It is shown that in many-electron systems quantum transfer amplitudes and thus transfer probabilities may be strongly influenced by fast fluctuating fields, in particular, caused by simultaneous electron transfers. Corresponding mutual interplay of many electron jumps, arising at the fundamental level of quantum phases, results in long-correlated (1/f type) conductance fluctuations. However, this could not be theoretically captured if neglect the real discreteness of quantum energy spectra and use the continuous spectrum approximation when building kinetic theory. Basing on first principles, the estimates of low-frequency fluctuations of tunneling conductance are presented.

I. INTRODUCTION

The electrical 1/f-noise (flicker noise) was discovered in 1925. Up to now low frequency 1/f-noise is known in different physical, chemical, biological, cosmic, geophysical and social systems [1-5], but there is no conventional explanation of this phenomenon. In electronics frequently just the 1/f conductance fluctuations limit qualities of real devices. What is important, 1/f-noise is much higher sensitive to mesoscopic structure of conductors as well as to external influences than conductivity itself, hence, it may bring a rich and delicate information on microscopic mechanisms of conductivity.

Traditionally 1/f-noise is thought be reducable to some "slow" random processes with a broad distribution of large life-times (relaxation times). The standard idea is that conductivity determined by "fast" mechanisms is modulated by "slow" varying parameters (number of charge carriers, local disorder, occupation of electron traps, Coulomb potentials, lattice magnetic moments, etc., for some hypotheses see [1-4,6-14,17-18]). But concrete origin of long life-times remains mysterious even in best experimentally investigated situations such as the electron mobility 1/f fluctuations in semiconductors [1,5,10] (besides, this theory could scarcely likely say why 1/f-noises in liquid and solid metals are comparable [5], or why hypothetical activation energies required to fit observed 1/f spectra much exceed actual energies determined from the response to charge injections [7]).

The problem looks sharply in case of "bad" (narrow-band, variable-range hopping, tunnel, magnetically disturbed) conductivity like in doped semiconductors [4,7,11], oxides [6,9], amorphous silicon [11-14], materials with colossal magnetoresistance [15-16,19,20,32], cermets [36], etc. In these systems long-range Coulombian (or magnetic) interactions are important stimulating the attempts to reduce 1/f-noise to slow random charge redistributions [4,11-14,17-18], but it is hard to explain why real 1/f spectra do not saturate below 100 Hz [7,17,18]. The interesting concept of self-organized criticality [31] also connects low-frequency noise with large spatial and temporal scales.

The alternative way to understanding 1/f-noise is developed since 1982 [5,21-25,28-30]. It attributes 1/f-noise not to "slow" processes but, in opposite, to "fast" microscopic kinetic processes responsible for resistivity. The logics is simple: if producing stochastic behaviour any dynamical system constantly forgets its own history, hence, it is indifferent to a number of kinetic events happened in the past, hence, it has neither stimulus nor possibility to follow some absolutely certain "probability of event per unit time". Therefore, a dynamical (Hamiltonian) system what produces relaxation (irreversibility) and noise makes this, generally speaking, without keeping definite (well time-averagable) "probabilities per unit time" and consequently produces fluctuations of the relaxation (dissipation) and noise rates. These fluctuations do not destroy detailed balance and do not cause any compensating reaction, hence, have no characteristic time scale (see [25] for more details and examples). In this theory the long-living statistical correlations associated with 1/f spectrum are the manifestation of pure freedom of random flow of kinetic events, not of long memory [21-25]. In this respect, as N.S.Krylov argued [26] in 1950, we all should get rid of the erroneous opinion that statistical correlations always reflect some actual causality.

The kinetic theory misses this 1/f-noise just because assumes quite certain "probabilities per unit time" although this ansatz is masked by ansatzes like "molecular chaos", "random phases", "thermodynamical limit", "continuous spectra", etc. But the correctly performed derivation of gas kinetics from Hamiltonian dynamics [28,25] demonstrates violation of "molecular chaos" and 1/f fluctuations of diffusivities and mobilities of gas particles. Thus 1/f-noise arises even in the system where Wittigly nothing like slow processes and giant life-times do exist.

While a gas is characterized by strong but short (well time-space separated) kinetic events (collisions), many systems possess weak but long-lasting (bad separated) interactions. For example, phononic systems (dielectric crystals). Nevertheless, as was shown in [30], 1/f type fluctuations of dissipation (inner friction) and of Raman light scattering observed in these systems (quartz) also can be deduced just from the basic (phonon-phonon) kinetics, if only take into account that any particular interaction interplays parametrically with other simultaneous interactions. Thus, it is wrong to suppose "elementary" kinetic events be statistically independent, in contrary to what we see in kinetics.

The purpose of the present paper is to show that similar situation may realize in quantum conducting many-electron systems. Concretely, that the very evolution of quantum amplitude and probability concerning one particular electron transfer may depend on fast fluctuating potentials and fields (electric, magnetic, etc.) induced by other simultaneously
occurring transfers, i.e. many elementary kinetic events (quantum jumps) are essentially interplaying. As far as we
know, this effect still was not under consideration. It is expected be especially strong if characteristic transfer time
(the time of evolution of quantum transfer probability right up to certainty) noticeably exceeds correlation time of
fluctuating fields. Such a condition seems natural if any conducting electron feels displacements of many other
electrons (or changes of many magnetic moments) by means of long-range interactions.

Our first principal statement is that short-correlated random fields if influencing the formation of quantum transfer
amplitudes lead to strong uncertainty of transfer probabilities, thus well defined "probabilities per unit time" do not
exist. As a consequence, long-correlated conductance fluctuations arise whose principal scale properties coincide with
thats of 1/f-noise. Hence, to explain 1/f-noise we have no need in extremely slow charge redistributions, instead
the very fast ones are sufficient. The second statement is that these effects can not be caught if neglect the actual
quantum discreteness of energy spectra.

For simplicity, here to comprehend the main idea we concentrate on the case of tunnel conduction (some aspects
of the discreteness in tunnel junctions were touched in [37] but only relating to low temperatures without accounting
for Coulombian effects). More extended argumentation including rigorous analysis of Hamiltonian models will be
published separately.

II. CHARACTERISTIC TIMES OF TUNNEL CONDUCTIVITY

Let us consider characteristic time scales relating to electron tunneling between metallic sides. Under a small voltage
$U$ applied to tunnel junction the mean charge transported during time $\Delta t$ can be phenomenologically expressed as

$$\Delta Q = e \cdot \frac{U e}{\delta E} \cdot \frac{\Delta t}{\tau_{\text{trans}}}$$

Here $\delta E$ is the mean separation of electron energy levels, $U e/\delta E$ is the number of levels effectively contributing to
electric current, and $\tau_{\text{trans}}$ is the mean transmission time required for one electron jump from a given level at one side
to wherever at opposite side. Though quantum jumps can realize in a moment, that moment is random and may lie
approximately equally anywhere in the interval $0 < \Delta t < \tau_{\text{trans}}$, with $\tau_{\text{trans}}$ being the time necessary to accumulate
the quantum transmission probability up to a value $\sim 1$. Clearly, $e/\tau_{\text{trans}}$ serves as the mean current per level. Here
from one gets the tunnel conductance as

$$G = \frac{\Delta Q}{U \Delta t} = e^2 \nu \gamma, \ \nu = 1/\delta E, \ \gamma = 1/\tau_{\text{trans}}$$

(1)

with $\nu$ being electron density of states and $\gamma$ mean jump probability per unit time.

Of course, any real junction possesses a finite capacity $C$ and thus finite characteristic time

$$\tau_{\text{rel}} \approx RC \equiv C/G$$

Its physical role can be identified (as usually in RC-circuits) as the relaxation time of junction charging and the
correlation time of thermal voltage fluctuations (if only suppose that Coulombian interaction between sides manifests
itself in stochastic form). Compare the above defined times:

$$\tau_{\text{trans}}/\tau_{\text{rel}} \approx e^2 \nu / C = E_C/\delta E$$

(2)

with $E_C = e^2/C$ being the Coulomb energy. Clearly, (2) is merely number of levels promoting charge relaxation.

Now let us demonstrate that this ratio may essentially exceed unit,

$$\tau_{\text{trans}}/\tau_{\text{rel}} \gg 1$$

(3)

even if Coulombian effects are weak in the trivial sense $E_C \ll T$. For certainty, consider a flat junction with area
$S$, side thicknesses $w$, and with $d$ and $\epsilon$ being the thickness and dielectricity of isolating barrier, respectively, and
use formula $C \approx \epsilon S/4\pi d$. The estimate for the metallic density of states is

$$\nu \approx \mu / N_e \approx Sw/(\hbar v_F a^2), \ \text{where} \ \mu \ \text{is Fermi energy,} \ N_e \ \text{number of metallic electrons,} \ v_F \ \text{Fermi velocity and} \ a \ \text{is atomic size (} Sw/N_e = a^3 \ \text{is volume per one electron), thus}$$

$$e^2 \nu / C \approx \frac{e^2}{\hbar c} \frac{c}{v_F} \frac{4\pi dw}{ea^2} \approx \frac{dw}{a^2}$$

where $c$ is the speed of light (introduced to see that $c/v_F$ overcompensates the fine structure constant) and typical
value $\epsilon \sim 20$ is substituted.
Obviously, the inequality (3) is satisfied if either \( d \) or \( w \) (moreover if both) a few times exceeds \( a \), i.e. practically always (in this sense Coulombian effects never could be neglected). Therefore the quantum probability of some particular electron jump grows inevitably under influence of relatively fast varying inter-side voltage \( u(t) \), \( u(t) \sim \sqrt{T/C} \), produced by many other jumps at the same time factually happening in both directions between other levels. Not only a moment of jump realization is random but the jump probability itself turns be random, that is different kinetic events become entangled. This may be named quantum (Coulombian) interaction of electron transfers. It can not be completely described by one-electron language, but the insert of fluctuating voltage gives the quasi one-particle approximation which, to some extent, may substitute for a real picture.

**III. TIME RATIOS CONNECTED WITH ENERGY DISCRETENESS**

The standard kinetic scheme deals with the tunnel Hamiltonian

\[
H = H_0 + H_{\text{run}} , \quad H_{\text{run}} = \sum_{kq} g_{kq} (b_q^+ a_k + a_k^+ b_q)
\]

and with three ansätze attracted to avoid formal difficulties brought in by the discreteness of electron energy levels, namely: i) energy spectrum in sides is so dense that the continuous limit is possible, \( \sum_k \to \int ... \nu(E) dE \); ii) the Fermi’s golden rule \( p_{kq}(\Delta t) \to 2\pi \Delta t (\delta_{kq}^2 / \hbar) \delta(E_{kq}) \), where \( p_{kq}(t) \) is jump probability and \( E_{kq} \) energy difference between states; iii) \( \delta_{kq}^2 \) is sufficiently smooth function of \( E_{kq} \).

Here \( \Delta t \) is a time interval necessary to adequately evaluate jump probabilities for kinetic equations. This scheme requires the restriction \( \Delta t \ll \tau_{\text{gold}} \), where \( \tau_{\text{gold}} = 2\pi \hbar / \delta E \). However, if we wanted to account for effects of the voltage fluctuations \( u(t) \) we would need at least in \( \Delta t \) comparable with \( \tau_{\text{trans}} \), i.e. in the additional condition \( \tau_{\text{trans}} / \tau_{\text{gold}} < 1 \). But in an adequate model just the opposite relation must be expected,

\[
\tau_{\text{trans}} / \tau_{\text{gold}} > 1
\]

It is easy seen if note that \( \tau_{\text{trans}} / \tau_{\text{gold}} = \delta E / \Delta E \), \( \Delta E = 2\pi \hbar / \tau_{\text{trans}} \), where \( \delta E \) is the energy uncertainty associated with instability of intra-side electron states. Hence, if the desirable condition was valid, then close states would be undistinguishable, in other words, electron spectra in sides would undergo mutual rebuilding because of too good transparency of tunnel barrier (see also [27] on the relation \( R / R_0 > 1 \)).

To analyse the ratio (4) more carefully, estimate \( \tau_{\text{trans}} \). The mean transported charge \( \Delta Q \) is expressed by

\[
\Delta Q = e(\Delta N_+ - \Delta N_-) = e \sum_{kq} [f(E_{kq}^-) - f(E_{kq}^+)] p_{kq} , \quad \Delta N_\pm = \sum_{kq} f(E_{kq}^\pm) [1 - f(E_{kq}^\mp)] p_{kq}
\]

where \( \Delta N_\pm \) is number of electrons tunneling in left (right) direction, \( E^\pm \) are energy levels in the sides, \( f(E) = 1 / [1 + \exp(E - \mu) / T] \), and

\[
p_{kq} = p_{kq}(\Delta t, U) \approx 4 \delta_{kq}^2 \sin^2(E_{kq} \Delta t / 2\hbar) / E_{kq}^2 , \quad E_{kq} = E_{kq}^+ - E_{kq}^- - eU
\]

is tunneling probability evaluated by ordinary perturbation theory. The corresponding low field conductance

\[
G = [\Delta Q / U \Delta t]_{U \to 0} = e^2 \sum_{kq} [-\partial f(E_{kq}) / \partial E_{kq}] p_{kq}(\Delta t, 0) / \Delta t
\]

turns into Eq.1 with

\[
\gamma = 1 / \tau_{\text{trans}} = \sum_q p_{kq}(\Delta t, 0) / \Delta t \approx 2\pi g^2 \nu / \hbar , \quad (E_{kq}, E_{q} \approx \mu)
\]

and \( g \) being characteristic magnitude of \( g_{kq} \). Combining this relation with Eq.1 one obtains

\[
\frac{\tau_{\text{trans}}}{\tau_{\text{gold}}} = \frac{e^2}{2\pi G} = \frac{R}{R_0} \approx \left( \frac{\delta E}{2\pi g} \right)^2 , \quad R_0 = \frac{2\pi \hbar}{e^2} \sim 20 \text{kOhm}
\]

Hence, the condition for really weak interaction is the smallness of transfer matrix elements as compared with the energy level spacing.
IV. DISCRETENESS, PHASE DECOHERENCE AND FLUCTUATIONS OF QUANTUM TRANSFER PROBABILITIES

Since the wish \( \tau_{\text{trans}} / \tau_{\text{gold}} < 1 \) was invoked by the continuous spectrum approximation, we may suspect that the discreteness must be involved in an evident form to describe the fluctuations of transfer probabilities, while the perturbation theory is still applicable. Consider the picture when quantum transfers from a fixed level "k" at one side to any level "q" at another side are influenced by a fast fluctuating field (FFF), here the voltage noise \( u(t) \), in its turn induced by electron jumps beyond our attention. Thus we use quasi one-electron picture, basing on the ansatz that in many-electron surroundings with a sufficiently rich energy spectrum \( \text{FFF} \) behaves as a random process. The modern theory of quantum chaos gives powerful support for this statement [33,34] (although taking in mind that in specific systems coherent collective charge oscillations are possible instead of stochasticity and relaxation).

In this section we may consider thermodynamical equilibrium taking \( U = 0 \). First introduce the randomly accumulating diffusive phase shift

\[
\varphi(t) = \frac{e}{\hbar} \int_0^t u(t') dt'
\]

Instead of (5), the standard Shroedinger equations for transfer amplitudes yield

\[
p_{kq} \approx |A_{kq}|^2, \quad A_{kq} \equiv \frac{g_{kq}}{\hbar} \int_0^{\Delta t} \exp(iE_{kq}t/\hbar) Z(t) dt, \quad Z(t) = \exp[i\varphi(t)]
\]

(8)

Clearly, now \( p_{kq} \) become random values governed (parametrically excited or damped) by the phase shift in its turn produced by fluctuating potential difference between "in" and "out" states. Further, introduce the phase correlation function, the phase correlation time and corresponding energetic measure of quantum coherence by

\[
K(t_1 - t_2) = \langle Z(t_1)Z^*(t_2) \rangle, \quad \tau_{\text{phase}} = \int_0^\infty |K(\tau)| d\tau, \quad \varepsilon_{\text{coh}} = 2\pi \hbar / \tau_{\text{phase}}
\]

(9)

where angle brackets denote averaging with respect to FFF.

As demonstrated below, expectedly \( \tau_{\text{phase}} \) is rather short as compared with \( \tau_{\text{trans}} \) and naturally limited by \( \tau_{\text{rel}} \). Therefore under the integral in (8) \( Z(t) \) can be treated as complex shot noise. Consequently, at \( \Delta t >> \tau_{\text{rel}} \) the transfer amplitudes \( A_{kq} \) behave approximately as complex Brownian paths, and regardless of details of \( Z(t) \)'s statistics we have reasons to write

\[
\langle p_{kq}^2 \rangle = \left\langle \left| A_{kq} \right|^4 \right\rangle \approx 2 \left\langle \left| A_{kq} \right|^2 \right\rangle^2 = 2 \langle p_{kq} \rangle^2, \quad \langle p_{kq}, p_{kq} \rangle \approx \langle p_{kq} \rangle^2
\]

(10)

Here the Malakhov’s cumulant brackets are used,

\[
\langle x, y \rangle \equiv \langle xy \rangle - \langle x \rangle \langle y \rangle
\]

This is our first finding: if considered at time intervals of order of the actual transition time the quantum amplitudes and probabilities may become 100% uncertain due to the phase decoherence (phase diffusion) forced by FFF. The second is that the mean probabilities grow linearly with time:

\[
\langle p_{kq} \rangle \approx \Delta t (g/\hbar)^2 \int K(\tau) \exp(iE_{kq}\tau/\hbar) d\tau \propto \Delta t \quad (E_{kq} = E_q^+ - E_q^-)
\]

i.e. FFF remove a need in the golden rule to have uniform growth of probabilities even at \( \Delta t >> \tau_{\text{gold}} \). Notice for the next that according to this formula now \( \varepsilon_{\text{coh}} \) (instead of \( 2\pi \hbar / \tau_{\text{gold}} \)) serves as the width of energy region available from a given level ("k" or "q").

But most interesting subject is the summary transition rate \( \gamma = p / \Delta t \), with \( p = p_k \)

\[
p_k \equiv \sum_q p_{kq} = \int_0^{\Delta t} \Gamma(t_1 - t_2)Z(t_1)Z^*(t_2) dt_1 dt_2, \quad \Gamma(\tau) = \sum_q \frac{\langle g_{kq} \rangle^2}{\hbar} \exp(i\tau E_{kq}/\hbar)
\]

(11)

Here the kernel \( \Gamma(\tau) \) represents the discreteness. Its analytical properties are of much importance for all the theory. In the continuous approximation provided ansatizes i)-iii) one would have \( \Gamma(\tau) = \delta(\tau) / \tau_{\text{trans}} \) , and \( p \) would be definitely nonrandom, \( p = \Delta t / \tau_{\text{trans}} \). In reality, because of the discreteness this kernel is quite non-local and does
not decay completely at arbitrary long time. For visuality only, let us choose equidistant spectrum at the right-hand side, \(E^+_q - E^-_k = n\delta E + \varepsilon\) (\(n\) is integer, \(\varepsilon \sim \delta E\)), then

\[
\Gamma(\tau) = \frac{1}{\tau_{\text{trans}}} \exp(i\varepsilon \tau/\hbar) \sum_n \delta(\tau - n\tau_{\text{gold}}) \tag{12}
\]

Now the third important point is easy seen: if \(\varepsilon_{\text{coh}} > \delta E\) then the mean probability \(\langle p \rangle \approx \Delta t/\tau_{\text{trans}}\) practically coincides with what is given by usual kinetics, even in spite of formal violation of the golden rule (in this sense, FFF effectively expand applicability of usual scheme).

And the main fourth point is that the phase decoherence produced by FFF if combined with the discreteness results in randomness of the summary quantum jump probabilities. Indeed, due to the possibility to consider \(A_{kq}'s\) as Brownian walks we obtain

\[
\langle p, p \rangle \approx \int \ldots \int_0^{\Delta t} \Gamma(t_1 - t_2)\Gamma(t_3 - t_4)K(t_1 - t_4)K(t_3 - t_2)dt_1...dt_4 \tag{13}
\]

Here under the same condition \(\varepsilon_{\text{coh}} > \delta E\) only the regions \(t_1 \approx t_4, t_3 \approx t_2\) are significant but many delta functions from (12) contribute. The resulting transfer probability variance is

\[
\langle p, p \rangle \approx \frac{\Delta t^2}{\tau_{\text{trans}}\tau_{\text{gold}}} \int |K(\tau)|^2 d\tau \approx \frac{\tau_{\text{phase}}}{\tau_{\text{gold}}} \langle p \rangle^2 \tag{14}
\]

(we took into account that "width" of delta functions determined by the inverse width of whole energy band is wittingly smaller than \(\tau_{\text{phase}}\)).

If inequality (4) was inverted because of too small \(\delta E\) the expression (13) would formally turn into zero, but in fact it always remains nonzero under any realistic slightly non-equidistant spectrum. Nevertheless, we may predict that violation of (4), as well as of (3), leads to supression of the probability fluctuations. Oppositely, large \(\delta E > \varepsilon_{\text{coh}}\) means most strong fluctuations with variance \(~\langle p \rangle^2~\) or greater. But thats are accompanied by decreasing correlations between transfers from different levels:

\[
\langle pk_1, pk_2 \rangle \approx \langle pk_1 \rangle \langle pk_2 \rangle \frac{1}{\tau_{\text{gold}}} \int \exp[i(E_{k_1} - E_{k_2})\tau/\hbar] |K(\tau)|^2 d\tau \tag{15}
\]

Evidently, while \(\varepsilon_{\text{coh}} > \delta E\) the currents from many levels fluctuate concordanly, at \(\varepsilon_{\text{coh}} < \delta E\) even close levels inject independently one on another. Besides, even \(\langle p \rangle\) become sensible to the energy shift \(\varepsilon\) in (12), therefore, this extreme case should be carefully analysed with accounting for realistic (uncommensurable) discrete energy spectra (and better out of the frame of the quasi one-particle picture).

To end this section, estimate the phase correlation time \(\tau_{\text{phase}}\) (what may be named also phase decoherence time). Notice that \(K(\tau)\) is nothing but characteristic function of the phase. In principle, it is determined just by the transfer statistics, in a complicated self-consistent picture. Since the latter now is out of our look, we confine ourselves by rough reasonings. For instance, at \(E_C << T\) it is likely natural to treat \(u(t)\) as the Ornstein-Uhlenbeck random process, then

\[
K(\tau) \approx \exp \left[ \frac{T}{C} \left( \frac{e}{\hbar} \frac{\tau_{\text{rel}}}{\tau} \right)^2 \right] \left[ 1 - \frac{\tau}{\tau_{\text{rel}}} - \exp(-\frac{\tau}{\tau_{\text{rel}}}) \right]
\]

The corresponding decoherence time is \(\tau_{\text{phase}} \sim (\hbar/e\sqrt{C/T}) / (e/e/\tau_{\text{rel}}) \). Perhaply, this is low bound for it (too rigid in the sense that it does not include the conductance \(G\)). At \(E_C \sim T\) (what qualifies small junctions) the charge quantization is essential and the better model for \(\varphi(t)\) is infinitely divisible random walk (on this subject see [37]) formed with rare increments by \(\Delta \varphi = (e/\hbar)(\pm e/\hbar/\tau)\), where \(\theta\) is random duration of charged stay distributed with some probability density \(W(\theta)\) and total time fraction \(\sim 1\). The suitable expression is \(W(\theta) = \tau_{\text{rel}}^{-1} \exp(-\tau/\tau_{\text{rel}})\) what corresponds to the characteristic function

\[
\Xi(\eta, \tau) \equiv \left\{ \exp[i\eta \int_0^\tau u(t)dt] \right\} \approx \exp \left\{ \left| \tau \right| / \tau_{\text{rel}} \int_0^\infty [\cos(\eta e/\tau/\theta) - 1] W(\theta) d\theta \right\}
\]

(her the Levy-Khinchin representation [37] was applied). Taking \(\eta = e/\hbar\) and assuming \(2\pi R_0 > 1\), obtain \(K(\tau) = \Xi(e/\hbar, \tau) \approx \exp(-|\tau|/\tau_{\text{rel}})\). Thus, in this oppositely extreme case \(\tau_{\text{phase}}\) may be of order of \(\tau_{\text{rel}}\) (likely representing upper bound for \(\tau_{\text{phase}}\)) and possibly larger than \(\tau_{\text{gold}}\).
V. LOW-FREQUENCY CONDUCTANCE FLUCTUATIONS

In this section we return to externally driven junction, so the total voltage will be \( U(t) = U + u(t) \) with the applied voltage \( U \) interpreted as its average value.

Of course, fluctuations of electron jump probabilities eventually result in more or less analogous conductance fluctuations. The conductance is implied as \( \Delta Q / U \Delta t \) where \( \Delta Q \) is the conditional quantum average value of transported charge taken under a fixed realization of \( u(t) \). The charge transport depends also on occupancies of energy levels. If \( U = 0 \) then \( \Delta Q \), by its definition, must turn into zero regardless of \( u(t) \). This means the existence of some natural statistical correlations between voltage noise and occupancies whose accurate description would need in a self-consistent many-electron picture. To avoid this difficulty, let us work up the storage

\[
\Delta Q = e \sum_{kq} [f(E_{k}^{-}) - f(E_{k}^{+})] |A_{kq}|^2
\]

making time integration by parts in the amplitudes as if \( u(t) \) was absent, to transform (16) into

\[
\Delta Q = e^2 U \int_0^{\Delta t} \Lambda(t_1 - t_2) \exp[-ieU(t_1 - t_2)/\hbar] Z(t_1)Z^*(t_2) dt_1 dt_2
\]

where new kernel is introduced,

\[
\Lambda(\tau) = \sum_{kq} \left( \frac{g_{kq}}{\hbar} \right)^2 \frac{f(E_{k}^{-}) - f(E_{k}^{+})}{E_{k}^{+} - E_{k}^{-}} \exp[i\tau(E_{k}^{+} - E_{k}^{-})/\hbar]
\]

This form is consistent with the detailed balance \( \Delta Q(U = 0) = 0 \) and hence may serve for estimates.

In general, if electron spectra in sides are much wider than \( T \) then at \( U < T/e \) we believe that, approximately, i) conductivity obeys Ohmic law and ii) its relative fluctuations do not depend on \( U \). Then Eqs. 17-18 can be simplified by means of linearization into

\[
\Delta Q \approx e^2 U \sum_k [-\partial f(E_k^-) / \partial E_k^-] p_k
\]

with \( p_k \) defined by (11). Naturally, the conductance fluctuations essentially depend on relation between decoherence and level spacing. At sufficiently small phase decoherence time ("large" junction, widely correlated jumps from different levels), the Eq.15 helps easy obtain

\[
\langle G, G \rangle \sim \frac{\delta E}{T} \left( \frac{\langle G \rangle}{T} \right)^2, \quad (\tau_{\text{phase}} << \tau_{\text{gold}})
\]

At large decoherence time ("small" junction, non-correlated jumps) conductance fluctuations are very sensible to concrete structure of energy spectra in sides, first of all to the degree of their relative commensurability. Omitting the details, the result is that the relative conductance variance may prove to be anywhere in the interval

\[
\frac{\tau_{\text{phase}}}{\tau_{\text{gold}}} \cdot \frac{\delta E}{T} \left( \frac{\langle G, G \rangle}{\langle G \rangle^2} \right) < 1, \quad (\tau_{\text{phase}} > \tau_{\text{gold}})
\]

i.e can achieve as great magnitude as \( \sim 1 \). In this case, accurate accounting for fluctuations of the occupancies may be especially important.

VI. COMPARISON WITH EXPERIMENT

The good experimental illustration for permissibly exciting properties of quantum transfers was brought in [36] where 1/f-noise in the cermet (granular composite) Ni-nanoparticles (25%-Al2O3)was investigated. In this system the parameters of a typical elementary tunnel junction formed by neighbouring metal particles are \( \delta E \approx 0.2 \text{meV} \), \( d \approx 2 \text{nm} \), \( C \approx 5 \cdot 10^{-6} \text{cm} \), \( E_C \sim T \) (at room temperature), and \( R \approx 30 \text{MOhm} \), which mean that \( \tau_{\text{gold}} \approx 3 \cdot 10^{-11} \text{s} \), \( \tau_{\text{rel}} \approx 1.5 \cdot 10^{-10} \text{s} \), \( \tau_{\text{trans}} / \tau_{\text{rel}} = E_C / \delta E \approx 200 \) and \( \tau_{\text{trans}} \approx 3 \cdot 10^{-8} \text{s} \). Both the inequalities (3) and (4) are well satisfied, thus giving us all grounds (see also Discussion) to suspect that 1/f-noise could be attributed to quantum Coulombian interactions, in the above mentioned sense (possibly with a contribution by electron-phonon processes).
where the correlation function $K$ decreasing factor what follows from the Eqs.17-18 is already in the frame of quasi one-electron picture, if consider full non-linear dependence (17) of $\Delta U$ charge transporting levels, $D \sim \delta E/eU$ transfer from one side embraces $\approx \delta E/eU > \delta E$ in metal granules. When applied voltage per elementary junction exceeds (20)-(21) should be multiplied by a decreasing factor (in previous section) fails, i.e. linearized expression (19) becomes invalid if applied to fluctuations, and the estimates clear confirmation of our aproach to 1/f-noise in this system. At the same time, this means that the hypothesis ii) may begin already as a weak $\Delta U$ maintained conductance fluctuations. Moreover, $\delta E$ fluctuations and described by 1/f spectrum, instead of $\delta E/eU > \delta E$ into "non-Ohmic" linear one (although mean conductance obeys Ohm law up to $\sim T/\delta E \sim 100$ times larger current).

The experimental appearance of the factor $\delta E/eU$ is evident manifestation of the role of discreteness and it itself gives clear confirmation of our aproach to 1/f-noise in this system. At the same time, this means that the hypothesis ii) (in previous section) fails, i.e. linearized expression (19) becomes invalid if applied to fluctuations, and the estimates (20)-(21) should be multiplied by a decreasing factor $D(U)$ ($D(0) = 1$, $D(U) < 1$). Indeed, as we underlined, any transfer from one side embraces $\approx \varepsilon_{coh}/\delta E$ levels at another side. If this number is smaller than the number of charge transporting levels, $eU/\delta E$, then the latters act as uncorrelated quantum channels. Hence, $\Delta Q$ variance may become linear function of $U$ being proportional to the number of channels, what corresponds in small junctions to $D \sim \delta E/eU$ at $eU > \delta E$.

Definite formal support for this euristic reasoning (literally appropriate at zero temperature only) could be found already in the frame of quasi one-electron picture, if consider full non-linear dependence (17) of $\Delta Q$ on $U$. The decreasing factor what follows from the Eqs.17-18 is

$$D(U) = S(eU)/S(0), S(E) \equiv \int K_{\Lambda}(\tau) \exp(-iE\tau/h) \left[ \int K(\theta) K(\tau-\theta)d\theta \right] d\tau$$

where the correlation function $K_{\Lambda}(\tau)$ is defined by

$$K_{\Lambda}(\tau) = \frac{1}{2\Delta t} \int_{-\Delta t}^{\Delta t} \Lambda(t + \frac{\tau}{2}) \Lambda^*(t - \frac{\tau}{2}) dt$$

In case of small junctions the averaging over an ensemble of probable electron spectra must be added what automatically takes place in cerments. In theory, we need in some statistics of energy levels (see [34] on the known variants). At present, notice only that the absence of a rigid measure of local spacing of energy levels (except the mean value $\delta E$ as related to whole spectra only) might results naturally in a weak dependence of $K_{\Lambda}(\tau)$ on $\tau$ characterized by logarithm $\ln(\tau_{gold}/\tau)$ which implies just the inverse proportionality $D(U) \sim \varepsilon_{coh}/eU$ at $eU > \varepsilon_{coh}$.

**VII. DISCUSSION AND RESUME**

The principal peculiarity of above considered fluctuations of quantum transfer probabilities and corresponding conductance fluctuations is that their relative measure seems independent on the time under observation, $\Delta t$. It looks as if conductance undergoe fluctuations with non-decaying statistical correlations. Outwardly thats resemble static fluctuations investigated in the theory of disordered conductors (so-called universal conductance fluctuations) [34]. But by essence these are different things: one starts from the decoherence just when other finishes at it.

Though our analysis was limited by times $\Delta t \sim \tau_{trans}$, there is a feeling that similar statistics expands to arbitrarily longer time scales. Formal proof that it is really true, attributing to Hamiltonian models of quantum channels (with "weak links" like tunnel barriers) subjected to FFF (Coulombian or magnetical), will be done separately. For the present, semi-formal arguments are: i) all what was obtained results from the trivial rule that in general one should manipulate with quantum amplitudes to find quantum probabilities as a final product; ii) all what was obtained results from "fast" noise and phase decoherence, but not from some causal correlations which could not be continued to longer time. We believe that a self-consistent analysis connecting FFF and electron transport statistics will result in some slow (nearly logarithmic) dependence of $\langle G, G \rangle$ on the observation time reflecting non-static character of the fluctuations and described by 1/f spectrum, instead of $\delta f$ having the same formal dimensionality (such dependence may begin already as a weak $\Delta t$ dependence of the correlator (23)).

An accurate approach to a tunnel junction would deal with non-equilibrium steady state governed by the Hamiltonian.
\[ H = H_0 + H_{\text{tun}}, \quad H_0 = H_- + H_+ + H_C - U \Delta Q \]

where \( H_C \) describes inter-side Coulombian interaction, \( \Delta Q \) is the operator of transported charge, and \( H_\pm \) describe two sides with their leads serving also as thermal baths ensuring relaxation to equilibrium occupancies. Like in practice, in theory it is rather hard to "solder leads", especially if wishing to avoid appeals to standard kinetic schemes. But, regardless of technical difficulties, we may state that none processes in leads could influence the formation of non-diagonal (inter-side) elements of the operator \( \Delta Q \) since that's determined by the tunneling itself and charging of tunnel barrier only. Therefore, the better relaxation in leads the more grounds we have to extend the Eq.16 to arbitrary time intervals (now with \(|A_{kj}|^2\) representing the number of passing electrons). The comparatively non-principal corrections to be performed are accounting for fluctuations of occupancies and including conductance fluctuations into higher-order statistics of FFF (voltage noise).

To resume, we demonstrated that if do not neglect the actual quantum discreteness when constructing kinetic models of transport processes then possible strong sensibility of quantum transfer amplitudes and probabilities to fast fluctuating fields becomes visible (in particular, created by the transfers themselves) which may result in fundamental 1/f type low-frequency fluctuations of transport rates. Hence, the now reigning quantum kinetic models ask for evident general comments. All come from the well known Pauli's kinetic master equations. In [35] Van-Hove developed its formal groundation under so called \( \lambda^2t \)-limit. But, with no doubts, this theory does not foresee anything like 1/f noise. What is the matter, the question arises naturally (some comments on this issue were suggested in [24]). We hope that our above consideration highlights the possible principal answer: in fact the Van-Hove's formalism supposes the limit \( \delta E \to 0 \) (i.e. the continuous spectrum idealization) be performed before the limit \( g \to 0 \) with \( g = \lambda \) representing (as above) the magnitude of weak interactions. Thus the influence of quantum discreteness onto statistics of quantum jumps (between eigenstates of unperturbed Hamiltonian) is lost. We think that it will be not hopeless to properly improve the present kinetics (see Introduction).

ACKNOWLEDGEMENTS

I acknowledge Dr. Yu.V.Medvedev and participants of his seminar at the Department of kinetic properties of disordered and nonlinear systems in DonPTI NAS of Ukraine for support, stimulating criticism and helpfull discussions.

REFERENCES

1. P.Dutta, P.Horn, Rev.Mod.Phys., 53(3), 497 (1981).
2. F.N.Hooge, T.G.M.Kleinenpenning, L.K.J.Vandamme, Rep.Prog.Phys., 44, 481 (1981).
3. M.B.Weissman, Rev.Mod.Phys., 60, 537 (1988).
4. M.B.Weissman, Rev.Mod.Phys., 65, 829 (1993).
5. G.N.Bochkov and Yu.E.Kuzovlev. UFN, 141, 151 (1983), English translation: Sov.Phys.-Usp., 26, 829 (1983).
6. B.Raquet, J.M.D.Coey, S.Wirth and S. Von Molnar, Phys.Rev., B59, 12435 (1999).
7. J.G.Massey and M.Lee, Phys.Rev.Lett., 79, 3986 (1997).
8. M.J.C. van den Homberg, A.H.Verbruggen, P.F.A.Alkemade, S.Radelaar, E.Ochs, K.Armbruster-Dagge, A.Seeger and H.Stoll, Phys.Rev., B 57, 53 (1998).
9. A.Ghosh, A.K.Raychaudhuri, R.Streekala, M.Rajeswari and T.Venkatesan, Phys.Rev., B 58, R14666 (1998).
10. X.Y.Chen, P.M.Koenrad, F.N.Hooge, J.H.Woler and V.Aninkevicius, Phys.Rev., B 55, 5290 (1997).
11. G.M.Khera and J.Kakalios, Phys.Rev., B 56, 1918 (1997).
12. M.Gunes, R.E.Johanson and S.O.Kasap, Phys.Rev., B 60, 1477 (1999).
13. K.M.Abbkemeier, Phys.Rev., B 55, 7005 (1997).
14. G.Snyder, M.B.Weisman and H.T.Hardner, Phys.Rev., B 56, 9205 (1997).
15. A.Lisauskas, S.I.Khartsev and A.M.Grishin, Studies of 1/f-noise in \( La_{1-x}M_xMnO_3 \) (\( M=\text{Sr,Pb} \)) epitaxial thin films, in J.Low.Temp.Phys., as MOS-99 Proceedings.
16. A.Lisauskas, S.I.Khartsev, A.M.Grishin and V.Palenskis, Electrical noise in ultra thin giant magnetoresistors, in Mat.Res.Soc.Proc. Spring-99 Meeting.
17. B.I.Shklovskii and A.L.Efros. Electronic properties of doped semiconductors, Springer-Verlag. Berlin, 1984.
18. V.I.Kozub, Solid State Commun., 97, 843 (1996).
19. V.Podzorov, M.Uehara, M.E.Gershenson and S.-W.Cheong, lanl arXiv [cond-mat/9912064].
20. M.Viret, L.Ranno and J.M.D.Coey, Phys.Rev., B 55, 8067 (1997).
21. Yu.E.Kuzovlev and G.N.Bochkov. On the nature and statistics of 1/f-noise. Preprint No.157, NIRFI, Gorkii, USSR, 1982.
22. Yu.E.Kuzovlev and G.N.Bochkov.Izv.VUZov.-Radiofizika, 26, 310 (1983), transl. in Radiophysics and Quantum Electronics (RPQEC, USA), No 3 (1983).
23. G.N.Bochkov and Yu.E.Kuzovlev. Izv.VUZov.-Radiofizika, 27, 1151 (1984), transl. in Radiophysics and Quantum Electronics (RPQEC, USA), No 9 (1984).
24. G.N.Bochkov and Yu.E.Kuzovlev. On the theory of 1/f-noise. Preprint N 195, NIRFI, Gorkii, USSR, 1985.
25. Yu.E.Kuzovlev, lanl arXiv [cond-mat/9903356].
26. N.S.Krylov. Works on the foundations of statistical mechanics. Princeton U.P., Princeton, 1979 (Russian original book published in 1950).
27. Yu.A.Genenko and Yu.M.Ivanchenko, Teor.Mat.Fiz., 69, 142 (1986).
28. Yu.E.Kuzovlev, Zh. Eksp. Teor. Fiz, 94, No.12, 140 (1988), transl. in Sov.Phys.-JETP, 67 (12), 2469 (1988).
29. Yu.E.Kuzovlev, Phys.Lett., A 194, 285 (1994).
30. Yu.E.Kuzovlev, Zh. Eksp. Teor. Fiz, 111, No.6, 2086 (1997), transl. in JETP, 84(6), 1138 (1997).
31. P.Bak. Self-organized criticality: why nature is complex. Springer, N-Y, 1996.
32. J.M.D.Coey, M.Viret and S. von Molnar, Adv.Phys., 48, 167 (1999).
33. G.Casati and B.Chirikov. Fluctuations in quantum chaos. Preprint, Budker Inst. of Nuclear Physics SB RAS, 1993.
34. C.W.J.Beenakker, Rev.Mod.Phys., 69, N 3, 731 (1997).
35. L.Van Hove, Physica, 21, 517 (1955).
36. J.V.Mantese, W.I.Goldburg, D.H.Darling, H.G.Craighead, U.J.Gibson, R.A.Buhrman and W.W.Webb, Solid State Commun., 37, 353 (1981).
37. W.Feller, Introduction to probability theory and its applications, John Wiley, N-Y, 1966.
38. Yu.A.Genenko and Yu.M.Ivanchenko, Phys.Lett., 126, 201 (1987).