations. This fact in combination with another generic feature of Kondo phenomenon—its supersensitivity to a strength of the tunneling coupling (and therefore its supersensitivity to the mechanical displacement of a quantum dot)—makes Kondo NEM coupling a promising phenomenon for practical applications.

An example of such a device is schematically shown in Fig. 1. A nanoisland is mounted on the metallic cantilever, which may vibrate under an external force. The contact between the source and drain electrodes is a combination of a time-dependent tunneling bridge between the source S and the island and a metallic bridge formed by a vibrating cantilever connecting the island with the drain D (see Ref. [2] for experimental realization).

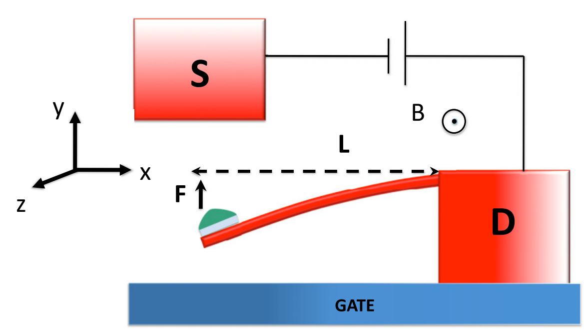

FIG. 1 (color online). Shuttle with a cantilever. The nanoisland (green) is separated from the wire (drain) by a tunnel barrier (light blue) of constant width. The width of the barrier to the source changes during the cycle. Shuttle oscillations are stimulated by initial conditions, e.g., bending the cantilever by the gate voltage at time $t = 0$.

We consider the configuration where a cantilever is displaces in the y direction, $\vec{u} = (0, y, 0)$ in a magnetic field $\vec{B} = (0, 0, B)$. In this case the Laplace force $\vec{F}$ acts on the cantilever in the same direction y, $\vec{F}_L = L \cdot \vec{I} \times \vec{B} = (0, F, 0)$. Here L is the length of the cantilever. Additionally, small electromotive force $\vec{F}_{\rm emf} = (f, 0, 0)$ acts on the electrons in the cantilever. In the limit of a strong Coulomb blockade in the nanoisland, the Kondo screening accompanies the tunneling "source-island-cantilever," and a unique possibility arises to study the contribution of a purely quantum many-particle Kondo effect on the classical oscillation of a shuttle (cantilever + island). The study of a "Kondo force" in shuttling is the main subject of this Letter.

We study two coupled subsystems: a tunneling contact "source-moving island-moving cantilever" treated as a purely quantum system in a framework of the Anderson-Kondo model, and a macroscopic wire with an attached island oscillating under an external constraining force. We work in the Kondo limit, where the nanoisland is represented by its spin $\vec{S}$, so that internal degrees of freedom are the spin-flip processes. The source-drain transport is a combination of quantum tunneling source-island-cantilever and Ohmic transport "island-drain." In this "time-dependent Schrieffer-Wolff" limit (see below), the Hamiltonian of the quantum subsystem is

$$H = H_{\rm lead} + H_{\rm ex} + \delta H, \quad H_{\rm lead} = \sum_{\alpha=l,r}\sum_{k\sigma}\xi_k c^\dagger_{\alpha k\sigma}c_{\alpha k\sigma},$$

$$H_{\rm ex} = \sum_{\alpha\alpha'} J_{\alpha\alpha'}\vec{s}_{\alpha\alpha'}\cdot\vec{S}, \qquad \delta H = \frac{eV_{\rm bias}}{2}(\hat{N}_l - \hat{N}_r). \qquad (1)$$

Here the indices l, r stand for the electronic states in the source and cantilever, respectively, $\xi_k = \varepsilon_k - \mu$ are the excitation energies of lead electrons, $\hat{N}_\alpha = \sum_{k\sigma}\hat{n}_{\alpha k\sigma}$ are the corresponding electron density operators, $\vec{s}_{\alpha\alpha'} = \frac{1}{2}\sum_{kk'}c^\dagger_{\alpha k\sigma}\vec{\tau}_{\sigma\sigma'}c_{\alpha' k'\sigma'}$, $\vec{S} = \frac{1}{2}d^\dagger_m\vec{\sigma}_{mm'}d_{m'}$ are the spin operators for the electrons in the leads and in the nanoisland, respectively, $\vec{\tau}$ and $\vec{\sigma}$ are the vectors of Pauli matrices acting on the states in the leads and dot. At small bias voltage $V_{\rm bias}$ the source and the cantilever are supposed to be in the adiabatically stationary state of thermal equilibrium. The parameters are $J_{\alpha\alpha'} = 4v^*_\alpha v_{\alpha'}/E_c$, where v_α is the tunneling amplitude between the nanoisland and the metallic lead α, E_c is the Coulomb blockade energy. The exchange couplings J_{ll} and J_{lr} are time dependent due to the dependence of the tunneling amplitude between the source and the moving nanoisland on the island position $v_l = v_l[\vec{u}(t)]$. The time dependence of this amplitude is a set of pulses corresponding to electron injection from the metallic reservoir to the shuttle periodically approaching the bank S [10]. We confine our treatment to the simplest case of $S = 1/2$ (odd occupation of a nanoisland in the neutral state) and the single channel tunneling between the nanoisland and the leads.

The oscillations of the cantilever with the attached nanoisland are determined by the classical Newton equations

$$\ddot{\vec{u}} + \frac{\omega_0}{Q_0}\dot{\vec{u}} + \omega_0^2\vec{u} = \frac{1}{m}\vec{F}, \qquad (2)$$

where $\omega_0 = \sqrt{k/m}$ is the oscillator frequency of the free cantilever, Q_0 is a quality factor of the NEM device.

Our aim is to study the spin and charge transport by means of a shuttle oscillating in accordance with Eq. (2) in the presence of many-particle Kondo screening described by the Hamiltonian (1). The coupling between the classical and quantum subsystem is realized via the parameters $J_{ll}(\vec{u})$, $J_{lr}(\vec{u})$, $\vec{F}(\vec{u})$, where the time dependence $\vec{u}(t)$ should be calculated self-consistently. Meanwhile, J_{rr} does not depend on displacement $\vec{u}$ (see Fig. 1).

The cotunneling Hamiltonian may be rationalized by means of the Glazman-Raikh rotation, which in our situation is time dependent:

$$\begin{pmatrix} c_{lk\sigma} \\ c_{rk\sigma}\end{pmatrix} = \begin{pmatrix}\cos\vartheta_t & -\sin\vartheta_t \\ \sin\vartheta_t & \cos\vartheta_t\end{pmatrix}\begin{pmatrix}\psi_{1k\sigma} \\ \psi_{2k\sigma}\end{pmatrix} \equiv U_t\begin{pmatrix}\psi_{1k\sigma} \\ \psi_{2k\sigma}\end{pmatrix}, \qquad (3)$$

with $\tan\vartheta_t = |v_r/v_l(t)|$. Use of this transformation for diagonalization of the Schrödinger operator $\mathcal{L} = -i\hbar d/dt + H(t)$ results in generation of an additional term H_B proportional to $-i\hbar U_t^{-1}\partial_t U$ in the transformed Hamiltonian (see, e.g., Ref. [13]),

$$H' = H_{\rm lead} + H_B + H_{\rm ex} + \delta H, \qquad (4)$$

$$H_{\rm lead} = \sum_{a=1,2}\sum_{k\sigma}\xi_k\psi^\dagger_{ak\sigma}\psi_{ak\sigma},$$

$$H_B = i\hbar\frac{d\vartheta_t}{dt}\sum_{k\sigma}(\psi^\dagger_{1k\sigma}\psi_{2k\sigma} - \psi^\dagger_{2k\sigma}\psi_{1k\sigma}), \qquad (5)$$

$$H_{\rm ex} = \frac{J}{4}\sum_{kk',\sigma\sigma',m,m'}\psi^\dagger_{1k\sigma}\vec{\tau}_{\sigma\sigma'}\psi_{1k'\sigma'}d^\dagger_m\vec{\sigma}_{mm'}d_{m'},$$

$\delta H = \frac{eV_{\rm bias}}{2}[(\hat{N}_2 - \hat{N}_1)\cos 2\vartheta_t + \sum_{k\sigma}(\psi^\dagger_{1k\sigma}\psi_{2k\sigma} + {\rm H.c.})\sin 2\vartheta_t]$. The term H_B may be treated as an additional gauge potential in the lead Hamiltonian describing a Berry-like phase [13] generated by shuttle motion [14]. Only the even partial wave ψ_1 survives in the cotunneling term $H_{\rm ex}$ with the time-dependent effective indirect exchange coupling $J(t) = J_{ll}(t) + J_{rr}$ (see, e.g., Ref. [15]). This time dependence may be parametrized [9] with the assumption that the source-island tunneling amplitude is an exponential function of a distance y between the source and the moving nanoisland, while the tunneling nanoisland-cantilever is constant: $v_r = v_0$, $v_l = v_1\exp[y(t)/\lambda]$. The spatial coordinates are counted off the equilibrium position of the cantilever, so that v_1 is exponentially small, $v_1/v_0 \sim \exp(-y_0/\lambda)$. Here λ is the confinement radius (tunnel length) of the electron wave function within the island, y_0 is the distance between the source and the island at equilibrium.

We suppose that the shuttling mode is slow enough and the electron transport is adiabatic; i.e., the exchange couplings J_{ll}, J_{lr} depend parametrically on time via the displacement coordinates $\vec{u}(t)$. Then to find the tunneling current one may trace the time dependence of local occupations of the left and right banks (source and nanoisland) near the point of tunneling contact. The current operator is

$$\hat{I} = \frac{e}{2}\frac{d}{dt}(\hat{N}_r - \hat{N}_l), \tag{6}$$

where in the case of an immovable nanoisland only the even mode 1 contributes to the current. In our case both modes 1 and 2 are involved in the tunneling transport due to the term H_B in Eq. (4). After the Glazman-Raikh rotation the current operator transforms into

$$\hat{I} = \frac{d}{dt}\hat{\mathcal{Q}}_t + \frac{d}{dt}\hat{q}_t, \tag{7}$$

where

$$\begin{aligned} \hat{\mathcal{Q}}_t &= \frac{e}{2}\cos 2\vartheta_t(\hat{N}_1 - \hat{N}_2), \\ \hat{q}_t &= -\frac{e}{2}\sin 2\vartheta_t \sum_{k\sigma}(\psi^\dagger_{1k\sigma}\psi_{2k\sigma} + \psi^\dagger_{2k\sigma}\psi_{1k\sigma}). \end{aligned} \tag{8}$$

Here the operator $\hat{\mathcal{Q}}_t$ controls the time-dependent electron occupation in the source lead, and the operator $\hat{q}_t$ is responsible for all tunneling and cotunneling processes including the admixture of odd components $\psi_{2k\sigma}$ to the tunneling charge transport induced by the gauge field H_B [13,16].

The time-dependent Glazman-Raikh angle defined by (9) results in adiabatic time dependence of the Breit-Wigner factor

$$\sin^2 2\vartheta_t = \frac{4\Gamma_l\Gamma_r}{(\Gamma_l + \Gamma_r)^2} = \frac{1}{\cosh^2\frac{[y(t)-y_0]}{\lambda}}. \tag{9}$$

Using the Friedel-Langreth sum rule [17], one may write [18]

$$N_1 - N_2 = \frac{\delta_t}{\pi}, \tag{10}$$

where $\delta_t = \delta_\uparrow + \delta_\downarrow$ is a total time-dependent Friedel phase. At the unitary limit, $\delta_{\uparrow,\downarrow} = \pm\pi/2$.

We are interested in the Kondo effect contribution to the tunneling current. This contribution is characterized by the spin dependent scattering phase shift $\delta_\sigma(\varepsilon)$ in the source lead, which approaches the unitarity limit $\pi/2$ at $T \to 0$ and $\varepsilon \to \varepsilon_F$. In the adiabatic limit $\hbar\omega_0 \ll k_B T_K^{\min}$ under conditions $(k_B T, g\mu_B B, |eV_{\text{bias}}|) \ll k_B T_K^{\min}$ the phenomenological Fermi liquid Hamiltonian H_{Noz} may be used [19] (here k_B is the Boltzmann constant, μ_B is the Bohr magneton, and g is the Landé factor). In this Hamiltonian both scattering and interaction are scaled by the time-dependent Kondo temperature $T_K(t)$ taking minimal value $T_K^{\min}$ at maximal distance from the source.

In order to get the full tunnel current in the adiabatic approximation, we (i) calculate a linear response with respect to both bias $eV_{\text{bias}} \ll k_B T_K$ and $\hbar d\vartheta_t/dt \leq \hbar\omega_0\vartheta_{\max} \ll k_B T_K$, (ii) take into account cancellations arising due to emergent $SU(2)$ symmetry associated with channels [14,20], and (iii) perform averaging with the adiabatic Hamiltonian (4) and (5) at zero bias and zero temperature. The finite temperature and bias effects are accounted for by the Nozières method [19,21]. As a result, the tunnel current $\bar{I}_t = \bar{I}_0(t) + \bar{I}_{\text{int}}(t)$ consists of two parts: the Friedel phase contribution

$$\bar{I}_0(t) = \frac{e}{2\pi}\cos 2\vartheta_t \frac{d\delta_t}{dt} \tag{11}$$

and the "Ohmic" current [22]

$$\bar{I}_{\text{int}}(t) = \frac{e^2}{\hbar}V_{\text{bias}}\frac{d}{dt}\left[\sin 2\vartheta_t \int_{-\infty}^{t} dt' \sin 2\vartheta_{t'}\Pi^R_{t-t'}\right], \tag{12}$$

where

$$\Pi^R_t = -\frac{i}{2}\sum_{k\sigma}\sum_{\alpha\neq\gamma}[G^R_{\alpha k\sigma}(t)G^K_{\gamma k\sigma}(-t) + G^K_{\alpha k\sigma}(t)G^A_{\gamma k\sigma}(-t)],$$

and $G^\Lambda_{\alpha=1,2}$ are bold retarded ($\Lambda = R$), advanced ($\Lambda = A$), and Keldysh ($\Lambda = K$) Green's function [22].

Let us first rewrite the Friedel part of the tunneling current (11) via the parameters characterizing the Kondo tunneling in the low-temperature strong coupling limit at $T \ll T_K^{\min}$, where

$$k_B T_K(t) = D_0 \exp\left[-\frac{\pi E_c}{4(\Gamma_l + \Gamma_r)}\right], \tag{13}$$

D_0 is the ultraviolet cutoff for the Kondo problem with a scale of the band energy in the source, $\Gamma_\alpha = \pi\rho_0|v_\alpha|^2$, ρ_0 is the density of electronic states at the Fermi level ε_F. In our adiabatic regime T_K parametrically depends on time, following the time dependence of $\Gamma_l(t)$. The Hamiltonian H_{Noz} [19,21], [21] establishes the relations between δ_σ, B, and T_K near the unitary limit, such as $\delta_t = 2|eV_{\text{bias}}|/[k_B T_K(t)] \ll 1$. The magnetic field enters only into the relative Friedel phase $\delta_\uparrow - \delta_\downarrow = \pi - 2g\mu_B B/[k_B T_K(t)]$. We neglect the influence of magnetic field on T_K, since we work in the limit $g\mu_B B \ll k_B T_K^{\min}$. Alternatively, a nonuniform magnetic field negligible at the dot and gradually increasing along the cantilever could be assumed in the model.

In the adiabatic limit the Friedel phase δ_t and Glazman-Raikh angle ϑ_t are not independent, but are connected through (9) and (13)

$$\frac{1}{\delta_t}\frac{d\delta_t}{dt} = \frac{\pi E_c}{4\Gamma_0}\sin 2\vartheta_t \frac{d\vartheta_t}{dt}, \tag{14}$$

with $\Gamma_0 = \pi\rho_0|v_0|^2 \ll E_c$. Thus, the Friedel contribution to tunnel current can be expressed in terms of shuttle velocities as follows:

$$\bar{I}_0(t) = \frac{\dot{y}}{\lambda}\frac{eE_c}{8\Gamma_0}\frac{eV_{\text{bias}}}{k_B T_K(t)}\frac{\tanh\left(\frac{y-y_0}{\lambda}\right)}{\cosh^2\left(\frac{y-y_0}{\lambda}\right)}. \tag{15}$$

Here, the time dependence of the tunnel current is predetermined by the time dependence of tunnel integrals for the nanoisland moving in the y direction, i.e., by the function $y(t)$ and its derivative $\dot{y}$. Moreover, one can see that even in the case of possible instability, large amplitude oscillations are exponentially suppressed. The typical behavior of $I_0(t)$ is shown in Fig. 2. The nonsinusoidal form of current is associated with time dependence of both tunnel width and Kondo temperature.

The second term $\bar{I}_{\rm int}(t)$ given by Eq. (12) leads to Ohmic contribution to the current with unitary conductance $G_0 = e^2/h$ (see discussion below):

$$\bar{I}_{\rm int}(t) = G_0 V_{\rm bias} \sin^2 2\vartheta_t \sum_\sigma \sin^2 \delta_\sigma. \tag{16}$$

The force in the right-hand side of the Newton equation (2) is a sum of the driving force F_0, the Lorentz force F_L, and electromotive (emf) force $F_{\rm emf}$:

$$F(y, t) = F_0(t) + \bar{I}_t BL + F_{\rm emf}. \tag{17}$$

The emf force can be estimated as $F_{\rm emf} \sim \dot{y}(BL)^2 G_0$ [23]. Because of sequential geometry of the electric circuit, the current $\bar{I}_t = \bar{I}_0(t) + \bar{I}_{\rm int}(t)$ is the tunneling current defined by (15) and (16). In the limit of small bias voltage $|eV_{\rm bias}| \ll k_B T_K^{\rm min}$, electrons in the source and the cantilever are supposed to be in an adiabatically stationary state of thermal equilibrium. Then the parametrization (15) is valid and with accuracy to small parameters $O\{[eV_{\rm bias}/(k_B T_K^{\rm min})]^2, [(g\mu_B B)/(k_B T_K^{\rm min})]^2\}$ the Lorentz force may be written as

$$F_L = F_{\rm ad}[y(t)] - \dot{y}\frac{dF_{\rm ad}}{dy}\frac{\hbar\pi E_c}{16\Gamma_0 k_B T_K^{(0)}}, \tag{18}$$

where $F_{\rm ad} = 2BLG_0 V_{\rm bias} \cosh^{-2}\frac{[y(t)-y_0]}{\lambda}$ and $T_K^{(0)}$ is a Kondo temperature at equilibrium position. Small correction to the adiabatic Lorentz force in Eq. (18) may be

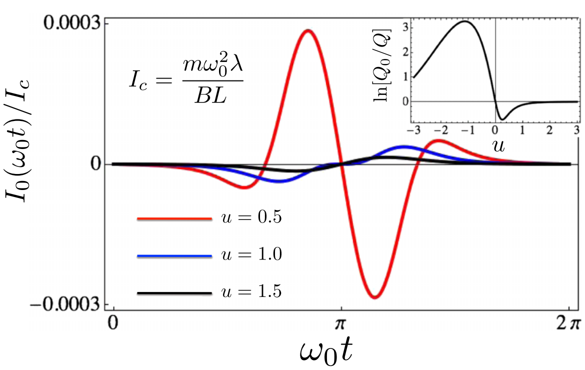

FIG. 2 (color online). Time dependence of the current I_0 for different values of asymmetry parameter $u = y_0/\lambda$. For all three curves, shuttle oscillates with amplitude $y_{\rm max} = \lambda$, $\hbar\omega_0/(k_B T_K^{\rm min}) = 10^{-3}$, $|eV_{\rm bias}|/(k_B T_K^{\rm min}) = g\mu_B B/(k_B T_K^{\rm min}) = 0.1$, with $T_K^{(0)} = 2$ K, $\lambda/L = 10^{-4}$. Inset: $\ln[Q_0/Q]$ as a function of u, $Q_0 = 10^4$.

considered as a first term in the expansion over a small nonadiabatic parameter $\omega_0\tau \ll 1$, where τ is the retardation time associated with inertia of the Kondo cloud. Using this interpretation, one gets τ

$$\tau = \frac{\hbar\pi E_c}{16\Gamma_0 k_B T_K^{(0)}} = \frac{1}{2}\left|\frac{Q^{-1}(B) - Q^{-1}(-B)}{\omega(B) - \omega(-B)}\right|, \tag{19}$$

where $Q(B)$ and $\omega(B)$ are the quality factor and the oscillator's frequency at finite magnetic field B, respectively.

Equation (19) allows one to obtain information about dynamics of the Kondo clouds from the analysis of the experimental investigation of the mechanical vibrations. The retardation time associated with dynamics of the Kondo cloud is parametrically large compared with the time of formation of the Kondo cloud $\tau_K = \hbar/T_K$ and can be measured owing to small deviation from adiabaticity [24]. We would also like to emphasize a supersensitivity of the quality factor to the change of the equilibrium position of the cantilever characterized by the parameter y_0. The plot $\ln[Q_0/Q]$ is presented in the inset of Fig. 2. From this plot one can see that both suppression $Q > Q_0$ and enhancement $Q < Q_0$ of the dissipation of nanomechanical vibrations (depending on the direction of the magnetic field and the equilibrium position of the cantilever) can be stimulated by Kondo tunneling. The latter demonstrate potentialities for the Kondo induced electromechanical instability which will be a subject for separate analysis.

Equations (15), (18), and (19) represent the central results of this Letter. On the one hand, we have shown that the electric current associated with the Kondo effect results in a magnetic field dependent Q factor allowing us to fine-tune the nanomechanical resonator. On the other hand, the non-Ohmic part of the current provides information about retardation effects related to the motion of the Kondo cloud. Thus, the measurement of the Kondo forces in a single electron transistor is complementary to conductance measurements information.

In conclusion, we have shown that the Kondo phenomenon in single electron tunneling gives a very promising and efficient mechanism for electromechanical transduction on a nanometer length scale. Measurement of the nanomechanical response on Kondo transport in a nanomechanical single-electronic device enables one to study the kinetics of the formation of Kondo screening and offers a new approach for studying nonequilibrium Kondo phenomena. The Kondo effect provides a possibility for superhigh tunability of the mechanical dissipation as well as supersensitive detection of mechanical displacement.

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[21] The Nozières [19] fixed point Fermi liquid Hamiltonian $H_{\rm Noz} = \sum_{k\sigma}\tilde{\xi}_k \phi^\dagger_{k\sigma}\phi_{k\sigma} - \sum_{kk'\sigma}[(\tilde{\xi}_k + \tilde{\xi}'_k)/2\pi\rho_0 k_B T_K]\times \phi^\dagger_{k\sigma}\phi_{k'\sigma} + (1/\pi\rho_0^2 k_B T_K)\sum_{kk'}\phi^\dagger_{k\uparrow}\phi_{k'\uparrow}\sum_{kk'}\phi^\dagger_{k\downarrow}\phi_{k'\downarrow}$ is written in terms of $\phi_{k,\sigma}$ fermions with a bandwidth $k_B T_K$ and spectrum $\tilde{\xi}_k$. It is used for accounting for inelastic contributions to the transport (see Ref. [15]).
[22] See also Ref. [15] for the details of the derivation of electric conductance for the Kondo problem using the Kubo formula.
[23] The smallness of the electromotive force is guaranteed by the upper bound on the external magnetic field applied to the cantilever $\Phi/\Phi_0 \cdot L/\lambda < E_c/\Gamma_0 \cdot |eV_{\rm bias}|/(k_B T_K)$, where $\Phi = B \cdot \mathcal{S}_\lambda$ is a flux through the area $\mathcal{S}_\lambda \sim \lambda^2$ and $\Phi_0 = h/e$ is a flux quantum. The value of the field corresponding to this bound is roughly estimated as $B_u \sim 10$ T.
[24] Parametric largeness of $\tau \gg \tau_K$ justifies neglecting corrections of the order of $\omega_0\tau_K \ll \omega_0\tau$ to the scattering phase and the Nozières Hamiltonian.