# Unusual Quasiparticles and Tunneling Conductance in Quantum Point Contacts in $\nu = 2/3$ Fractional Quantum Hall Systems

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Understanding topological matter in the fractional quantum Hall (FQH) effect requires identifying the nature of edge state quasiparticles. FQH edge state at the filling factor $\nu = 2/3$ in the spin-polarized and unpolarized phases is represented by the two modes of composite fermions (CF) with the parallel or opposite spins described by the chiral Luttinger liquids. Tunneling through a quantum point contact (QPC) between different or similar spin phases is solved exactly. With the increase of the applied voltage, the QPC conductance grows from zero and saturates at $e^2/2h$ while a weak electron tunneling between the edge modes with the same spin transforms into a backscattering carried by the charge $q = e/2$ quasiparticles. These unusual quasiparticles and conductance plateau emerge when one or two CF spin-polarized modes in the QPC tunnel into a single mode. We propose experiments on the applied voltage and temperature dependence of the QPC conductance and noise that can shed light on the nature of edge states and FQH transport.

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The quantum Hall effect (QHE) occurs in strong magnetic fields in a two-dimensional electron gas, giving the Hall conductance quantized in terms of $e^2/2\pi\hbar$, where $e$ is the electron charge and $\hbar$ is the Planck constant. When electron correlations emerge, the fractional quantum Hall effect (FQHE) exhibits a fractionally quantized conductance. FQHE excitations are fractionally charged quasiparticles that obey Abelian or non-Abelian anyon statistics [1–7]. FQHE with its edge states is the breeding ground for topological effects, quantum phase transitions, and qubits for topological quantum computing [8–18].

Conductance quantization can arise from the edge states [19] in the Landauer-Buttiker approach [20]. For the FQHE, Wen [21] found that for Laughlin [22] states with filling factors $\nu = n\Phi_0/B = 1/m$, where $m$ is odd, $n$ is the electron density, $\mathbf{B}$ is the magnetic field, and $\Phi_0 = 2\pi\hbar c/e$ is the flux quantum, the edges can be described as the chiral Luttinger liquids. While in quantum wires [23–25] conductance is independent of the Luttinger liquid parameters, the picture of chiral Luttinger liquids explains the FQHE conductance [21,26] and tunneling of $\nu = 1/m$ edge states through the quantum point contact (QPC) [27–29]. In [27] the thermodynamic Bethe ansatz and the integrability of the Luttinger model with impurity interaction were used to find the nonequilibrium temperature-dependent conductance and noise.

A more complex situation emerges for hierarchical FQHE states [5,30,31], such as $\nu = 2/3$ state. Based on an idea that it is a hole conjugate of the $\nu = 1/3$ state, [32,33] suggested that the edge at $\nu = 2/3$ consists of an outer downstream integer mode (from the underlying $\nu = 1$ state) and an inner, counterpropagating $-1/3$ edge mode. This conjugated state and two other models, two independent $\nu = 1/3$ edge channels and a pseudospin singlet state were discussed for a bilayer system in [34]. In [35,36] it was shown that disorder scattering between edge modes leads to two decoupled modes, a single downstream charge mode with conductance $2e^2/3h$ and an upstream neutral mode. Such charge-neutral (or spin) separation, as argued in [21], can also arise solely due to long-range Coulomb interactions, and stems from the nature of composite fermion (CF) [37,38] states on the edge of $\nu = 2/3$ system [39]. A reconstructed edge due to soft confining potential was suggested in [40,41]. In recent years, different approaches to edge structure of the $\nu = 2/3$ states and their transport properties continue to be actively investigated experimentally and theoretically [42–48]. With different approaches to the structure of the edge, there have been two approaches to transport. Conductance quantization implies equilibration between the chemical potentials of the reservoirs and the outgoing edges. The first approach, incoherent transport model, assumes a short coherent length and suppression of quantum interference between the channels [49–53]; the second approach is a quantum solution taking interference into account [27,54]. Experimentally, studies of tunneling through quantum point contacts (QPC) at $\nu = 2/3$ in [55–57] have shown a conductance plateau at $e^2/3h$, however, recent work [58] demonstrated its appearance within 2% of $e^2/2h$.

 

In this Letter, we find exact quantum mechanical solutions to two models of tunneling through QPC that emerge in several $\nu = 2/3$ edge configurations. One of those is associated with the transition between spin-polarized and unpolarized $\nu = 2/3$ FQH phases [39,59,60], which occurs due to the crossing of the two CF energy levels with opposite spin polarization [16,61,62]. It was shown [63–66] that the two different spin phases can be induced in the neighboring regions of the FQH liquid by electrostatic gates. The effect of boundary between the regions leads to the model with two CF edge modes of the spin-polarized phase coupled by tunneling through QPC to one CF edge mode of the same spin in the unpolarized phase. The other configuration leading to the same model is the QPC tunneling between two spin-polarized FQH regions when two CF edge modes are well separated in space due to a potential profile in one of the regions, so that only one mode can tunnel to the neighboring region. In this model for both configurations, using hybridization to account for the interference of modes in the region with the two CF states, consideration can be reduced to tunneling between two single modes. The solution shows that with an increase of the applied voltage or temperature, the QPC tunneling conductance grows from zero and saturates at $e^2/2h$ plateau while a weak electron tunneling between the edge modes with the same spin transforms into the backscattering carried by the charge $q = e/2$ quasiparticles. This charge also arises in the shot noise. This solution similarly emerges in a model describing tunneling of one mode into one mode, which describes tunneling between unpolarized region and a polarized region with two CF edge modes well separated in space due to potential profile or tunneling between two such polarized regions. In this case no hybridization and no interference occur, and the results confirm the $e^2/2h$ conductance we obtained in [66] for the contact between spin-polarized and unpolarized regions using strong coupling boundary conditions. Testing the predicted nonlinear applied voltage dependence of the tunneling current, its temperature dependence and noise properties can distinguish coherent and incoherent transport models and uncover the physics of edges in both polarized and unpolarized FQHE at $\nu = 2/3$.

*Edge states model*—Our key assumption is the separation of the charge and neutral or spin edge modes due to the long-range Coulomb interaction [21]. Such separation occurs both in polarized and unpolarized $\nu = 2/3$ phases [39], and is crucial for experimental observation of their Hall resistivity quantization at $R = 3h/2e^2$ [59,60,66]. While the separation in the polarized phase can also appear due to disorder-induced scattering between edge modes [35], this cannot be its origin in the unpolarized phase with opposite spins of edge modes. The similarity between the $\nu = 2/3$ plateau in the polarized and unpolarized phases [65,66] is therefore in favor of the Coulomb origin of charge-neutral (spin) separation. We treat the polarized and unpolarized edges on equal footing, using chiral Luttinger liquids description with action

$$S = \frac{1}{4\pi}\int dt \int dx\big[-3\partial_x\varphi_c(\partial_t + v_c\partial_x)\varphi_c + \partial_x\varphi_n(\partial_t - v_n\partial_x)\varphi_n\big], \tag{1}$$

where $v_c$ and $v_n$ are velocities of the charge and the neutral modes, $\varphi_c$ and $\varphi_n$ are their bosonic operators defining the corresponding densities via $\rho_{c,n} = (1/\sqrt{2}\pi)\partial_x\varphi_{c,n}$. We use the notation $\varphi_{pc}$ and $\varphi_{uc}$ for charge modes, $\varphi_{pn}$ and $\varphi_{us}$ for neutral (spin) modes. The neutral operators describe the difference in the occupation numbers of $\Lambda$ levels in the polarized phase and the doubled spin density in the unpolarized phase. We assume no edge reconstruction. Technical details of description of edge modes are discussed in the Supplemental Material [67].

*Tunneling and charge current*—We begin with tunneling through a QPC between the polarized and unpolarized phases. Tunneling is spin conserving and is carried out by interacting electrons. In the weak coupling limit, a rare electron tunneling is carried out by uncorrelated whole electrons, but for strong coupling the process transforms into correlated electron tunneling. The point contact electron tunneling at $x = 0$ between polarized and unpolarized phases is described by Hamiltonian

$$\mathcal{H}_T = -\sum_{a=1,2}(U_a\xi_u\xi_{pa}e^{i(\Psi_a(t)-Vt)} + \text{H.c.}). \tag{2}$$

Here $\{\xi_b, \xi_c\} = 2\delta_{b,c}$, $\xi$'s are Majorana fermions accounting for the fermion statistics of different CF edge modes and composing the corresponding Klein factors, and

$$\Psi_a(t) = \Phi_{pa}(0,t) - \Phi_{u1}(0,t) = \frac{1}{\sqrt{2}}\big[3\varphi_{pc}(0,t) \mp \varphi_{pn}(0,t) - 3\varphi_{uc}(0,t) + \varphi_{us}(0,t)\big], \tag{3}$$

where $\Phi_{pa}$ and $\Phi_{u1}$ are the CF fields, the tunneling amplitudes $U_a$ can be chosen real and positive, and $V$ is the applied voltage. We assume coherent propagation of charge and neutral (spin) modes along the edge and allow for the possibility of interference. The simpler model of tunneling of one mode into one mode is described by Hamiltonian (2), in which tunneling amplitude $U_2$ equals zero and the Klein factors are omitted.

The tunneling charge current from Eq. (2) is given by

$$J_T = -\partial_t\widehat{Q_P}(t) = i\left[\int dx\rho_{pc}, \mathcal{H}_T\right] = \frac{\delta}{\delta(Vt)}\mathcal{H}_T. \tag{4}$$

Its average and fluctuations are defined by the operators $\Psi_a(t)$, $a = 1,\ 2$, whose evolution in the absence of tunneling is characterized by their second-order correlators. The latter are found from the Gaussian action Eq. (1) for $\varphi_{pc}$, $\varphi_{pn}$ and similarly for $\varphi_{uc}$ and $\varphi_{ps}$ in the form

$$\langle\Psi_a(t)\Psi_b(0)\rangle = (3+\delta_{a,b})g(t); \tag{5}$$

$g(t)$ is the single point correlator at $x=0$ of a normalized chiral boson field, with action as in Eq. (1) for $\varphi_{pn}$,

$$\langle\varphi_{pn}(x,t)\varphi_{pn}(0,0)\rangle = -\ln\left\{\delta\left[i\left(t+\frac{x}{v_n}\right)+\alpha\right]\right\} \equiv g\left(t+\frac{x}{v_n}\right). \tag{6}$$

Here $\alpha = 1/D$, $D$ is the energy cutoff in both edges, and $\delta\to 0$ should be taken in the final results. In terms of the normalized chiral boson fields $\phi_j(x,t)$, $j=0,a$, the operators $\Psi_a(t)$, $a=1,2$ can also be represented as

$$\Psi_a(t) = \phi_a(0,t)+\sqrt{3}\phi_0(0,t). \tag{7}$$

The fermionization

$$\psi_a(x,t)=\frac{\xi_{pa}}{\sqrt{2\pi\alpha}}e^{i\phi_a(x,t)}, \quad \psi_0(x,t)=\frac{\xi_u}{\sqrt{2\pi\alpha}}e^{i\sqrt{3}\phi_0(x,t)} \tag{8}$$

allows us to rewrite the tunneling Hamiltonian (2) as

$$\mathcal{H}_T = -\sum_{a=1,2}(2\pi\alpha U_a\psi_0^+(0,t)\psi_a(0,t)e^{-iVt}+\text{H.c.}), \tag{9}$$

which describes the tunneling process of electrons from the two noninteracting chiral channels into the FQHE edge of $\nu=1/3$ filling factor. Combining these interfering channels into a single tunneling channel and applying bosonization of its chiral fermion field

$$\psi_T(x,t)=\frac{1}{\sqrt{\sum_a U_a^2}}\sum_a U_a\psi_a(x,t)\equiv\frac{\xi_T}{\sqrt{2\pi\alpha}}e^{i\phi_T(x,t)}, \tag{10}$$

we obtain the tunneling Hamiltonian $\mathcal{H}_T$ and current $J_T$:

$$\mathcal{H}_T = -\tilde{t}\cos(\phi_T(0,t)-\sqrt{3}\phi_0(0,t)-Vt) \tag{11}$$

$$J_T = -\tilde{t}\sin(\phi_T(0,t)-\sqrt{3}\phi_0(0,t)-Vt), \tag{12}$$

where $\tilde{t}^2 = 2\sum_a U_a^2$, and the single Klein factor is omitted since its drops out from any perturbative order in $\tilde{t}$ due to the charge conservation. In Hamiltonian (11) $\phi_0 = \sqrt{3/2}(\varphi_{pc}-\varphi_{us})$ is a combination of the charge edge modes of the different $\nu=2/3$ FQH regions and $\phi_T$ is constructed from their neutral and spin modes, which in the case $U_2=0$ can be explicitly expressed as $\phi_T = (\varphi_{pn}-\varphi_{us})/\sqrt{2}$. Therefore the tunneling current $J_T$ describes charge and spin tunneling into the spin-unpolarized FQH phase. The same scheme also applies to the QPC tunneling between two spin polarized FQH regions when two CF edge modes are well separated in space due to potential profile in one of the regions, so that only one mode $\Phi_{Ip} = 3\varphi_{Ipc}-\varphi_{Ipn}$ can tunnel to the neighboring region. The modes $\Phi_{Ip}$, $\varphi_{Ipc}$ and $\varphi_{Ipn}$ then enter our equations instead of $\Phi_{u1}$, $\varphi_{uc}$ and $\varphi_{us}$, correspondingly. The calculations of the tunneling current is defined by the scaling dimension 2 of the tunneling operators in Eq. (11) which manisfests itself in the power dependence $t^{-4}$ of their correlators, according to Eq. (5). We thus demonstrated this scaling power to emerge in the description of the complex edge at $\nu=2/3$ for several cases of tunneling between polarized or unpolarized FQH phases.

Tunneling from the hybridization-induced single noninteracting chiral channel in Eqs. (11), (12) can be viewed, by analogy to [68], as tunneling between the two counterpropagating primary one-component edges of equal filling factor $\nu'=1/2$ with densities $\rho_{d,n}(x,t) = \pm(1/2\pi)\sqrt{\nu'}\partial_x\phi_{d,n}(x,t)$ and electron annihilation operators $e^{i(1/\sqrt{\nu'})\phi_{d,n}(x,t)}$, $\phi_{d,n}$ are normalized right and left moving chiral fields. Thus, Eq. (2) maps onto and is equivalent to the problem of tunneling between $\phi_d(x,t)$ and $\phi_n(x,t)$ described by Hamiltonian

$$\mathcal{H}_T = -\tilde{t}\cos\left(\frac{1}{\sqrt{\nu'}}\phi_d(0,t)-\frac{1}{\sqrt{\nu'}}\phi_n(0,t)-Vt\right). \tag{13}$$

The current $J_T$ for $\mathcal{H}_T$ of Eq. (13) coincides with $J_T$ in Eqs. (4), (12). Its average [27,28] at finite temperatures $T$ is given by

$$\langle J_T\rangle = \frac{1}{4\pi}\left(V-\frac{\Gamma^2}{2}\int d\omega\frac{f(\frac{\omega-V}{T})-f(\frac{\omega+V}{T})}{\omega^2+\Gamma^2}\right), \tag{14}$$

$f(\epsilon/T)$ is the Fermi function, and $\Gamma = 2D^2/(\pi\tilde{t})$ [69]. Equation (14) is evaluated via the digamma-function $\psi(x)$ [70], describing the dependence of $\langle J_T\rangle$ on the applied voltage and temperature in the whole range of parameters:

$$\langle J_T\rangle = \frac{1}{4\pi}\left(V-\Gamma\,\mathrm{Im}\,\psi\left[\frac{1}{2}+\frac{\Gamma+iV}{2\pi T}\right]\right). \tag{15}$$

In Fig. 1 conductance $\langle J_T\rangle/V$ is plotted in terms of variables $1/\tilde{V} = \Gamma/V$ and $1/\tilde{T} = \Gamma/2\pi T$. At $T=0$

$$\langle J_T\rangle = \frac{1}{4\pi}\left[V-\Gamma\tan^{-1}\left(\frac{V}{\Gamma}\right)\right]. \tag{16}$$

In order to elucidate different mechanisms of the tunneling, we study limiting cases of Eq. (16). At low $V$

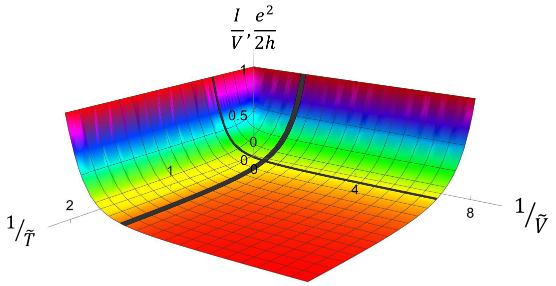


FIG. 1. Conductance $\langle J_T \rangle / V = I/V$ dependence on dimensionless variables $1/\tilde{V} = \Gamma/V$ and $1/\tilde{T} = \Gamma/2\pi T$. Thick and thin black curves show the dependence of conductance on $1/\tilde{T}$ at $1/\tilde{V} = 2$ and on $1/\tilde{V}$ at $1/\tilde{T} = 0.5$, correspondingly.

limit of model Eq. (13), $V \ll \Gamma$, the current $\langle J_T^l \rangle$ is cubic in $V$:

$$\langle J_T^l \rangle = \frac{1}{3} \frac{V}{4\pi} \left( \frac{V}{\Gamma} \right)^2. \tag{17}$$

Therefore, we observe that in the coherent approach to transport, the behavior of tunneling current and conductance at small voltages is nonlinear in the applied voltage for $T \ll V$, as opposed to the picture of incoherent transport [51–53] that results in Landauer-Buettiker current linear in voltage and voltage-independent conductance.

In the high-voltage regime, $V \gg \Gamma$, Eq. (15) gives the following result for the average current denoted $J_T^h$:

$$\langle J_T^h \rangle = \frac{V}{4\pi} - \frac{\Gamma}{8}. \tag{18}$$

Here the first term defines conductance $G = e^2/2h$. The second term gives the reduction of the tunneling current $\langle J_T \rangle$ due to the quasiparticle backscattering current $J_{bsc}$

$$J_{bsc} = \langle J_T \rangle - \langle J_T^h \rangle = \frac{\Gamma}{8}, \tag{19}$$

where $\langle J_T \rangle = V/4\pi$.

In the Supplemental Material [67], we show an alternative calculation of the tunneling currents using the strong coupling boundary conditions. We also calculate the quasiparticle charge, and neutral and spin tunneling currents for tunneling through QPCs in several configurations. We note that QPC conductance $G = e^2/2h$ has been also discussed [58] in terms of the incoherent equilibration model. However, neither the nonlinear current-voltage characteristics nor the temperature dependence following from Eq. (15) arise in that approach. $G = e^2/2h$ has been also discussed for a special tunneling configuration, in which tunneling between two $\nu = 2/3$ regions proceeds through the quantum dot with $\nu = 1$ [71].

*Quasiparticle charge*—The reduction of tunneling current due to quasiparticle backscattering is described and their charge is determined by taking into account a sudden change of strong coupling boundary conditions at $t = t_0$. For tunneling, e.g., between the $p1$ edge mode in the polarized region and $u1$ edge mode with the same spin in the unpolarized region through the QPC at $x = 0$ the strong coupling boundary conditions $\tilde{t} \to \infty$ are given by

$$\frac{1}{2} \sum_{\pm} \left( \tilde{\Phi}_{p1}(\pm 0, t_0) - \tilde{\Phi}_{u1}(\pm 0, t_0) \right) = 2\pi n. \tag{20}$$

Here both $\Phi_{\alpha,1}(x,t)$, $\alpha = p, u$ fields are extended to finite $x$ from their $x = 0$ expressions in Eq. (3), as chiral right moving fields. A sudden variation of the boundary conditions changes $n = 0$ to $n = 1$; both $n$ minimize $-\tilde{t} \cos[\Phi_{p1}(0,t) - \Phi_{u1}(0,t)]$. The boundary conditions also must keep continuous the two dual fields

$$\eta_1(x,t) = \Phi_{p1}(x,t)\theta(-x) + \Phi_{u1}(x,t)\theta(x) \tag{21}$$

$$\eta_2(x,t) = \Phi_{u1}(-x,t)\theta(-x) + \Phi_{p1}(x,t)\theta(x), \tag{22}$$

where $\theta(x)$ is the Heaviside step function. The jump from $n = 0$ to $n = 1$ leads to a jump in the dual fields

$$(\eta_1(x,t_0) - \eta_2(x,t_0))|_{x=-0}^{+0} = -4\pi. \tag{23}$$

Since $(\eta_1(x,t_0) + \eta_2(x,t_0))|_{x=-0}^{+0} = 0$, the corresponding jumps in these fields are given by

$$\Delta\eta_2(+0,t_0) = 2\pi = -\Delta\eta_1(+0,t_0). \tag{24}$$

This leads to the change $\Delta(3\varphi_{pc}(x,t_0) - \varphi_{pn}(-x,t_0)) = 2\sqrt{2}\pi\theta(x)$. Therefore, using the continuity condition $\Delta(\varphi_{pc}(x,t_0) + \varphi_{pn}(-x,t_0)) = 0$, we find

$$\Delta\varphi_{pc}(x,t_0) = \frac{\pi}{\sqrt{2}}\theta(x), \quad \Delta\varphi_{pn}(-x,t_0) = -\frac{\pi}{\sqrt{2}}\theta(x). \tag{25}$$

Hence the changes in the charge density and the density of the neutral mode are given by

$$\delta\rho_{pc} = \frac{1}{\sqrt{2}\pi} \partial\Delta\varphi_{pc} = \frac{1}{2}\delta(x), \qquad \delta\rho_{pn} = \frac{1}{2}\delta(x). \tag{26}$$

Changes in tunneling densities of the spin mode in the unpolarized region and relations between changes of densities for charge, neural, and spin modes on both sides of the QPC are presented in the Supplemental Material [67].

The transferred charge (26) enables [72] finding the charge of backscattering quasiparticles, $q = e/2$. Surprisingly, this charge differs from $q = e/3$ that one expects from [35] for $\nu = 2/3$. This is due to the absence of tunneling between two edge states of opposite spin polarizations in the polarized and unpolarized regions, or, for tunneling between the regions with the same spin polarization, due to absence of tunneling from at least one channel in one of the regions, owing to the potential profile.

*Shot noise*—The charge of the tunneling quasiparticles can be measured studying the shot nonequilibrium noise [73–78]. The shot noise of the tunneling current [27,68,79–81] is given by

$$S = \int_{-\infty}^{+\infty} dt(\langle J_T(t) J_T(0)\rangle - \langle J_T\rangle^2). \tag{27}$$

From [27] it follows that

$$\begin{aligned} S &= -\frac{\nu}{2(1-\nu)} V^2 \partial_V \left(\frac{\langle J_T\rangle}{V}\right) \\ &= \frac{1}{2}\frac{\Gamma}{4\pi}\left(\tan^{-1}\left(\frac{V}{\Gamma}\right) - \frac{v/\Gamma}{1+(\frac{V}{\Gamma})^2}\right). \end{aligned} \tag{28}$$

At small voltages and weak backscattering $V \ll \Gamma$

$$S^l = \frac{1}{3}\frac{V}{4\pi}\left(\frac{V}{\Gamma}\right)^2. \tag{29}$$

At large voltages $V \gg \Gamma$

$$S^h = \frac{1}{2}\frac{\Gamma}{4\pi}\frac{\pi}{2}. \tag{30}$$

The Schottky formula gives the charge of tunneling quasiparticles $q$ in the limits of weak electron tunneling and weak quasiparticle backscattering as the ratio of the corresponding values of noise to the currents, $S^l/\langle J_T\rangle$ and $S^h/\langle J_{bsc}\rangle$. The resulting charge (writing $e$ explicitly) is

$$S^l/\langle J_T\rangle = q = e, \quad S^h/\langle J_{bsc}\rangle = q = e/2. \tag{31}$$

Thus, the result for shot noise confirms the one from the analysis of the boundary conditions: in our models for QPC tunneling at $\nu = 2/3$, the quasiparticle charge is $e/2$. Notably, the arising noise signal is much stronger than the noise in the incoherent equilibration approach.

*Conclusion*—The coherent model of transport in the $\nu = 2/3$ FQHE involving interference of chiral Luttinger liquid edge modes leads to an exact solution of the problem of tunneling through the QPC between $\nu = 2/3$ FQHE regions with different or similar spin phases. With the increase of the applied voltage, the QPC conductance grows from zero and saturates at $e^2/2h$ while a weak electron tunneling through the QPC between the same spin edge modes transforms into the backscattering carried by the fractional charge $q = e/2$ quasiparticles. Unusual new quasiparticles and fractional conductance emerge in the QPC with one or two CF modes scattering into one mode as is the case in tunneling between polarized and unpolarized phases and can occur in tunneling between similar phases due to engineering of the potential profile. All these cases are characterized by the same power time dependence of the correlators of the tunneling operators. Using the fermionization method, we have shown that tunneling of the two CF modes from one of the sides of the QPC is renormalized with the account of interference between the two modes, and is equivalent to tunneling of a single mode.

Recent experiments on QPC tunneling at $\nu = 2/3$ have shown signatures of $e^2/2h$ plateau. Besides our model this conductance can be discussed using the incoherent equilibration transport approach. However, then the nonlinear current-voltage characteristics and temperature dependence of tunneling current and noise, which are predicted in the present work and can be tested experimentally, do not emerge in the incoherent model.

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