Signature of directed chaos in the conductance of a nanowire

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We study the conductance of chaotic or disordered wires in a situation where equilibrium transport decomposes into biased diffusion and a counter-moving regular current. A possible realization is a semiconductor nanostructure with transversal magnetic field and suitably patterned surfaces. We find a non-trivial dependence of the conductance on the wire length. It differs qualitatively from Ohm’s law by the existence of a characteristic length scale and a finite saturation value.

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According to Ohm’s law, the resistance of a wire is proportional to its length. This is a straightforward consequence of the diffusive motion of electrons in the disordered potential of a normal material. However, unlike the time when Ohm arrived at his fundamental observation, conductors can be tailor-made today with almost complete control over the microscopic structure. Therefore it is important to understand the consequences of non-diffusive electron dynamics on the electronic conductance or other transport properties. This question has been studied in much detail for semiconductor nanostructures in which the motion of electrons is ballistic rather than diffusive. In such systems disorder is negligible and consequently all transport properties are determined by the shape of the sample, as in a billiard model. For example, in the ideal case of a perfectly clean nanowire with parallel walls the resistance should be zero independent of the length, and indeed this remarkable prediction has been confirmed experimentally. Beside ballistic systems, also the effects of anomalous diffusion on the electronic or thermal conduction properties have attracted a lot of attention.

In the present paper we study the electronic conductance of a wire in the case of a different and very profound modification of the microscopic dynamics. We consider systems where directed chaos leads to biased diffusion in the absence of any potential gradient. Directed chaos means that the time-averaged velocity of chaotic trajectories is non-zero due to broken time-reversal symmetry and due to the specific phase space structure. This effect may occur in various types of systems including, e.g., cold atoms in suitably pulsed optical potentials or chains of electronic billiards in a transversal magnetic field. It has been investigated both, theoretically and experimentally, in a number of recent publications. The interest is due to some intriguing and potentially very useful properties. For example, quantities like velocity average, velocity dispersion or scattering delay times are intrinsically dependent on the transport direction. However, all previous studies focused on the ratchet-like directed transport in effectively infinite periodic systems with directed chaos. In contrast, we address here for the first time the typical electronic setup of a finite sample which is coupled to two electron reservoirs. We show that directed chaos has profound consequences also in this context where the transport velocity is not directly measurable; instead the conductance becomes the most basic and most relevant quantity. Due to biased diffusion, a new length scale \( \lambda_{\text{ch}} \) appears and rules the asymptotic decay of the conductance with the sample length, see Eq. (10) below. Moreover, chaotic trajectories can propagate through samples of arbitrary length and consequently the conductance approaches a non-zero constant given in Eq. (9). Note that it is not possible to reduce the description of directed chaos to a diffusion equation with bias. In our explicit result for the conductance, Eq. (17), we must account also for the detailed structure of the underlying mixed phase space.

To be specific, we investigate the prototypical model first introduced in [12]. We consider a two-dimensional electron gas (2DEG) confined to a quasi one-dimensional channel (Fig. 1a). One wall of the channel is straight. Electrons are specularly reflected, but no back scattering occurs in the transport direction. The other wall has a rough surface causing strong and essentially random scattering. The details of this roughness are not crucial (see below). Directed chaos is induced by a perpendicular magnetic field which breaks time-reversal invariance. We stress that a realization of our model does not require more than a novel combination of elements which are all well understood and experimentally accessible in the context of a mesoscopic 2DEG, namely transversal magnetic fields of moderate strength, negligible bulk scattering and surfaces which are either disordered or manufactured with a precisely defined geometry.

For our purpose it is sufficient to treat the electrons as independent classical particles, i.e., we assume a semi-classical regime with many transversal modes in the channel, \( N \sim m^*v_F b/h \gg 1 \). We use dimensionless units in which channel width and Fermi velocity are unity, \( b = v_F = 1 \). The strength of the magnetic field is parameterized by the cyclotron radius \( r_c = m^*v_F / e B \). Besides the length of the sample this is the only free parameter in our problem. We will assume that the magnetic field is
ordered surface with a correlation length that is below
served. A physical realization of this behavior is a dis-
is a periodic array of semicircular scatterers with small
shown in the insets of Fig. 1a. In one case the surface
preserved. In [12] we considered the two extreme cases
in the second case the direction
from one direction, say to the left. On the average,
chaotic or random orbits which are reflected from both
walls (Fig. 1a). The regular orbits are transporting con-
triparallel walls with different surfaces. The upper wall re-
fl ects specularly while at the lower wall the reflection angle is
essentially random. Possible physical realizations of this ran-
domness are magnified in the insets (see text). Typical tra-
jectories are either regular (dashed line) or random (full line)
and transport in opposite directions. In (b) the transversal
Poincaré sections for \( v_x > 0 \) and \( v_x < 0 \) are displayed.
not too strong, \( r_c \geq 1 \). In this case there are no pinned
orbits in the bulk of the channel and the phase space contains
only two types of electron trajectories: regular orbits skipping along the clean channel boundary, and
chaotic or random orbits which are reflected from both
walls (Fig. 1a). The regular orbits are transporting con-
tinuously in one direction, say to the left. On the average,
the chaotic orbits are transporting in the opposite direction,
thus compensating for the regular transport and
making the system unbiased as a whole. However, the
transport velocity \( \dot{x} \) of a chaotic electron is fluctuating,
in fact the dynamics is diffusive with a superimposed drift
along the channel. The average drift velocity can be ob-
tained by application of a phase space sum rule [3]. For
our model we find for the long-time velocity average of
almost all chaotic trajectories
\[
v_{\text{ch}} = \frac{1}{2} \frac{\theta - \sin \theta \cos \theta}{\pi(1 - \cos \theta) - \sin \theta + \theta \cos \theta}
\] (1)

with \( \theta(r_c) = \arccos(1 - 1/r_c) \) [12]. This result is inde-
pendent of the precise modelling of the rough channel
surface as long as the phase space structure of Fig. 1b is
preserved. In [12] we considered the two extreme cases
shown in the insets of Fig. 1a. In one case the surface is
a periodic array of semicircular scatterers with small
radius \( R \to 0 \) [14], i.e., the dynamics of the system is
deterministic and there is no disorder whatsoever. In
the second case the direction \( \varphi \) of the trajectory was randomized upon every scattering from the rough
surface. Specifically it was chosen with probability density
\( P(\varphi) = \frac{1}{2} \sin \varphi \) from the interval \([0, \pi]\) such that the
invariant measure on the energy shell \( dx \, dy \, d\varphi \) was pre-
served. A physical realization of this behavior is a dis-
ordered surface with a correlation length that is below
the Fermi wavelength. While this second system is non
deterministic when approximated classically, its trans-
port properties are essentially the same as for the case of
deterministic directed chaos. In particular, in both
cases an ensemble of chaotic trajectories spreads diffu-
sively around the moving center of mass, \( \langle \Delta x^2 \rangle = D_{\text{ch}} t \)
[12]. The precise value of the diffusion constant depends
on the detailed modeling of the rough boundary. Ana-
lytical results for \( D_{\text{ch}} \) are available in the case of random
scattering [12] and therefore we shall use this version of
the model in the numerical calculations below.

The electronic conductance is obtained within the
framework of the Landauer-Büttiker formula [1], \( G(L) =
(2e^2/h) T(L) \). We approximate the transmission semi-
classically, \( T \sim N(t) \), and discuss quantum corrections
at the end of this letter. \( t(L) \) denotes the total classi-
cal probability that an electron is transmitted through a
sample of length \( L \) if it enters the system at \( x = 0 \) from
the left with random initial conditions \( P(y, \varphi) = \frac{1}{2} \cos \varphi
(y \in [0, 1], \varphi \in [-\frac{\pi}{2}, \frac{\pi}{2}]) \). According to our convention
(regular orbits are skipping to the left) these initial con-
ditions are within the chaotic component of phase space.
Similarly, we define the probability \( t'(L) \) that an electron
is transmitted if it enters at \( x = L \) from the right with
an analogous distribution (but \( \varphi \in [\frac{\pi}{2}, \frac{3\pi}{2}] \)). The total
transmission probability must be the same for the two
distinct transport directions,
\[
t(L) = t'(L).
\] (2)

There are various ways to arrive at this fundamental iden-
tity. For example, one observes that \( t \neq t' \) would result
in the accumulation of particles in one of the reservoirs
even when the system is in thermal equilibrium. Alterna-
tively, a microscopic derivation can be based on the fact
that the scattering map of Hamiltonian systems is area
preserving. Although Eq. (2) is not specific for directed
chaos, this identity has very interesting consequences in
the present context. While \( t(L) \) is entirely due to chaotic
trajectories whose properties are not immediately acces-
sible, \( t'(L) \) can be decomposed into conditional probabili-
ties for regular and chaotic trajectories,
\[
t'(L) = \mu_{\text{reg}} t'_{\text{reg}}(L) + \mu_{\text{ch}} t'_{\text{ch}}(L),
\] (3)
where
\[
\mu_{\text{reg}} = \frac{1}{2} \int_{\text{reg}} d\varphi \cos \varphi = \frac{1}{2} \frac{\theta - \sin \theta \cos \theta}{1 - \cos \theta}
\] (4)
is the the relative area of the regular component in the
transversal Poincaré section (Fig. 1b), and \( \mu_{\text{ch}} = 1 - \mu_{\text{reg}} \).
Due to the lack of back scattering along the regular tra-
jectories we have \( t'_{\text{reg}}(L) \equiv 1 \). Moreover, it is clear that
\( t'_{\text{ch}} \to 0 \) for \( L \to \infty \). This is so because for trajectories
which are moving through arbitrarily long samples against the average chaotic flow the time-averaged ve-
locity cannot converge to \( v_{\text{ch}} \). Hence these trajectories

FIG. 1: (a) A prototypical example for directed chaos is a
2DEG in a perpendicular magnetic field which is confined by
two parallel walls with different surfaces. The upper wall re-
fl ects specularly while at the lower wall the reflection angle is
essentially random. Possible physical realizations of this ran-
domness are magnified in the insets (see text). Typical tra-
jectories are either regular (dashed line) or random (full line)
and transport in opposite directions. (b) The transversal
Poincaré sections for \( v_x > 0 \) and \( v_x < 0 \) are displayed.
must be of measure zero in phase space. Taken together, the mentioned facts yield a remarkable result which is illustrated in the right inset of Fig. 2 for long systems: the probability that a chaotic trajectory transmits from \( x = 0 \) to \( x = L \) is given by the relative phase space area occupied by the counter-moving regular trajectories,

\[
t(\tau) \to \mu_{\text{reg}} \quad (\tau \to \infty). 
\]  

The goal is now to understand quantitatively how \( t_{ch}'(L) \) decays from \( t_{ch}'(0) = 1 \) to zero as the length of the system increases. The results of numerical simulations for a wide range of values \( r_c \) are shown in Fig. 2. The data suggest an exponential behavior for long wires. This can be understood after replacing the microscopic dynamics by the Fokker-Planck equation (FPE) for biased diffusion. At \( x = 0 \) and \( x = L \) we assume absorbing boundary conditions. The probabilities to reach these exits from a point \( 0 \leq x \leq L \) within the sample are then

\[
p_l'(x) = \frac{e^{-x/\lambda_{ch}} - e^{-L/\lambda_{ch}}}{1 - e^{-L/\lambda_{ch}}}, 
\]

\[
p_l(x) = \frac{1 - e^{-x/\lambda_{ch}}}{1 - e^{-L/\lambda_{ch}}}, 
\]

respectively [13]. Here

\[
\lambda_{ch} = D_{ch}/v_{ch} 
\]

denotes the Peclet length of the biased diffusion process.

In order to contribute to the transmission from the right to the left end of a wire of length \( L \), a chaotic particle should first be transmitted through a segment of length \( l < L \) and then, starting from \( x = L - l \), be absorbed at \( x = 0 \). Based on this argument we propose as a recursion relation for the chaotic transmission

\[
t_{ch}'(L) = t_{ch}'(l)p_l'(L - l). 
\]

Note that this relation cannot be valid for arbitrary \( l \). For example, \( l = 0 \) leads to \( t_{ch}'(L) = 0 \) as for a diffusing particle starting at one of the absorbing boundaries the probability to reach the opposite end is identically zero. The reason for this limitation can be understood as follows. Upon replacing the original dynamical system by a 1D FPE we have discarded all information about the momentum of the electron. Strictly speaking it is then impossible even to define a transmission probability since this requires to enter the system with a given direction. Hence, for Eq. (9) to be valid, the length scale for momentum correlations should be negligible compared to the distance from the boundaries, \( \lambda_{ch} \ll x \) and \( \lambda_{ch} \ll l \). Under this assumption we find to leading order \( t_{ch}'(L) = t_{ch}'(l) e^{-(L-l)/\lambda_{ch}} \) with the solution

\[
t_{ch}'(L) = c \exp(-L/\lambda_{ch}). 
\]

The dashed line in Fig. 2 shows this exponential for \( r_c = 50 \) (with suitably chosen prefactor). Indeed the asymptotic decay of the transmission probability is reproduced. However, for short systems the behavior is clearly not exponential. Moreover, even in the asymptotic regime \( L \to \infty \) the FPE approach is not accurate enough to predict the prefactor \( c \) of the exponential decay [17]. This is no surprise. According to Eq. (10), a different prefactor corresponds to an additive constant in the system length \( L \) and this type of error must be expected after discarding the correlation length of the momentum.

![FIG. 2: The chaotic transmission (+) is compared to Eq. (10) for \( r_c = 200, 50, 10, 6, 3, 2 \) (top-bottom). Each data point represents \( 10^6 \) trajectories. For \( r_c = 50 \) the asymptotic exponential is shown with a dashed line. Left inset: Dependence of the Peclet length on the cyclotron radius. Right inset: For \( r_c = 10 \) the total transmission (●) is compared to Eq. (19) and to the asymptotic constant \( \mu_{\text{reg}} = 0.294 \).](image_url)

We conclude that a satisfactory theory for \( t_{ch}'(L) \) cannot be based on the FPE alone. As an alternative model let us consider a generalization of the persistent random walk [14] to a biased walk (BPRW): A particle moves with velocity \( \pm v_F \) and is reflected in direction-dependent obstacles at distances \( L_1 \). For a segment of length \( L = nL_1 \) we define the transmission and reflection probabilities by \( t_n = 1 - r_n = t(nL_1) \) (left to right) and \( t'_n = 1 - r'_n = t'_n(nL_1) \) (right to left).

A multiple-scattering expansion allows to express these probabilities in terms of a single segment. We find

\[
t_n = \frac{1 - r'_1}{(t'_1/t_1)^n - r_1'/r_1}, 
\]

\[
t'_n = \frac{1 - r_1}{(t_1/t'_1)^n - r_1/r'_1}. 
\]

In our case the chaotic transport is biased to the right. Therefore we assume \( t_1 > t'_1 \) and find in the limit \( n \to \infty \)

\[
t_\infty = 1 - r_1/r'_1 
\]

and

\[
t'_n = t_\infty (t_1/t'_1)^{-n} \quad (n \to \infty). 
\]
However, there is no direct connection between the parameters of the BPRW and the dynamics of our original model. In order to close this gap we must make use of the information about the underlying phase space structure, Eq. [3], and the result obtained within the FPE approach, Eq. [10]. The former implies \( t_\infty = \mu_{\text{reg}} \) or equivalently \( r_1/r'_1 = \mu_{\text{ch}} \). Further the comparison of Eqs. [10] and [12] yields

\[
\begin{align*}
    c \exp(-nL_1/\lambda_{\text{ch}}) &= \mu_{\text{reg}} (t_1/t'_1)^{-n}, \\
\end{align*}
\]  

Now it is easy to read off the correct prefactor of the asymptotic exponential decay, \( c = \mu_{\text{reg}} \), which was used for the dashed line in Fig. 2. On the other hand we infer \( \exp(L_1/\lambda_{\text{ch}}) = t_1/t'_1 \) which is substituted into Eqs. [11] and [12]. The final result for the chaotic transmission probability from right to left is then

\[
\begin{align*}
    t'_\text{ch}(L) = & \frac{\mu_{\text{reg}}}{\exp(L/\lambda_{\text{ch}})} - \mu_{\text{ch}}, \\
\end{align*}
\] 

while the total transmission is given by

\[
\begin{align*}
    t(L) = \frac{\mu_{\text{reg}}}{1 - \mu_{\text{ch}} \exp(-L/\lambda_{\text{ch}})}. \\
\end{align*}
\] 

These two contradicting statements can be reconciled as follows. In the presence of directed chaos the localization length \( \xi \) diverges exponentially with the number of transversal modes \( (\ln \xi \sim N \sim h^{-1}) \) [13]. Thus, even for moderate \( N \), the localization length will easily exceed the sample length \( L \) or the coherence length of the given material. In this regime we can safely ignore localization. Other quantum effects like weak localization or tunneling are not expected to change our results qualitatively although they may lead to small corrections \( \sim N^{-1} \) in the relevant parameters \( \lambda \) and \( \mu_{\text{reg}} \). A numerical analysis of such effects will be attempted elsewhere.

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[16] Finite \( R \) results in a correction to the first term in the denominator of Eq. [4], see [12].
[17] Extrapolation of Eq. [10] to \( L = 0 \) yields \( c = 1 \). It is also possible to repeat the argumentation leading to Eq. [10] for \( t(L) \). Using Eq. [13], one finds in this way \( c = \mu_{\text{reg}}/\mu_{\text{ch}} \). Both results disagree with the correct value \( c = \mu_{\text{reg}} \).
[18] For the BPRW we derive explicit expressions for transport velocity \( v \) and diffusion constant \( D \). Within the FPE we must then expect a decay of the transmission probability with length scale \( \lambda = D/v \). In general this disagrees with the value \( \lambda = L_1/\ln(t_1/t'_1) \) obtained from Eq. [13]. A continuum limit \( (L_1, r_1, r'_1 \to 0 \) for fixed \( t_\infty \) is not sufficient to remove this discrepancy.