Shot-Noise Signatures of Charge Fractionalization in the $\nu = 2$ Quantum Hall Edge

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We investigate the effect of interactions on shot noise in $\nu = 2$ quantum Hall edges, where a repulsive coupling between copropagating edge modes is expected to give rise to charge fractionalization. Using the method of nonequilibrium bosonization, we find that even asymptotically the edge distribution function depends in a sensitive way on the interaction strength between the edge modes. We compute shot noise and the Fano factor from the asymptotic distribution function, and from comparison with a reference model of fractionalized excitations, we find that the Fano factor can be close to the value of the fractionalized charge.

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In contrast to three spatial dimensions, where excitations of an interacting many particle system often carry the same quantum numbers as in the noninteracting case, interactions in 1D systems completely change the character of the excitation spectrum [1,2]. A prototype model for this physics is the Luttinger model, where electrons are no longer well defined quasiparticles and where electronic excitations decompose into spin and charge parts moving with different velocities [1,3].

An important example of interacting 1D systems is the edge states of incompressible quantum Hall liquids [4,5], where as a result of strong interactions, charge fractionalization can occur [6–12] and manifests itself in shot noise [13–16]. For the case of the filling fraction $\nu = 2$, there are two chiral edge modes copropagating at different velocities v_1 and v_2. In the presence of a short range interaction v_{12} between them, a pulse of charge e injected into edge mode one at a first quantum point contact (QPC1) decomposes into a charge pulse and a neutral pulse. In the charge pulse, a charge $e^* = (e/2)\sin 2\theta$ [where $\tan 2\theta = v_{12}/(v_1 - v_2)$ parametrizes the strength of interactions] travels on mode two and $e/2 + \sqrt{e^2/4 - (e^*)^2}$ on mode one [17]. In the neutral pulse, there is a charge $-e^*$ on mode two and a charge $e/2 - \sqrt{e^2/4 - (e^*)^2}$ on mode one. In this way, by exciting edge channel one via a partially transmitting QPC1, high frequency charge noise is generated on edge mode two [7]. At a QPC2, allowing for partial transmission of channel two, both charges $\pm e^*$ traveling within the charge (neutral) pulse give rise to low frequency shot noise with a Fano factor e^*/e [17].

Alternatively, one can look at this problem by using the concept of energy relaxation [18–21]. Interactions play a crucial role in the thermalization process that drives a system through states described by the Gibbs equilibrium ensemble. Generically, the dynamics is only constrained by two integrals of motion, total energy and total particle number. Integrable models like the $\nu = 2$ quantum Hall edge have infinitely many integrals of motion, and therefore it is not clear if an equilibrium state can ever be reached [22]. If the two edge modes are driven out of equilibrium with respect to one another, the system relaxes towards a nonthermal steady state [18–21,23], whose distribution function determines shot noise at a QPC2. The corresponding Fano factor depends on the strength of the interaction between the edge modes, and in general neither agrees with the fractional charge e^* introduced above nor with the result for two equilibrated edge modes. For the special case of a half open QPC1, however, the Fano factor is close to e^*/e, suggesting an interpretation in terms of charge fractionalization. Some of our results were obtained independently in Refs. [8,24]. In Ref. [8], a setup similar to that in Fig. 1 was analyzed perturbatively in the transmission probability a of QPC1, capturing only the initial stage of relaxation. A nonperturbative analysis is presented in Ref. [24], and the nonanalytic dependence of noise on a in the limit $a \ll 1$ is emphasized. If the integrability of the

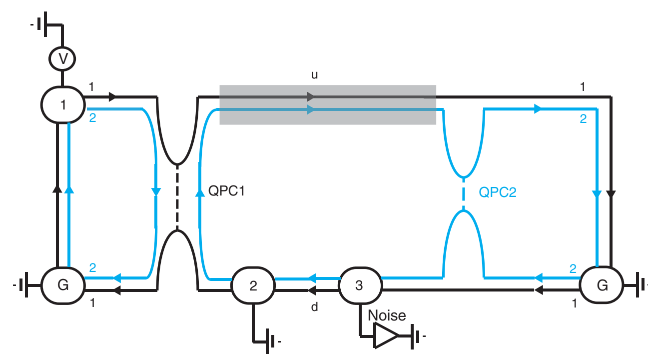

FIG. 1 (color online). Sketch of a $\nu = 2$ Hall bar with a QPC1, where inner modes (2, light blue lines) are fully reflected, while partial transmission of outer modes (1, black lines) is possible. At a QPC2, the opposite situation is realized. The shaded area is the interaction region, where partial energy relaxation takes place. The upper edge is biased with voltage V at contact one; current noise is measured at contact three.

$\nu = 2$ edge is broken, the system eventually relaxes to a thermal state [25].

We consider the setup (Fig. 1) where a Hall bar is pinched by two QPCs. The outer edge mode is labeled 1 and the inner one 2. The top and bottom edges originate at zero temperature from reservoirs at voltages $V_1 = V$ and $V_2 = 0$. At the QPC1, the outer modes are partially transmitted with probability a, while the inner ones are fully reflected; as a consequence, only the outer mode becomes noisy. After the QPC1, the two edge modes interact over some distance (shaded area in Fig. 1) before reaching the QPC2. Here, the outer modes are fully transmitted while the inner ones are partially reflected with probability p. Current noise is then measured at contact three. Using the recently developed nonequilibrium bosonization technique [26–29] within a quantum-quench model [30–32], we compute the shot noise at the QPC2, with particular emphasis on its dependence on the strength of the interaction between the edge modes.

The edges and the QPC2 are described by the following Hamiltonian ($\hbar = k_B = 1$):

$$\mathcal{H}_\eta = 2\pi \int_x [v_1 \rho_{1\eta}^2(x) + v_2 \rho_{2\eta}^2(x) + v_{12}\rho_{1\eta}(x)\rho_{2\eta}(x)],$$
$$\mathcal{H}_{\rm QPC2} = t_2 \psi_{2u}^\dagger(x)\psi_{2d}(x) + \text{H.c.} \quad (1)$$

Here, $\mathcal{H}_\eta$ describes chiral modes and $\eta = u, d$ labels the upper and lower edges. The local interaction needs to satisfy the stability criterion $v_{12}^2/4 \leq v_1 v_2$ [33]. $\mathcal{H}_{\rm QPC2}$ describes tunneling of electrons at the QPC2, with t_2 the tunneling amplitude. The fields $\rho_{i\eta}(x)$ in Eq. (1) describe density fluctuations and are related to bosonic displacement fields by $\rho_{i\eta}(x) = \partial_x \phi_{i\eta}(x)/2\pi$; here, i labels different edge modes. The bosonic fields satisfy $[\phi_{i\eta}(x), \phi_{j\xi}(y)] = \imath\pi\delta_{i\eta,j\xi}\text{sgn}(x-y)$, and the fermionic field is represented as $\psi_{i\eta}(x) = (2\pi\alpha)^{-1/2} e^{\imath\phi_{i\eta}(x)}$, with α denoting a short distance cutoff on the scale of the magnetic length. For later reference, we decompose the bosonic fields as $\phi_{i\eta}(x) = \varphi_{i\eta}(x) + \varphi_{i\eta}^\dagger(x)$ and $\varphi_{i\eta}(x) = \sum_{q>0}\sqrt{2\pi/qL}\, e^{-q\alpha/2} e^{\imath s_\eta q x} b_{i\eta}(s_\eta q)$, where $s_\eta = \pm 1$, respectively, for right (u) and left (d) movers, and $b^\dagger$ and b are canonical bosonic operators.

Following Ref. [20], we do not model the QPC1 explicitly but instead consider its effect on the downstream electron distribution of mode ($1u$) in a noninteracting setting and model the distribution as a "double step" function

$$f(\epsilon) = a\theta(-\epsilon + \mu_1) + (1-a)\theta(-\epsilon + \mu_2), \quad (2)$$

where $\mu_1 = (1-a)eV$ and $\mu_2 = -aeV$ ($eV > 0$) are chosen such that the average density in mode ($1u$) corresponds to zero bias. As a consequence of this choice, there is no density shift in mode ($2u$).

Next, we consider the effects of the intermode interaction on the distribution function (2). Instead of switching on the interaction right after the QPC1, we use the model of a quantum quench, where the interaction v_{12} is suddenly turned on for times $t > 0$ everywhere in space. Because of the chirality of the edge states, the quantum quench faithfully models the effect of a position dependent interaction; see Refs. [17,20]. The interacting Hamiltonian can be diagonalized by means of a Bogoliubov transformation M. For copropagating states ($v_1 v_2 > 0$), M can be represented by the following matrix:

$$M = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix}, \quad (3)$$

allowing us to express $\mathcal{H}$ in terms of new fields $\beta_{i,q} = \sum_j M_{ij} b_{j,q}$. The mixing angle θ expresses the strength of the interaction through the relation $\tan 2\theta = v_{12}/(v_1 - v_2)$. At this point, the new operators evolve in the Heisenberg picture as $\beta_{iq}(t_0) = e^{-\imath q\tilde{v}_i t_0}\beta_{iq}(t=0)$, with new velocities $\tilde{v}_{1(2)} = v_{1(2)}\cos^2\theta + v_{2(1)}\sin^2\theta \pm (1/2)v_{12}\sin 2\theta$.

As a final step, we undo the Bogoliubov transformation in order to express the $\beta_{iq}(t_0)$ in terms of the original basis. As a result, we obtain a relation between the bosonic operators at $t_0 > 0$ and those at $t = 0$:

$$b_{1q}(t_0) = u_q(t_0) b_{1q} + s_q(t_0) b_{2q},$$
$$b_{2q}(t_0) = s_q(t_0) b_{1q} + v_q(t_0) b_{2q}, \quad (4)$$

where $b_{iq} \equiv b_{iq}(t=0)$. Now, all the time dependence is encoded in the coefficients

$$u_q(t_0) = \cos^2\theta e^{-\imath q\tilde{v}_1 t_0} + \sin^2\theta e^{-\imath q\tilde{v}_2 t_0},$$
$$v_q(t_0) = \cos^2\theta e^{-\imath q\tilde{v}_2 t_0} + \sin^2\theta e^{-\imath q\tilde{v}_1 t_0}, \quad (5)$$
$$s_q(t_0) = (1/2)\gamma_\theta(e^{-\imath q\tilde{v}_1 t_0} - e^{-\imath q\tilde{v}_2 t_0}),$$

where $\gamma_\theta = \sin 2\theta$. To leading order in the tunneling amplitude t_2, the current noise at the QPC2 can be expressed in terms of greater (lesser) Green functions $G_{i,\eta}^{>(<)}(\epsilon)$ [34] as

$$S_{\omega\to 0} = \frac{2e^2}{h}\frac{|t_2|^2}{2\pi}\int_\epsilon G_{2u}^<(\epsilon)G_{2d}^>(\epsilon) + G_{2d}^<(\epsilon)G_{2u}^>(\epsilon), \quad (6)$$

with $G^<(\epsilon) = G^>(-\epsilon)$. Using the boson representation of electron operators, we can compute $G_{2u}^{>(<)}(\tau)$ of the fully interacting edge mode. Because of the nonequilibrium distribution of edge mode ($1u$), calculating the expectation value of a product of bosonic exponents is highly nontrivial. Here, we discuss the results for the "long time limit" of the Green function, in which the system reaches a nonequilibrium steady state:

$$G_{2u}^<(\tau) = \langle \psi_{2u}^\dagger(t_0+\tau, x_0)\psi_{2u}(t_0, x_0)\rangle$$
$$= G_0^<(\tau)\langle e^{\sum_q \lambda_{1u}^\star(q,t_0,\tau) b_{1u,q}^\dagger} e^{-\sum_q \lambda_{1u}(q,t_0,\tau) b_{1u,q}}\rangle,$$
$$G_0^<(\tau) = \frac{1}{2\pi}\frac{1}{(-\imath\tilde{v}_1\tau + \alpha)^{\sin^2\theta}}\frac{1}{(-\imath\tilde{v}_2\tau + \alpha)^{\cos^2\theta}}. \quad (7)$$

Here, $G_0^<(\tau)$ is the equilibrium Green function of edge mode 2 in the presence of interactions. All the information about nonequilibrium effects is contained in the average over bosonic coherent states in Eq. (7), where $\lambda_{1u}(q, t_0, \tau) = \imath(2\pi/qL)^{1/2} e^{\imath q x_0 - q\alpha/2}[s_q(t_0 + \tau) - s_q(t_0)]$. As emphasized in Ref. [26], nonequilibrium effects make the theory non-Gaussian, and higher order cumulants appear in the evaluation of the above expectation value. In order to compute the expectation value over the nonequilibrium state, we refermionize the bosonic operators, introducing new fermionic operators [2]:

$$\begin{aligned} b_{1u,q}^\dagger &= \imath(2\pi/qL)^{1/2}\sum_k c_{1u,k+q}^\dagger c_{1u,k}, \\ b_{1u,q} &= -\imath(2\pi/qL)^{1/2}\sum_k c_{1u,k-q}^\dagger c_{1u,k}. \end{aligned} \tag{8}$$

Since the bosonic operators describe free particle-hole excitations, the c operators are also free and therefore can be connected to the incoming states via a scattering matrix. Then, the expectation values of products of Fermi operators can be evaluated using an appropriate fermionic density matrix ρ_{1u}. The crucial step now consists in noticing that the computation of higher order cumulants is similar to the problem of full counting statistics, and using Klich's trace formula [27,35], it can be expressed in terms of a Fredholm determinant of the Toeplitz type, normalized to its zero temperature, equilibrium value

$$\bar{\Delta}_\tau(\delta) = \frac{\det[1 + (e^{-\imath\delta_\tau} - 1)f(\epsilon)]}{\det[1 + (e^{-\imath\delta_\tau} - 1)\theta(-\epsilon)]}, \tag{9}$$

where $f(\epsilon)$ is given by Eq. (2). The scattering phase $\delta_\tau = -\sum_q (2\pi/qL)^{1/2}[\lambda(q, t_0, \tau) + \lambda^*(q, t_0, \tau)] = 2\pi(e^*/e)\,\omega_\tau(t_0, x_0)$ contains information about the interedge interaction, and the window function

$$\begin{aligned} \omega_\tau(t_0, x_0) = {} & \theta[x_0 - \tilde{v}_1(t_0 + \tau)] - \theta[x_0 - \tilde{v}_1 t_0] \\ & + \theta[x_0 - \tilde{v}_2 t_0] - \theta[x_0 - \tilde{v}_2(t_0 + \tau)]. \end{aligned} \tag{10}$$

As a function of t_0, $\omega_\tau(t_0, x_0)$ represents two unit square pulses of opposite signs, with widths τ, and with a separation equal to $x_0(\tilde{v}_1^{-1} - \tilde{v}_2^{-1})$. Since $\delta_\tau = 2\pi(e^*/e)\omega_\tau(t_0, x_0)$, these pulses can be identified with charges $\pm e^*$ passing an observer at position x_0. In the case of two separated pulses, the expectation value of bosonic coherent states factorizes into a product of two single pulse determinants having the same scattering phase $\delta_{\tau,\mathrm{single}} = 2\pi(e^*/e)[\theta(-t_0) - \theta(-t_0 - \tau)]$, and we can rewrite Eq. (7) as $G_{2u}^<(\tau) = G_0^<(\tau)\bar{\Delta}_\tau^2(\delta_{\mathrm{single}})$. The determinant Eq. (9) can be evaluated numerically by treating t_0 and ϵ as conjugated variables and by carefully defining a regularization scheme [27]. Finally, the lesser Green function $G_{2d}^<(\epsilon) = \theta(-\epsilon)/\tilde{v}_1^{\sin^2\theta}\tilde{v}_2^{\cos^2\theta}$ is easily evaluated due to its equilibrium nature. Fourier transforming Eq. (7) into energy space, we can compute the distribution function at the QPC2; as a consequence of interactions, the distribution function is broadened from a single step (see Fig. 2). However, it does not have the same functional form as a Fermi distribution but rather describes a nonequilibrium steady state. The distribution obtained by only retaining the Gaussian term in the cumulant expansion clearly deviates from the full one, making evident the necessity for including higher order terms. The nonequilibrium distribution also deviates from an equilibrium Fermi distribution with effective temperature $T^* = eV\sqrt{(3/2)a(1-a)}/\pi$, obtained by assuming that the two edge modes fully equilibrate and that each of them carries half the energy flux injected into the upper edge via the QPC1. Using Eq. (6), we can finally evaluate the low frequency noise; in doing so, we relate the reflection probability p to the microscopic Hamiltonian through $p = |t_2|^2/2\pi\tilde{v}_1^{\sin^2\theta}\tilde{v}_2^{\cos^2\theta}$. In Fig. 3, we display the dependence of low frequency noise on the transmission a of the QPC1, normalizing the noise by its value at $a = 1/2$. One clearly sees that it deviates both from the standard free fermion dependence $a(1-a)$ and from the effective equilibrium result with $S_{\mathrm{eq}} = 4epI\log 2\sqrt{(3/2)a(1-a)}/\pi$. To put the strength of the noise at the QPC2 into perspective, we define a reference noise expected for noninteracting electrons tunneling through both the QPC1 and QPC2 along a single edge, obtained by using the distribution Eq. (2) in Eq. (6):

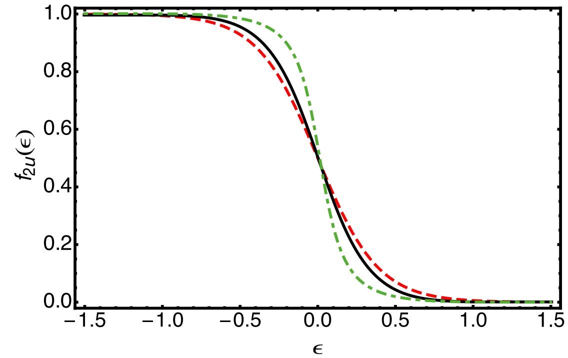

FIG. 2 (color online). Steady state distribution of an edge mode ($2u$) asymptotically away from the QPC1. The solid black line represents the nonequilibrium distribution obtained from Eqs. (7) and (9) by considering all cumulants. The dash-dotted green line represents the distribution obtained by retaining only the Gaussian term. The dashed red line represents the fully equilibrated distribution at effective temperature $T^* = eV\sqrt{(3/2)a(1-a)}/\pi$. The mixing angle is $\theta = 0.47$, and the transmission probability of the QPC1 is $a = 1/2$.

$$S_{\mathrm{ref}}(\omega \to 0) = 4epIa(1-a) \quad \text{with} \quad I = \frac{e^2}{h}V. \tag{11}$$

Since the distribution Eq. (2) gives rise to both a particle and a hole current, the prefactor in Eq. (11) is 4 instead of the usual 2 (see the inset of Fig. 4). Defining a Fano factor $F = S/S_{\mathrm{ref}}$, we can make contact with the concept of fractional charges described in the introduction. Assuming that for fractional

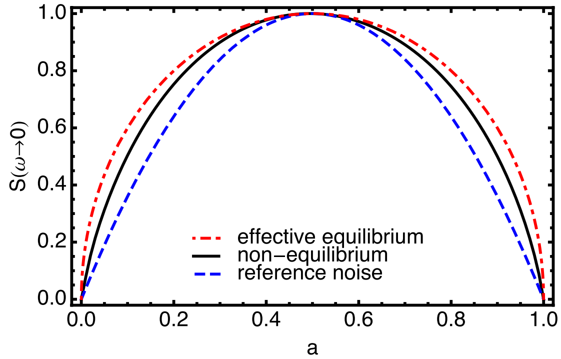

FIG. 3 (color online). Shot noise after the QPC2 as a function of a, normalized to its $a = 1/2$ value, for a mixing angle $\theta = 0.47$. The solid black line represents the full nonequilibrium result. The dashed blue line represents the reference noise of noninteracting electrons. The dash-dotted red line represents noise in a fully equilibrated thermal state.

charges the tunneling probability p in Eq. (11) is renormalized to $(e^*/e)p$ [17], the Fano factor is given by $F = \sin 2\theta/2$. In Fig. 5, the Fano factor is shown as a function of the mixing angle for the specific transmission $a = 1/2$ of the QPC1. For this value of a, there is a surprisingly good agreement between the value $e^*/e = (1/2)\sin 2\theta$ and F of the full nonequilibrium noise, suggesting that the Fano factor can indeed be interpreted as being due to the formation of fractionalized charges in the $\nu = 2$ quantum Hall edge.

We find that the zero frequency noise power depends in a singular way on a in the limit $a \ll 1$; see also Ref. [24]. To obtain the noise in this limit, the functional determinant can be approximated by its long time asymptotics (valid for $eV\tau \gg 1$) $\bar{\Delta}_\tau(\delta) \simeq \exp[-|\tau|/(2\tau_\phi)]$, where the dephasing rate $\tau_\phi^{-1} = -(eV/2\pi)\log[1 - 4a(1-a)\sin^2(\pi\gamma_\theta/2)]$. Knowledge of $\bar{\Delta}_\tau(\delta)$ for large times allows us to accurately calculate the distribution function of mode $(2u)$ for energies $\epsilon \ll eV$. However, for $a \ll 1$, the distribution function only deviates from a step function on the scale aeV, such that the long time asymptotics allows an exact calculation of the distribution function. Using Eq. (6) and taking the $a \ll 1$ limit, we find $S(\omega \to 0) \simeq 8pa\log(1/a)\sin^2(\pi\gamma_\theta/2)eV(e^2/h\pi^2)$. This nonanalyticity in a explains the divergence in S with x_0 found in Ref. [8] when calculating S perturbatively in a.

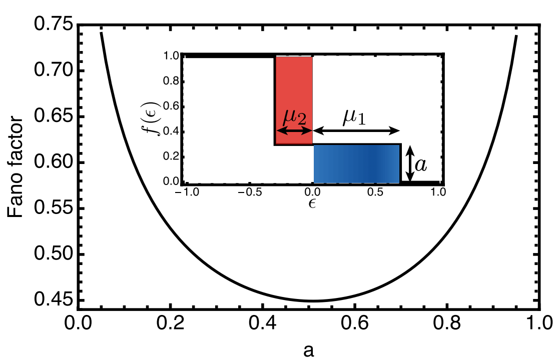

FIG. 4 (color online). Fano factor $F = S/S_{\text{ref}}$ as a function of transparency of the QPC1 for a mixing angle $\theta = 0.47$. At $a = 1/2$, the Fano factor is $F = 0.45$. Inset: Double step distribution [Eq. (2)]. The upper red area describes a hole current I_h and the lower blue area a particle current $I_p = I_h = (e^2/h)Va(1-a)$, impinging on the QPC2. From this, we obtain a reference noise $S_{\text{ref}} = 2ep(I_h + I_p)$; see also Eq. (11).

A useful way to characterize the nonlinear dependence of experimentally measured shot noise on the transmission probability a of the QPC1 is by fitting it to a function proportional to $[a(1-a)]^d$ [36]. For the reference noise of Eq. (11), d is trivially equal to unity. For "thermal" noise with effective temperature T^*, one finds $d = 0.5$. For the full nonequilibrium noise, we find that its dependence on a can be well fitted by the above power law and that d varies from $d = 0.85$ for $\theta = \pi/16$ to $d = 0.68$ for $\theta = \pi/4$; see Fig. 5. In this way, from knowledge of d, the mixing angle θ can be inferred, without using the Fano factor.

In summary, due to the joint effect of interactions and nonequilibrium, the distribution function of an originally unbiased, zero temperature mode $(2u)$ interacting with a noisy mode $(1u)$ evolves towards a nonthermal steady state that depends on the interaction strength in an characteristic way. Comparing the shot noise and Fano factor from our numerically exact calculation with a simple model of charge fractionalization, we find that the Fano factor can

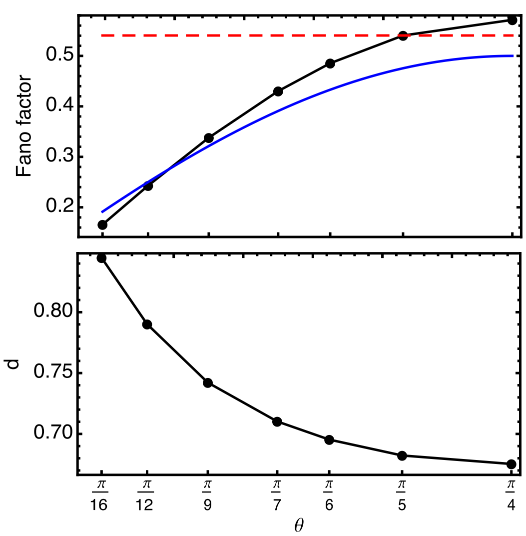

FIG. 5 (color online). Upper panel: Fano factor as a function of the mixing angle for transmission $a = 1/2$ of the QPC1. The dashed red line represents the fully equilibrated edge; F is independent of interactions. The black dots represent the full nonequilibrium situation. The solid blue line represents the reference model of a diluted system of fractional charges [$F = (1/2)\sin 2\theta$]. Lower panel: The dependence of the full nonequilibrium noise is calculated numerically and fitted by a function proportional to $[a(1-a)]^d$ for different values of θ. The black lines connecting the dots are a guide to the eye.

indeed be interpreted in terms of charge fractionalization in the $\nu = 2$ quantum Hall edge.

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