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Transient and Sharvin resistances of Luttinger liquids

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Although the intrinsic conductance of an interacting one-dimensional system is renormalized by the electron-electron correlations, it has been known for some time that this renormalization is washed out by the presence of the (non-interacting) electrodes to which the wire is connected. Here, we study the transient conductance of such a wire: a finite voltage bias is suddenly applied across the wire and we measure the current before it has enough time to reach its stationary value. These calculations allow us to extract the Sharvin (contact) resistance of Luttinger and Fermi liquids. In particular, we find that a perfect junction between a Fermi liquid electrode and a Luttinger liquid electrode is characterized by a contact resistance that consists of half the quantum of conductance in series with half the intrinsic resistance of an infinite Luttinger liquid. These results were obtained using two different methods: a dynamical Hartree-Fock approach and a self-consistent Boltzmann approach. Although these methods are formally approximate we find a perfect match with the exact results of Luttinger/Fermi liquid theory.

I. INTRODUCTION

The theory of quantum transport describes how a quantum device is connected to the voltage sources and measuring apparatus of the macroscopic world¹. This description has not always been obvious, and the correct formula for the conductance of a quantum circuit was strongly debated at the beginning of quantum nanoelectronics (formerly mesoscopic physics). In its original paper², Landauer identified the conductance with $g_I = g_F T / (1 - T)$ where $g_F = e^2/h$ is the quantum of conductance and T is the probability for an electron to be transmitted through the device (in one dimension, ignoring spin). It took several years for the correct expression $g = g_F T$ —now known as the Landauer formula— to emerge and it was not established until the seminal experiment of Wees *et al.*³ which measured the quantification of conductance in unit of g_F in a point-contact geometry. The expression $g = g_F T$ for the conductance is in itself rather spectacular: it predicts that even in the total absence of scattering ($T = 1$, perfect transmission), a quantum circuit has a finite resistance $1/g_F$, hence that some energy must be dissipated through Joule effect. The solution to this paradox comes from the concept of Sharvin —or contact— resistance: The expression $g = g_F T$ should be understood as

$$\frac{1}{g} = \frac{1}{2g_F} + \frac{1}{g_I} + \frac{1}{2g_F}, \quad (1)$$

i.e. the measured resistance $1/g$ corresponds to an intrinsic resistance of the circuit $1/g_I$ in series with the universal Sharvin resistance $1/(2g_F)$ of each contact. More recently, the resistance entering in the RC relaxation time of a quantum capacitor (a situation with one unique electrode) has been measured explicitly⁴ and corresponds to the Sharvin resistance $1/(2g_F)$ ⁵. The purpose of this article is to generalize the concept of Sharvin resistance, now well understood in Fermi liquid electrodes, to interacting one-dimensional electronic systems, i.e. Luttinger

liquids⁶.

Luttinger liquids (L) can be understood formally using the bosonization formalism⁶⁻⁹. In a simple physical picture, the excitation spectrum of a Luttinger liquid is made of plasmons¹⁰, ripples of the Fermi sea, that obey bosonic statistics. These plasmons have a linear dispersion relation $\omega = v_L k$, where the Luttinger liquid or plasmon velocity v_L is renormalized with respect to the Fermi velocity v_F . Likewise, the conductance g_L also gets renormalized^{11,12} with respect to the Fermi liquid (F) result g_F ,

$$\frac{v_L}{v_F} = \sqrt{1 + \frac{U}{\pi v_F}}, \quad (2a)$$

$$\frac{g_L}{g_F} = \frac{v_F}{v_L}, \quad (2b)$$

where U characterizes the strength of the electron-electron interaction^{12,13}. It was soon understood^{14,15}, however, that the intrinsic renormalized conductance of the Luttinger liquid could not be observed directly: in presence of Fermi liquid electrodes, the effect of interaction is washed out and one recovers the Fermi liquid result g_F .

In this article, we extend these results to interacting electrodes. We study the —transient— response of a Luttinger liquid connected to two Fermi liquid reservoirs when a voltage bias is abruptly switched on across the wire. This technique provides a shortcut for the study of the Sharvin resistance: at short times, the system does not yet know about the presence of the Fermi liquid electrodes, so that one observes a plateau that corresponds to the conductance between two Luttinger liquids. At longer times one should recover the non-interacting Fermi liquid conductance^{14,15}. The interest in the transient response of Luttinger liquids have been mostly theoretical so far¹⁶⁻²², with a focus on the phenomena on spin-charge separation. However, the recent experimental progress in manipulating quantum nanoelectronic circuit at high frequencies (10 GHz and beyond) while retaining low tem-

peratures (a few tens of a mK) now make this type of measurement feasible in the lab^{23–28}. Experiments at terahertz frequencies are also becoming possible^{29,30}.

II. MODEL

We consider an infinite (quasi-) one-dimensional wire described by a Hamiltonian of the form $\hat{H}(t) = \hat{H}_0 + \hat{H}_b(t) + \hat{H}_{\text{int}}$. $\hat{H}_0 = \sum_{\langle ij \rangle, \sigma} \gamma_{ij} c_{i\sigma}^\dagger c_{j\sigma}$ contains nearest neighbor hoppings $\gamma_{ij} = -1$, while $c_{i\sigma}^\dagger$ and $c_{i\sigma}$ are the creation and destruction operators on site i with spin σ . $\hat{H}_b(t) = \sum_{i\sigma} V_b(t) \theta(i_b - i) c_{i\sigma}^\dagger c_{i\sigma}$ corresponds to a time dependent bias voltage $V_b(t)$ where the potential drop happens between site i_b and $i_b + 1$ ($\theta(x)$ is the Heaviside function). We will consider various shapes for the voltage pulse $V_b(t)$ in the regime where the voltage varies quickly with respect to the propagation time through the interacting region, yet slow with respect to the bandwidth of the model (in order to avoid spurious effects associated with energy scales comparable with the Fermi energy E_F). At time $t < 0$ the system is at equilibrium and at zero temperature. The interacting term is of the Hubbard form $\hat{H}_{\text{int}} = U \sum_i s(i) (c_{i\uparrow}^\dagger c_{i\uparrow} - n_0) (c_{i\downarrow}^\dagger c_{i\downarrow} - n_0)$ where the function $s(i) = (\tanh[(i - i_L)/d] - \tanh[(i - i_R)/d])/2$ characterizes the transition between the central interacting region $[i_L, i_R]$ and the non-interacting electrodes over a width d , while n_0 is the equilibrium density. A sketch of the system is shown in Fig. 1. We consider two (formally) approximate methods to study this problem: a time dependent Hartree-Fock approach which is equivalent to the Random Phase Approximation, hence already known to capture the salient features of the one dimensional plasmons⁶ and a much simpler self-consistent Boltzmann approach that provides analytical expressions in a number of situations.

III. TIME DEPENDENT HARTREE-FOCK APPROACH

The first method is a time dependent mean-field approach where \hat{H}_{int} is replaced by its mean-field value \hat{H}_{HF} at time t . The problem becomes effectively diagonal in spin, so that from now on we ignore the spin degree of freedom and consider the spinless problem (the spinfull conductance can be recovered by multiplying our results by a factor 2). Note that this time dependent Hartree-Fock approach is very different from its static counterpart; in particular it captures screening effects at the Random Phase Approximation level³¹. The interaction term takes the form $\hat{H}_{\text{HF}} = U \sum_i s(i) c_i^\dagger c_i [\langle c_i^\dagger(t) c_i(t) \rangle - n_0]$, where the notation $c_i(t)$, $c_i^\dagger(t)$ refers to the Heisenberg representation. We solve this problem using the method developed in Refs. [32–34] within the Keldysh formalism which is implemented as a top layer of the Kwant package³⁵. It

amounts to solving a set of one-particle Schrödinger equations

$$i\partial_t \psi_{\alpha E}(i, t) = \sum_j H_{ij}(t) \psi_{\alpha E}(j, t), \quad (3)$$

with $H_{ij} = \gamma_{ij} + V_b(t) \theta(i_b - i) + U s(i) (n(i, t) - n_0)$ and the initial condition $\psi_{\alpha E}(i, t = 0) = \psi_{\alpha E}^{\text{st}}(i)$ where $\psi_{\alpha E}^{\text{st}}(i)$ is a scattering state α at energy E of the time-independent Hamiltonian \hat{H}_0 at $t = 0$. The local density of electrons $n(i, t)$ is given by,

$$n(i, t) = \sum_\alpha \int \frac{dE}{2\pi} f_0(E) |\psi_{\alpha E}(i, t)|^2 \quad (4)$$

where f_0 is the Fermi function. The new twist with respect to Refs. [32–34] is the presence of the self-consistent term $U s(i) (n(i, t) - n_0)$ which transforms the set of initially independent Schrödinger equations into coupled ones. Since up to 1000 different energy values are necessary to discretize the integral of Eq. (4), such a set of coupled partial integro-differential equations is in general intractable. We leverage the fact that the dynamics of Eq. (4) is much slower than the one of Eq. (3) to develop a doubly adaptative timestepping scheme: we construct a linear extrapolation of $n(i, t)$ which is used by a Runge-Kutta method with Dormand and Prince stepsize control³⁶ to integrate the one-particle Schrödinger equations (3). At a second larger (and also adaptive) time step, the integral Eq. (4) is recalculated and the extrapolation of $n(i, t)$ is updated. In all the simulations below, the overall error in the calculated observables is of the order 10^{-4} .

Figure 1 shows a typical simulation for the propagation of a voltage pulse through the (blue) interacting region. The left panel shows two simulated charge densities $n(i, t)$, one with interactions (solid line) and one without interactions (dashed line). The injected Gaussian pulse $V_b(t) = V_b e^{-(t-t_0)^2/2\tau^2}$ propagates faster in the interacting case. The right panel of Fig. 1 shows the dependence of the plasmon velocity v_L on the interaction strength U for three different Fermi energies. We find a perfect match between the numerical data (symbols, extracted from a linear fit of the numerical simulations as shown in the blue lines of the left panel) and the Luttinger-liquid prediction Eq. (2a) (solid lines). The fact that the dispersion relation of the Luttinger liquid is perfectly reproduced by a mean-field approach is in itself nontrivial⁶.

We note that Equation (2) together with the Luttinger liquid approach of the one dimensional Hubbard model is only valid for moderate interactions. When U becomes of the order of the bandwidth (e.g. $U \sim 2$ in our case), the plasmon physical picture ceases to be relevant and one needs to resort to other approaches such as the Bethe ansatz^{37,38}. Likewise, the results presented here do not apply in the strongly interacting limit.

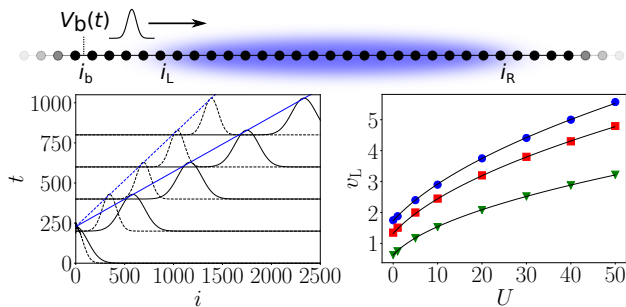


FIG. 1. (Color online) Top panel: sketch of the system under consideration including the position of the voltage drop (electrode/wire interface) i_b and interacting blue region. Transport simulation results for the Plasmon velocity v_P in one-dimensional quantum wires. Left panel: electronic density $n(i, t)$ (in arbitrary units) as a function of site i for different values of time t (as indicated on the y-axis) after injection of a Gaussian voltage pulse of width $\tau = 100$. The interacting region is $|i_R - i_L| = 1000$ sites wide. Solid lines: $U = 10$, dashed lines: $U = 0$ (no interaction). The blue lines are linear fits from which we extract the velocity. Right panel: extracted Plasmon velocity v_L vs. interaction strength U for three different Fermi energies $E_F = -1, -1.5$ and -1.9 . Symbols are from the simulations and lines correspond to Eq. (2a).

A. Transient conductances

We now turn to the transient conductances of our system: at $t = 0$ we abruptly raise the voltage bias across the wire from zero to V_b and measure the current I flowing through the system. The results are presented in Fig. 2 and Figure 3 where solid (resp. dashed) lines mark the result with (without) interaction. The solid line of Fig. 2 shows the current I_1 inside the interacting region as a function of time after a voltage drop in the interacting region. We observe a well developed plateau until the pulse reaches the boundary of the interacting region, followed by a second, stationary, plateau. The duration of the first plateau is simply given by the propagation time $|i_R - i_L|/v_P$ in the interacting region. We define the (dimensionless) transient conductance g as I/I_F , where the current is evaluated on a plateau and $I_F = V_b/2\pi$ is the non-interacting value. Note that since the system is not stationary, the transient conductance varies inside the system, in particular depending on the measurement being performed inside the wire versus inside the electrodes. The value of the transient conductance is illuminating: the short time value is given by the Luttinger liquid theory Eq. (2b) while the long time (stationary) value is given by the Fermi liquid value $I_1/I_F = 1$, see the inset of Fig. 2. It is by itself rather remarkable that our mean-field approach captures these two values exactly. The transient conductance I_2/I_F measured inside the electrode is however very different: it is given by a new formula

$$\frac{1}{g_T} = \frac{1}{2} \left[\frac{1}{g_L} + \frac{1}{g_F} \right], \quad (5)$$

which corresponds to half a quantum of resistance in series with half the Luttinger liquid intrinsic resistance. This is our first hint that $1/(2g_L)$ should be understood as the Sharvin resistance of a Luttinger liquid. Figure 3 is very similar to Fig. 2 but the voltage drop happens at the electrode-wire interface. There, only the T transient conductance is observed. In reality the position of the drop is determined by the interplay between the electrostatic capacitance and the quantum capacitance of the circuit (see [32] for a discussion). In practice, the drop is most likely to happen at the electrode, making Fig. 3 the experimentally most relevant case.

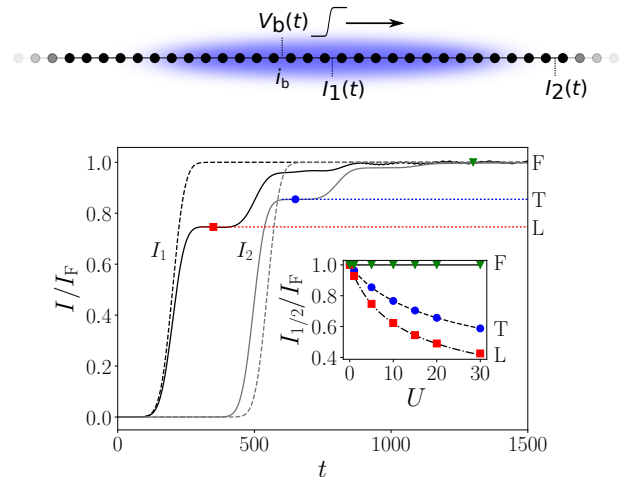


FIG. 2. (Color online) Numerical simulation of the current I as a function of time t in a one-dimensional quantum wire in presence (solid lines, $U = 5$) and absence (dashed lines, $U = 0$) of interactions. The electrical potential V_b is switched on smoothly at time $t_0 = 200$ and stays constant after. The subscript on $I_{1/2}$ indicates the position where the current is actually measured, as shown in the sketch. Inset: transient conductances as a function of U , see text (symbols) and theoretical predictions (dashed line: transient Eq. (5), dashed dotted: Luttinger liquid Eq. (2b) and solid line: Fermi liquid).

IV. SELF CONSISTENT BOLTZMANN APPROACH

To proceed, we introduce a second, even simpler approach, in the spirit of the classical theory of surface plasmons¹⁰. We introduce the semi-classical probability distribution $f(x, k, t)$ for an electron to be at position x (we use x instead of i as we will be working in the continuum limit) and momentum k at time t . In our ballistic sample, $f(x, k, t)$ satisfies a collisionless Boltzmann equation³⁹,

$$\partial_t f = -v_k \partial_x f - F(x, t) \partial_k f, \quad (6)$$

where $v_k = dE/dk$ is the non-interacting velocity corresponding to the dispersion relation $E(k) = -2 \cos(k)$

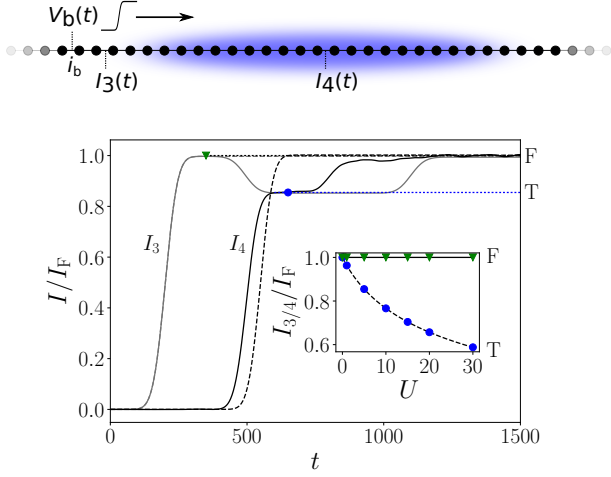


FIG. 3. (Color online) Same as Fig. 2 but the voltage drop happens at the electrode wire interface.

of \hat{H}_0 . The classical force $F(x, t)$ arises from the electric field of the bias potential $V_b(x, t)$ and from the field generated by the pulse itself through Coulomb repulsion $F(x, t) = \partial_x V_b(x, t) + U \partial_x [s(x)n(x, t)]$. The density $n(x, t)$ is obtained through

$$n(x, t) = \int_{-\pi}^{\pi} \frac{dk}{2\pi} f(x, k, t) \quad (7)$$

and we arrive at a close set of equations.

A. General solution

Interestingly, the above self-consistent Boltzmann equation reproduces exactly the plasmon velocity and conductance of the Luttinger liquid Eq. (2a) and (2b) as well as the transient regime Eq. (5) that will be discussed below. The low energy spectrum of the Boltzmann equation can be obtained by expanding $f(x, k, t) = f_0[E(k)] + f_1(x, k, t)$ to first order in f_1 around the Fermi function f_0 . After integration over k , we introduce the density of right $n_+(x, t) = \int_0^{\pi} \frac{dk}{2\pi} f_1$ and left $n_-(x, t) = \int_{-\pi}^0 \frac{dk}{2\pi} f_1$ movers and arrive at

$$\partial_t n_{\pm} \pm \partial_x \left[v_F n_{\pm} + \frac{U}{2\pi} s(x) n_1 - \frac{V_b}{2\pi} \right] = 0, \quad (8)$$

with $n_1 = n_+ + n_-$. Equation (8) has the form of a Liouville equation, hence describes ballistic propagation of different modes. In particular for a constant interaction $s(x) = 1$, it can be solved exactly and its general solution reads (in terms of the initial condition at $t = 0$),

$$\begin{aligned} n_{\pm}(x, t) = & \frac{1}{4v_F v_L} \sum_{\eta=\pm 1} \eta \left[(v_F + \eta v_L)^2 n_{\pm}(x \mp \eta v_L t, t=0) \right. \\ & - (v_F^2 - v_L^2) n_{\mp}(x \mp \eta v_L t, t=0) \\ & \left. + 2v_F (v_F \pm \eta v_L) \int_0^t \frac{dt'}{2\pi} \frac{\partial}{\partial x} V_b(x - \eta v_L(t-t'), t') \right] \end{aligned} \quad (9)$$

and indeed propagates with the plasmon velocity v_L given by Eq. (2a). This theory can be generalized to the multi-channel situation, and one recovers the same theory as what can be derived from the bosonization approach^{12,40}. The multi-channel theory has recently been verified experimentally in a wave guide geometry⁴¹ and is consistent with experiments in the integer quantum Hall regime²⁸.

B. Derivation of the transient conductances from the Boltzmann approach

Let us now derive the values of the transient conductances from the Boltzmann approach. In a transient regime where the densities $n_+(x)$ and $n_-(x)$ are *locally* stationary (they may still propagate in a remote region), Eq. (6) admit two constants of motion, the current I and the renormalized density N ,

$$I = v_F [n_+(x) - n_-(x)] \quad \forall x \quad (10a)$$

$$N = \left(1 + \frac{U s(x)}{\pi} \right) [n_+(x) + n_-(x)] \quad \forall x \quad (10b)$$

Let us now consider the general situation of a junction between one Luttinger liquid on the left [characterized by an interaction parameter U_1 and the corresponding conductance g_{L1} , see Eq. (2)] in contact with a second Luttinger liquid on the right (U_2 and g_{L2}). At $t = 0$, $n_+(x, t=0) = n_-(x, t=0) = 0$ and we suddenly raise a sharp potential step $V_b(x) = V_b \theta(x_1 - x) \theta(t)$ with x_1 deep in the left electrode. Deep in the left and right electrode (where $s(x) = \text{constant}$), we can use Eq. (9) and obtain the structure of the solution; in a second step, these solutions are matched using Eqs. (10a) and (10b). After some algebra, we arrive at $I = g V_b$ with,

$$\frac{1}{g} = \frac{1}{2g_{L1}} + \frac{1}{2g_{L2}}, \quad (11)$$

which is our chief analytical result. Eq. (11) calls for a number of comments.

(i) First, it is valid for an arbitrary function $s(x)$, i.e. the conductance only depends on the nature of the left and right electrodes and *not* on the intermediate region.

(ii) In particular, in the case of a Luttinger liquid sandwiched in between non-interacting electrodes [$s(0) = 1$, $s(\pm\infty) = 0$], one recovers the established fact that the electron-electron interaction does not renormalize the conductance^{14,15}.

(iii) For an infinitely long and homogeneously interacting wire [$s(x) = 1$], we recover the Luttinger liquid result $g = g_L$.

(iv) For a non-interacting electrode in contact with an interacting one, Eq. (11) reduces to Eq. (5).

V. DISCUSSION AND OUTLOOK

Equation (11) is strikingly similar to Eq. (1) that defines the Sharvin resistance of a non-interacting elec-

trode. In fact, Eq. (11) can be interpreted as the generalization of the concept of Sharvin resistance to Luttinger-liquid electrodes, attributing the contact resistance $1/(2g_L)$ to each electrode. We have studied the perfectly transmitting situation (e.g. the absence of impurities that could lead to scattering) that corresponds to the absence of intrinsic resistance. The generalization to the presence of impurities leads to power law behaviors in the $I(V_b)$ which is beyond the scope of the Boltzmann approach^{9,42}. The concept of Sharvin resistance, however, can be used beyond the present situation. Indeed, what defines an electrode in practice depends on where the energy relaxation takes place; an electrode is essentially any system whose size is comparable to its energy relaxation length. Hence, we predict that the renormalization of the conductance Eq. (11) should be amenable to observation in d.c. experiments involving interacting wires of sufficient length. For instance, a geometry close to the one used in [41] should allow for the determination of the interacting Sharvin resistance with $U \sim d/a_B$ (d distance to the gate, a_B effective Bohr radius). The above results might also be measured directly in fast transient experiments, a regime whose experimen-

tal study has only recently began. Our identification of the interacting Sharvin resistance also implies that the relaxation time of a quantum capacitor connected to a Luttinger liquid is controlled by $1/(2g_L)$ ⁴³. This last experiment could be performed, for instance, by extending the measurements of Ref. [4] to the fractional quantum Hall regime.

Finally, let us note that besides the presented context of nanoelectronics, our approach could also prove useful in other areas such as cold atoms, see e.g. [44] and [45], where Feshbach resonances allows to tune the interactions over a wide range.

ACKNOWLEDGMENTS

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