Transport in finite incommensurate Peierls-Fröhlich systems

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We show that the conductance of a one-dimensional, finite charge-density-wave (CDW) system of the incommensurate type is not renormalized at low temperatures and depends solely on the leads. Within our formalism, we argue that a similar behavior (perfect conductance) should occur for a wide class of one-dimensional strongly correlated finite systems where interactions are current dependent. The universal conductance is related to the presence of an (anomalous) chiral symmetry. The fundamental role played by the finiteness of the sample and the adiabaticity of the contacts to Fermi-liquid leads is evidenced.

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I. INTRODUCTION

Transport in strongly correlated, low-dimensional mesoscopic systems has been a subject of intense investigation. The theoretical tools for studying one-dimensional (1D) systems (exact solutions, bosonization, renormalization group, etc.) are already quite well-known in the literature and some important concepts were developed in the last two decades. For instance, it is understood that 1D, interacting gapless systems at low energies fall into the Luttinger liquid universality class. The thermodynamics and the dynamical correlation functions of this liquid are also well established. Its transport properties however are less understood. For some time it was believed that the linear conductance of a Luttinger liquid at zero temperature is renormalized by the interaction strength. Recent advances in technology allowed the fabrication and manipulation of nanostructures, such as quantum wires, where this theoretical prediction could be tested. The first result found was negative, namely, the measured conductance was equal to the quantum $e^2/h$ for each propagating channel; later, some deviation from this value was observed. Soon it was understood that, for a finite Luttinger liquid junction or wire coupled to leads, the conductance is dominated by the noninteracting electron gas in the leads, i.e., it should not be renormalized by the interaction in the wire. Why this argument fails to explain particularly the results of Ref. remains an open issue. On the other hand, one expects that processes leading to backscattering, such as disorder, could be involved.

In contrast to these findings, one would imagine that a junction formed by a gapped system, like a half-filled Mott-Hubbard insulator or a Peierls-Fröhlich CDW should show an exponentially suppressed conductance at sufficiently low temperatures. Recently, Ponomarenko and Nagaosa demonstrated that a Hubbard-Mott insulator where the order parameter (gap) is dynamical has in fact perfect conductance, $G = 2e^2/h$ (the factor of 2 stands for spin degeneracy), i.e., transport is also controlled by the leads. In this work we show that a similar behavior occurs for an incommensurate CDW system adiabatically connected to Fermi-liquid leads. This result agrees with that obtained in Ref. where transport of charge in disordered mesoscopic CDW heterostructures was studied within the Keldysh formalism. In contrast to the elaborate calculation of Ref. here we present very simple physical arguments based on the existence of a chiral symmetry for the system as a whole (CDW plus leads) when the phase of the CDW order parameter is dynamic. This symmetry becomes anomalous upon the application of an external electric field. Outside the sample, the resulting chiral current is associated to the transport of free electrons; inside the CDW region, the current has an extra component related to the dynamics of the lattice degrees of freedom. As long as these degrees of freedom are not frozen or pinned by impurities, the sample conductance will depend solely on the physics outside the CDW region.

One is tempted to try to understand the underlying reason for such universal behavior in a system-independent manner. We argue that perfect conductance should be a general property of all one-dimensional Fermi-liquid/finite-system/Fermi-liquid structures, as long as they display a (anomalous) chiral symmetry. This occurs in turn when the finite system possesses this symmetry and is adiabatically connected to the Fermi-liquid reservoirs. Our work, in some sense, makes more specific some general ideas on this subject which have recently appeared in the literature.

II. THE PEIERLS-FRÖHLICH SYSTEM

The model effective Lagrangian density that describes a 1D incommensurate charge-density-wave system can be...
divided into three terms\(^\ddagger\) (\(\hbar = c = 1\)), namely,
\[
\mathcal{L}_{\text{CDW}} = \mathcal{L}_{\text{el}} + \mathcal{L}_{\text{el-ph}} + \mathcal{L}_{\text{ph}},
\]
where
\[
\mathcal{L}_{\text{el}}[\psi] = i\bar{\psi}(\partial_\mu + ieA_\mu)\gamma^\mu \psi, \tag{2}
\]
\[
\mathcal{L}_{\text{el-ph}}[\psi, \Delta] = \Delta \bar{\psi}P_L\psi + \Delta^\dagger P_R\bar{\psi}, \tag{3}
\]
and
\[
\mathcal{L}_{\text{ph}}[\Delta] = \frac{1}{2v} (\partial_t \Delta \partial_t \Delta - v^2 \partial_x \Delta \partial_x \Delta) - \frac{\omega_{\text{ph}}^2}{2v} \Delta^\dagger \Delta. \tag{4}
\]

Let us specify our conventions. We follow the usual space-time notation \(x^\mu \equiv (x^0, x^1) = (t, x)\) and \(\partial_\mu \equiv \partial/\partial x^\mu\), with the metric tensor \(g_{\mu\nu} = \text{diag}(1, -1)\). \(v\) denotes the CDW velocity in units such that the Fermi velocity \(v_F = 1\). The Dirac fermion field \(\psi\) has right- and left-moving components (the electron modes), \(\psi = (\psi_L, \psi_R)^T\), spanning a two-dimensional vector space where the following gamma matrices act:
\[
\begin{align*}
\gamma^0 &\equiv \sigma_3, \\
\gamma^1 &\equiv -i\sigma_2, \quad \text{and} \\
\gamma^5 &\equiv \gamma^0 \gamma^1 = \sigma_3.
\end{align*}
\]
Here, \(\sigma_1, \sigma_2, \sigma_3\) are the standard Pauli matrices. We also have \(\bar{\psi} \equiv \psi^\dagger \gamma_0\), whereas the right- and left-moving projectors are defined as \(P_{L,R} \equiv (1 \pm \gamma_5)/2\). (For the sake of simplicity, throughout this work we assume that the fermions are spinless.) The field \(\Delta\) and its adjoint \(\Delta^\dagger\) represent lattice harmonic, optical phonon modes of frequency \(\omega_{\text{ph}}\) and momenta \(\pm 2p_F\), where \(p_F\) is the Fermi momentum. By assumption, \(\hbar/p_F\) is incommensurate with the lattice constant.

In order to implement the adiabatic transition between the Peierls-Frohlich system (hereafter named junction) and the asymptotic Fermi-liquid leads we replace \(\Delta(x, t)\) by \(h(x)\Delta(x, t)\), where \(h(x)\) is a smooth function that vanishes outside a region \(\Omega\) of length \(L\); inside this region, \(h(x)\) jumps to the saturation value \(1\). This choice allows us to describe the complete system (junction plus leads) with a single Lagrangian density. Problems related to wave function matching at the interfaces are thus avoided.

When the field \(\Delta\) is frozen, chiral symmetry is explicitly broken even at the classical level and the fermions acquire a mass. The gap in the spectrum of excitations yields an insulating behavior and charge transport through the junction region is exponentially suppressed. For \(T \ll \Delta\), an elementary calculation yields the electric conductance
\[
G = \frac{4e^2}{\hbar} \exp\left(-\frac{2\Delta}{\Delta_L}\right), \tag{5}
\]
where \(\Delta_L \equiv \hbar p_F / L \ll \Delta\). Notice that exponentially small conductances are also the prediction of a quasiclassical theory of normal-metal/CDW junctions developed in Ref. \(^4\). If, on the other hand, we allow for \(\Delta\) to be a dynamical field, classical global chiral symmetry is restored. We show below that, in this situation, charge transport in the CDW junction is strongly enhanced.

III. BOSONIZATION

Let us bosonize the fermionic fields following the standard procedure, which includes the chiral anomaly in a consistent way. The universal rules are\(^3\)
\[
i\bar{\psi}\gamma^\mu \psi \leftrightarrow \frac{1}{2} \partial_\mu \phi \partial^\mu \phi, \tag{6}
\]
\[
j^\mu_\phi \equiv \bar{\psi}\gamma^\mu \psi \leftrightarrow \frac{\beta}{2\pi} \epsilon^\mu_\phi \partial^\nu \phi, \quad j^\lambda_\phi \equiv \bar{\psi}\gamma^\mu \gamma^5 \psi \leftrightarrow \frac{\beta}{2\pi} \partial^\mu \phi, \tag{7}
\]
where \(\beta = 2\sqrt{\pi}\) and \(\phi\) is a massless boson field. The constant \(C\) depends on how the theory (path integral) is renormalized. After the rules stated above, the bosonized Lagrangian density becomes
\[
\mathcal{L}_{\text{boson}} = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi + \frac{\hbar}{2} \left(C^* \Delta^\dagger e^{i\beta \phi} + C \Delta e^{i\beta \phi}\right) + \frac{c^2}{2\pi} \epsilon^{\mu\nu} \partial_\mu A_\nu \phi + \frac{\hbar^2}{2v} \partial_\mu \Delta^\dagger \partial_\mu \Delta \\
- \frac{v}{2} \partial_x (h\Delta^\dagger) \partial_x (h\Delta) - \frac{\omega_{\text{ph}}^2}{2v} \Delta^\dagger \Delta. \tag{8}
\]

The equations of motion for \(\phi\) and \(\Delta\) are, respectively,
\[
\partial_\mu \partial^\mu \phi + \frac{i\beta h}{2} \left(C^* \Delta^\dagger e^{-i\beta \phi} - C \Delta e^{i\beta \phi}\right) = -\frac{e\beta}{2\pi} E(x, t), \tag{9}
\]
and
\[
\frac{1}{2v} \hbar^2 \partial_t^2 \Delta - \frac{v\hbar}{2} \partial_x^2 (h\Delta) + \frac{\omega_{\text{ph}}^2}{2v} \Delta = \frac{C\hbar^2}{2} e^{-i\beta \phi}, \tag{10}
\]
where \(E(x, t) = \partial_t A_1(x, t) - \partial_x A_0(x, t)\). If we multiply Eq. (11) by \(\Delta^\dagger\) and subtract from the resulting expression its adjoint, we find, after using Eq. (10), that
\[
\partial_\mu j^\mu_\phi = -\frac{e\beta}{2\pi} E(x, t), \tag{11}
\]
where the total axial current components are
\[
j^\mu_0 = \frac{i\beta h^2}{2v} \left(\Delta^\dagger \partial_t \Delta - (\Delta \partial_t \Delta^\dagger)\right) + \partial_t \phi \tag{12}
\]
and
\[
j^\mu_1 = \frac{i\beta h^2}{2} \left[\left(h^\dagger \partial_t (h\Delta^\dagger) - \partial_x (h^\dagger) (h\Delta)\right) + \partial_x \phi\right]. \tag{13}
\]
Equation (12) is the anomalous divergence of the chiral current associated to the classical global symmetry
\[
\phi \to \phi + \theta, \quad \Delta \to \Delta e^{-i\beta \theta}. \tag{14}
\]
(In terms of the original fermionic fields, \(\psi \to e^{i\theta \gamma_5} \psi\)). It is important to notice that the chiral current depends on the particular region we are considering. Outside the junction, where \(h(x) = 0\), the current is due to electrons only. Inside the junction, both electrons and lattice degrees of freedom contribute to the chiral current.)
IV. SOLUTION TO THE FIELD EQUATIONS

We can understand why the conductance does not depend on the properties of the CDW junction by solving the equations of motion (10) and (11) for the fields $\phi$ and $\Delta$. This line of reasoning was introduced by Maslov and Stoner to arrive at a similar conclusion in the case of a Luttinger liquid wire connected to Fermi-liquid reservoirs. We thus assume that the electric field is zero outside the region $\Omega$, which we take extending from $-L/2$ to $+L/2$. The electric field $E(x,t)$ is switched on at an instant $t=t_+$ and eventually reaches the stationary value $E(x)$. To facilitate the discussion, let us assume that the electric field $E(x)$ and the profile $h(x)$ are even functions of $x$. As a result, the field $\phi$ will have the same property.

For late enough times, the charge current $I = eJ^y = -(e\beta/2\pi)\partial_t \phi$ should be stationary (time independent) in the region around $\Omega$. The (causal) solution of Eq. (14) compatible with these constraints is $\phi(x,t) = -k(t/2) + \phi_0$, where the plus (minus) sign corresponds to $x < -L/2$ ($x > L/2$) and

$$k = \frac{2\pi I}{e\beta}.$$

We now propose the following asymptotic ansatz to extend this solution to the interior of $\Omega$ for late times:

$$\phi(x,t) = \phi_a(x,t) \equiv f(x) - kt$$

and

$$\Delta(x,t) = \Delta_a \equiv D(x)e^{i\beta kt},$$

where the functions $f(x)$ and $D(x)$ satisfy the coupled differential equations

$$\partial_x^2 f + \frac{i\beta \hbar}{2} \left( C^*D^*e^{-i\beta f} + CDe^{i\beta f} \right) = \frac{e\beta}{2\pi} E(x)$$

and

$$\frac{h^2 \partial^2 f}{2v} + \frac{e\hbar}{2} \partial_x^2 (hD) - \frac{\omega^2 \hbar^2}{2v} D = \frac{C^*h}{2}e^{-i\beta f},$$

The correct large-distance asymptotics are implemented assuming that $f(x) = -kx + \phi_0$ for $x < -L/2$ and $f(x) = kx + \phi_0$ for $x > L/2$. The consistency of the ansatz will be checked in the next section. For now, we remark that it is not necessary to know the exact forms of $f(x)$ and $D(x)$ to calculate the conductance of the junction. The reason is the following. Using the anomalous divergence of Eq. (12), together with Eqs. (17) and (18), we arrive at

$$\partial_x \left\{ \frac{i\beta \hbar}{2} [D^*\partial_x(hD) - D\partial_x(hD^*)] + \partial_x f \right\} = \frac{e\beta}{2\pi} E(x).$$

Integrating this expression between points $a$ ($a < -L/2$) and $b$ ($b > L/2$) and since $h(a) = h(b) = 0$ outside $\Omega$, we get

$$2k = \partial_x f(b) - \partial_x f(a) = \frac{e\beta}{2\pi} [V(a) - V(b)],$$

where $V(x)$ is the electric potential. Recalling Eq. (10), we arrive at the relation

$$I = \frac{e^2}{2\pi} [V(a) - V(b)],$$

which leads to the linear conductance (restoring the $\hbar$ factor),

$$G = \frac{e^2}{\hbar}.$$ 

The conductance is perfect, that is, it is not renormalized by the dynamics inside the junction, only depending on the properties of the electron gas in the leads. An interpretation of this result in terms of the dynamics of the fields $\phi$ and $\Delta$ is in order. Notice that in the asymptotic regime, the bosonized electron field and the phonon phase “rotate” coherently, allowing for a perfect matching between charged wave packets traveling in and out of the junction region. The time evolution represented by Eqs. (12) and (13) is very similar to the chiral transformation of Eq. (15). Thus, the action cost of this time evolution is rather small and favorable. The lattice is left in an excited oscillating state, with a frequency controlled by the bias and the Fermi velocity in the leads, while a perfect electric current is established through the system.

V. ASYMPTOTIC IN AND OUT BEHAVIORS

Notice that $f(x)$ is a continuous and differentiable function that grows linearly for $|x| > L/2$ and is bounded for $|x| < L/2$. Therefore, if $t > t_+$ and for large enough $t_+\phi_a(x,t)$ will be negative in the interval $[l(t),r(t)]$ and positive outside, where $l(t) = -t + \phi_0/k < -L/2$ and $r(t) = t - \phi_0/k > L/2$. The points $l(t)$ and $r(t)$ move in opposite directions, away from $\Omega$. They propagate through the leads, with Fermi velocity, the information about the switching on of the electric field in the region $\Omega$ at earlier times. Thus, to be accurate, the asymptotic solution for the field $\phi$ should be represented by

$$\phi_a(x,t) = [f(x) - kt] \theta(kt - f(x)),$$

where $\theta(z)$ is the step function. It is easy to see by direct substitution that $\phi_a$ and $\Delta_a$ satisfy the time dependent equations (11) and (13) with the stabilized electric field profile $E(x)$; hence, at a given instant $t > t_+$, the asymptotic electric current is given by Eq. (23) inside $[l(t),r(t)]$, and is zero outside.
In order to check the consistency of the asymptotic solution proposed for the field $\phi$ with a periodic time dependence for $\Delta$, let us solve Eq. (11) for $\Delta(x, t)$ in terms of $\phi(x, t)$. Inverting Eq. (11), we can write that

$$\Delta(x, t) = -\frac{C_0}{2} \int dx' \int_0^t dt' \ G(x; t, x', t') h(x') \times e^{-i2\phi(x', t')},$$

where $G$ is the retarded Green’s function of the differential operator

$$\left[ -\frac{h^2}{2}\frac{\partial^2}{\partial t^2} - \frac{v h}{2} (h'' + 2h'\partial_x + h\partial_x^2) + \frac{h^2\omega_p^2}{2v} \right].$$

(27)

The use of the retarded prescription is justified as follows. The effective action obtained by path integrating over the phonon degrees of freedom in Eq. (11) depends explicitly on the Feynman phonon propagator; therefore, the associated effective field equation for the bosonized $\phi$ field is neither real nor causal. This happens because this equation refers to in-out vacuum expectation values of the $\phi$ field. Alternatively, one can use a close-time path formalism\cite{14} to construct real and causal effective equations for the in-in vacuum expectation values of the Hermitian field $\phi$. It can be shown that the in-in effective equations can be obtained by replacing, in the corresponding in-out expressions, the Feynman propagator by the retarded one. The resulting effective equations are equivalent to Eqs. (10) and (13) with $\Delta$ solved by means of a retarded prescription [c.f. Eq. (24)].

Since the operator (27) is time independent, it is useful to pass to a Fourier representation, namely,

$$\tilde{G}(x, x'; t - t') = \int \frac{d\omega}{2\pi} \tilde{G}_\omega(x, x') e^{-i\omega(t - t')}.$$ (28)

$G_\omega(x, x')$ will be an analytic function with poles located in the lower half plane, but with a vanishingly small negative imaginary part. For instance, for the limiting case $h(x) = \theta(L/2 - |x|)$, the poles are placed at $\omega_n = \pm \sqrt{\omega_p^2 + v^2\kappa_n^2 - \kappa_n}$, where $\kappa_n = (2n + 1)\pi/L$, $n = 1, 2, \ldots$. This pole structure implies that, for large $t$, the main contribution to Eq. (28) comes from large $t'$. Thus, the asymptotic behavior of $\Delta(t)$ can be obtained by inserting in Eq. (28) the asymptotic solution of Eq. (25), which in turn can be simplified to $\phi(x', t') = f(x') - kt$, as $x'$ is limited to a finite interval. Integrating in $t'$, we get the asymptotic $\Delta$ behavior proposed in Eq. (13).

We will now look for a time-dependent electric field $E(x, t)$ such that $\phi = \phi_a$ for all times. For this purpose, we extend $\phi_a$ backwards in time and try to check if the induced functions $\Delta(x, t)$ and $E(x, t)$ represent a sensible switching on process. For early enough times $t < t_-$, $f(x) - kt$ will be positive everywhere, implying that $\phi_a(x, t) = 0$. Inserting this result into Eq. (11) yields

$$E(x, t) = -\frac{i\pi h}{e} (C^*\Delta - C\Delta).$$

(29)

On the other hand, the retarded prescription in Eq. (24) shows that at $t < t_-$, the integrand should be evaluated at $t' < t < t_-$, i.e., when $\phi = 0$:

$$\Delta(x, t) = -\frac{C_0}{2} \int dx' \int_0^t dt' \ G(x; x', t - t') h(x').$$

(30)

We see that at early times $\Delta$ is time-independent. Since the retarded Green’s function is real, the phase of $\Delta$ is uniform and equal to that of $C^*$. Inserting Eq. (30) into Eq. (29), we conclude that the induced value of the electric field at early times ($t < t_-$) is zero, representing therefore a realistic switching on process.

VI. UNIVERSALITY OF FREE BOSONIZATION RULES AND LANDAUER CONDUCTANCE

We will now consider a generic finite system with current interactions, adiabatically connected to Fermi-liquid leads; we will follow the path integral approach to bosonization (see, e.g., Ref. 17), which is based on Fujikawa’s anomalous determinant.\cite{18}

Firstly, using the method of Ref. 19, we can easily check the universality of the free bosonization rules for our system. The fermionic partition function is

$$Z_F[A] = \int \bar{D}\psi D\tilde{\psi} \exp \left\{ i \int dx \left[ \bar{\psi} (i\partial_\mu - eA_\mu)\gamma^\mu \psi \right] + iI[j^\mu] \right\}.$$ (31)

The functional $I[j^\mu]$ denotes a current-dependent interaction term, which we suppose to be localized in $\Omega$. This interaction part can be represented in terms of a functional Fourier transform,

$$\exp \{ iI[j^\mu] \} = N \int \mathcal{D}a_\mu \exp \left\{ -i \int dx h(x)a_\mu j^\mu + iS[a_\mu] \right\}.$$ (32)

The constant $N$ is chosen such that $I[0] = 0$. Defining $h(x)$ as in the previous sections, we see that, by construction, $I[j^\mu] = 0$ for currents localized outside $\Omega$. As before, the role of the smooth function $h(x)$ is to implement the adiabatic contact of the interacting finite system to the noninteracting leads.

We can insert Eq. (32) into Eq. (31) and bosonize the fermions as in a free theory with an external source $eA_\mu + h(x)a_\mu$.
\[
Z_F[A] = \int \mathcal{D}\phi \mathcal{D}a_\mu \exp \left\{ i \int \tilde{d}^2 x \left[ \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - (eA_\mu + h a_\mu) \frac{1}{\sqrt{\pi}} \epsilon^{\mu\nu} \partial_\nu \phi \right] + iS[a_\mu] \right\}. \tag{33}
\]

Integrating over \(a_\mu\), we obtain the bosonized action
\[
S_{\text{boson}} = \int \! d^2 x \left( \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - eA_\mu \frac{1}{\sqrt{\pi}} \epsilon^{\mu\nu} \partial_\nu \phi \right) + i \left[ \frac{1}{\sqrt{\pi}} \epsilon^{\mu\nu} \partial_\nu \phi \right]. \tag{34}
\]

This is the universality of the free bosonization rules; it implies the universality of Landauer conductance at \(T = 0\) for a finite system with current interactions localized in the junction. Indeed, the bosonized equation of motion corresponding to Eq. (34) can be written as an anomalous divergence,
\[
\partial_\mu \left[ \partial^\mu \phi + \frac{1}{\sqrt{\pi}} \epsilon^{\mu\nu} \frac{\delta I}{\delta j^\nu(x)} \right] = -\frac{e\beta}{2\pi} E(x,t). \tag{35}
\]

Here, the asymptotic behavior will also have the form of Eq. (23), but with \(f(x)\) replacing \(\phi\) in Eq. (25) for a stationary electric field \(E(x)\). Notice that the contribution to the axial current coming from the interaction is localized. Therefore, the integration of Eq. (23) over the spatial coordinate, in analogy to Eq. (24), will lead to a perfect conductance again.

An analogous result is obtained if we consider a finite CDW system with an additional (local) current interaction of the form shown in Eq. (33). We can follow the same steps that led from Eq. (31) to Eq. (34). The only difference is that we must replace the massless free fermionic part in Eq. (31) by the corresponding CDW Lagrangian of Eq. (1) and, of course, include the path integration over the dynamical field \(\Delta\). The equivalence between the partition function for free fermions, with mass \(m\) and parameter \(\Delta(x)\), and the corresponding bosonized version can then be used to complete the steps showing the universality of the free bosonization rules. When the bosonized interaction term \(I[1/\sqrt{\pi} \epsilon^{\mu\nu} \partial_\nu \phi]\) is added to the bosonized CDW action [c.f. Eq. (2)], the anomalous chiral current will be modified by a term localized in the junction region. This kind of term, as we have seen, does not change the transport properties of the system. For instance, these conclusions hold when forward-scattering impurities are present in the CDW junction; if impurities are also present in the Fermi-liquid leads, however, some renormalization of the conductance is expected, in agreement with the results of Ref. [1].

**VII. SUMMARY**

We have shown that one-dimensional structures of the Fermi-liquid/finite-system/Fermi-liquid type display perfect Landauer conductance quantization at low temperatures, provided the finite system part presents an anomalous chiral symmetry. An important ingredient behind this behavior is the adiabaticity of the contacts between the finite system and the Fermi-liquid leads or reservoirs, which permits the extension of the chiral symmetry to the system as a whole. In this case, an anomalous chiral current is always present when a bias voltage is applied; outside the sample, this current is associated to the transport of free fermions. These general properties cause the charge transport through the system to be dominated by the reservoirs.

The natural language used to study this problems is bosonization. In 1+1 dimensions, it offers a simple way to take into account the chiral anomaly. In this framework, we have shown that the universality of Landauer conductance is intimately connected to the universality of the free bosonization rules.

The general relation between universality in transport properties of gapless systems and the existence of conserved chiral charges was first proposed in Ref. [2]. In that work, however, the role of finite sample size and the adiabaticity of the contacts was not apparent. We hope to have contributed to clarify this point with our discussion. In particular, we have studied a finite CDW system of the incommensurate type. The dynamical character of the phonon field is responsible for the restoration of the anomalous chiral symmetry, which is absent in the case of static phonon fields. This system, when adiabatically connected to reservoirs, presents a nonrenormalized, Fermi-liquid like conductance at low temperatures.

We estimate that this conductance quantization should be experimentally observable up to temperatures \(k_B T\) of order \(h\omega_p\) (the optical phonon frequency). For higher temperatures, fluctuations could destroy the CDW phase coherence. In that case, the conductance would involve a thermal activation mechanism, as in a band gap material. We have seen that the addition of forward scattering impurities to the CDW junction will not modify its perfect conductance. On the other hand, it is important to remark that, in practice, any strong inhomogeneities, leading to backscattering, will break global chiral symmetry. In a similar way, the pinning of the phonon field phase by impurities will cause the conductance to be suppressed. The effect of commensurability of the CDW wavelength with the underlying lattice constant will be discussed elsewhere.

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