Indications of Universal Excess Fluctuations in Nonequilibrium Systems

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The fluctuation in electric current in nonequilibrium steady states is investigated by molecular dynamics simulation of macroscopically uniform conductors. At low frequencies, appreciable decomposition of the spectral intensity of current into thermal and excess fluctuations provides a simple picture of excess fluctuations behaving as shot noise. This indicates that the fluctuation–dissipation relation may be violated in a universal manner by the appearance of shot noise for a wide range of systems with particle or momentum transport.

KEYWORDS: nonequilibrium steady state, fluctuation–dissipation relation, shot noise, electric conduction

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In equilibrium states, the fluctuation of an observable is related universally to a linear response function by the fluctuation–dissipation relation (FDR). In nonequilibrium steady states (NESSs), the FDR is often violated, and ‘excess fluctuation’ (XF) appears. XF plays crucial roles in many fields of physics, including single electron tunneling, the squeezing of photons, the measurement of fractional charge, and the determination of the fundamental limits of quantum interference devices. However, unlike equilibrium fluctuation, it is not yet well understood whether universal properties exist in XF.

Experimentally, the FDR violation is hardly observable in heat conduction because a convection current or a phase transition is induced for large temperature difference (which drives heat conduction) before XF becomes detectable. In contrast, the violation has been widely observed in systems with particle (or momentum) transport, such as electric conductors and photoemitting devices. We therefore consider such systems.

Among such systems are simple systems, including mesoscopic conductors, which are applicable to macroscopic conductors far from equilibrium, because of the strong decoherence. We therefore use molecular dynamics (MD) simulation on a model system, the size of which is 1\(x\)10\(^3\)\(x\)10\(^3\). In the \(x\)-direction, we impose the periodic boundary condition. The boundaries in the \(y\)-direction, we constant out of the system thereby keeping the system in a NESS. Impurities are immobile and play a role of providing random potential.

We assume short-range interactions among all particles. Since interaction potential is well characterized by its scattering cross section, detailed forms of the potential are
expected to be irrelevant when studying general nonequilibrium properties. Therefore, we here take a simple form, \(k_0(\max[0,d_j])^{5/2}\). Here, \(k_0\) is a constant and \(d_j\) is the overlap of the potential ranges. \(R_j\) is the radius of the potential range (\(R_e, R_p\), and \(R_l\) for an electron, phonon, and impurity, respectively), and \(r_j\) is the position of the \(j\)-th particle. We can change the strength of scattering by varying \(R_j\) (and particle number density).

This model corresponds to a perfect sample because the total number of carriers does not change and because impurities do not move. This system is macroscopically uniform although the translational invariance is broken by impurities and the thermal walls for phonons. Furthermore, the model and results are also applicable to systems that have a mass flow of neutral particles.22

We use units in which \(m_e, R_e, e,\) the Boltzmann constant, and a reference energy are unity. Regarding the other parameters, the main result, eq. (5), is insensitive to their values, as will be shown later. We here fix \(R_p = 1, T_0 = 1,\) and \(k_0 = 4000;\) the other parameters are varied to illustrate the possible universality of the result.

To investigate nonequilibrium states of this model, we perform MD simulation using Gear’s fifth-order predictor–corrector method.21 The time-step width is set to \(10^{-3}\). The initial position of each particle is randomly arranged so as not to be in contact with the other particles, and the initial velocities of the electrons and phonons are given by the Maxwell distribution with temperature \(T_0\). We calculate various quantities after the system reaches a NESS.

The electric field applied to the system is composed of a time-independent field \(E\), which is varied in a wide range, and a time-dependent field \(\varepsilon f(t)\), which is small. The electric field induces electric current \(I(t) = e_n L_0 V_x(t)\), where \(V_x\) is the velocity in the \(x\)-direction (i.e., along the electric field) of the center of mass of electrons. We take \(\varepsilon \neq 0\) only when we calculate the differential response function \(\mu(t - \tau; E)\) of a NESS, which is defined by

\[
\langle \delta I(t) \rangle_{E,0} = \int_{-\infty}^{\infty} dt \mu(t - \tau; E) L_0 \varepsilon f(t) + O(\epsilon^2),
\]

for \(t > \tau\) and by \(\mu(t - \tau; E) = 0\) for \(t < \tau\). Here, \(\delta I = I - \langle I \rangle_{E,0}\), and \(\langle \cdot \rangle_{E,0}\) denotes the average at the NESS in the electric field \(E + \varepsilon f(t)\). The convolution theorem yields \(\tilde{\mu}(\omega; E) = \lim_{\varepsilon \to 0} \langle \delta I(t) \rangle_{E,0} / L_0 \varepsilon f(\omega)\), where the tilde denotes the Fourier transform. Note that \(\tilde{\mu}(\omega; E)\) differs from that in an equilibrium state, \(\tilde{\mu}(\omega; 0)\).

We are mainly interested in the current fluctuation that is characterized by the spectral intensity \(S_I(\omega; E)\) of \(I(t)\) for \(\varepsilon = 0\). By the Wiener–Khinchine theorem, \(S_I(\omega; E) = \tilde{\mu}^2(\omega; E)\) is equal to the Fourier transform of the autocorrelation function \(\langle \delta I(t) \delta I(0) \rangle_{E,0}\) of current. In equilibrium states (\(E = 0\)), the FDR, \(S_I(\omega; 0) = 2T_0 \varepsilon \mu(\omega; 0)\), holds for all \(\omega\). Here, \(T_0\) is the temperature of the conductor, which is equal to \(T_0\) when \(E = 0\). We plot both sides of this relation in Fig. 1(a), and confirm that it holds in our simulation.

When larger \(E \neq 0\) is applied, \(\langle I \rangle_{E,0}\) becomes nonlinear with \(E\), as shown in the inset of Fig. 1(a). In such NESSs, we find that the FDR is violated, i.e., for any \(T\) that is independent of \(\omega\),

\[
S_I(\omega; E) \neq 2T_0 \varepsilon \mu(\omega; E) \quad \text{for some } \omega.
\]

This is demonstrated in Fig. 1(b), which shows \(S_I(\omega; E)\) and \(2T_0 \varepsilon \mu(\omega; E)\) for \(E = 0\). The inset shows (\(I\))_{E,0} versus \(E\). (b) \(S_I(\omega; E) = 2T_0 \varepsilon \mu(\omega; E)\). and \(2T_0 \varepsilon \mu(\omega; E)\) for \(E = 0.06\) (nonlinear response regime). In these simulations, \(m_e = 1\), \(R_p = 0.5\), \(L_0 = 750\), \(L_x = 125\), \(n_e = n_p = 0.016\), and \(n_i = 2/375\). The data points are the averages of five samples (impurity configurations) and the error bars are the standard deviations among them.
where \( I_0 \) is a certain crossover value of the current. In the latter region \( (|I|, |I_0|) \), \( I_{\text{exs}} \) takes the form of shot noise, where \( W \) is the Fano factor.

We have thus found that the dominant mechanism that breaks the FDR is the appearance of shot noise. To confirm that this observation holds widely for the model considered here, we also study \( I_{\text{exs}} \) in the following cases: (i) another impurity density, \( n_i = 0.016 \), (ii) other linear dimensions \( L_x \) (along \( E \)) of the system, \( L_x = 375, 300, 187.5, \) and 150, and (iii) the values of the other parameters are changed significantly (e.g., \( n_s = 0.008, m_p = 10, \) and \( R_L = 2 \)).

(iv) The thermal walls for phonons are set away from the boundaries for electrons, as shown in the top-left inset of Fig. 3(b).

Figure 3 shows the results in case (iv). In this case, the local phonon temperature \( T_p \) around the boundaries for electrons is markedly different from \( T_0 \), as shown in Fig. 3(a), where \( T_p = m_p/(v_p^E - v_p^{E,0})^2 \) (\( v_p^E \) is the local phonon velocity in the \( x \)-direction). Despite this fact, \( I_{\text{exs}} \) is well-fitted again by eq. (5), as shown in Fig. 3(b), if we define thermal fluctuation again by eq. (3) using \( T_p \).

Furthermore, we have found, although the data are not shown here, that eq. (5) also holds well in cases (i)–(iii) (except when the densities are so high that a liquid-solid phase transition takes place). Note in particular that the validity of eq. (5) in case (iii) suggests that it holds independently of details of the models, because case (iii) naturally includes, for example, the case where the \( R_L \) are specific functions of \( n_s \).

The above observations strongly indicate the robustness of eq. (5). Note that this possible universality is visible only when thermal fluctuation in nonequilibrium states is appropriately defined as eq. (3). In fact, we have found (although the data are not shown here) that the possible universality is obscured if we use \( T_x(E) \) instead of \( T_0 \) in thermal fluctuation.

Using the results in case (ii), we also investigate the \( L_x \) dependences of \( W \) and \( I_0 \). We evaluate \( W \) and \( I_0 \) by fitting the numerical results of \( S_{\text{exs}} \) for large \( |I|, |I_0| \) to the asymptotic form of eq. (5). In Fig. 4, we show \( W_L \) and \( I_0 \) versus \( L_x \). We see that \( W_L \) is almost independent of \( L_x \), i.e., \( W \sim 1/L_x \). This agrees with the partial result for macro-
scopic conductors in ref. 9 (however, see note 24), and coincides with the results for long mesoscopic conductors.\(^7,8\) Furthermore, we observe that \(I_0\) is almost independent of \(L_x\), although the error bars are somewhat large.

By combining the present results with the results on simple systems,\(^3\)–\(^7\)–\(^17\) we conjecture that the FDR is violated not in a random and system-dependent manner but in a universal manner by the appearance of shot noise, for a wide range of systems from mesoscopic to macroscopic. All details of individual systems are absorbed into \(W\), \(I_0\), and the differential response function \(\text{Re} \tilde{\mu}\) (by which thermal fluctuation is defined). The origin of current fluctuation in the present model is the chaotic behavior of interacting many particles in classical systems, while that in the simple systems\(^3\)–\(^17\) is essentially the probabilistic nature of quantum or thermal-activation processes of noninteracting particles. Despite such a marked difference, \(S_{\text{ext}}\) takes an identical form in all systems. This observation may be used as a touchstone in nonequilibrium thermodynamics or statistical mechanics beyond the linear response theory.

Note that the present results could never be obtained by a naive perturbation expansion, in powers of the driving force \(E\), about an equilibrium state. For example, the relation \(S_{\text{ext}} \propto \langle |\langle f\rangle| \rangle\) suggests that such a power series would not converge for large \(E\). Using MD simulation, we have successfully investigated such a “non-perturbative regime”. Our results may be confirmed experimentally, for example, in high-quality doped semiconductors, which may be prepared by modulation doping, at room temperature.

In conclusion, we have presented a study of excess fluctuations in a nonequilibrium system and found that the fluctuation–dissipation relation is violated in a manner that may be universal. We hope that our work will stimulate further research that will test the correctness of this conjecture for wider classes of systems.

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18) Although the fluctuation theorem [see, for example, D. J. Evans and D. J. Searles: Adv. Phys. 51 (2002) 1529] would formally hold in NESSs, it cannot be used to treat or predict XF discussed in refs. 2–17 and here.
19) In some mesoscopic conductors, transmittance \(T\) varies as a function of applied voltage, and so does \(W (= 1 - T)\). Although \(S_I\) then becomes a nonlinear function of \(|\langle f\rangle|\), it is also called shot noise because the physics involved is the same as that in the other cases where \(T\) is constant or small.
20) Although 1/f noise also appears in mesoscopic conductors, it is sufficiently small in good samples for shot noise to be observed.\(^5\)–\(^11\)
22) T. Yuge and A. Shimizu: Prog. Theor. Phys. Suppl. 178 (2009) 64. Note that Umklapp scattering is unimportant in electron transport in semiconductors because their average wave number is much smaller than the sizes of reciprocal lattice vectors.
23) In the two-sided Welch’s test, the null hypothesis that \(S(\omega \geq 0; E = 0.06) = 2 \tau, \text{Re} \tilde{\mu}(\omega \geq 0; E = 0.06)\) is rejected at the 99.9% confidence level.
24) The derivation of the scaling law in refs. 9 and 11 is rigorous in the macroscopic regime if shot noise is dominant. However, similarly to other works, refs. 9 and 11 did not show that shot noise is indeed dominant in macroscopic conductors.