Finite-frequency noise for edge states at a filling factor $\nu = 2/5$

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Received 21 May 2012
Accepted for publication 27 June 2012
Published 30 November 2012
Online at stacks.iop.org/PhysScr/T151/014025

Abstract
We investigate the properties of the finite-frequency noise in a quantum point contact geometry for the fractional quantum Hall state at a filling factor $\nu = 2/5$. The results are obtained in the framework of Wen’s hierarchical model. We show that the peak structure of the colored noise allows us to discriminate among different possible excitations involved in the tunneling. In particular, the optimal values of voltage and temperature are found in order to enhance the visibility of the peak associated with the tunneling of a 2-agglomerate, namely an excitation with charge double the fundamental one associated with the single quasiparticle.

PACS numbers: 71.10.Pm, 72.70.+m, 73.43.−f

(Some figures may appear in color only in the online journal)

1. Introduction

The fractional quantum Hall effect [1] is one of the most remarkable examples of a strongly correlated electron system. Since its discovery it has been the subject of an intense study from both the theoretical and the experimental point of view, with the aim of providing a unified picture of the plethora of different observed values of the filling factor $\nu$ [2]. A suitable theoretical framework for the description of these states is provided by the theory of edge states [6]. For the Laughlin series [3] at $\nu = 1/(2n + 1)(n \in \mathbb{Z})$, a chiral Luttinger liquid theory described in terms of a single bosonic mode was proposed [4]. For the more general Jain series [5] with $\nu = p/(2np + 1)(p \in \mathbb{Z})$, a possible extension has been proposed by Wen and Zee [7] considering $|p| \geq 1$ additional hierarchical fields, propagating with a finite velocity along the edge.

Noise experiments in the quantum point contact (QPC) geometry [8] have been crucial for demonstrating the existence of the peculiar fractionally charged excitations predicted by the theoretical models [3]. In particular, it was proved that, for the Jain sequence, the elementary quasiparticle (qp) charge is given by $e^* = e/(2np + 1)$ [9–11]. More recently, some intriguing transport experiments have been performed on a QPC at very low temperature and extremely weak backscattering for various filling factors, showing a change in the power law of the backscattering current as a function of temperature and an increasing of the effective tunneling charge, as extracted from the noise [12–14]. In order to explain both these unexpected behaviors, we observed that propagating neutral modes could crucially modify the scaling dimensions of operators, leading to a crossover in the relevance of the tunneling process according to measurements [15–19].

The symmetrized noise at a finite frequency [20] represents an important tool for characterizing the predicted crossover phenomenology, complementary and alternative to current and zero-frequency noise. From the experimental point of view, measurements of colored noise have been carried out for a QPC in a 2D electron gas at zero magnetic field [21]. Nevertheless, a great deal of effort was devoted to extending the observations to more interesting cases including the FQH regime. From a theoretical point of view, the colored noise was investigated for the Laughlin sequence [22, 23] and also for more exotic states such as $\nu = 5/2$ [24, 25]. In these cases, much of the relevant physics occurs at frequencies close to the Josephson frequencies multiple of $\omega_0 = e^* V / h$ ($e^*$ is the charge of the elementary excitation) that is in the range of GHz for external bias $V$ in the range of $\mu V$. These values, although high, should be experimentally observable.
In this paper, we will discriminate among the different tunneling charges involved in the transport, being their contributions resolved at different frequencies. We will focus on the edge state at \( \nu = 2/5 \), where only two bosonic fields, one charged and one neutral, are involved [6]. We will demonstrate the necessity to consider voltages larger than the neutral modes cutoff frequency in order to efficiently detect the presence of the 2-agglomerate.

The paper is organized as follows. In section 2 we recall Wen’s model for the description of the FQH state at \( \nu = 2/5 \). In section 3 we provide the expression for the backscattering current and noise in terms of the tunneling rates. In section 4 we analyze the scaling behaviors for various energy regimes. In section 5 we present the finite-frequency noise peak structure for \( \nu = 2/5 \) as a function of the bias voltage in order to determine the optimal values of the physical parameters to observe the 2-agglomerate peak. Section 6 is devoted to the conclusions.

2. The model

We consider infinite edge states of a Hall bar in the Jain sequence at filling factor \( \nu = p/(2np + 1) \) [5]. We focus on the case \( \nu = 2/5 \) \((n = 1 \text{ and } p = 2)\), described by the Lagrangian density \([6, 17] \( h = 1 \)
\[
\mathcal{L} = \frac{5}{8\pi} \partial_x \varphi^c \overline{\partial_t} \varphi - \frac{1}{8\pi} \partial_x \varphi^c \overline{\partial_t} \varphi^c + \frac{1}{8\pi} \partial_x \varphi^c \overline{\partial_t} \varphi^c - \frac{1}{8\pi} \partial_x \varphi^c \overline{\partial_t} \varphi^c + \frac{1}{8\pi} \partial_x \varphi^c \overline{\partial_t} \varphi^c, \tag{1}
\]

where \( \nu_c \) is the propagation velocity of the charged bosonic mode \( \varphi^c \) and \( \nu_n \) is that of the neutral bosonic mode \( \varphi^n \). The electron number density of the edge depends on the charged field alone, via the relation \( \rho(x) = \partial_x \varphi^c(x)/2\pi \). Owing to this fact the charge mode velocity is affected by Coulomb interactions, leading to the reasonable assumption \( \nu_c \gg \nu_n \) [15, 26, 27]. From (1) one can note that the neutral mode co-propagates with respect to the charged one as is typical of the states with \( p > 0 \) [6]. The commutators of the bosonic fields are \([\varphi^c(x), \varphi^c(y)](\omega) = i\pi \nu_c/\pi \text{sgn}(x - y) \text{ with } \nu_c = 2/5 \) and \( \nu_n = 2 \).

We consider a generic \( m \)-agglomerate annihilation operator with charge \( e_m = (en\nu_c)/2 = me^\ast \) with \( e^\ast = e/5 \) being the charge of the elementary excitation, namely the single qp \( e \) (the electron charge). It can be written as
\[
\Psi^{(m,j)}(x,t) = \frac{\mathcal{F}^{(m,j)}}{\sqrt{2\pi a}} e^{i\frac{me^\ast}{2a}(\tau - \nu t)} e^{i\varphi(x,t)}, \tag{2}
\]

where \( a \) is a finite-length cutoff, \( m \in \mathbb{N} \) is related to the charge of the excitation and the additional quantum number \( j \in \mathbb{Z} \) plays a role analogous to the isospin and is necessary in order to provide the proper fractional statistics [18]. Note that for a given qp excitation the integers \( m \) and \( j \) must have the same parity. The Klein factors \( \mathcal{F}^{(m,j)} \) are ladder operators necessary in order to change the particle numbers of the \( m \)-agglomerates and to provide the proper statistical properties between excitations [16, 18]. From the long-time limit of the two-point imaginary time Green function \( \mathcal{G}^{(m,j)}(\tau) = \langle T_{\tau} \Psi^{(m,j)}(0,\tau) \Psi^{(m,j)^\dagger}(0,0) \rangle \) calculated at \( x = 0 \) and at zero temperature [28], one can extract the scaling dimension of the \( m \)-agglomerate
\[
\Delta^{(m,j)} = \frac{1}{2} \left[ \nu_c \left( \frac{m}{2} \right)^2 + \nu_n \left( \frac{j}{2} \right)^2 \right]. \tag{3}
\]

For energies higher than the typical neutral mode bandwidth \( \omega_n = \nu_n/a \), the neutral mode contribution to the dynamics saturates and one has the effective scaling dimension [15, 16]
\[
\bar{\Delta}^{(m,j)} = \frac{1}{2} \nu_c \left( \frac{m}{2} \right)^2 \tag{4}
\]

that only depends on the charged sector of the theory. From now on, the corresponding charged mode bandwidth \( \omega_c = \nu_c/a \) is assumed to be the greatest energy scale of the system.

Note that the above scaling dimension could be strongly affected by interaction with the external environment [15, 19, 29–34]; nevertheless, in this paper, for the sake of simplicity, we only consider the standard unrenormalized case.

3. Transport properties

The tunneling through the QPC at \( \chi = 0 \) of a generic \( m \)-agglomerate between the \( R \) and \( L \) edges of the Hall bar is described by the Hamiltonian \( H^{(m)} = \sum_m \Psi_{\chi}^{(m)}(0) \Psi_{L}^{(m)}(0) + \text{ h.c.} \), where \( \Psi_{\chi}^{(m)}(\Psi_{L}^{(m)}) \) is the \( m \)-agglomerate annihilation operator. From equation (3) it is easy to note that the more relevant single qp’s \((m = 1)\) are those with \( j = \pm 1 \) \((\Delta^{(1,\pm1)} = 3/10)\), while among all the \( 2 \)-agglomerate \((m = 2)\) we consider that with \( j = 0 \) \((\Delta^{(2,0)} = 1/5)\). All the other operators with the same charge \((\text{the same } m)\) and different \( j \) are less relevant in the renormalization group sense and give a negligible contribution to the transport properties with respect to the considered one [15]. From now on, for notational convenience, we will omit the index \( j \) where not needed.

The tunneling amplitudes \( t_m \) depend, in general, on the geometry of the constriction, and can be energy dependent [8]. In the following, they will be assumed as \( m \)-dependent constants in order to make the discussion clearer and to restrict the number of free parameters of the theory. The total tunneling Hamiltonian will consist of the sum over all possible \( m \)-agglomerates \( H_T = \sum_m H^{(m)}_T \).

In the following we will focus on the weak tunneling regime. At lowest order in the tunneling Hamiltonian, the backscattering current of the \( m \)-agglomerate \( I^{(m)}_b \) and the finite-frequency symmetrized noise \( S^{(m)}_{\text{sq}}(h\omega) \) can be written in terms of the tunneling rates [20]. By using the detailed balance relation, the current is [16] \((k_B = 1)\)
\[
I^{(m)}_b(h\omega) = me^\ast \left( 1 - e^{-h\omega/\kappa_B T} \right) \Gamma^{(m)}(h\omega), \tag{5}
\]

with tunneling rate
\[
\Gamma^{(m)}(E) = |t_m|^2 \int_{-\infty}^{+\infty} dt e^{imEt} G^{\geq}_{m,R}(0,-t)G^{\leq}_{m,L}(0,t). \tag{6}
\]

Here, \( \omega_n = e^\ast V \) is the Josephson frequency associated with the single qp with \( V \) being the bias. The correlators \( G^{\geq}_{m,j}(0,t) = \langle \Psi_{L}^{(m,j)}(0,t) \Psi_{\chi}^{(m,j)^\dagger}(0,0) \rangle = G^{\leq}_{m,j}(0,t) \ast \) are the...
two-point Green functions of the \( m \)-agglomerate operators on the edge \( l = R, L \).

The noise spectral density is

\[
S_B^{(m)}(\omega) = \int_{-\infty}^{+\infty} dt \, e^{i\omega t} S_B^{(m)}(t),
\]

(7)

where

\[
S_B^{(m)}(t) = \left[ \langle \delta I_B^{(m)}(t) \delta I_B^{(m)}(0) \rangle + \langle \delta I_B^{(m)}(0) \delta I_B^{(m)}(t) \rangle \right],
\]

(8)

with \( \delta I_B^{(m)} \) the fluctuactions of the current with respect to the average.

Recalling the definition of the tunneling rate in (6), one has

\[
S_B^{(m)}(\omega) = (me^*)^2 \sum_{t,\nu=\pm} \left[ S^{(m)}(\epsilon/\alpha m + \eta \omega_n) \right],
\]

(9)

or, equivalently, in terms of the backscattering current

\[
S_B^{(m)}(\omega) = me^* \sum_{t,\nu=\pm} \coth \left( \frac{\epsilon/\alpha m + \eta \omega_n}{2T} \right) \times I_B^{(m)}(\epsilon/\alpha m + \eta \omega_n).
\]

(10)

The above result is consistent with the non-equilibrium fluctuation–dissipation theorem [20, 35]. Note that from (9) one can easily restore the well-known result for the zero-frequency limit [16, 36, 37].

At the lowest perturbative order the total noise will be

\[
S_B(\omega) = \sum_m S_B^{(m)}(\omega),
\]

(11)

being the contributions of the different \( m \)-agglomerate independent. Note that, analogously, the total backscattering current is given by \( I_B = \sum_m I_B^{(m)} \). The simple relations in (10) and (11) are due to the Poissonian statistics of the tunneling processes at lowest order in \( |t_0|^2 \) and to the independency of the sources of noise.

4. Scaling behavior

Let us now focus on the evaluation of the tunneling rate in (6), starting from the zero-temperature limit. Here, the bosonic Green’s functions are [38–40]

\[
\langle \phi^0(0, t) \phi^0(0, 0) \rangle = -v_s \ln(1 + i\omega_n t),
\]

(12)

with \( s = c, n \), leading to the tunneling rate [15]

\[
\Gamma^{(m)}(E) = \left| t_{0n} \right|^2 \frac{2\pi a^{\alpha s}}{\alpha c \omega_n} \cdot \left[ \frac{mE}{\omega_n} - \frac{mE}{\omega_n} \right] \coth \left( \frac{mE}{\omega_n} \right) \times \Theta(mE).
\]

(13)

Here, \( F_1[\alpha; a; b; c] \) is the Kummer confluent hypergeometric function [41], and \( \alpha = v_m m^2/2 \) and \( \delta = v_n j^2/2 \) are the charged and the neutral exponent, respectively.

The rate shows two different regimes. At low energies \( (E \ll \omega_n) \) the rate scales as

\[
\Gamma^{(m)}(E) \approx E^{\Delta^{(m, \beta) - 1}},
\]

(14)

receiving contributions from both the charged and the neutral modes. In the same limit, for frequencies close to the Josephson resonance \( \omega \to m\omega_n \), one has (cf (9))

\[
S_B^{(m)}(\omega \to m\omega_n) \approx (\omega - m\omega_n)^{\Delta^{(m, \beta) - 1}}.
\]

(15)

As stated before, for the single qp the scaling dimension is given by \( \Delta^{(1, 1)} = 3/10 \), while for the next excitation, the 2-agglomerate, the scaling is driven by the charged mode contribution only in such a way that \( \Delta^{(2, 0)} = 1/5 \) (cf (3)). These behaviors indicate the 2-agglomerate as the dominant excitation in the current [15]. All other excitations, with higher charges and different isospin, have higher scaling dimensions and can be safely neglected.

Note that the peculiar values of the scaling dimensions imply a divergent power-law behavior of the total noise in correspondence to the 2-agglomerate Josephson frequencies \( \omega_n, \omega_c \) and a dip for the single qp \( \omega_n \).

For applied voltages higher than the neutral mode cutoff \( \omega_h > \omega_n \), the neutral modes saturate leading to a lower effective dimension \( \Delta^{(1, 1)} = 1/20 \) for the single qp (cf (4)). On the other hand, the 2-agglomerate scaling is unaffected. In this case one has a two-peak structure in correspondence to the Josephson frequencies of the two excitations \( \omega_n, \omega_c \), respectively.

At finite temperature the above behaviors will be smoothed with more remarkable changes close to the Josephson resonances \( \omega = \omega_h \) for \( T > |\omega - m\omega_n| \). To quantitatively determine the temperature influence on the noise, one has to consider the finite-temperature expressions for Green’s functions in (12) for the bosonic fields [15, 16]:

\[
\langle \phi^0(0, t) \phi^0(0, 0) \rangle = v_s \ln \left[ \frac{|\Gamma(1 + T/\omega_n - i\pi T)|^2}{\Gamma(1 + T/\omega_n)} \right]
\]

(16)

with \( s = c, n \).

The tunneling rate can still be analytically evaluated for temperatures lower than the bandwidths, namely \( T \ll \omega_n, \omega_c \), leading to

\[
\Gamma^{(m)}(E) = \left| t_{0n} \right|^2 \frac{(2\pi a^{\alpha s})}{\alpha c \omega_n} \cdot \frac{2}{T} \cdot \frac{e^{\frac{mE}{2T}}}{\alpha + \delta} \times B \left( \alpha + \delta - i \frac{mE}{2\pi T}, \alpha + \delta + i \frac{mE}{2\pi T} \right). \]

(17)

where \( B(\alpha; b) \) is the Euler beta function [41] \( (\alpha = v_m m^2/2 \) and \( \delta = v_n j^2/2 \). At higher temperatures the rates have to be evaluated numerically [16].

Close to the Josephson frequencies, the noise in (10) reduces to

\[
S_B^{(m)}(\omega \to m\omega_n) \approx TG^{(m)}(T)
\]

(18)

written in terms of the \( m \)-agglomerate linear conductance

\[
G^{(m)}(T) = (me^*)^2 \frac{\Gamma^{(m)}(0)}{T}.
\]

(19)

The power-law behavior for the height of the peaks as a function of temperature is therefore given by \( S_B^{(m)}(\omega \to m\omega_n) \approx T^{u + \delta - 1} \) for \( T \ll \omega_n \) and \( S_B^{(m)}(\omega \to m\omega_n) \approx T^{u - 1} \) for \( T \gg \omega_n \).

The visibility of the 2-agglomerate peak is guaranteed in this regime for \( S_B^{(2)}(2\omega_n) < S_B^{(2)}(2\omega_n) \).
5. Results

As stated before, for the composite edge states belonging to the Jain sequence, important information on the carrier charges involved in the tunneling can be extracted from knowledge of the total noise at a finite frequency in (11). Each $m$-agglomerate contribution leads to a resonance at $\omega = m\omega_0$. This enables us to resolve them and to obtain indications of their power-law behavior. Depending on their scaling dimensions, they show peaks (for $\Delta < 1/4$), dips (for $1/4 < \Delta < 1/2$) or a monotonic increasing as a function of frequency (for $\Delta > 1/2$). Note that, for notational convenience, we generically indicated with $\Delta$ both the low-energy scaling dimension in (3) and the effective scaling above the neutral mode bandwidth in (4) depending on the considered regime. In figure 1 we show the behavior of the spectral noise for the considered $v = 2/5$ case at low temperature ($T = 10$ mK).

The ratio $S_B(\omega)/2eI_B$ is plotted to recover the proper value of the Fano factor in the zero-frequency limit [16, 25]. To further simplify the description, we focus only on the first two contributions to the noise, namely single qp and 2-agglomerate, neglecting the contributions from other possible excitations. Various curves correspond to different values of the applied bias, namely different $\omega_0$. Note that in the figure the frequency is rescaled with respect to $\omega_0$.

For $\omega_0 > \omega_n$ (dotted magenta and short-dashed blue curves) it is easy to note a pronounced peak in correspondence to $\omega = \omega_0$ because of the single-qp contribution. Despite the fact that the temperature is very low, thermal effects lead to a smoothening of the peak. A second contribution, less pronounced, is observable at $\omega = 2\omega_0$, a signature of the presence of 2-agglomerate that has a less divergent power-law behavior. This secondary peak becomes more and more evident with increasing $\omega_0$.

For bias voltages such that $\omega_0 < \omega_n$ (long-dashed green and solid red curves) the neutral mode contribution increases the scaling dimension turning the single-qp peak into an extremely broad dip that completely hides the 2-agglomerate contribution. This is a clear demonstration that the ideal condition to observe the peak pattern for the $v = 2/5$ case is $\omega_0 > \omega_n$. A remark is in order concerning the role of temperature. Indeed, the smoothening induced by thermal effects could completely wash away the observed structure under the condition $T > \omega_n$; therefore very low temperatures are needed to see this structure.

6. Conclusions

In this paper we analyzed the finite-frequency noise for the FQH state at $v = 2/5$. We considered the presence of two more relevant excitations: the single qp, with charge $e/5$, and the 2-agglomerate, with charge $2e/5$. The finite-frequency noise has the unique possibility to resolve spectroscopically the contributions of different excitations looking at different Josephson resonances. We showed that the peak associated with the 2-agglomerate is more evident at a Josephson frequency higher than the neutral mode cutoff, where the tail of the single-qp contribution decreases faster. We also commented on the evolution of the height of the single-qp peak as a function of temperature in order to determine the optimal condition for the visibility of the considered peak pattern.

Acknowledgments

We thank M Carrega, T Martin and F Portier for useful discussions. We acknowledge support from the EU-FP7 via ITN-2008-234970 NANOCTM.

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Figure 1. Spectral noise $S_B(\omega)/2eI_B$ as a function of the normalized frequency $\omega/\omega_0$. Different values of the bias are given by $\omega_0 = 30$ mK (solid red), $\omega_0 = 100$ mK (dashed green), $\omega_0 = 200$ mK (short dashed blue) and $\omega_0 = 300$ mK (dotted magenta). Other parameters are $\omega_n = 5$ K, $\omega_n = 50$ mK, $|t_2|/|t_1| = 1$ and $T = 10$ mK.
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