Enhancements of Andreev conductance induced by the photon/vibron scattering

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We analyze the subgap spectrum and transport properties of the quantum dot embedded between one superconducting and another metallic reservoirs and additionally coupled to an external boson mode. Emission/absorption of the bosonic quanta induces a series of the subgap Andreev states, that eventually interfere with each other. We discuss their signatures in the differential conductance both, for the linear and nonlinear regimes.

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I. INTRODUCTION

The bosonic modes, like photons \([1]\) or vibrational degrees of freedom \([2]\), can strongly affect electron tunneling through the nanoscopic systems \([3]\). When a level spacing of nanoobject is large in comparison to the boson energy \(\omega_0\) and a line-broadening is sufficiently narrow, a series of the side-peaks \([4]\) may appear due to emission/absorption of the bosonic quanta. Such features (spaced by \(\omega_0\)) have been really observed in measurements of the differential conductance in several nanojunctions \([4, 5]\).

Similar bosonic modes are currently studied also in the systems, where the quantum dots/impurities are coupled with superconducting reservoirs \([9, 10]\). Since the proximity effect spreads electron pairing onto these quantum dots, the bosonic features manifest themselves in a quite peculiar way. They could be observed by the Josephson \([9, 10]\) and the Andreev spectroscopies \([12, 17]\), in photon-assisted subgap tunneling \([18]\), transient phenomena \([19]\), or in prototypes of the nano-refrigerators operating due to the multi-phonon Andreev scattering \([20]\).

First of all, in a subgap regime (assuming \(\omega_0\) smaller than energy gap \(\Delta\) of superconductor) the bosonic features are expected to be more numerous than in the normal state. This is a consequence the proximity effect, mixing the particle and hole excitations. Secondly, it has been shown numerically \([12, 16, 20]\) that the linear (zero-bias) Andreev conductance exhibits the bosonic features spaced by a half of \(\omega_0\). To our knowledge, this intriguing theoretical result was neither clarified on physical arguments nor checked experimentally. Verification would be feasible by the tunneling spectroscopy using e.g. low-frequency vibrations of some heavy molecules or slowly-varying ac electromagnetic field. Let us emphasize, that such low-energy boson mode need not be related with any pairing mechanism of the superconducting reservoir.

The purpose of our paper is to provide a simple analytical argument, explaining the reduced frequency \(\omega_0/2\) of the bosonic features in the linear Andreev conductance versus the gate-voltage. We also study in detail the multiple subgap states originating from the boson emission/absorption processes. We analyze their signatures both in the quantum dot spectrum and the tunneling transmission. The latter quantity can be probed by the (low-temperature) differential conductance as a function of the source-drain bias. We predict that the multiple Andreev states could be seen with a period, dependent on the gate voltage.

For calculations we consider the setup displayed in figure 1. It can be practically realized in a single electron transistor (SET) using e.g. the carbon nanotube suspended between the external electrodes (like in Refs \([5, 6]\)). Another possibility could be the scanning tunneling microscope (STM), where the conducting tip (N) probes some vibrating quantum impurity (QD) hosted in a superconducting (S) substrate \([21]\). In both SET and STM configurations such boson mode can be eventually related to external ac field.

In what follows we introduce the Hamiltonian and discuss the method for treating the bosonic mode. We next investigate the bosonic signatures in the QD spectrum and in the subgap Andreev conductance. For clarity, we focus on the limit \(\Gamma_N \ll \omega_0\) whereas the second coupling \(\Gamma_S\) can be arbitrary. In the last section we address the correlation effects.

FIG. 1: (color online) A scheme of the quantum dot between the metallic (N) and superconducting (S) electrodes and coupled to the monochromatic boson (phonon or photon) mode.
II. MICROSCOPIC MODEL

For microscopic description of the tunneling scheme shown in Fig. 1 we use the Anderson impurity model

$$\hat{H} = \hat{H}_N + \hat{H}_S + \hat{H}_{mol} + \hat{H}_T.$$  

where the energy level is lowered by the polaronic shift $\xi = -\lambda^2/\omega_0$ and the effective potential $U = -2\lambda^2/\omega_0$.

Boson operators are shifted $\hat{d}_\sigma^{(t)} = \hat{d}_\sigma^{(t)} - \frac{\lambda}{\omega_0} \sum_\sigma \hat{d}_\sigma^{(t)}$ whereas fermions are dressed with the polaronic cloud $\hat{X} = e^{-(\lambda/\omega_0) (\hat{a}^\dagger - \hat{a})}$. (4)

Reservoirs $\hat{H}_\beta$ are invariant on the unitary transformation (2) but the operator $\hat{X}$ appears in the hybridization term $\hat{H}_T$. For simplicity we absorb it into the effective coupling constants $\Gamma_\beta = 2\pi \sum_k |V_{k\beta}|^2 \langle \mathcal{X}^\dagger \mathcal{X} \rangle$ which can be defined for the wide band limit.

The effective single particle excitation spectrum is given by the Green’s function

$$G_\sigma(\tau_1, \tau_2) = -i \left\langle \hat{T}_\tau \hat{d}_\sigma(\tau_1) \hat{d}_\sigma^\dagger(\tau_2) \right\rangle_{\hat{H}_T},$$  

where $\hat{T}_\tau$ denotes the time ordering operator. Since trace is invariant on the unitary transformations $\left\langle \ldots \right\rangle_{\hat{H}} = \left\langle \ldots \right\rangle_{\hat{H}_T}$ it is convenient to compute the statistical averages with respect to $\hat{H}$. In particular, (3) can be expressed as

$$G_\sigma(\tau_1, \tau_2) = -i \left\langle \hat{T}_\tau \hat{d}_\sigma(\tau_1) \hat{d}_\sigma^\dagger(\tau_2) \right\rangle_{\hat{H}_T} \sim \left\langle \hat{T}_\tau \hat{X}(\tau_1) \hat{X}^\dagger(\tau_2) \right\rangle_{\hat{H}_{bas}}$$  

because the fermionic and bosonic degrees of freedom are separated by the Lang-Firsov transformation. From a standard procedure [1, 22] one obtains

$$\left\langle \hat{T}_\tau \hat{X}(\tau_1) \hat{X}(\tau_2) \right\rangle_{\hat{H}_{bas}} = \exp \left\{ - (\lambda/\omega_0)^2 \times \left[ (1 - e^{-\omega_0(\tau_1 - \tau_2)}) (1 + N_p) + (1 - e^{\omega_0(\tau_1 - \tau_2)}) N_p \right] \right\}$$  

with the Bose-Einstein distribution $N_p = [e^{\beta\omega_0} - 1]^{-1}$. Fourier transform of the Green’s function (7) is found as

$$G_\sigma(\omega) = \sum_l g_\sigma(\omega - l\omega_0) e^{-\lambda^2/2 N_p} (1 + N_p) \left[ \frac{\lambda}{\omega_0} \right]^2 \frac{2\lambda^2/2 N_p (1 + N_p)}{2 \lambda^2/2 N_p (1 + N_p)},$$  

where $I_l$ denote the modified Bessel functions and $g_\sigma(\tau_1, \tau_2) = -i \left\langle \hat{T}_\tau \hat{d}_\sigma(\tau_1) \hat{d}_\sigma^\dagger(\tau_2) \right\rangle_{\hat{H}_{bas}}$ is the fermionic part of (3). In the ground state (8) simplifies to

$$\lim_{T \to 0} G_\sigma(\omega) = \sum_l g_\sigma(\omega - l\omega_0) e^{-g^2} g^l \frac{g^l}{l!}$$  

with the adiabatic parameter $g = (\lambda/\omega_0)^2$. 

Due to the proximity induced on-dot pairing the single particle Green’s function $G_{\uparrow}(\tau_1, \tau_2)$ is mixed with the (anomalous) propagator

$$F(\tau_1, \tau_2) = -i \left\langle \hat{T}_r \hat{d}^\dagger_1(\tau_1) \hat{d}^\dagger_1(\tau_2) \right\rangle_{\hat{H}} \ .$$

This important fact has been remarked in the previous considerations of $dc$ Josephson current \[11\] and it also plays significant role for the Andreev spectroscopy (see the next section). The boson part of the anomalous propagator \[10\] takes the following form

$$\left\langle \hat{T}_r \hat{X}^\dagger(\tau_1) \hat{X}(\tau_2) \right\rangle_{\text{bos}} = \left\langle \hat{T}_r \hat{X}^\dagger(\tau_1) \hat{X}(\tau_2) \right\rangle_{\hat{H}_{\text{fer}}} \ .$$

Their Fourier components obey the Dyson equation

$$\begin{bmatrix}
 g(\omega) & f(\omega) \\
 f^*(-\omega) & -g^*(-\omega)
\end{bmatrix}^{-1} = \begin{bmatrix}
 \omega - \tilde{\varepsilon} & 0 \\
 0 & \omega + \tilde{\varepsilon}
\end{bmatrix} - \Sigma_{\text{QD}}^0(\omega) - \Sigma_{QD}^{\text{corr}}(\omega),$$

where $\Sigma_{\text{QD}}^0$ is the selfenergy matrix of uncorrelated molecular dot and the second contribution $\Sigma_{QD}^{\text{corr}}$ is due to the effective Coulomb interaction $\hat{U}$. In the wide-band limit the selfenergy $\Sigma_{\text{QD}}^0(\omega)$ can be expressed as

$$\Sigma_{\text{QD}}^0(\omega) = -i \frac{\Gamma_N}{2} \begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} - \frac{\Gamma_S}{2} \gamma(\omega) \begin{bmatrix} 1 & \frac{\Delta}{\omega} \\ \frac{\Delta}{\omega} & 1 \end{bmatrix}$$

with

$$\gamma(\omega) = \begin{cases}
\frac{\omega}{\sqrt{\Delta^2 - \omega^2}} & \text{for } |\omega| < \Delta, \\
\frac{\omega}{\sqrt{\omega^2 - \Delta^2}} & \text{for } |\omega| > \Delta.
\end{cases}$$

We investigated the effective spectral function $\rho(\omega) = -\pi^{-1} \text{Im}G(\omega + i0^+) \text{ at zero temperature}$, focusing on the intermediate electron-boson coupling $g \sim 1$. Figures\[2\] show the QD spectrum for $\hat{U} = 0$, neglecting the correlation effects $\Sigma_{QD}^{\text{corr}}$. Influence of the Coulomb potential $\hat{U}$ is discussed in section V.

\[2\] illustrates evolution of the bosonic features with respect to the superconductor gap $\Delta$. In the normal state (for $\Delta = 0$) such lorentzian peaks are located at $\omega = \tilde{\varepsilon} + i\omega_0$ (with integer $l \geq 0$) and their broadening is $\Gamma_N + \Gamma_S$. For finite $\Delta$ all peaks split into the lower and upper ones due to the induced on-dot pairing. In the extreme limit $\Delta \gg \Gamma_S$ the selfenergy $\Sigma_{QD}^0(\omega)$ becomes static

$$\lim_{\Delta \gg \Gamma_S} \Sigma_{QD}^0(\omega) = -\frac{1}{2} \begin{bmatrix} i\Gamma_N & \Gamma_S \\ \Gamma_S & i\Gamma_N \end{bmatrix}$$

therefore the effective quasiparticle energies evolve to $\omega_0 \pm \sqrt{\tilde{\varepsilon}^2 + (\Delta / 2)^2}$ and their broadening shrinks to $\Gamma_N$.

Focusing on such superconducting atomic limit \[10\] we show in Fig. 3 the subgap bosonic peaks with respect to $\tilde{\varepsilon}$. In the SET configuration the energy level $\tilde{\varepsilon}$ would be tunable by applying the gate voltage. In particular, these peaks may overlap with each other when $\tilde{\varepsilon} \approx \omega_0 / 2$ as reported earlier in the Refs.\[12\] \[16\] \[20\]. This effect can be deduced analytically from

$$l\omega_0 + \frac{\sqrt{\tilde{\varepsilon}^2 + (\Gamma_S / 2)^2}}{\Gamma_S} = l\omega_0 - \sqrt{\tilde{\varepsilon}^2 + (\Gamma_S / 2)^2} .$$

\[17\]
The neighboring peaks ($l' = l + 1$) overlap when $\omega_0/2 = \sqrt{\varepsilon^2 + (\Delta/2)^2}$. For small $\Gamma_S$ such situation takes place at $\varepsilon \approx \frac{1}{2}\omega_0$. Other crossings would be eventually possible for the higher-order multiplications of $\omega_0/2$.

Figure 4 displays the subgap spectrum $\rho(\omega)$ as a function of the coupling $\Gamma_S$. From (17) we conclude that for $\varepsilon = 0$ the bosonic peaks overlap at $\Gamma_S = \omega_0$. Energy of these crossing points is $\omega = (\frac{1}{2} + l)\omega_0$. Here (for $g = 1$) we observe four such crossings, but for stronger electron-boson couplings a number of the in-gap states and their crossings would increase.

IV. ANDREEV CONDUCTANCE

Under nonequilibrium conditions the charge current can be transmitted at small voltage $|eV| < \Delta$ via the Andreev scattering, engaging the in-gap states. This anomalous transport channel occurs when electrons from the metallic lead are converted into the Cooper pairs (propagating in superconducting electrode) with the holes reflected back to N electrode. The resulting current $I_A(V)$ can be expressed by the Landauer-type formula [24]

$$I_A(V) = \frac{2e}{h} \int d\omega T_A(\omega) [f_{FD}(\omega - eV) - f_{FD}(\omega + eV)]$$

with the Fermi-Dirac function $f_{FD}(\omega) = \left[\frac{e^{\omega/k_BT} + 1}{e^{\omega/k_BT} - 1}\right]^{-1}$ and the Andreev transmittance [24]

$$T_A(\omega) = \Gamma_N^2 |F(\omega)|^2.$$  

The Andreev transmittance occurs when $\omega$ coincides with the subgap quasiparticle states. In our present case we thus expect a number of such enhancements due the bosonic features. Let’s remark that $T_A(-\omega) = T_A(\omega)$ implies the Andreev conductance $G_A(V) = \frac{dI_A(V)}{dV}$ to be an even function of the bias $V$.

V. CORRELATION EFFECTS

In various experimental realizations of the quantum dots (such as self-assembled InAs islands [25], carbon nanotubes [26, 27] or semiconducting nanowires [28, 29]) attached to the superconducting leads the energy gap $\Delta$ was safely smaller than the repulsion potential $U$. For this reason, in the subgap Andreev spectroscopy the correlations hardly contributed any Coulomb blockade. Instead of it, they can eventually induce the singlet-doublet quantum phase transition [30, 31] and/or the Kondo physics [32]. In this paper we consider the strongly asymmetric coupling $\Gamma_N \ll \Gamma_S$ and focus on the deep subgap regime $\Gamma_N, \Gamma_S \ll \Delta$, therefore the Kondo-type effects [30–32] would be rather negligible.

Analysis of such singlet-doublet transition for the vibrating quantum dot has been previously addressed [13] using the NRG technique. We revisit the same issue here, determining the differential Andreev conductance (unavailable for the NRG calculations [15]), because this quantity could be of interest for experimentalists. For the sake of simplicity, we analyze the correlation effects
in the superconducting atomic limit $\Delta \gg \Gamma_S$. Hamiltonian of the molecular quantum dot \cite{3} can be additionally updated with the pairing terms $\frac{1}{2}\Gamma_S (\hat{d}_l^\dagger \hat{d}_l^\dagger + \hat{d}_l \hat{d}_l)$ originating from the static off-diagonal parts of the self-energy matrix \cite{10}.

In the absence of the boson field (i.e for $\lambda = 0$) the effective quasiparticle energies are given by $\pm U/2 \pm E_d$, where $E_d = \sqrt{(\varepsilon + U/2)^2 + (\Gamma_S/2)^2}$. In the realistic situations only two branches $\pm (U/2 - E_d)$ appear in the subgap regime, whereas the other high energy states $\pm (U/2 + E_d)$ overlap with a continuum beyond the gap. The quantum phase transition (QPT) from the singlet $|0\rangle + \varepsilon |\uparrow\downarrow\rangle$ to doublet $|\sigma\rangle$ configuration occurs at $U/2 = E_d$ \cite{33}. In order to estimate quantitatively the Andreev conductance we use the off-diagonal Green’s function $f(\omega)$ \cite{34,34}, restricting to its subgap part

$$f_{sub}(\omega) \simeq \left[ \frac{\alpha \omega}{\omega + \frac{\omega^2}{2} - \left( \frac{U}{2} - E_d \right)} - \frac{\alpha \omega}{\omega + \frac{\omega^2}{2} + \left( \frac{U}{2} - E_d \right)} \right]$$

with the usual BCS coefficient $\omega = \Gamma_S / 4E_d$ and the spectral weight $\alpha = \left[ \exp \left( \frac{U}{2k_B T} \right) + \exp \left( \frac{E_b}{k_B T} \right) \right] / Z$, where $Z = 2 \exp \left( \frac{U}{2k_B T} \right) + \exp \left( \frac{E_b}{k_B T} \right) + \exp \left( \frac{E_{d}}{k_B T} \right)$. The missing part of spectral weight $1 - \alpha$ belongs to the high-energy states (outside the gap). At zero temperature this subgap weight changes abruptly from $\alpha = 1$ (in the singlet state when $U/2 < E_d$) to $\alpha = 0.5$ (in the doublet state when $U/2 > E_d$).

In figure 6 we plot the Andreev conductance obtained for the half-filled quantum dot $\varepsilon = -U/2$ (QPT occurs then at $U = \Gamma_S$). We notice the subgap conductance enhancements around $|eV| = U/2 - E_d$. Yet, exactly at the QPT, both the singlet and doublet contributions cancel each other. Formally, this is due to the odd (asymmetric) structure of the Green’s function \cite{20}.

The superconducting atomic limit solution can be generalized onto $g \neq 0$ case in a straightforward way. The unitary transformation \cite{2} implies $\varepsilon \rightarrow \tilde{\varepsilon}, U \rightarrow \tilde{U}$ and following the steps \cite{10} we can determine the off-diagonal Green’s function. At zero temperature, we find

$$F_{sub}(\omega) \simeq \frac{\alpha \omega}{\omega + \frac{\omega^2}{2} - \left( \frac{U}{2} - E_d \right) + s \omega_0}$$

with $s \equiv \text{sign}(\tilde{U} - E_d)$.

Figure 7 shows the Andreev conductance obtained for the half-filled quantum dot using $g = 1, \omega_0 / \Gamma_N = 10, \Gamma_S / \Gamma_N = 20, T = 0$. The bosonic side-peaks give rise to additional subgap branches, similar to what has been reported for the spectral function \cite{15}. Right at the QPT, the zero-bias conductance again vanishes $G(0) \rightarrow 0$ and we observe only the higher order maxima at $|eV| = l \omega_0$ (with $l \geq 1$). Away from the QPT, the Andreev conductance shows the usual maxima at $|eV| = |\tilde{U}/2 - E_d| + l \omega_0$ whose spectral weights depend on $\tilde{U}$ and $l$.

VI. SUMMARY

We have investigated the subgap spectrum and transport properties of the quantum dot coupled between the metallic and superconducting electrodes in presence of the external boson mode $\omega_0$. We have found that the induced Andreev states eventually cross each upon varying the gate potential (through $\varepsilon$) or due to the correlations (via quantum phase transition from the singlet to doublet configurations). We have explored their signatures in the measurable charge transport. The tunneling conductance of such multilevel `molecule’ shows a series of characteristic enhancements, depend on: the gate voltage with
frequency \( \omega_0/2 \) (which can be deduced from Eqn. [17]), the bias \( V \) applied between external leads (Fig. 5), and the correlations (Fig. 7). External boson reservoir can thus substantially affect the anomalous Andreev current and it can be probed experimentally using the low-energy vibrational modes or the slowly-varying \( ac \) fields [1].

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