Letter

Shot noise and universal Fano factor as a characterization of strongly correlated metals

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Shot noise measures out-of-equilibrium current fluctuations and is a powerful tool to probe the nature of current-carrying excitations in quantum systems. Recent shot-noise measurements in the heavy-fermion strange metal $YbRh_2Si_2$ exhibit a strong suppression of the Fano factor (F)—the ratio of the current noise to the average current in the dc limit. This system is representative of metals in which electron correlations are extremely strong. Here we carry out the first theoretical study on the shot noise of diffusive metals in the regime of strong correlations. A Boltzmann-Langevin equation formulation is constructed in a quasiparticle description in the presence of strong correlations. We find that $F = \sqrt{3}/4$ in such a correlation regime. Thus, we establish the aforementioned Fano factor as universal to Fermi liquids, and we show that the Fano factor suppression observed in experiments on $YbRh_2Si_2$ necessitates a loss of the quasiparticles. Our work opens the door to systematic theoretical studies of shot noise as a means of characterizing strongly correlated metallic phases and materials.

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Introduction. It is standard to describe metallic systems with electron correlations in terms of quasiparticles. These are elementary excitations that carry the quantum numbers of a bare electron, including charge e. In strange metals near quantum criticality [1,2], however, the current carriers are expected to lose [3,4] a well-defined quasiparticle interpretation and hence the notion of a discrete charge. This issue is especially pronounced in quantum critical heavy-fermion metals [2,5–7], for which a beyond-Landau description involving Kondo destruction [8–10] has received considerable experimental support [11–16].

How to directly prove that quasiparticles are lost in correlated metals is largely an open question. One established means of such a characterization is in terms of the ratio of the thermal and electrical conductivities (κ and σ, respectively). The quasiparticle description requires that the Lorenz number, $L \equiv \kappa / T\sigma$, obeys the Wiedemann-Franz law [17,18]. Given that charge-neutral excitations such as phonons also contribute to the thermal current, alternative means of characterizing the absence of quasiparticles are much called for. Here we address this issue in terms of shot noise [19]—the out-of-equilibrium fluctuations of the electrical current.

When electron correlations are strong, the shot noise of diffusive metals has not been theoretically considered. Here we show that, with a suitable requirement on a hieararchy of length scales, the shot-noise Fano factor (F), defined as the ratio of average current fluctuations in the static limit to the average current, has a universal value $F = \sqrt{3}/4$ in Fermi liquids in the presence of Landau parameters and quasiparticle weight reduction. By extension, the quasiparticle description fails for strongly correlated metals when their Fano factor disobeys $F = \sqrt{3}/4$.

Shot noise has proven invaluable in understanding several correlated electronic systems and materials, such as quantum Hall liquids, superconductors, and quantum dots, and has offered a window into the nature of elementary charge carriers in their respective ground states [20]. For example, measurement of the shot-noise Fano factor has been pivotal in uncovering the $\frac{1}{3}$ charge fractionalization in the Laughlin state [21,22] and charge $2e$ Cooper pairs in (fluctuating) superconductors [23,24]. Furthermore, the Fano factor has been widely used to isolate dominant scattering mechanisms in mesoscopic systems. Typically, in a diffusive Fermi gas, $F = \frac{1}{3}$ when scattering is dominated by impurities [25,26]. In a similar Fermi-gas-based approach, when the inelastic electron-electron scattering rate is included and higher than the elastic scattering rate, but in the absence of significant electron-phonon scattering, F was shown to equal $\frac{\sqrt{3}}{4}$ [27,28].

Recently, shot-noise measurements in mesoscopic wires of the heavy fermion compound $YbRh_2Si_2$ showed a large suppression of the Fano factor well below that of a Fermi-gas-based diffusive metal [29]. This was found to occur at low enough temperatures (<10 K) where equilibration of electrons via a bosonic bath such as phonons is minimal and cannot account for the reduced shot-noise signal. Since the only other source of inelastic scattering is strong electron interactions, a natural question arises: can strong correlation effects in a Fermi liquid (FL) account for the observed shot-noise suppression or is it a signature of strange metallicity near the quantum critical point? This issue is particularly important, given that, even for the FL state of heavy-fermion metals, the effect of interactions is pronounced and can induce orders-of-magnitude renormalization of both the quasiparticle weight and effective interactions.

Here we report on the first theoretical study about the shot noise of strongly correlated diffusive metals. We show that, in the presence of strong correlations, the Fano factor of a diffusive FL F is equal to $\frac{\sqrt{3}}{4}$. In particular, it is independent of quasiparticle weight or Landau FL parameters. This includes, for example, the case when the quasiparticle residue $z_{\mathbf{k}}$ approaches an infinitesimally small (but nonzero) value or when the Landau parameters are as large as in the heavy-fermion metals (typically about $\sim 10^3$). Since the existence of quasiparticles is central to the robustness of the Fano factor prediction, the experimentally observed suppression [29] strongly indicates a loss of quasiparticles.

To this end, we derive the Boltzmann-Langevin transport equation for a diffusive metal in a regime suitable for addressing strong correlations. This allows us to analyze the role of strong interactions on the Fano factor. First, we notice that charge conservation constrains the Boltzmann equation to be independent of the quasiparticle weight. This holds even in the presence of anisotropic effects when $z_{\mathbf{k}}$ is strongly momentum dependent. As a consequence, the current noise, the average current, and the Fano factor determined from the Boltzmann-Langevin equation are independent of $z_{\mathbf{k}}$. Second, we calculate the shot noise and demonstrate that the shot noise and the average current get renormalized identically by the Landau parameters. As a result, the Fano factor remains robust to the introduction of arbitrarily strong interactions within FL theory. Also inherent in this cancellation is our observation that the conductances entering both the shot noise and the average current are equal and determined by the same quasiparticle lifetime; i.e., there is a symmetry of the scattering rate between single and multiparticle operators. This symmetry exists because in FLs the interactions are instantaneous (no frequency dependence) and the electron scattering processes are Poissonian, i.e., independent of one another. In the remainder of the Letter, we construct the Boltzmann-Langevin equation in the strongly correlated regime of a FL. We closely follow the diagrammatic analysis of Betbeder-Matibet and Nozieres [30]. In particular, we analyze the role of interactions via coupling through the scalar and vector potentials. We then calculate the Fano factor through the shot noise and the average current, and we discuss its relevance to experiments in $YbRh_2Si_2$ before presenting our conclusions.

Boltzmann-Langevin equation for interacting Fermi liquids. Interacting electrons in the presence of randomly distributed impurities are governed by the Hamiltonian $H = H_0 + H_I + H_{\text{imp}}$. Here $H_0 = \sum_{\mathbf{k},\sigma} \epsilon_{\mathbf{k}} c^\dagger_{\mathbf{k}\sigma} c_{\mathbf{k}\sigma}$ and $H_I = \frac{1}{2}\sum_{\substack{\mathbf{k},\mathbf{k}',\mathbf{q}\\ \sigma\sigma'}} V(\mathbf{q}) c^\dagger_{\mathbf{k}+\mathbf{q}\sigma} c^\dagger_{\mathbf{k}'-\mathbf{q}\sigma'} c_{\mathbf{k}'\sigma'} c_{\mathbf{k}\sigma}$ describe the noninteracting electron dispersion and the Coulomb interaction, respectively. In addition, $H_{\text{imp}} = \sum_{i,\mathbf{k},\mathbf{q},\sigma} U(\mathbf{q}) e^{-i\mathbf{q}\cdot\mathbf{R}_i} c^\dagger_{\mathbf{k}+\mathbf{q}\sigma} c_{\mathbf{k}\sigma}$ marks electron scattering from dilute impurities. The bare electron operator is denoted by $c_{\mathbf{k}\sigma}$, with $\mathbf{k}$ and σ representing momentum and spin. The Coulomb and impurity matrix elements are denoted by $V(\mathbf{q})$ and $U(\mathbf{q})$, respectively. Next we consider an external field with Fourier components $\lambda(\mathbf{k})$ that

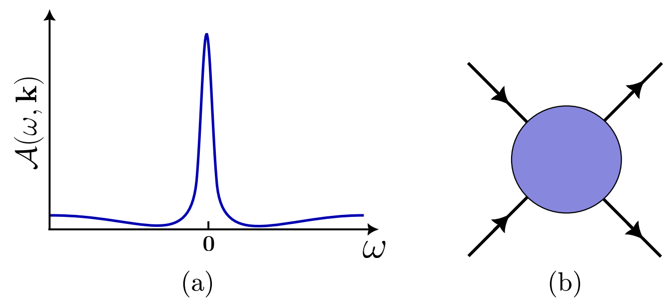

FIG. 1. (a) Plot of the spectral function $A(\omega, \mathbf{k})$ with $\mathbf{k}$ located at the Fermi momentum $\mathbf{k}_{\text{F}}$. The area of the peak is determined by $z_{\mathbf{k}}$. (b) A representation of the irreducible vertex contributing to the total vertex in the Bethe-Salpeter equation in Fig. 2(b).

couples to the interacting electrons, which contributes

$$H'(\mathbf{q}_0, \omega_0) = \sum_{\mathbf{k},\sigma} \lambda_\sigma(\mathbf{k}) c^\dagger_{\mathbf{k}+\mathbf{q}_0/2\sigma} c_{\mathbf{k}-\mathbf{q}_0/2\sigma} e^{-i\omega_0 t} + \text{H.c.} \quad (1)$$

to the total Hamiltonian, where $(\omega_0, \mathbf{q}_0) \equiv q_0$ is the energy and momentum transfer between the electrons and the external field.

To see how the external field modifies the local electronic density to linear order, we write a semiclassical total density as $n_\sigma(k, \mathbf{r}, t) = n^0_\sigma(\mathbf{k}) + \delta n_\sigma(k, \mathbf{q}_0) e^{i(\mathbf{q}_0\cdot\mathbf{r}-\omega_0 t)} +$ H.c. Here $k \equiv (\mathbf{k}, \Omega)$ is the average momentum and energy of the incoming particles. We have further defined variation of the density as the expectation value $\delta n_\sigma(\mathbf{k}, q_0) = \langle\psi(q_0)|\alpha^\dagger_{\mathbf{k}-\frac{1}{2}\mathbf{q}_0\sigma}\alpha_{\mathbf{k}+\frac{1}{2}\mathbf{q}_0\sigma}|\psi(q_0)\rangle$, where $\alpha_{\mathbf{k}\sigma}$ is the quasiparticle operator, and $|\psi(q_0)\rangle = |\psi_0\rangle + |\delta\psi(q_0)\rangle$ is the ground state $|\psi_0\rangle$ corrected by $|\delta\psi\rangle$ due to the external field. The density response $\delta n_\sigma(\mathbf{k}, q_0) = \frac{1}{2\pi i}\int d\Omega \chi_\sigma(k, q_0)$ can be evaluated using the diagram in Fig. 2(a), and it is given by

$$\chi_\sigma(k, q_0) = z^{-1}_{\mathbf{k}} \Lambda_\sigma(k, q_0) G(k^-) G(k^+). \quad (2)$$

Here $z_{\mathbf{k}}$ is the quasiparticle residue and is obtained by converting the quasiparticle operators in terms of the physical electron operators. $G(k^\pm) \propto z_{\mathbf{k}}$ are the interacting propagators

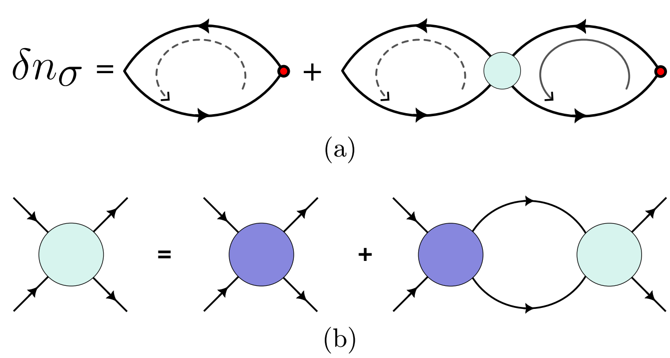

FIG. 2. Diagrammatic representations of the key processes contributing to the Boltzman equation. (a) Density response. The dashed (solid) arrow indicates the sum over the internal frequency (frequency, momenta, and spin). The smaller/red (larger/cyan) solid disk represents the external field λ that elicits the density response (total four-fermion vertex, Γ). (b) The Bethe-Salpeter equation for Γ in terms of the irreducible vertex.

for the shifted wave vector and energies $k^{\pm} \equiv (\mathbf{k} \pm \frac{1}{2}\mathbf{q}_0, \Omega \pm \frac{1}{2}\omega)$. Scatterings among electrons renormalize λ_σ [red disk in Fig. 2(a)] to Λ_σ, which is given by

$$\Lambda_\sigma(k, q_0) = \lambda_\sigma(\mathbf{k}) + \sum_{\mathbf{k}'\sigma'\Omega'} \lambda_{\sigma'}(\mathbf{k}')\Gamma_{\sigma,\sigma'}(k, k', q_0)G(k^-)G(k^+), \tag{3}$$

where $\Gamma_{\sigma,\sigma'}$ is the fully dressed four-fermion vertex. In order to include renormalizations from both electron-electron and electron-impurity scatterings, we utilize the Bethe-Salpeter equation [see Fig. 2(b)] to obtain $\Gamma_{\sigma,\sigma'}(k, k', q_0) = \hat{\Gamma}_{\sigma,\sigma'}(k, k', q_0) + \Delta\Gamma_{\sigma,\sigma'}(k, k', q_0)$, where

$$\Delta\Gamma_{\sigma,\sigma'}(k, k', q_0) = \sum_{\substack{\mathbf{k}''\sigma'' \\ \Omega''}} \hat{\Gamma}_{\sigma,\sigma''}(k, k'', q_0) \times G(k''^-)G(k''^+)\Gamma_{\sigma'',\sigma'}(k'', k', q_0) \tag{4}$$

and $\hat{\Gamma}_{\sigma,\sigma'}(k, k', q_0)$ (Fig. 1) is the sum of all irreducible vertex diagrams including Coulomb and impurity scatterings.

Eliminating the vertex part $\Lambda_\sigma(k, q_0)$ [Eq. (3)] between the density response [Fig. 2(a)] and the Bethe-Salpeter equation [Fig. 2(b)] in favor of the irreducible vertex $\hat{\Gamma}_{\sigma,\sigma'}(k, k', q_0)$, we obtain the Boltzmann transport equation for the quasiparticles,

$$dn_{\boldsymbol{p}}(\boldsymbol{x}, t)/dt + I(n_{\boldsymbol{p}}) = 0, \tag{5}$$

with $d/dt = \partial_t + \dot{\boldsymbol{x}} \cdot \partial_{\boldsymbol{x}}\dot{\boldsymbol{p}} \cdot \partial_{\boldsymbol{p}}$, where $\dot{\boldsymbol{x}} = \partial\epsilon_p/\partial\boldsymbol{p} = \boldsymbol{v}$, $\dot{\boldsymbol{p}} = -\partial\epsilon_p/\partial\boldsymbol{x} = e\boldsymbol{E} - \sum_{\boldsymbol{p}'} f_{\boldsymbol{p},\boldsymbol{p}'}\partial_{\boldsymbol{x}}\delta n_{\boldsymbol{p}'}$, and

$$\epsilon_{\boldsymbol{p}}(x) = \epsilon^0_{\boldsymbol{p}} - e\boldsymbol{E}\cdot\boldsymbol{x} + \sum_{\boldsymbol{p}'} f_{\boldsymbol{p},\boldsymbol{p}'}\delta n_{\boldsymbol{p}'} \tag{6}$$

denotes the total quasiparticle energy. $\epsilon^0_{\boldsymbol{p}}$ corresponds to the noninteracting fermion energy at global equilibrium. $\boldsymbol{v}$ denotes the quasiparticle velocity, and $\boldsymbol{E}$ refers to the static electric field acting on the quasiparticles. $I(n_{\boldsymbol{p}}) = I_{\rm im} + I_{ee}$ consists of electron-impurity and electron-electron collision integrals. Note that the vertex parts can be solved exactly in the static long-wavelength limit of $\mathbf{q}_0, \omega_0 \to 0$ with constant $|\mathbf{q}_0|/\omega_0$ [30]. Since the Ward identities constrain $\Lambda_{\mathbf{k}} \propto z_{\mathbf{k}}^{-1}$, it is clear from the expression of $\chi_\sigma(\mathbf{k}, \Omega, \mathbf{q}_0, \omega_0)$ that the density response, and as a consequence the Boltzmann equation, is independent of the quasiparticle residue.

While the Boltzmann transport equation is useful to calculate the nonequilibrium average electronic properties such as the electrical current, it is insufficient to describe their fluctuations. To do this, we introduce a Langevin source term $\delta J^{\rm ext}$ to the Boltzmann equation [31], which allows the room for quasiparticle fluctuations $n_{\boldsymbol{p}} \to n_{\boldsymbol{p}} + \delta n^{\rm fl}_{\boldsymbol{p}}$ to give the Boltzmann-Langevin equation for a strongly correlated Fermi liquid,

$$(\partial_t + \boldsymbol{v}\cdot\partial_{\boldsymbol{x}} + e\boldsymbol{E}\cdot\partial_{\boldsymbol{p}})\delta n^{\rm fl}_p(x, t) + \delta I = -e\delta\boldsymbol{E}\cdot\boldsymbol{v}\partial_{\epsilon_p}n_p + \frac{\partial n_p}{\partial\epsilon_p}\boldsymbol{v}\cdot\sum_{p'} f_{p,p'}\partial_{\boldsymbol{x}}\delta n^{\rm fl}_{p'} + \delta J^{\rm ext}(p, x, t), \tag{7}$$

where $\delta\boldsymbol{E}$ is the field fluctuation induced by quasiparticle fluctuations and is determined self-consistently through the Maxwell equation [26] $\nabla\cdot\delta\boldsymbol{E} = 4\pi\delta\rho^{\rm fl}$, where $\delta\rho^{\rm fl} = e\sum_p \delta n^{\rm fl}_p$ is the charge fluctuation. δI represents the change of collision integral due to fluctuating quasiparticle distribution $\delta n^{\rm fl}_p$.

We note that the Landau parameters appear explicitly in Eqs. (5)–(7) through the kinetic terms and the renormalizations to the quasiparticle energy ($\epsilon_{\mathbf{p}}$). Additionally, Eqs. (5) and (7) implicitly depend on the Landau parameters through the collision integrals (I), quasiparticle fluctuations ($\delta n^{\rm fl}_p$), and field fluctuations (δE). Here, $\delta\boldsymbol{E}$ is sensitive to the Landau parameters because it is generated by quasiparticle fluctuations $\delta n^{\rm fl}_p$ that depend on the interactions among the quasiparticles.

The fluctuations have zero mean value but have finite correlations. $\delta J^{\rm ext}$ denotes the extraneous flux of particles in $\boldsymbol{p}$ state and equals

$$\delta J^{\rm ext}(p, x, t) = \sum_{\boldsymbol{p}'} \delta J(p'p, x, t) - \delta J(pp', x, t). \tag{8}$$

It is the difference between flux from all $\boldsymbol{p}'$ states to the $\boldsymbol{p}$ state and flux from the $\boldsymbol{p}$ state to all $\boldsymbol{p}'$ states. We assume that the different fluxes are correlated when and only when the initial and final states are identical, thereby following a Poisson distribution of the form

$$\langle\delta J(p_1p_1', x_1, t_1)\delta J(p_2p_2', x_2, t_2)\rangle = LA\delta_{p_1,p_2}\delta_{p_1',p_2'}\delta(x_1 - x_2)\delta(t_1 - t_2)J(p_1p_1', x_1, t_1), \tag{9}$$

where $J(p_1p_1', x_1, t_1)$ is the mean flux of particles. L and A are the length and the cross section of the system. The presence of the Dirac δ functions in space-time reflects the fact that the duration and the spatial extent of collisions is much smaller than the electron-electron scattering lifetime and the scattering length respectively. We use Eq. (7) to calculate the shot noise.

Steady state. Consider a diffusive correlated metallic wire with length L and cross section A ($L \gg \sqrt{A}$). We model the correlated metallic wire as a strongly interacting diffusive Fermi liquid with an applied voltage. An applied voltage drives the metal into a nonequilibrium steady state. The static electric potential energy $e\phi(x) = -e\boldsymbol{E}\cdot\boldsymbol{x}$ serves as an s-wave perturbation to the quasiparticle distribution at every position along the system. The kinetics for the nonequilibrium system could be described by the Boltzmann transport equation (5) for the quasiparticles. The two ends of the wire stay in their own equilibrium: $n_p(\pm L/2, t) = f(\epsilon \pm eV/2, T_{\rm bath} = 0)$, which serve as the boundary conditions to the Boltzmann equation, where f is the Fermi-Dirac distribution function.

In the steady state, $\partial n_{\boldsymbol{p}}/\partial t = 0$. In order to solve Eq. (5), we expand the Boltzmann equation in terms of $\delta\bar{n}(\boldsymbol{x}, \boldsymbol{p}) = n_{\boldsymbol{p}}(\boldsymbol{x}) - n_0(\boldsymbol{x}, \epsilon)$ in the regime where $E_{\rm F} \gg 1/\tau_{\rm im} \gg 1/\tau_{ee} \gg D/L^2$. The first inequality corresponds to the condition for dilute impurities, and the following inequalities denote the strong correlation regime. The distribution $n_0(\boldsymbol{x}, \epsilon)$ stands for local equilibrium distribution and ϵ refers to the true quasiparticle energy defined in Eq. (6). $\delta\bar{n}(\boldsymbol{x}, \boldsymbol{p})$ corresponds to the departure from local equilibrium. This should be contrasted with the quasiparticle excitation $\delta n_{\boldsymbol{p}}(x) = n_{\boldsymbol{p}}(\boldsymbol{x}) - n_0(\epsilon^0_{\boldsymbol{p}})$ in Eq. (6), where $n_0(\epsilon^0_{\boldsymbol{p}})$ refers to the noninteracting fermion distribution at global equilibrium without any position dependence. The two quantities are connected by $\delta\bar{n}(\boldsymbol{x}, \boldsymbol{p}) = \delta n_{\boldsymbol{p}} - \frac{\partial n_0}{\partial\epsilon}\sum_{\boldsymbol{p}'} f_{\boldsymbol{p},\boldsymbol{p}'}\delta n_{\boldsymbol{p}'}(\boldsymbol{x})$. Only the local equilibrium distribution

$n_0(\boldsymbol{x}, \epsilon)$ as a function of real quasiparticle energy ϵ could make the collision integral $I = I_{\rm im} + I_{ee}$ vanish, and only $\delta\bar{n}(\boldsymbol{x}, \boldsymbol{p})$, rather than $\delta n_{\boldsymbol{p}}(\boldsymbol{x})$, determines the current [30,32].

When quasiparticle scattering is strong such that the electron-electron scattering length is small compared with the system size, $l_{ee} \ll L$, it drives the quasiparticles into local equilibrium, with a general position-dependent Fermi-Dirac form $n_0(\epsilon, x) = f[\epsilon, T(x)]$ such that $I(n_0) = 0$. Same as the equation in a Fermi gas [27,28], the explicit solution to the Boltzmann equation has the form

$$T(x) = \frac{\sqrt{3}}{2\pi} eV \sqrt{1 - \left(\frac{2x}{L}\right)^2}, \tag{10}$$

with $V = EL$. The details of the derivation are shown in the Supplemental Material (SM) [33]. The local equilibrium distribution function $n_0 = f[\epsilon, T(x)]$ is plotted at zero and finite environment temperature in Fig. S1 in the SM [33].

Shot noise. To calculate the shot noise, we model the mean particle flux as $J(p'p, x, t) = W(p'p)n_{p'}(x, t)[1 - n_p(x, t)]$, where $W(pp')$ is the scattering rate, and in the isotropic case $W(pp') = \delta(\epsilon_p - \epsilon_{p'})/(LAN_{\rm F}\tau_{\rm im})$. The extraneous flux of quasiparticles could be connected with the quasiparticle fluctuation and thereby connects with the current fluctuation through the Boltzmann-Langevin equation (7). We leave the details of the derivations to the SM [33] and list the final results of the current noise,

$$S = 2\int_{-\infty}^{\infty} dt \langle \delta I(t)\delta I(0)\rangle = \frac{\sqrt{3}}{2}\,\mathrm{GeV}, \tag{11}$$

where $T(x)$ is the local temperature in Eq. (10). $G = e^2 DN_{\rm F}A/L = ne^2\tau_{\rm im}A/(m^*L)$ represents the conductance of the FL.

We find that under electron correlation, the Landau parameters *only renormalize the conductance*. Therefore, the Fano factor $F = S/2eI$ is still $\sqrt{3}/4$ compared with the Fermi gas results in the hot electron regime, regardless of the interaction strength. This is despite the Landau parameter showing up in the quasiparticle energy, the distribution function, and its associated quasiparticle fluctuations [cf. Eqs. (5)–(7)], which is analogous to the robustness of the Wiedemann-Franz law against interactions in a Fermi liquid [18].

Discussion. Several remarks are in order. First, in our calculations, we assume that the corrections to the quasiparticle density and energies are linear in the external perturbation (linear response). The role of nonlinearities in the response and the Fano factor dependence of the interaction parameters will be captured by higher-order terms. Second, our calculations also assume intermediate mesoscopic length scales in accordance with experimental values. If the wire is too long, electron-phonon scattering must be considered, which acts to suppress the shot noise and the Fano factor. If the length of the wire is much smaller than the electron-electron scattering length L_{ee}, inelastic scattering does not contribute to the shot noise, and the Fano factor is automatically independent of the interaction strength and Landau parameters. Third, when retardation effects become important, as for example in the limit of lower electron densities, the Poissonian nature of the interaction fails. In this case, one must revisit the issue of the Fano factor dependence on the interaction parameters. In experiments on $YbRh_2Si_2$, the mesoscopic wire is quasi-three-dimensional where the FL theory continues to hold. However, when the transverse dimensions of the wire are sufficiently reduced, the FL theory eventually gives way to Luttinger liquid physics. In this case, the Fano factor generally depends on the dimensionless Luttinger liquid parameters since the shot noise responds to an effective charge ge, with g being the interaction parameter [19,34]. Finally, a large mean free path compared to the Fermi wavelength ($k_{\rm F}l \sim 1000$; see the SM [33]) indicates that disorder-related corrections are negligible.

In Table I of the SM [33], we have contrasted our work on the shot noise of correlated metals with earlier works [26–28,35] on simple metals. The contrast emphasizes that the combination of factors considered here identifies the shot-noise Fano factor as a universal ratio, akin to the Wiedemann-Franz law for Fermi liquids, especially in the strong-correlation regime.

To conclude, we have shown that the Fano factor of a strongly correlated Fermi liquid, defined by the ratio of its average current fluctuations to its average current, has a universal value of $F = \frac{\sqrt{3}}{4}$. More specifically, our results demonstrate that the Poissonian nature of the instantaneous Coulomb interaction and the charge conservation dictate a Fano factor that is independent of either the Landau parameters or (however small) the quasiparticle residue. This has important consequences for recent shot-noise experiments in the heavy-fermion material $YbRh_2Si_2$ where a strong suppression of the Fano factor was observed even when the effect of electron-phonon coupling is negligible [29]. The existence of a quasiparticle interpretation is a sufficient requirement for our analysis to hold. Thus, any suppression of the Fano factor below the universal ratio $F = \frac{\sqrt{3}}{4}$ strongly suggests the loss of quasiparticles. In addition, our theory sets the stage to examine shot noise in other many-body systems, such as where electrons are coupled to collective bosons as has recently been considered [36–38] or at the type of beyond-Landau (Kondo destruction) quantum criticality that has been developed for the heavy fermion strange metals like $YbRh_2Si_2$ [3,8–10]. More generally, our work points to the shot-noise Fano factor as a powerful characterization of strongly correlated metallic phases and materials.

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