PHYSICAL REVIEW B VOLUME 62, NUMBER 16 15 OCTOBER 2000-II

Detection of quantum noise

U. Gavish, Y. Levinson, and Y. Imry
Department of Condensed Matter Physics, Weizmann Institute of Science, Rehovot 76100, Israel

(Received 8 August 2000)

We discuss the detection of quantum fluctuations in the light of the relationship between time-dependent correlators and measurable properties [L. Van Hove, Phys. Rev **95**, 249 (1954)]. Considering the interaction between the fluctuating electron system and a resonant circuit or a photon mode, we prove that zero-point fluctuations (ZPF) *are not observable by a passive detector*, corroborating the results of Lesovik and Loosen (Pis'ma Zh. Éksp. Teor. Fiz. [JETP Lett. **65**, 269 (1997)]). By a passive detector we mean one which is itself effectively in the ground state, and cannot transfer energy to the ZPF whose detection is attempted. We find that the ZPF can, on the other hand, *be observed from deexcitation of an active detector*. We also make the connection between these statements and the recent discussion of whether decoherence can be caused by the ZPF. The distinction is made between decoherence via making a real excitation in the environment and effects due to its polarization by *virtual* excitations.

There is much recent interest in the observability of quantum noise, and especially of the zero-point fluctuations (ZPF). While it is well known that these fluctuations do manifest themselves, e.g., in the Debye-Waller factor, the Lamb shift, and the Casimir force, other more direct aspects of these fluctuations are still being debated. Questions such as which correlation function of the ZPF can be measured[2] (see also Ref. 31) and whether the ZPF can be amplified,[3] are discussed in the literature. One is sometimes led to believe that the random fields produced by charge and current fluctuations in the ZPF,[4] can be directly observed, and can cause, for example, decoherence of conduction-electrons wave functions.[5] We are going to show that that is not the case.

In the classical case the current correlator of an electron system (which we will refer to as the ''antenna'') $C_C(t'-t)\equiv\langle j(t)j(t')\rangle$ is real and symmetric, i.e., $C_C(t)=C_C(-t)$. Its Fourier transform is also real and symmetric, i.e., $S_C(\omega)=S_C(-\omega)$. [In what follows all Fourier transforms are defined as $S(\omega)=(1/2\pi)\int_{-\infty}^{+\infty}dt C(t)\exp(i\omega t)$.] If the current fluctuations are coupled to the electromagnetic field in vacuum, the radiated power at frequency ω is easily seen to be proportional to $S_C(\omega)$.[6,7] Alternatively, one can[2] couple the antenna (inductively) to some resonant circuit (with frequency ω_0) and measure the power induced by the antenna in the circuit. If the circuit is at zero temperature the measured signal is proportional to $S_C(\omega_0)$ and this is similar to radiation by the antenna into the vacuum. If the temperature of the circuit (or the coupled electro-magnetic (EM) field) is $T_0\neq 0$, there is also a backflow of energy from the circuit (or the field) to the antenna, and the measured signal is the net energy flow Q from the antenna to the circuit (or the field).

In the quantum case one has to replace $j(t)$ by the current operator $\hat{j}(t)=\exp(i\hat{H}t)\hat{j}\exp(-i\hat{H}t)$, where $\hat{H}$ is the Hamiltonian of the antenna. The operators $\hat{j}(t)$ for different times do not commute, and because of this the quantum correlator $C_Q(t'-t)\equiv\langle\hat{j}(t)\hat{j}(t')\rangle$ is generally not real and not symmetric, instead $C_Q(t)=C_Q(-t)^*$. Its Fourier transform is also nonsymmetric, $S_Q(\omega)\neq S_Q(-\omega)$. This can be seen from the explicit expressions[1,7]

$$C_Q(t)=\sum_i P_i\langle i|\hat{j}(0)\hat{j}(t)|i\rangle, \tag{1}$$

$$S_Q(\omega)=\hbar\sum_{if} P_i|\langle f|\hat{j}|i\rangle|^2\delta(E_i-E_f-\hbar\omega), \tag{2}$$

where $|i\rangle$ are the states of the antenna with energies E_i and populations P_i. In the case when the antenna is in equilibrium at a temperature T, one finds[1]

$$S_Q(\omega)=S_Q(-\omega)e^{-\hbar\omega/k_BT}, \tag{3}$$

which means that the classical symmetry holds only for low frequencies $\hbar|\omega|\ll k_BT$. In the time domain this means that the classical symmetry becomes valid only for late times $|t|\gg\hbar/k_BT$. Because the quantum correlator[1] $C_Q(t)$ is not real and not symmetric *it is not a directly measurable quantity*. This also happens often in nonequilibrium situations.

$S_Q(\omega)$ has the following important physical significance, generalizing the Born scattering results of Ref. 1. It is proportional to the energy emission rate[8] into the vacuum (i.e., the state of the EM field where all $N_\omega=0$) for $\omega>0$ and the absorption rate for $\omega<0$ and a given photon, $N_{|\omega|}=1$.

The customary way in the quantum case is[9,2] to consider the symmetrized correlator $C_S(t'-t)\equiv(1/2)\langle\hat{j}(t)\hat{j}(t')+\hat{j}(t')\hat{j}(t)\rangle$, which is real and symmetric like the classical one. However $C_S(t)$ is *not*, as we will show (see also Ref. 11), the measured ''radiating correlator,'' since it contains the ZPF. For example, if the antenna is in equilibrium at a temperature T, it follows from the fluctuation-dissipation theorem[9] that for $\omega>0$ one has $S_S(\omega)\sim[N_T(\omega)+(1/2)]\hbar\omega$, where $S_S(\omega)$ is the Fourier transform of $C_S(t)$ and $N_T(\omega)=[\exp(\hbar\omega/k_BT)-1]^{-1}$ is the Planck function. This means that $S_S(\omega)\neq 0$ when $T=0$, i.e., when the antenna is in its ground state. Since being in the ground state the antenna cannot radiate energy, $S_S(\omega)$ cannot be considered as the correlator measured by detecting the radiation. Attention to this point was already attracted in the paper of Lesovik and Loosen.[2] Very recently, during the final preparation of this paper, the distinction between the emission and absorption parts of the noise spectrum was made in Ref. 12 as well.

To further clarify the above and generalize it to a finite N_ω, consider a quantum oscillator representing the resonant circuit (capacitance C and inductance L), which interacts with a quantum antenna, the interaction being $V=\alpha\hat{x}\hat{j}$, where $\hat{x}$ is the coordinate of the oscillator ($\dot{x}$ is the current in the circuit) and α is the coupling constant (mutual inductance). Using a quantum kinetic equation[13] one can calculate the energy flow Q from the antenna to the oscillator, the increase of the energy of the oscillator δE due to this flow, and the increase of the oscillator displacement $\delta\langle x^2\rangle$ (see Ref. 2). One finds in this way $\delta E=Q\tau$ and $\delta\langle x^2\rangle=2a^2(\delta E/\hbar\omega_0)$ with

$$Q=\alpha^2(2\pi a^2\omega_0^3/\hbar)[S_Q(\omega_0)(N+1)-NS_Q(-\omega_0)]. \quad (4)$$

Here τ is the relaxation time of the oscillator, $a^2=(\hbar c/2)(C/L)$ is the square of the ZPF amplitude of the oscillator, and N (similar to N_ω above) is the average number of quanta in the oscillator. When the oscillator is in equilibrium at a temperature T_0, $N\equiv N_{T_0}(\omega_0)$. The result for Q follows also from the above mentioned properties of $S_Q(\omega)$,[1] where the terms with the factor N are the induced probabilities. It follows from Eq. (4) and the detailed balance Eq. (3) that if the antenna and the circuit are in equilibrium with different temperatures the energy flow is always from the hotter to the colder system. The flow Q, if it is positive, is the energy dissipated in the circuit and QT_0 is the entropy generation rate.

The relation between $\delta\langle x^2\rangle$ and E can also be obtained from the quantum virial theorem.[14] Written in terms of Q, this relation confirms the conjecture of Ref. 2 and allows us to state that the measured current noise spectral density (for positive ω) is (for an analogous result for a two-level model, see Ref. 11, Eqs. 1.5.33-34):

$$S_M(\omega)=S_Q(\omega)[N_{T_0}(\omega)+1]-N_{T_0}(\omega)S_Q(-\omega), \quad (5)$$

where $N_{T_0}(\omega)$ is the Planck function with the temperature of the measuring device. Thus, $S_M(\omega)$ generalizes the above-mentioned well known $S_Q(\omega)$ to the case of finite T_0. The measured correlator in the time domain is

$$C_M(t)=2\int_0^\infty d\omega S_M(\omega)\cos\omega t. \quad (6)$$

We emphasize that the statement that $S_M(\omega)$ is the measured noise power spectrum is valid for an arbitrary state [used to perform the averaging in Eqs. (1),(2)] of the antenna. This includes nonequilibrium states in mesoscopic systems,[15] for example, current-carrying ones, where shot noise[10] is relevant, as well as states encountered in the quantum optics[16] context.

A zero temperature circuit (passive detector) measures $S_Q(\omega)$ for $\omega>0$. If the antenna is in its ground state, it follows from Eq. (2) that $S_Q(\omega)=0$ for positive ω and one can see now that the signal of a passive detector is zero, i.e., a passive detector does not respond to the ZPF of the antenna current. Only an active detector (circuit with $T_0>0$) ''responds'' to the ZPF, which means that the antenna being in the ground state absorbs quanta from the detector *deexciting* it. The last statement is clearer from the representation (see Ref. 2):

$$S_M(\omega)=S_Q(\omega)+N_{T_0}(\omega)\,\mathrm{Im}\,\chi(\omega), \quad (7)$$

where

$$\begin{aligned}\mathrm{Im}\,\chi(\omega)&=S_Q(\omega)-S_Q(-\omega)\\ &=\hbar\sum_{if}(P_i-P_f)|\langle f|\hat{j}|i\rangle|^2\delta(E_i-E_f-\hbar\omega)\\ &=(1/2\pi)\int_{-\infty}^{+\infty}dt\langle[j(0),j(t)]\rangle\exp(i\omega t) \quad (8)\end{aligned}$$

is the absorption coefficient of the antenna, related to its response

$$\chi(\omega)=\frac{i}{\pi}\int_0^\infty dte^{i\omega t}\langle[\hat{j}(0),\hat{j}(t)]\rangle. \quad (9)$$

(Here $[,]$ means a commutator.) When the population of the antenna is not inverted (i.e., $P_f<P_i$ if $E_f>E_i$) the absorption $\mathrm{Im}\,\chi(\omega)>0$ for negative ω. Note that the equilibration between the antenna and the circuit happens only due to the absorptive part of the response $\mathrm{Im}\,\chi(\omega)$.

The deexcitation of an active detector can be regarded as an indirect observation of the ZPF. Indirect observation of the ZPF is also possible in pumped systems, like a parametric amplifier.[17] It is well known that the ZPF appear in various other physical effects such as the Lamb shift, the Debye-Waller exponent, and the Casimir force. How they influence a linear amplifier is considered, for example, in Ref. 3.

We now consider electronic dephasing when $T\to 0$ in mesoscopic systems. Recently, Mohanty *et al.*[18] have published extensive experimental data indicating that contrary to general theoretical expectations,[15,19,20] the dephasing rate in films and wires does not vanish as $T\to 0$. Serious precautions[22] were taken to eliminate experimental artifacts. It was speculated that such a saturation of the dephasing rate when $T\to 0$, might follow from interactions with the zero-point motion of the environment. These speculations have received apparent support from calculations in Ref. 5. However, the latter were severely criticized in Refs. 23 and 24 and were in disagreement with experiments in Ref. 25 and with results on the zero-point motion of the ions in Ref. 26. In fact, it is clear that since dephasing must be associated with a change of state of the environment,[19] it cannot happen as $T\to 0$. In that limit neither the electron nor the environment has any energy to exchange. Below, we review the proof[21,27] of this qualitative statement, based on the detailed balance condition, Eq. (3), for the structure factor. While demonstrating unequivocally that zero-point motion does not dephase, the proof does show what *further* physical assumptions can, in fact, produce an apparently finite dephasing rate when T is very small, but not in the strict $T\rightarrow 0$ limit.

It is useful to apply a recently derived semiclassical expression[28] for the dephasing rate:

$$1/\tau_\phi \propto \int\int d\mathbf{q} d\omega |V_q|^2 S_e(q,\omega) S_{env}(-q,-\omega), \quad (10)$$

where V_q is the Fourier transform of the interaction between the electron under discussion and the particles of the environment (the other conduction electrons, in this case), $S_e(q,\omega)$ is the van Hove dynamic structure factor (Fourier transform of the density correlator) of the diffusing electron (a Lorentzian with width Dq^2, in the classical limit), and $S_{env}(-q,-\omega)$ is the same for the environment. What Eq. (10) means is that the dephasing rate is given by a summation over all the $(\vec{q},\omega)$ channels of exchange between the electron and the environment. One[21,27] now uses the above expression Eq. (10) and applies the very general detailed-balance equilibrium relationship, as in Eq. (3)

$$S(q,\omega) = S(-q,-\omega) e^{-\hbar\omega/k_B T}, \quad (11)$$

to either $S_e(q,\omega)$ or $S_{env}(-q,-\omega)$. It is immediately seen that the integrand of Eq. (10) is a product of two factors one of which vanishes for $\omega > 0$ and the other for $\omega < 0$, as $T \rightarrow 0$. Thus the integral and the dephasing rate vanish, in general, when $T\rightarrow 0$. The dephasing of a conduction electron at the Fermi level constitutes a very sensitive passive detector for the fluctuations of the environment. However, it detects nothing when the environment is in its ground state as well. One may note, however, that if $S_{env}(-q,-\omega)$ has an approximate δ-function peak at small ω, due to an abundance of low-energy excitations, relatively strong dephasing will follow at the correspondingly low temperatures. A particular model for those, invoking defect dynamics, was suggested in Ref. 27.

Thus, the ''standard model'' of disordered metals (in which the defects are strictly frozen) gives, as expected, an infinite τ_ϕ at $T=0$. On the other hand, there may be other physical ingredients that can make τ_ϕ relatively short at very low temperatures (but *diverge at the $T\rightarrow 0$ limit*), *without* contradicting any basic law of physics. We reemphasize that this does *not* imply dephasing by zero-point fluctuations, which has been repeatedly, and wrongly, claimed in the literature. The failure of the semiclassical approximation used in these considerations was clarified in Ref. 21.

The reduction of the $T=0$ persistent current in a mesoscopic ring with a resonant transmitter, reported in Ref. 4, is a valid result for the model considered. It is due, however, to the reduction of the tunneling amplitude between the resonant level and the ring,[29] by a reversible polarization of the environment. Modeling the latter, for example, by a set of independent harmonic oscillators linearly coupled to the electron coordinate, it is found that the tunneling amplitude between the resonant level and the ring is reduced. The oscillators are polarized differently when the electron is in the wire and in the resonant level. The tunneling amplitude between the latter two states is reduced by the product of the overlaps between the two shifted states of each oscillator.[30] Under some conditions, this total overlap can even vanish.[29] However, only *virtual* excitations of the environment oscillators are produced, but no real excitation, needed[19] to cause ''decoherence.'' Such a reversible polarization of the environment by two partial waves is different from a real decoherence. It occurs only as long as the environment interacts with the electron. The polarization disappears if the coupling of the electron with the environment is switched off. This happens, for example, when the electron leaves the medium and is far enough from it. An electron wave whose energy is just at the Fermi level will emerge after spending an arbitrarily long time in an environment which is in a nondegenerate ground state, without creating any real excitation in it. This wave will therefore retain full coherence with another partial wave which has not interacted at all with the environment. This is very different from the case where a real excitation of the environment is produced by one of the partial waves. That will imply that the two partial waves will leave the environment in two orthogonal states. This orthogonality lives forever after the coupling has been switched off[19] and it causes a real dephasing of the electronic interference. Whether the persistent current reduction due to the above polarization is a true ''decoherence'' effect, was left without a clear answer in Ref. 4. We would like to make here the unambiguous statement that this reduction has nothing to do with ''$T=0$ decoherence.''

There is no doubt that the ZPF are measurable, using the right experimental arrangement. However, one of the main points of this paper is that one should *not* think naively that the EM fields produced by the fluctuating charges and currents in the ZPF (Refs. 4 and 5) are directly measurable by a passive detector. In quantum physics, one does not directly measure the time dependence of operators, but rather expectation values. If the Hamiltonian is time independent then the expectation value of the current in any stationary state, including the ground state, is time independent. The current dynamic correlators [as in Eq. (1)] are generally nonzero and time dependent also in the ground state. We demonstrated here, however, that they are not measurable through signals that they send to passive detectors, but can very well be measured for example via the nonzero absorptive part of their Fourier transform $S(\omega)$. The sign of the variable ω has an important physical significance (absorption vs emission) and the symmetrized $S(\omega)$ is, in general, not the relevant quantity for typical experimental setups in the quantum domain. Care is needed in using too vividly the picture[4] of flows of matter and energy between subsystems when the total closed system is in its nondegenerate ground state.

This research was supported by grants from the German-Israeli Foundation (GIF), the Israel Science Foundation, and the Israeli Ministry of Science and the French Ministry of Research and Technology. Instructive discussions with D. Cohen, S. Haroche, J-M Raimond, D.E. Khmelnitskii, the late R. Landauer, Z. Ovadyahu, A. Schwimmer, B.Z. Spivak, and A. Stern are gratefully acknowledged. Some of this work was done while U.G. and Y.I. were visiting the Ecole Normale Superieure in Paris; they thank S. Haroche for his hospitality and the Chaires Internationales Blaise Pascal for support.

[1] L. Van Hove, Phys. Rev. **95**, 249 (1954).
[2] G.B. Lesovik and R. Loosen, Pis'ma Zh. Éksp. Teor. Fiz. **65**, 280 (1997) [JETP Lett. **65**, 295 (1997)].
[3] J.R. Tucker and M.J. Feldman, Rev. Mod. Phys. **57**, 1055 (1985).
[4] P. Cedraschi, V.V. Ponomarenko, and M. Büttiker, Phys. Rev. Lett. **84**, 346 (2000).
[5] D.S. Golubev and A.D. Zaikin, Phys. Rev. Lett. **81**, 1074 (1998); Phys. Rev. B **59**, 9195 (1999); Phys. Rev. Lett. **82**, 3191 (1999).
[6] S.M. Rytov, Yu.A. Kravtsov, and V.I. Tatarski, *Principles of Statistical Radiophysics* (Springer, Berlin, 1987-1989), Vol. 3.
[7] One should take the $\vec{q}$ Fourier component of the current density, $\vec{q}$ being the wave vector of the emitted light. Taking $\vec{q}=0$ correspond to the dipole approximation. The component of $\vec{j}$ along the photon's polarization is understood (Ref. 8).
[8] G. Baym, *Lectures on Quantum Mechanics* (Addison-Wesley, New York, 1997), pp. 271-276.
[9] L.D. Landau and E.M. Lifshitz, *Statistical Physics* (Pergamon Press, Oxford, 1980), Pt. 1, Sects. 118, 124, and 125.
[10] V.A. Khlus, Zh. Éksp. Teor. Fiz. **93**, 2179 (1987) [JETP **66**, 1243 (1987)]; G.B. Lesovik, Pis'ma Zh. Éksp. Teor. Fiz. **49**, 513 (1989) [JETP Lett. **49**, 592 (1989)]; Th. Martin and R. Landauer, Phys. Rev. B **45**, 1742 (1992); M. Büttiker, *ibid.* **46**, 12 485 (1992).
[11] C.W. Gardiner, *Quantum Noise* (Springer Verlag, Berlin, 1991), Chap. 1.
[12] R. Aguado and L.P. Kouwenhoven, Phys. Rev. Lett. **84**, 1986 (2000).
[13] Y. Levinson, Zh. Éksp. Teor. Fiz. **57**, 660 (1970) [Sov. Phys. JETP **30**, 362 (1969)].
[14] According to that theorem, $\delta\langle x^2\rangle$ is given by $\frac{1}{2}\delta E$ divided by the force constant of the oscillator.
[15] Y. Imry, *Introduction to Mesoscopic Physics* (Oxford Unversity Press, 1997).
[16] See, for example, S. Haroche, Phys. Today **51** (7), 36 (1998).
[17] I.R. Senitzky, Phys. Rev. A **48**, 4629 (1993).
[18] P. Mohanty, E.M. Jariwala, and R.A. Webb, Phys. Rev. Lett. **78**, 3366 (1997).
[19] A. Stern, Y. Aharonov, and Y. Imry, Phys. Rev. A **41**, 3436 (1990); in *Quantum Coherence in Mesoscopic Systems*, edited by G. Kramer, Vol. 254 of *NATO Advanced Studies Institute Series B: Physics* (Plenum, New York, 1991), p. 99.
[20] B.L. Altshuler, A.G. Aronov, and D.E. Khmelnitskii, J. Phys. C **15**, 7367 (1982).
[21] D. Cohen and Y. Imry, Phys. Rev. B **59**, 11 143 (1999).
[22] R.A. Webb, P. Mohanty, and E.M. Jariwala, in *Quantum Coherence and Decoherence*, Proceedings of ISQM, Tokyo 1998, edited by K. Fujikawa and Y.A. Ono (North-Holland, Amsterdam, 2000).
[23] I.L. Aleiner, B.L. Altshuler, and M.E. Gershenson, Waves Random Media **9**, 201 (1999).; Phys. Rev. Lett. **82**, 3190 (1999).
[24] B.L. Altshuler, M.E. Gershenson, and I.L. Aleiner, Physica E **3**, 58 (1998).
[25] Yu.B. Khavin, M.E. Gershenson, and A.L. Bogdanov, Phys. Rev. Lett. **81**, 1066 (1998); Phys. Rev. B **58**, 8009 (1998).
[26] Interestingly, this issue appears to have been settled already in 1988. See, for example, J. Rammer, A.L. Shelankov, and A. Schmid, Phys. Rev. Lett. **60**, 1985 (1988).
[27] Y. Imry, H. Fukuyama, and P. Schwab, Europhys. Lett. **47**, 608 (1999).
[28] Y. Imry, in *Quantum Coherence and Decoherence* (Ref. 22).
[29] A.J. Leggett *et al.*, Rev. Mod. Phys. **59**, 1 (1987); A. Stern, Ph.D. Thesis, Tel-Aviv University (1990).
[30] The overlap reduction of the tunneling amplitude by coupling to the environment is largest when the latter is modeled by a single set of harmonic oscillators (Refs. 29 and 4). In a more realistic model (Ref. 21) where the harmonic oscillators bath is different for different parts of the sample, that overlap will be much closer to unity.
[31] G. B. Lesovik, Phys. Usp. **41**(2), 145 (1998).