Pump-probe scheme for electron-photon dynamics in hybrid conductor-cavity systems

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Recent experiments on nanoscale conductors coupled to microwave cavities put in prospect transport investigations of electron-photon interplay in the deep quantum regime. Here we propose a pump-probe scheme to investigate the transient dynamics of individual electron-photon excitations in a double quantum dot-cavity system. Excitations pumped into the system decay via charge tunneling at the double dot, probed in real time. We investigate theoretically the short-time charge transfer statistics at the dot, for periodic pumping, and show that this gives access to vacuum Rabi oscillations as well as excitation dynamics in the presence of double dot dephasing and relaxation.

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Introduction.— Hybrid quantum electrodynamic (QED) structures that incorporate nanoscopic conductors in superconducting circuits have attracted a lot of interest over the last few years. Several experiments have demonstrated coupling between both single [1–3] and double [4] quantum dots (DQDs) and a superconducting microwave cavity. Hybrid conductor-cavity systems have large potential for applications in quantum information technology [10–12] and can be used as a characterization tool for nanoscale systems [12–13]. Moreover, they put in prospect transport investigations as well as applications of electron-photon interactions in the deep quantum limit [16–20]. This includes DQD masers [21–22], entanglement detection in Cooper-pair splitters [22], testbeds for Franck-Condon physics [23], detection schemes for Majorana fermions [24–27], photon emission statistics [28] and efficient heat engines [29].

A key feature of nanoscopic conductor-cavity system which to date has received little attention is the dynamics of electron-photon excitations. Besides being of fundamental interest and of importance for conductor-cavity applications, knowledge of the dynamics is instrumental in identifying and characterizing parasitic effects responsible for excitation dephasing and relaxation [30]. However, an experiment investigating the dynamics would arguably require, in the same conductor-cavity system, both short time excitation manipulation and detection. Importantly, in circuit QED systems microwave photons can be generated and controlled in real time with large accuracy [31–32]. Moreover, time resolved counting of electrons in quantum dot systems has been demonstrated experimentally [33, 34].

In this letter we combine these two capacities and propose a pump-probe scheme for investigating the transient dynamics of electron-photon excitations in a hybrid conductor-cavity system (See Fig. 1). Single excitations are pumped into the system via an externally driven qubit coupled to the cavity. The excitations relax by electron tunneling at the DQD, probed by monitoring the dot occupation in real time. We analyze theoretically the short time charge transport statistics at the DQD for periodic pumping, and show that it provides information on vacuum Rabi oscillations, as well as excitation relaxation time scales. Moreover, the transport statistics clearly displays how DQD dephasing and relaxation alter the dynamics, key knowledge in the experimental efforts to reach the strong conductor-cavity coupling regime.

System.— The system under investigation consists of a DQD coupled to a superconducting microwave cavity, and tunnel coupled to a single lead electrode. The DQD, occupied with at most one excess electron, contains two active levels. Moreover, an externally controlled superconducting qubit, in Fig. 1 exemplified by a transmon [36], is coupled to the cavity. The DQD and the qubit are coupled to the cavity with the same strength $g_0$, and kept at resonance with the fundamental cavity mode, at frequency $\omega_0$. In addition the characteristic impedance $Z_0$ of the cavity is taken much smaller than the resistance quantum $R_Q = \hbar/e^2$. Under these conditions the DQD-cavity-qubit system is described by the generalized Tavis-Cummings (TC) Hamiltonian

$$\hat{H}_S = \hbar \omega_0 \hat{a}^\dagger \hat{a} + \frac{\hbar \omega_0}{2} \left( \hat{\sigma}_z + \hat{d}_e^\dagger \hat{d}_e - \hat{d}_g^\dagger \hat{d}_g \right) + \hbar g_0 \left[ \hat{a}^\dagger (\hat{\sigma}_- + \hat{d}_g^\dagger \hat{d}_e) + \hat{a} (\hat{\sigma}_+ + \hat{d}_e^\dagger \hat{d}_g) \right].$$ (1)

Here $\hat{a}^\dagger (\hat{a})$ is the photon creation (annihilation) operator, and $\hat{d}_{e/g}^\dagger (\hat{d}_{e/g})$ the electron creation (annihilation) operators for the excited/ground (bonding/anti-bonding)
Starting in the system ground state, exciting the qubit|±⟩ denotes the excited/ground state of the qubit, and ˆc_± denotes the creation operator for an electron at energy ε_±. We will take the detuning between levels in the dots, forming the DQD, to be zero. For this case the effective lead-DQD tunneling amplitudes are equal for |e⟩ and |g⟩. The tunnel Hamiltonian is then given by

\[ H_L = \sum_k \epsilon_k \hat{c}_k^\dagger \hat{c}_k, \]

with the Hamiltonian ˆH_L = \sum_k \epsilon_k \hat{c}_k^\dagger \hat{c}_k, where \hat{c}_k^\dagger denotes the creation operator for an electron at energy \epsilon_k. We will take the detuning between levels in the dots, forming the DQD, to be zero. For this case the effective lead-DQD tunneling amplitudes are equal for |e⟩ and |g⟩. The tunnel Hamiltonian is then given by

\[ H_L = \sum_k \epsilon_k \hat{c}_k^\dagger \hat{c}_k + H.c. \]

We consider the sequential tunneling regime Γ \ll \hbar \omega, with Γ = \pi \sum_k |t_k|^2 δ(ε_k \pm \hbar \omega_0/2) the DQD-tunneling rate. Applying the Born-Markov approximation a quantum master equation (QME) \[d\hat{\rho}/dt = L\hat{\rho}\] is derived for the time evolution of the reduced qubit-DQD-cavity system density operator \hat{\rho}. By taking \mu = 0 and \hbar \omega_0/2 \gg k_B T the occupation of the lead at the energies at which electrons tunnel into or out of the DQD does not depend on the temperature. The Liouvillian \mathcal{L} then becomes

\[ L\hat{\rho} = -i/\hbar [\hat{H}_S, \hat{\rho}] + \Gamma \left[ \mathcal{D}(\hat{\rho}, \hat{\gamma}^\dagger) + \mathcal{D}(\hat{\rho}, \hat{\gamma}) \right], \]

with \mathcal{D}(\hat{\rho}, \hat{\gamma}) = 2\hat{\gamma}^\dagger \hat{\gamma} \hat{\rho} - \hat{\gamma}^\dagger \hat{\gamma} \hat{\rho} - \hat{\rho} \hat{\gamma}^\dagger \hat{\gamma} \hat{\rho} \] for any operator \hat{\gamma}.

The system dynamics after an excitation at \[t = 0\] is illustrated in Fig. 2. The vacuum Rabi oscillations decay, with the excited state probabilities \[e^{-\Gamma t^2/4}\] for \[\Gamma \ll g_0\]. Hence, the system excitation is eventually transformed into an electron-hole pair excitation in the lead as |−e0⟩ → |−00⟩ → |−g0⟩, i.e. the DQD electron, in |e⟩, tunnels out into the lead followed by a back tunneling into |g⟩.

**Pump-probe scheme** – We consider an experimentally relevant pump-probe scheme, with a periodic pumping, or excitation, of the qubit. Importantly, the pump period \[\tau\] constitutes a versatile tool for investigating the different time scales of the dynamics. We focus on periodic steady-state operation, where the state of the system is entirely determined by the time \[t\] passed since the last pumping, i.e. \[\hat{\rho}(t + m\tau) = \hat{\rho}(t), \] with \[m = 0, 1, 2...\]. The density operator is then given by

\[ \hat{\rho}(t) = e^{\mathcal{L}t} \mathcal{R}_0 \mathcal{L}^{(\tau-t)} \hat{\rho}(t), \quad 0 < t < \tau, \]

where \[\mathcal{R}_\gamma = \hat{\gamma} \hat{\gamma} \mathcal{R}_\gamma^\dagger\]

The probe consists of a non-invasive charge detector (not shown in Fig. 1), monitoring the individual dot-lead tunneling events in time \[\tau\]. Since tunneling into (out of) the DQD creates a hole (an electron) excitation in the lead, we here consider the statistics of electron and hole transfers into the lead per pumping cycle. The full counting statistics for a time periodic system can be obtained along the lines of [41]. We introduce an electron (hole) counting field \[\chi_e (\chi_h)\] into the Liouvillian of Eq. 5, so that \[\mathcal{L} \to \mathcal{L}(\chi_e, \chi_h)\]. The cumulant generating function (CGF) over \[N \geq 1\] periods is

\[ S_N(\chi_e, \chi_h) = \ln\langle [\mathcal{R}_0 \mathcal{L}^{(\chi_e, \chi_h)}]^{N}\rangle_0, \]

where \[\langle \gamma \rangle_0 = \text{tr}[\hat{\gamma} \hat{\rho}(0)]\]. From the CGF one obtains the mean and variance of the number of electrons (holes) emitted per period as

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\[ S_N(\chi_e, \chi_h) = \ln\langle [\mathcal{R}_0 \mathcal{L}^{(\chi_e, \chi_h)}]^{N}\rangle_0, \]
in this long period regime the system acts as an ideal approximation. The QME then reduces to an ordinary and can be averaged out in the density matrix (secular probabilities). This picture is substantiated by the electron-photon dy-

an electron-hole pair excitation in the lead (see Fig. 2a). \( N \) per period, during the probability of emitting \( \Gamma \tau \) electrons and \( \chi \) periods, is given by \( \Gamma \tau \tau \Gamma \) and a period \( \Gamma \) describes one such event during \( \chi \) transfer cycles. Legend in b) same as in a). c) Electron (green lines) and hole (red lines) noise as a function of \( \Gamma \tau \) for \( N = 1 \). For \( N \gg 1 \) (black lines) electron and hole noise coincide. Inset: Average number of excitations \( n \) as a function of \( \Gamma \tau \). In both main panel and inset, the numerical (solid lines) and analytical, large \( \Gamma \tau \) (dashed lines) results, are shown. d) Full probability distribution \( P_N(n) \) for \( N = 50 \) for different \( \Gamma \tau \).

\[
\langle n_{e(h)} \rangle_N = \frac{1}{N} \frac{\partial S_N}{\partial (\chi_{e(h)})} \bigg|_{\chi_e,\chi_h=0} \quad \text{and} \quad \langle \delta n_{e(h)}^2 \rangle_N = \frac{1}{N} \frac{\partial^2 S_N}{\partial (\chi_{e(h)})^2} \bigg|_{\chi_e,\chi_h=0}.
\]

Moreover, the probability of emitting \( n_e \) electrons and \( n_h \) holes per period, during \( N \) periods, is given by \( P_N(n_e, n_h) = \int_0^\infty \frac{d\chi_e}{\sqrt{2\pi}} \frac{d\chi_h}{\sqrt{2\pi}} e^{-iN(n_e\chi_e+n_h\chi_h)+S_N(\chi_e,\chi_h)}/(2\pi)^2. \)

**Strong coupling regime.** To illustrate in a compelling way the proposed pump-probe scheme, we consider the strong coupling regime \( g_0 \gg \Gamma \) and a period \( \tau \gg 1/g_0 \) (Fig. 3). In this regime vacuum Rabi oscillations are fast and can be averaged out in the density matrix (secular approximation). The QME then reduces to an ordinary master equation in the TC-basis, with the dynamics determined by \( \Gamma \) and \( \tau \). For long periods \( \tau \gg 1/\Gamma \) the system will describe a complete cycle of transferring one energy quantum from the qubit, via a cavity photon, to an electron-hole pair excitation in the lead (see Fig. 2c). This picture is substantiated by the electron-photon dynamics, illustrated by the corresponding time dependent probabilities \( P_{eq} \) for the (set of) TC-eigenstates, shown in Fig. 2d. Here \( \sigma \) denotes an empty (U) or occupied (O) dot and \( q \) the number of excitations. We stress that in this long period regime the system acts as an ideal electron-hole pair pump, or equivalently, a pump of energy quanta \( \hbar \omega. \)

Decreasing the period \( \tau \) there are corrections to the complete transfer cycle. This is clear from the plot of the TC probabilities in Fig. 2c: higher energy states with \( q \gg 2 \), including two cavity photons, acquire a finite population. Importantly, the effect of finite \( \Gamma \tau \) is clearly manifested in the full counting statistics. A careful inspection of Eq. 4 shows that the corrections scale as \( x = e^{-\Gamma \tau}/4 \), much smaller than unity for large \( \Gamma \tau \). Solving the CGF to leading order in \( x \) we get

\[
\frac{S_N(\chi_e,\chi_h)}{N} = \frac{1}{N}(\chi_e + \chi_h) + \frac{x(e^{-2i(\chi_e+\chi_h)} - 1)}{8} + x \frac{\delta S}{N}.
\]

The first term describes the complete cycle, while the second term describes the first order corrections. The third term \( \delta S = (\cos(\chi_h) - 4)/(3 + 5e^{-i(\chi_e+\chi_h)} + 3e^{i(\chi_e+\chi_h)} - (e^{-2i(\chi_e+\chi_h)} - 1))/8 \) gives the short time behavior, going to zero for \( N \gg 1 \). Focusing on the \( N \gg 1 \) limit, the electron and hole transfer statistics become identical, with the number of transferred excitations \( n = n_e = n_h \), and we can write the full distribution (first order in \( x \))

\[
P_N(n) = (1 - x/8)\delta_{n1} + (x/8)\delta_{nN,1-N-2}.
\]

From Eq. 7, as well as from the plot of the numerically evaluated \( P_N(n) \), in Fig 3 it is clear that the effect of \( x \) is i) a non-zero probability for \( N \gg 2 \) excitations to be transferred, and consequently, ii) a reduced probability for the ideal cycle, transferring \( N \) excitations. This can be explained as follows: for non-zero \( x \) there is a finite probability that an excitation is pumped out of the system during a qubit ramp. The \( P_N(n) \) to leading order in \( x \) describes one such event during \( N \) cycles. When an excitation is pumped out of the system it does not generate an electron-hole pair. Moreover, during the subsequent period the system will remain in the ground state, where no electron-hole emission can occur, thus resulting in a total loss of two periods for missed excitations. We note that a similar mechanism was discussed in [42].

For arbitrary \( \Gamma \tau \), the numerically calculated probability distribution \( P_N(n) \) for \( N \gg 1 \) as well as the two lowest cumulants for the electron and hole transfer statistics are shown in Fig. 3. The \( P(n) \) shows that the “missed two excitations” effect, for \( \Gamma \tau \sim 1 \), turns into a “missed several pairs of excitations” effect, with a finite probability for missing an even number of excitations 2, 4, 6, .... The probability to miss an odd number of excitations is smaller, albeit increasing for decreasing \( \Gamma \tau \).

The average number of emitted electrons and holes are equal and independent of \( N \), \( \langle n_e \rangle_N = \langle n_h \rangle_N = \langle n \rangle \). In Fig. 3b \( (n) \) is plotted together with the analytical solution for large \( \Gamma \tau \), \( (n) = 1 - e^{-\Gamma \tau}/4 \), following from Eq. 6. For \( \Gamma \tau \gg 10 \) the two results coincide, while for smaller \( \Gamma \tau \) the analytical solution overestimates the number of emitted excitations.

In contrast to the average number, the electron and hole noise differ for \( N \sim 1 \) while for \( N \gg 1 \) we find \( \langle \delta n_e^2 \rangle_N = \langle \delta n_h^2 \rangle_N \equiv \langle \delta n^2 \rangle \). As shown in Fig. 3 for \( \Gamma \tau \gg 1 \), both the electron and hole noise, for any \( N \), are well captured by the large \( \Gamma \tau \) approximation, \( \langle \delta n^2 \rangle_N = \langle \delta n_e^2 \rangle_N = \langle \delta n_h^2 \rangle_N = \langle \delta n^2 \rangle N \gg 1 \).
When treated by considering two opposite cases of pumping peri-
vacuum Rabi oscillations are manifested as oscillations
line (in both panels) is strong coupling result (\(g\)).

\[
\frac{(N+1)/2N}{e^{-\Gamma \tau/4}} \text{and } \frac{(\delta n^2)_N}{= \frac{(3N+5)/6N}{e^{-\Gamma \tau/4}}},
\]

obtained from Eq. (6). The noise is thus exponentially
between the excited, \(|e\rangle\), and ground, \(|g\rangle\), states of the
DQD. As illustrated in Fig. 5 this has several consequences
for the average number of transferred excitations \(\langle n \rangle\) (similarly for the fluctuations, not shown).

Increasing the dephasing from \(\Gamma_D \ll g\) the amplitude of the
crude Rabi oscillations are first suppressed, reaching a
minimum around \(\Gamma_D \sim g\). However, increasing the dephasing
the oscillations are revived for \(\Gamma_D > g\).

In this strongly dephased regime the system can be seen
as a coherent qubit-cavity subsystem, displaying vacuum
Rabi oscillations with frequency \(2\sqrt{g}\), coupled incoherently
to the DQD. Moreover, with increasing \(\Gamma_D\) the
DQD relaxation and dephasing – In recent conductor-
cavity experiments[4, 5] DQD decoherence prevented
reaching the strong coupling regime. This raises the question
how the signatures of coherent electron-photon in-
teraction in transport quantities are altered in the presence of dephasing and relaxation. Dephasing and relaxation,
with rates \(\Gamma_D\) and \(\Gamma_R\) respectively, are accounted for by adding the terms
\[
\mathcal{L}_D[\rho] = \frac{\Gamma_D}{2} \mathcal{D}[\hat{d}_e^{\dagger} \hat{d}_e - \frac{\delta^2}{2}] \rho\]
and
\[
\mathcal{L}_R[\rho] = \frac{\Gamma_R}{2} \mathcal{D}[\hat{d}_g^{\dagger} \hat{d}_g] \rho\]
to the Liouvillian in Eq. (4).

The effect of dephasing is to suppress the coherence

Vacuum Rabi oscillations – Outside the strong coupling regime, i.e. for \(\Gamma \sim g\), the coherent electron-
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FIG. 4: (color online) Manifestation of vacuum Rabi oscillations. The average \(\langle n \rangle\) (left panel) and fluctuations \(\langle \delta n^2 \rangle\) (right panel) of number of transferred electron-hole pairs per
pumping period as a function of \(\Gamma \tau\) for different \(g_0/\Gamma\). Dashed
line (in both panels) is strong coupling result (\(g\)).

FIG. 5: (color online) Effects of dephasing and relaxation. Average number of excitations \(\langle n \rangle\) as a function of \(\Gamma \tau\) for
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different relaxation rates \(\Gamma_R\), with \(\Gamma_D = g_0\) (right panel).

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