Characteristic features in non-Markovian noise spectrum of transport current

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Based on the construction of an efficient scheme for non-Markovian noise calculation, we analyze the noise spectrum of transport current through interacting single quantum dot and double dots. We show that the non-Markovian correction is remarkable. It leads to a number of characteristic spectral structures, including that due to finite bandwidth and that sensitive to and enhanced by the magnitude of Coulomb interaction.

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Shot noise due to charge discreteness in mesoscopic transport has stimulated great interest in recent years. It provides much rich information beyond the average current [1, 2]. Conventionally, shot noise and higher cumulants of current in full counting statistics (FCS) are largely restricted to zero frequency, and the Born-Markov memoryless master equation approach is employed [3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16]. Memory effects on the non-Markovian memory kernel prescription [22] have been discussed recently [3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16].

In this work, instead of the infinite Coulomb interaction and even finite bandwidth, we investigate the circuit electron transport setup as readout devices for solid-state charge qubits [20].

Our study is based on the particle-number “n”-resolved quantum master equation (QME) approach [8, 14], implemented with the memory kernel prescription. Let us start with the conventional QME for the reduced system operator containing memory [22]:

$$\dot{\rho}(t) = -i\mathcal{L}\rho(t) - \int_0^t d\tau \Sigma(t, \tau)\rho(\tau),$$

with $\Sigma(t, \tau) = \Sigma(t - \tau) \equiv \langle \mathcal{L}'(\tau) e^{i\mathcal{L}(t-\tau)} \mathcal{L}'(\tau) \rangle_B$. Here, $\mathcal{L}' \equiv [H, \cdot ]$ is the system Liouvillian; $\mathcal{L}'(\cdot) \equiv [H'(\cdot), \cdot ]$ is the system–reservoirs tunnel coupling Liouvillian; $\langle \cdot \rangle_B$ denotes an average over the reservoirs bath degree of freedom of electrodes with different chemical potentials under bias voltage. The electron reservoirs bath Hamiltonian assumes $h_B = \sum \epsilon_{ak} a_k^\dagger a_k$. The system $H$ is rather general and written in terms of the creation $\{d_\mu^\dagger \equiv d_\mu^\dagger \}$ and annihilation $\{d_\mu \equiv d_\mu \}$ operators of electron in system orbitals. The tunnel coupling in the $h_B$–interaction picture is $H'(t) = \sum a_k \epsilon\mu e^{i\omega t} (t_a t_k d_\mu a_k^\dagger + H.c.) e^{-i\omega t} \equiv \sum_{\alpha\mu} \{\tilde{f}^\dagger_{\alpha\mu}(t) d_\mu + d_\mu^\dagger \tilde{f}^-_{\alpha\mu}(t)\}$, in the view point of dissipative dynamics. The coefficients of leads on central system is completely characterized by the reservoir correlation functions $C^\sigma_{\alpha\mu}(t; \tau) = \langle \tilde{f}^\sigma_{\alpha\mu}(t) \tilde{f}^-_{\alpha\mu}(\tau) \rangle_B$, where $\tilde{f}$ takes the opposite sign of $\sigma = \pm$. It results in the self-energy kernel $\Sigma(t)$ in Eq. (1) of

$$\Sigma(t) = \sum_{\alpha\mu\sigma} \left\{ \left\{ d_\mu^\dagger D^{(\sigma)}_{\alpha\mu}(t; \mathcal{L}) + H.c. \right\} \right\},$$

where $D^{(\sigma)}_{\alpha\mu}(t; \mathcal{L}) = \sum \epsilon_\mu e^{i\omega t} C^\sigma_{\alpha\mu}(t) d_\mu$. The above QME invokes the second–order Born approximation and is valid for sequential tunneling. For later use, let us introduce the half–Fourier transform that resolves $\Sigma(t)$ with the formal solution $\tilde{\rho}(\omega) = \Sigma(\omega + i(L - \omega))^{-1} \rho(t_0)$. The stationary state is then $\tilde{\rho} = [-i\omega \tilde{\rho}(\omega)]_{\omega=0}$. The resolution of Eq. (2) reads

$$\tilde{\Sigma}(\omega)\tilde{\rho} = \sum_{\alpha\mu\sigma} \left\{ d_\mu^\dagger \tilde{D}^{(\sigma)}_{\alpha\mu}(L - \omega)\tilde{\rho} - \dot{\tilde{\rho}}\tilde{D}^{(\sigma)}_{\alpha\mu}(L + \omega) \right\},$$

for the non-Markovian memory kernel prescription [22].
with $\tilde{D}^{(\sigma)}_{\mu}(\omega) = \sum_{\nu} \int_0^\infty dt \, e^{-\sigma \omega t} C^{(\sigma)}_{\mu \nu}(t) d^\nu_{\mu}$.

Rather than the above unconditional QME, a richer information contained equation for conditional state will be more desirable. This is the well established particle-number-resolved QME, which contains such transport information as current, shot noise, and even all the higher moments of current fluctuations \cite{14}. The key quantity in the “n”-resolved QME is $\rho^{(n)}$, the reduced system state conditioned on $n_a$ electrons passed through the tunnel junction between the α-th lead and the central system. The unconditional state is related to $\rho^{(n)}$ via $\rho = \sum_{n_a} \rho^{(n_a)}$. Following Refs. \cite{2,14,23}, especially the idea of decomposition of the Hilbert space presented in \cite{14}, the non-Markovian version can be readily formulated out as (setting $D^{(\sigma)}_{\mu} = \sum_{\alpha} D^{(\sigma)}_{\mu \alpha}$ and $\alpha' \neq \alpha$)

$$\dot{\rho}^{(n_a)}(t) = -i\mathcal{L} \rho^{(n_a)}(t) - \sum_{\mu} \int_0^t dt' \left\{ d^\nu_{\mu} D^{(\sigma)}_{\mu}(t - t'; \mathcal{L}) \times \rho^{(n_a)}(t) - D^{(\sigma)}_{\alpha \mu}(t - t'; \mathcal{L}) \rho^{(n_a)}(t) d^\sigma_{\alpha} \right\}$$

$$- D^{(\sigma)}_{\alpha \mu}(t - t'; \mathcal{L}) \rho^{(n_a+\sigma+1)}(t) d^\sigma_{\alpha} + \text{H.c.}.$$  

(4)

The conditional state is straightforwardly related to the distribution function for the tunneled electron numbers: $P(n_a, t) = \text{Tr}[\rho^{(n_a)}(t)]$. With the knowledge of $P(n_a, t)$, all transport properties can be obtained. First, for the current, through $I_{\alpha}(t) = -\frac{e}{\hbar} \sum n_a P(n_a, t)$, we have $I_{\alpha}(t) = -\int_0^t \, dt J^{(\gamma)}_{\alpha}(t, \tau) \rho(\tau)$ or $I_{\alpha}(t) = -J^{(\gamma)}_{\alpha}(t, \omega) \rho(\omega)$, with

$$J^{(\gamma)}_{\alpha}(t, \omega) = \sum_{\mu} \left\{ [\tilde{D}^{(\gamma)}_{\alpha \mu}(\omega) \hat{O}] d^\sigma_{\mu} + [\tilde{D}^{(\gamma)}_{\alpha \mu}(-\omega) \hat{O}] d^\sigma_{\mu} + d^\sigma_{\mu} [\hat{O} \tilde{D}^{(\gamma)}_{\alpha \mu}(\omega)] + d^\sigma_{\mu} [\hat{O} \tilde{D}^{(\gamma)}_{\alpha \mu}(-\omega)] \right\}.$$  

(5)

The stationary current reads $\tilde{I}_{\alpha} = -J^{(\gamma)}_{\alpha}(\omega, 0) \rho(\omega)$.

For the current noise spectrum, we have $S(\omega) = aS_{L}(\omega) + bS_{R}(\omega) - ab S_\text{c}(\omega)$. The capacitive coupling related parameters $a$ and $b$ link the total current circuit with the left and right junction currents as $I(t) = aI_L(t) - bI_R(t)$. Accordingly, $S_{L/R}(\omega)$ is the noise spectrum of $I_L/R(t)$, and $S_\text{c}(\omega)$ is the charge fluctuation spectrum on the central dots, satisfying $\tilde{Q}(t) = -[I_L(t) + I_R(t)]$. For $S_{L/R}(\omega)$, it follows the MacDonal’s formula $S_{\alpha}(\omega) = 2\omega \int_0^\infty dt \text{sin} \omega t \sum_{n_a} \langle n_a^2 \rangle \delta(n_a^2 \pi^2) \sum_{\gamma} \left\{ \text{Re} \left[ \sum_{\mu} \langle n_a^2 \rangle \text{Im} \langle \tilde{J}^{(\gamma)}_{\alpha \mu}(\omega) \rho(\omega) \hat{O} \rangle \right] \right\}$. With the help of Eq. (4) and the MacDonal’s formula, the current noise spectrum can be finally expressed as

$$S_{\alpha}(\omega) = 4\omega \text{Im} \left\{ \text{Tr} \left[ \tilde{J}^{(\gamma)}_{\alpha}(\omega) \hat{N}^{\dagger}(\omega) \rho_{\text{c}}(\omega) \right] \right\} + 2\omega \text{Re} \left\{ \text{Tr} \left[ \tilde{J}^{(\gamma)}_{\alpha}(\omega) \hat{N}(\omega) \rho_{\text{c}}(\omega) \right] \right\},$$

(6)

where $\hat{N}(\omega) = i(\hat{\Sigma}(\omega) + i(\mathcal{L} - \omega)^{-1} \tilde{J}^{(\gamma)}_{\alpha}(\omega) \rho(\omega) \hat{O} \hat{\omega}$.

For the charge noise, the current conservation leads to $S_{\text{c}}(\omega) = 2S_{LR}(\omega) + S_{L}(\omega) + S_{R}(\omega)$, with $S_{LR}(\omega) = 2\omega \text{Im} \left\{ \text{Tr} \left[ \tilde{J}^{(\gamma)}_{L}(\omega) \hat{N}^{\dagger}(\omega) \rho_{\text{c}}(\omega) \right] \right\}$ resulted from the MacDonald’s formula as described above.

To illustrate the non-Markovian reservoir coupling effect on noise spectrum, we choose a simplified Lorentzian spectral density, $\Gamma_{\alpha \mu}(\omega) = \Gamma_{\alpha \mu}^{\text{WS}}/[(\omega - \mu_\alpha)^2 + \Gamma_{\alpha \mu}^{\text{WS}}^2]$, if it is nonzero. The WBL is achieved when $\Gamma_{\alpha \mu}^{\text{WS}} \rightarrow \infty$. The reservoir correlation functions determine the rates of electron tunneling between the system and leads. They are related to the spectral density via the fluctuation–dissipation theorem \cite{24}: $C^{(\gamma)}_{\alpha \mu}(t) = \frac{1}{\Omega} \omega e^{-\omega^2 \mu_\alpha} \Gamma_{\alpha \mu}(\omega)/(1 + e^{\sigma(\omega - \mu_\alpha)})$, with $\mu_\alpha$ the chemical potential of the α–lead and $\beta$ the inverse temperature.

Before going to applications, two general remarks are worthwhile. (i) The key feature of non-Markovian master equation is the memory kernel in Eq. (1), which leads to the frequency-dependent tunneling rates. As an illustration, take a look at the term $\tilde{D}^{(\gamma)}_{\alpha \mu}(\epsilon_{\mu} - \omega) d^\sigma_{\alpha}$ in Eq. (3). It describes the tunneling from the α–lead to the system state of energy $\epsilon_{\mu}$, associated with an energy of $\omega$ absorbed (setting $h = 1$). On the other hand, the Markovian approximation is of $\omega$-independent tunneling rates \cite{21}. (ii) The frequency-dependent tunneling rates have profound effect on the noise spectrum. For $|\omega - \mu_\alpha| \gg W$, the tunneling rates become zero, leading to vanishing noise at high frequency limit. In contrast, the Markovian approach cannot predict this behavior \cite{21}, as the tunneling rate exhibits no frequency dependent. Even for WBL model where both the non-Markovian and Markovian treatments do not approach to the zero noise result, qualitative differences arise, as will see below.

**Preliminary Insight.** Consider the simplest system, a single dot with single spinless level, $H = \epsilon_0 a^\dagger a$, involving only two states, the empty (0) and the occupied (1). We also assume large bias and low temperature: $\mu_\alpha \gg \mu_\beta$ and the Fermi function $f_\alpha(\omega) = 1$ for $\omega < \mu_\alpha$, and zero otherwise. For the charge occupation fluctuation on the dot, $S_N(\omega) = \int_0^\infty dt \, e^{i\omega t} \langle N(t)N(0) + N(0)N(t) \rangle$, we find that (denoting $\gamma = |\Gamma_L/\Gamma_R|$)

$$S_N(\omega) = \begin{cases} 1, & 0 < \omega < \omega_{L0}, \\ 1 + \gamma/2, & \omega_{L0} < \omega < \omega_{R0}, \\ (1 + \gamma)/2, & \omega > \omega_{R0}, \end{cases}$$  

(7)

where $\omega_{L0} = |\mu_\alpha - \epsilon_0|$, and $S_N^0(\omega) = 4\Omega/(\Gamma^2 + \omega^2)$ is the result under Markovian approximation, with $\Gamma = \Gamma_L + \Gamma_R$. Therefore, in frequency domain, the non-Markovian effect is manifested as a discontinuity at frequencies $\omega_{L0}/(\Gamma^2 + \omega^2)$ reflecting the nature of resonance with the Fermi level \cite{13}. Figure 1 depicts the time domain counterpart of Eq. (1), i.e., the charge number correlation function $S_N(t)$ where nonexponential decay behavior is evident. Note that $S_N(t)$ is nothing but the probability for an electron remains in the dot at time $t$, assuming it was at $t = 0$. The nonexponential decay, which is typical for non-Markovian dynamics, essentially arises from the nonlocal time memory effect in Eq. (2). Throughout this work, we set $\mu_L = eV/2$ and $\mu_R = -eV/2$.

**Interacting Single Dot.** The relevant Hamiltonian reads $H = \sum_\mu \epsilon_\mu a^\dagger_\mu a_\mu + U n_1 a^\dagger_0 a_0$, where $\hat{n}_1 = a^\dagger_0 a_0$ and $\mu$ the
spin index, so that $C_{0\mu}^{(\pm)}(\omega) = C_{\alpha}^{(\pm)}(\omega)\delta_{\mu\nu}$. To display the effect of Coulomb interaction more transparently, we assume the dot level spin-degenerate, $\epsilon_1 = \epsilon_1 = \epsilon_0$. Also, we consider the transport in strong Coulomb blockade regime where $\epsilon_0 + U > \mu_L > \epsilon_0 > \mu_R$.

Let us start with region of $\omega < \omega_o$ and carry out the noise spectra analytically, which in terms of Fano factor $S(\omega)/(2I)$ read (denoting $I_{\text{eff}} = 2\Gamma_L + \Gamma_R$)

$$F_L(\omega) = \frac{2(4\Gamma_L^2 + \Gamma_R^2 + \omega^2)}{\Gamma_L^2 + \omega^2} - \frac{2\Gamma_{\text{eff}}\Pi_L}{\Gamma_L(\Gamma_L^2 + \omega^2)};$$

$$F_R(\omega) = \frac{2(4\Gamma_R^2 + \Gamma_L^2 + \omega^2)}{\Gamma_R^2 + \omega^2} - \frac{\Gamma_{\text{eff}}}{\Gamma_R(\Gamma_R^2 + \omega^2)} \times \left[ \Pi_R + \left( \omega^2 + 2\Gamma_L/\Gamma_{\text{eff}} - 6\Gamma_L^2 \right) \Phi_R \right];$$

$$F_c(\omega) = \frac{2\omega^2}{\Gamma_L^2 + \omega^2} \frac{\Gamma_{\text{eff}}\Phi\Gamma_{\text{eff}} - \Phi_{\text{eff}}}{\Gamma_L\Gamma_R(\Gamma_L^2 + \omega^2)};$$

where $\Pi_\alpha = \omega\Phi_\alpha\Phi + 2\Gamma_L\Gamma_R\Phi\widetilde{\Phi} + 2\Gamma_R\Phi_\alpha\Phi + \omega^2\Gamma_\alpha\Phi_\alpha/\Gamma_{\text{eff}}$, $\Phi_\alpha(\epsilon_0, \omega) = \chi_\alpha(\epsilon_0 - \omega) - \chi_\alpha(\epsilon_0 + \omega) \equiv \Phi_\alpha$, $\Phi = \Phi_L + \Phi_R$, $\Phi = \Phi_L - \Phi_R$, and $\chi_\alpha(x) \equiv \frac{1}{2\pi i} \text{Res} \left[ \frac{1}{x} + \frac{i}{2} (x - \mu_A) \right]$ ($\Psi$ is the digamma function). At zero frequency, $\Phi_\alpha(\epsilon_0, 0) = 0$, the first two terms in Eq. (8) vanish, and the Markovian result (the first term) is recovered [21]. The second terms in Eq. (8), which arise in fact from the renormalization of the dot level, reduce the tunneling rates and suppress the shot noise. They can be neglected in the high-frequency regime considered hereafter.

Increasing frequency usually results in non-Markovian effect more distinct, as the short timescale dynamics becomes more important. Consider the noise spectrum at intermediate frequencies. From Fig. 2, we notice a number of characteristic frequencies, i.e., $\omega_{L(R)} = (|\mu_{L(R)}| - \epsilon_0)$ and $\omega_{L(R)} = \epsilon_0 + U - |\mu_{L(R)}|$, at which sudden change of fluctuation magnitudes takes place. The processes involve electron transfer between dot and leads, associated with the energy absorption/emission of detection. The resonant step-structures are completely absent in Markovian treatment; much richer non-Markovian transient dynamics occur in time domain.

Markovian theory concludes also that in Coulomb blockade regime the noise spectrum does not depend on $U$ [21]. This result seems plausible, since noise is determined by the random occupation of the dot, which would be no longer relevant to the magnitude of $U$, provided the dot has already been in the strong Coulomb blockade regime. This argument is, however, valid only for long-time dynamics. With the increase of frequency, the energy nonconserved states will be involved, due to the energy uncertainty in the relevant short-time regime, and the noise becomes $U$-dependent.

In the WBL where the memory arises only from Fermi distribution, we obtain (i) $F_\alpha(\omega_{00} < \omega < \omega_{11}) = \frac{1}{2}(1 + \gamma)\Phi_\alpha$ and $F_\alpha(\omega > \omega_{11}) = \frac{1}{2} + \gamma\Phi_\alpha$, where $\gamma L = 1/\gamma R = \gamma$; and (ii) $F_c = 1 + \frac{1}{2} \gamma^2(1 + \gamma) = 1 + \gamma + \frac{1}{2} \gamma^2$ and $\frac{1}{2}(3 + 2\gamma + \gamma^2)$, for $\omega_0 < \omega < \omega_{L1}$, $\omega_{L1} < \omega < \omega_{R1}$, and $\omega > \omega_{R1}$, respectively. These results are plotted in the insets of Fig. 2. The noise can be either super-Poissonian or sub-Poissonian, depending on the ratio $\gamma \equiv \Gamma_L/\Gamma_R$. Strikingly, it differs from the Markovian results: $F_c^{\text{M}} = 1$ (i.e., Poissonian noise) and $F_c^{\text{M}} = 2$.

Remarkable non-Markovian effect is also manifested at high frequency limit. For the finite bandwidth model, the non-Markovian theory predicts a correct vanishing noise, while the Markovian treatment cannot give this prediction. For infinite wide-band model, the contradiction is not so sharp, however, interesting difference also emerges. Consider, for instance, $F_L$ and $F_R$. The Markovian treatment gives $F_L = F_R = 1$, i.e., the Poissonian behavior. However, the non-Markovian treatment results in $F_L = \frac{1}{2} + \gamma$ and $F_R = \frac{1}{2} + \gamma\gamma^{-1}$ for strong Coulomb interaction $\epsilon_0 + U > \mu_L > \epsilon_0 > \mu_R$; and $F_L(R) = \frac{1}{2}(1 + \gamma\pm 1)$ for weak Coulomb interaction $\mu_L > \epsilon_0$, $\epsilon_0 + U > \mu_R$. They can be either super-Poissonian or sub-Poissonian depending on the ratio $\gamma$. Finally, the results also show that the noise is enhanced by Coulomb interaction strength.

**Interacting Double Dots.** Now consider the transport through two coupled quantum dots, described by

$$H = \sum_{\alpha=L,R} \epsilon_\alpha d_{\alpha}^d d_{\alpha} + U N_L N_R + \Omega (d_{L}^d d_{R} + d_{R}^d d_{L}).$$

We as-

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**FIG. 1:** Charge correlation function $S_N(t)$, under non-Markovian (solid line) and Markovian (dashed line) treatments. Parameters (in unit of meV) are $T = 0.05$, $eV = 5$, $\epsilon_0 = 1$, $\Gamma = 1$, and $W = 15$.

**FIG. 2:** Noise spectrum of transport current through an interacting quantum dot, decomposed to (a) junction currents and (b) charge-number fluctuation components, respectively. Parameters (in unit of meV) are $U = 10$, $T = 0.05$, $eV = 6$, $\epsilon_0 = 1$, $\Gamma = 0.4$, with finite bandwidth of $W = 15$. The insets are the corresponding results of WBL.
sume that the intra-dot Coulomb interaction is infinitely large; thus only the inter-dot interaction term appears explicitly. Moreover, for simplicity, we consider spinless electron here. In the strong inter-dot Coulomb blockade regime, the charge configurations relevant to steady-state transport are the empty $|0\rangle$, left-dot occupied $|L\rangle$, and right-dot occupied $|R\rangle$ states. Similar to the single-dot case, the non-Markovian memory effect will also incorporate the double-dot occupied $|d\rangle$ state into dynamics. The characteristic frequencies in noise spectrum are now determined by the respective positions of the eigenenergies $E_e$ and $E_g$ of the coupled dots from the Fermi levels, i.e., $|E_{e,g} - \mu_0\rangle$ and $|E_{e,g} + U - \mu_0\rangle$. In addition, another characteristic frequency, the Rabi frequency $\Delta_{eg} = E_e - E_g = \sqrt{\epsilon^2 + 4\Omega^2}$, with $\epsilon = \epsilon_L - \epsilon_R$, will strongly affect the noise spectrum, particularly for a large $\Omega$. The resultant noise spectrum is depicted in Fig. 3.

For the double dots, interesting non-Markovian features also appear at high frequency limit. For the WBL, the Fano factor reads

$$F_\alpha(\omega) = \frac{\Gamma_\alpha - 4(2 + \gamma^{-1})\Omega^2 + 4\varepsilon^2_{ren}}{(8\Gamma_R\Omega^2)}$$  (9)

where $\varepsilon_{ren} = \varepsilon + \Delta E_L - \Delta E_R$, with $\Delta E_\alpha = \frac{1}{2}[(\chi_\alpha(\epsilon_0) - \chi_\alpha(\epsilon_0 + U))]$, accounts for the renormalization effect. Strikingly, it follows Eq. (9) that the noise can be highly super-Poissonian, if the coherent coupling $\Omega$ between the dots is small. This result can be understood as follows. At high frequency limit, $S_\alpha = \Gamma_\alpha$, while the steady-state current reads $I = 4\Gamma_R\Omega^2/[(\Gamma_\alpha^2 + 4(2 + \gamma^{-1})\Omega^2 + 4\varepsilon^2_{ren})]$. In other words, the noises of the left and right junction currents do not vanish, but the steady-state current will if the coupling $\Omega \rightarrow 0$. Differing remarkably, the Markovian treatment results in $S^M_\alpha = 2I$, thus $F^M_\alpha = 1$.

In summary, upon the establishment of the "n"-resolved QME with memory, we illustrated that a non-Markovian treatment is essential to correctly account for the frequency-dependent shot noise. Markovian approach is applicable only in the classical regime. It breaks down completely, even for the wide-band leads, as long as the measurement frequency becomes comparable to or higher than the characteristic excitation frequencies. This is the quantum regime where memory effect is strikingly prominent. The resultant characteristic features in the noise spectrum provide very useful access to some important system parameters of study. Support from GRC Hong Kong and NNSF China is acknowledged.

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