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CORRELATED FERMIONS and TRANSPORT in MESOSCOPIC SYSTEMS

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T. Martin, G. Montambaux and J. Trân Thanh Vân

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TRANSPORT OF ONE–DIMENSIONAL CORRELATED ELECTRONS

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Sevearl recent results concerning frequency dependent and nonlinear transport of one-dimensional correlated electrons are reviewed. The frequency dependent conductance of a Luttinger liquid is discussed. For zero-range interaction the conductance is similar to that of a non-interacting ideal quantum wire, however, with a renormalized zero-frequency limit. For a screened Coulomb potential the conductance as a function of the frequency shows features of the elementary excitations. They can be simulated in terms of a capacitance, an inductance and a resistance. The non-linear DC-conductance of the Luttinger liquid in the presence of a tunnel barrier shows Coulomb blockade if the interaction is of finite range. The capacitance scales in the same way as for the system without tunnel barrier. The linear DC-response of a tunnel barrier as a function of the frequency and the voltage of a monochromatic driving field shows oscillations which can be interpreted as signature of a frequency-induced depinning of a charge density wave.

I. INTRODUCTION

Recently, electrical transport processes in nano-structures became a subject of increasing experimental and theoretical efforts [1-3]. Especially frequency and time-dependent transport processes [4,5], and the influence of interactions [6] were thoroughly studied. Since they provide insight into the elementary excitations of these systems of geometrically confined interacting electrons, such investigations are of great fundamental interest. In addition, potential applications of nano-structures in future single electron devices, which will have to be operated at very high frequencies, require detailed knowledge of interaction effects in their AC behavior.

Investigations concentrated mainly on the linear AC-response of non-interacting electrons [7], the AC-conductance of a one-dimensional (1D) tunnel barrier [8], periodically driven quantum wells [9], photo-induced tunneling through a tunnel barrier [10], photo-assisted transport through double barrier structures [11,12], and through quantum point contacts [13,14]. In the latter works, a charging model in connection with a semi-classical approach was used, in order to model interaction effects. On the other hand, in the DC-transport through quantum dots, correlation effects were shown to be of great importance [15]. For 1D systems driven transport was studied in the presence of zero-range interaction [16] within the Luttinger model. The latter is a paradigmatic example of a correlated electron system. In view of the quantum nature of the systems, which has to be properly taken into account in a low-temperature transport theory, the study of the AC-transport in a *Luttinger model* should be of considerable interest [17].

We explore in this paper the *linear AC-transport* with a *long-range interaction.* We show that the AC-conductance reflects the elementary excitations of the electrons, in addition to the spatial properties of the electrical probe field. We deduce a quantum capacitance, and a quantum inductance from these results. We present also results for the *non-linear current-voltage characteristic* of a tunnel junction in the presence of finite range interaction. We show that Coulomb-blockade is closely related with the finite-range of the interaction, and that the capacitance of the system scales in the same way as that of the 1D system without tunnel

junction. The *photo-induced conductance* of a potential barrier, when the interaction is of zero-range, will be shown to exhibit features that are closely related to the charge density wave nature of the ground state.

II. AC–TRANSPORT OF A LUTTINGER LIQUID

The Luttinger liquid is a model for the *low-energy excitations* of 1D interacting electrons [18]. Its major importance is that the excitation spectrum can be calculated analytically, as well as many other properties, like the linear conductivity, even in the presence of perturbing potentials [17]. The main assumptions are: (1) linearization of the free electron dispersion relation near the Fermi level, (2) extension of the energy spectrum to include negative energies.

The Luttinger Hamiltonian is easily diagonalized by introducing Bosonic operators, b_k, $b_k^\dagger$. One finds $H=\sum_k \omega(k) b_k^\dagger b_k$, with the dispersion relation

$$\omega(k) = v_F \mid k \mid \sqrt{1+\frac{\hat{V}(k)}{\pi v_F}}, \tag{1}$$

for a general interaction $V(x-x')$ with the Fourier-transform $\hat{V}(k)$. Apparently, $v(k) \equiv v_F\sqrt{1+\hat{V}(k)/\pi v_F}$ plays the role of a (k-dependent) 'charge wave' velocity. For an interaction of zero-range, $\omega(k) = v_F \mid k \mid g^{-1}$ with the interaction strength g^{-1} given by $v(0)/v_F$.

We consider a 3D screened Coulomb interaction potential [19]

$$V(r) = \frac{e^2}{4\pi\varepsilon\varepsilon_0}\frac{e^{-\alpha r}}{r}, \tag{2}$$

with the inverse screening length α, and the relative dielectric constant ε. In a semiconductor heterostructure with narrow electron channels, say in the x-direction, the electrons may be confined within the lowest subband by suitably tuning the voltages at external gates. Then, the interaction of the electrons within the subband may be obtained from eq. (2) by averaging over the coordinates perpendicular to x using the density corresponding to the confining wave function (width d). One obtains

$$V_L(x) = V_0^{(L)}\frac{\alpha e^{-\alpha|x|}}{2} \qquad (\alpha d \gg 1), \tag{3}$$

$$V_C(x) = V_0^{(C)}\frac{e^{-\alpha\sqrt{x^2+d^2}}}{\sqrt{x^2+d^2}} \qquad (\alpha d \ll 1), \tag{4}$$

denoted by 'Luttinger limit' and 'Coulomb limit', respectively. The former tends for $\alpha \longrightarrow \infty$ to the zero-range interaction potential ($\delta(x)$) of the Luttinger model with $V_0^{(L)} \equiv e^2/\varepsilon\varepsilon_0(\alpha d)^2$ = const. The latter, with $V_0^{(C)} = e^2/4\pi\varepsilon\varepsilon_0$, becomes the 1D-equivalent of the unscreened Coulomb interaction, when $\alpha \to 0$. The corresponding Fourier-transforms are

$$\hat{V}_L(k) = V_0^{(L)}\frac{\alpha^2}{k^2+\alpha^2}, \qquad \hat{V}_C(k) = V_0^{(C)}\,2K_0(d\sqrt{k^2+\alpha^2}), \tag{5}$$

with $K_0(x)$ the Bessel function. The dispersion relations for $V_C(x)$ and $V_L(x)$ are shown in Fig. 1 for different coupling strengths.

The dispersion exhibits two asymptotic regions, $\omega(k) \approx v_F|k|g_{C,L}^{-1}$ for $|k|\lambda \ll 1$ and $\omega(k) \approx v_F|k|$ for $|k|\lambda \gg 1$. The length scale λ induced by the interactions is equal to d in the Coulomb limit and to α^{-1} in the Luttinger limit.

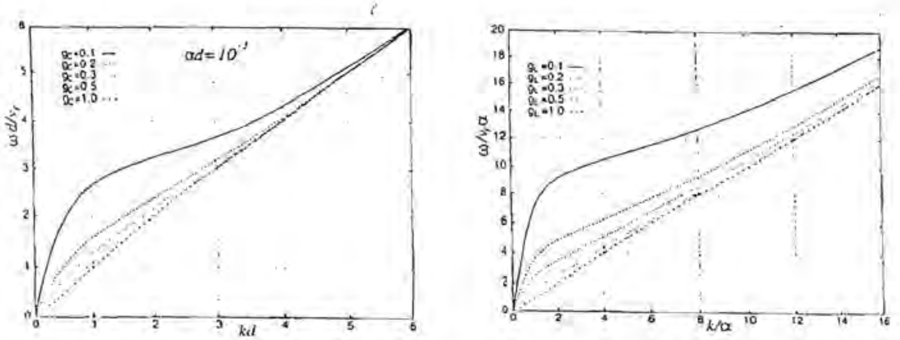

FIG. 1. Dispersion relation $\omega(k)$ with finite-range interaction. Coulomb limit (left) Luttinger limi (right) for different coupling strength.

The velocity of the 'charge sound wave' is strongly enhanced by the interaction as compared to v_F, for long wavelengths. For short wavelengths it becomes v_F. In the intermediate region, $k\lambda = O(1)$, there is a crossover between the two asymptotic dispersions such that $\omega'(k) \equiv d\omega(k)/dk$ has a minimum, ω_p, at some finite wave number k_p, i. e. $\omega''(k_p) = 0$. It corresponds to plasmon-like excitations with non-zero wave vector. In the Coulomb limit, starting from the observation (cf. Fig. 1) that $k_p d = O(1)$ one obtains for $g_C \ll 1$ (strong interaction)

$$\omega_p \approx \frac{v_F a}{d}\sqrt{\frac{2V_0^C}{\pi v_F}} \equiv \frac{v_F a}{g_0 d}. \tag{6}$$

The constant a depends only weakly on the interaction.

The dispersion relation for the Luttinger limit is qualitatively very similar. In the limit of strong interaction ($g_L \ll 1$)

$$\omega_p \approx \frac{v_F\,\alpha}{g_L}. \tag{7}$$

The corresponding wave vector is $k_p \approx 3^{1/4}\alpha/\sqrt{g_L}$. The flat part of $\omega(k)$ at intermediate wave numbers implies the presence of a peak in the derivative of the inverse of the dispersion, $k(\omega)$. The behavior of the excitation spectrum can be detected in Raman scattering [20]. We show now that this can be directly observed in the AC-conductance.

In order to calculate the AC-transport properties we have to couple the electrons with an external electric field $E(x,t)$ which depends explicitely on space and time. The linear relation between the average current density, and $E(x,t)$, defines the non-local conductivity $\sigma(x,t)$

$$\langle j(x,t)\rangle = \int_{-\infty}^{\infty} dx' \int_{-\infty}^{\infty} dt' \sigma(x-x', t-t') E(x', t'). \tag{8}$$

The latter is determined by the Kubo-formula, and it is directly related to the equilibrium retarded current-current correlation function

$$R(x,t) = -i\Theta(t)\langle[j(x,t), j(0,0)]\rangle. \tag{9}$$

In order to evaluate this, we use the Bosonic representation of the current density

$$j(x,t) = -\frac{e}{\sqrt{\pi}}\frac{\partial}{\partial t}\vartheta(x,t), \tag{10}$$

where $\partial\vartheta(x)/\partial x$ is the electronic density, which is given in terms of the Bosonic operators as

$$\vartheta(x) = -i\sum_{q\neq 0}\sqrt{\frac{v_F}{2L\omega(q)}}\,\mathrm{sgn}(q)\left(e^{iqx}b_q - e^{-iqx}b_q^\dagger\right). \tag{11}$$

After some straightforward calculations one gets for the *absorptive part* of the non-local quantum mechanical AC-conductivity,

$$\mathrm{Re}\,\sigma(x,x';\omega) = \frac{v_F e^2}{h}\int_0^\infty dk \cos k(x-x')\left(\delta(\omega(k)+\omega) - \delta(\omega(k)-\omega)\right). \tag{12}$$

The conductance, $\Gamma(\omega)$, is obtained from the conductivity by calculating the absorbed power when a monochromatic electric probe field is applied [7],

$$\Gamma(\omega) = \frac{v_F e^2}{h}\int_0^\infty dk L(k)\left(\delta(\omega(k)+\omega) - \delta(\omega(k)-\omega)\right). \tag{13}$$

Here, $L(k) \equiv \left|\int_{-\infty}^{\infty} dx e^{ikx}E(x)\right|^2/U^2$, with the voltage $U \equiv -\int_{-\infty}^{\infty} dx E(x)$. For a monotonic dispersion, (cf. Fig. 1), we find

$$\Gamma(\omega) = \frac{e^2 v_F}{h}\left(\frac{d\omega}{dk}\right)^{-1}_{\omega(k)} L(k(\omega)). \tag{14}$$

The conductance for the Luttinger limit is shown in Fig. 2.

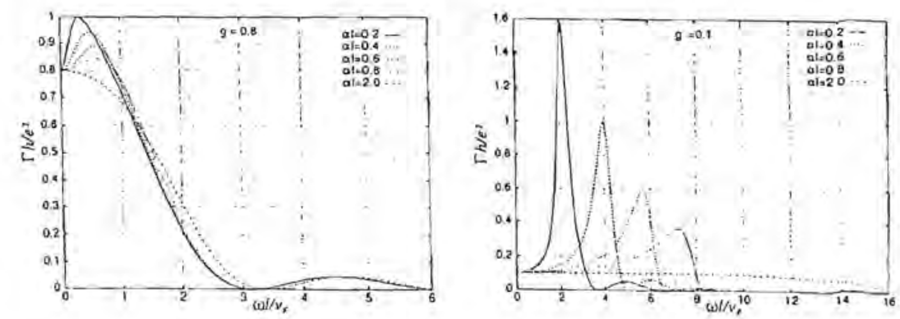

FIG. 2. AC-conductance $\Gamma(\omega)$ in the Luttinger limit for weak (left) and strong (right) coupling, and different screening.

For a δ-interaction $\Gamma(\omega) = (ge^2/h)L\,(g\omega/v_F)$, the same as without interaction [7], except for the renormalization of the prefactor and the Fermi velocity with g. For an interaction potential of finite range we get asymptotically $\Gamma(\omega) = (ge^2/h)L\,(g\omega/v_F)$, and $\Gamma(\omega) = (e^2/h)L\,(\omega/v_F)$, for $\omega \to 0$ and $\omega \to \infty$, respectively, due to the limiting behavior of $\omega(k)$ for small and large $|\,k\,|$.

The most important feature of (14) is that there is a separation between the internal properties of the system, represented by $dk(\omega)/d\omega$, and the external probe field, represented by $L(k)$ — the Fourier transformed of its spatial auto-correlation function [7]. The result (14) implies that internal properties of a mesoscopic quantum system, reflected by the maximum in the AC-conductance at $\omega \approx \omega_p$ (cf. Fig. 2) can be determined, provided that the spatial properties of the probe field are known. If the latter is constant within an interval ℓ, $L(k) = \sin^2 k\ell/(k\ell)^2$. For small $k\ell$, this is essentially a constant. There are zeroes for $k\ell = n\pi$. If $E(x)$ is a smooth function around $x = 0$, rapidly decaying for $|\,x\,| \to \infty$, $L(\omega)$ will also be a rapidly decaying function, without any particular structure. Then, the features of the dispersion of the elementary excitations are directly displayed by the AC-conductance provided that $\omega_p < \omega_1$. The frequency $\omega_1 = \omega(\pi/\ell)$ corresponds to the width in frequency of the Fourier-transform of the probe field. It is given by eq. (1) when $k = \pi/\ell$. In order to observe the resonance experimentally one should have $\ell/d < 1$. Near-field optical spectroscopy [21] should in principle be able to provide the technique for such a measurement.

We conclude by mentioning that the above results suggest to define a *capacitance* and an *inductance* of a quantum wire when the interaction is strong *and* long-range. We concentrate on the Coulomb limit. The Luttinger limit can be treated analogously [22].

Consider the classical capacitance, C, and inductance, L, of a charge and a current distribution via electrostatic and magnetic energy, respectively. For a 'wire' (length $\mathcal{L}(\to \infty)$, charge and current density $\propto \exp(-(y^2+z^2)/2d^2))$ they are $Q^2/2C \equiv (4\pi\varepsilon_0)^{-1}Q^2\mathcal{L}^{-1}\ln\mathcal{L}/d$, and $LI^2/2 \equiv (\mu_0/4\pi)I^2\mathcal{L}\ln\mathcal{L}/d$, respectively. Here, Q, I are the total charge and current, respectively. Then $C = 2\pi\varepsilon_0\mathcal{L}/\ln(\mathcal{L}/d)$ and $L = (\mu_0/2\pi)\mathcal{L}\ln(\mathcal{L}/d)$. The corresponding resonance frequency is $\omega_0^2 \equiv (LC)^{-1} = (\varepsilon_0\mu_0)^{-1}\mathcal{L}^{-2}$.

As the plasma frequency ω_p with wave vector $k_p \propto \pi/d$, ω_0 represents a resonance of a mode with wave number $k_R = \pi/\mathcal{L}$ of the classical wire. This suggests to replace $\mathcal{L}$ by d in ω_0 when attempting to translate the classical resonance frequency into ω_p. The result eq. (6) can then be reproduced by identifying additionally

$$\varepsilon_0 \to \frac{e^2}{h}\frac{g_C^2 K_0(\alpha d)^2}{v_F a^2} \approx \frac{e^2}{h}\frac{g_0^2 K_0(\alpha d)}{v_F a^2}, \qquad \mu_0 \to \frac{h}{e^2}\frac{1}{K_0(\alpha d)v_F}. \tag{15}$$

This choice is, of course, not unique. It is motivated by the fact that $e/\sqrt{\varepsilon_0}$ represents the interaction strength in the electrostatic energy. The product of ε_0 and μ_0 should be independent of αd. Thus, the extra factor $K_0(\alpha d)$ when replacing ε_0 requires a factor K_0^{-1} when replacing μ_0. The latter are necessary in order to compensate the logarithmic terms in C and L which result from the cutoff at $\mathcal{L}$ in the integrations for the electrostatic and magnetic energies. In the microscopic theory, the cutoff is provided by the screening length α^{-1}. Therefore, we choose to replace $\ln(\mathcal{L}/d)$ in the classical expressions for C and L by $K_0(\alpha d)(\propto -\ln\alpha d)$ for $\alpha d \to 0$. The prefactor e^2/hv_F appears for dimensional reasons.

The *quantum capacitance* and *quantum inductance* of the Luttinger 'wire' can then be defined as follows

$$C_q \equiv \frac{e^2}{h}\frac{2\pi}{v_F a^2}g_0^2 d, \quad L_q \equiv \frac{h}{e^2}\frac{1}{2\pi}\frac{d}{v_F}. \tag{16}$$

As the DC-*quantum contact conductance*, $\Gamma_q \equiv e^2 g_C/h$, C_q and L_q are independent of the wire length. They depend only on microscopic properties. While Γ_q vanishes as $|\ln\alpha d\,|^{-1/2}$ for $\alpha \to 0$ (Coulomb interaction), the latter stay finite. The above results imply that the AC-transport behavior of the Luttinger wire can also be simulated by a classical circuit of an inductance, a capacitance and resistances, and leads to a generalisation of the above results which accounts also for the finite width of the resonance in the AC-conductance [22]. This width is due to the bulk modes of the correlated electrons and is represented by a resistance which is conceptually different from Γ_q^{-1}.

III. CURRENT-VOLTAGE CHARACTERISTIC OF A TUNNEL JUNCTION

The Coulomb blockade effect in the non-linear current-voltage (I-U) characteristics of mesoscopic tunnel junctions [6,23] was studied intensively by theory and experiment [2] during the past decade. Due to the repulsive interaction between the electrons, tunneling is supressed for voltages below $U_C = e/2C$, if the temperature is smaller than $T_C \equiv E_C/k_B$ (k_B Boltzmann constant, e elementary charge, C capacitance of the circuit). The quantity $E_C \equiv eU_C$ is called the charging energy.

In the semi-phenomenological theory of Coulomb blockade [24] the tunnel junction is modeled by a capacitance and a tunnel resistance R_t. An external impedance $Z(\omega)$ is included into the circuit. Here, we present some microscopic results for a 1D tunnel junction [25]. The parameters introduced in the above mentioned theory by *ad-hoc* assumptions are deduced from the interaction between the electrons.

Two semi-infinite (1D) systems of interacting electrons treated within the above Luttinger approximation are coupled by a tunnel junction. The interaction potential between the electrons is again assumed to have a finite, non-zero range. The tunneling current as a function of a voltage applied across the junction is calculated using Fermi's golden rule with respect to the tunneling part of the Hamiltonian.

The result for the current-voltage characteristic $I(U)$ is

$$I(U) = \frac{e\Delta^2}{4}\left[1 - e^{-\beta eU}\right]\int_{-\infty}^{\infty} dt e^{ieUt}e^{-W_g(t)}. \tag{17}$$

(Δ tunnel matrix element, β inverse temperature). The thermal equilibrium correlation function is

$$W_g(t) = \int_0^{\infty} d\omega\frac{J(\omega)}{\omega^2} \times \left[(1 - \cos(\omega t))\coth\left(\frac{\beta\omega}{2}\right) + i\sin(\omega t)\right]. \tag{18}$$

It contains as the key quantity the spectral function $J(\omega)$ which is in our theory directly given in terms of the dispersion law of the charge excitations of the system (cf. eq. (1)) [25].

As usual, the tunnel resistance is given by the inverse of the tunneling probability Δ^{-2}. In addition, the charging energy is found to be the interaction potential at zero-distance, $E_C \propto V(x=0)$, and the dissipative resistance is given by the spatial average of the interaction potential, $R \propto \sqrt{\bar{V}}$. The spectrum of the elementary excitations determines completely the impedance of the circuit, $Z(\omega) \equiv J(\omega)/\omega - 2$ (Fig. 3). In order to explain the latter no additional 'environmental modes' are needed. The crucial point is that the above mentioned asymptotic behavior of $I(U)$ for large U appears to be directly related to the finite, non-zero range of the interaction. The latter implies that the dispersion relation of the elementary excitations of the electron system becomes the one of free electrons in the short wavelength limit. Our results show that in 1D a charging energy, and, in turn, a capacitance, cannot be defined if the electron-electron interaction is zero-range.

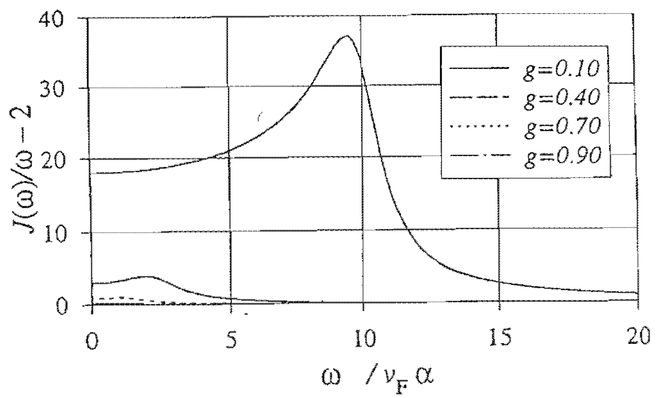

FIG. 3. The function $J(\omega)/\omega - 2$ of the Luttinger limit for different interaction parameters g (see text).

The capacitance, given here by $C \propto V^{-1}(0) \propto g_0^2 d$, scales in the same way as the one, C_q, found above by considering the frequency dependent conductance of the pure quantum wire (of interacting electrons). This was obtained by a completely different philosophy starting from the classical theory of antennas. In fact, in the Coulomb limit, $C/C_q = O(1)$. Following the earlier approach, one could consider the present problem as being the quantum equivalent of two (only very weakly coupled) wires. From the comparison of their AC-properties with those of the quantum system considered here, one would get the same scaling law for the capacitance.

The capacitance responsible for the Coulomb blockade is naturally given here by the total capacitance of the circuit, i. e. the two wires and the junction. Apparently, in our ideal 1D model, the junction as such does not contribute significantly towards the capacitance. It is only needed for detecting the Coulomb blockade induced by the latter.

IV. PHOTO-INDUCED TUNNELING

Stationary tunneling was investigated thoroughly in recent years for the above model [17,26,27]. The results provided access to a microscopic picture for the Coulomb blockade effect in nanometer-tunnel contacts, as discussed above. There are also indications that the Luttinger liquid is important for the understanding of the edge currents in fractional quantum Hall systems [16]. Until now, transport in the presence of an additional time-dependent potential was studied only by using "Fermi's golden rule" [16].

We applied the path integral formalism [28] to evaluate the current-voltage characteristics of a tunnel barrier in the Luttinger liquid for zero-range interaction, and subject to a voltage

$$V(x,t) = [V_\omega(t) + V_\Omega(t)]\Theta(x).$$

The "probe voltage" $V_\omega(t) = V_\omega \cos\omega t$ is of infinitesimally small amplitude, and $V_\Omega(t) = V_\Omega \cos\Omega t$ is an additional bias voltage, for instance due to a microwave field, of arbitrary amplitude. We show [28] that the *absorptive* AC-conductance with respect to the "probe voltage" in the presence of the "microwave voltage", $\Gamma_\Omega(\omega)$, can be written as a superposition of the linear AC-conductances of the tunnel contact in the absence of the "microwave" at $\omega + n\Omega$ (n integer),

$$\Gamma_\Omega(\omega) = \sum_{n=-\infty}^{\infty} J_n^2\left(\frac{eV_\Omega}{\Omega}\right)\frac{(\omega+n\Omega)}{\omega}\Gamma_0(\omega+n\Omega). \tag{19}$$

In the limit of a high barrier the AC-tunnel conductance (for $T \to 0$) is

$$\Gamma_0(\omega) = \frac{1}{R_t}\frac{1}{\Gamma(2/g)}\left(\frac{\omega}{\omega_c}\right)^{2/g-2} e^{-|\omega|/\omega_c}. \tag{20}$$

The frequency ω_c corresponds to the energy-cutoff of the Luttinger excitation modes.

In eq.(19), $J_n(z)$ is the Bessel function of order n. By evaluating (19) in detail, we find that the zero-temperature DC-conductance $\Gamma_\Omega(0)$ shows characteristic oscillations when the frequency of the driving voltage is changed (Fig. 4).

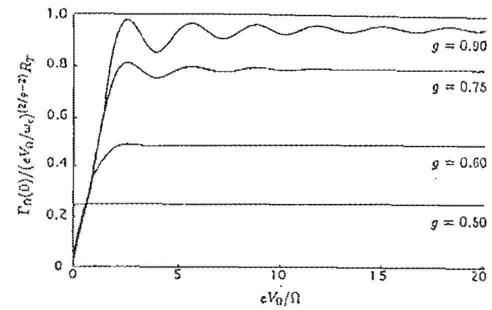

FIG. 4. The DC-photoconductance of a tunneling barrier in a Luttinger liquid as a function of the frequency of the "photon field", $\Gamma_\Omega(0)$, in units of $(eV_\Omega/\omega_c)^{2/g-2}/R_t$, at zero-temperature for various interaction parameters g (R_t tunneling resistance, V_Ω amplitude of driving voltage, ω_c cutoff frequency).

The conductance in the presence of the additional time-dependent bias voltage can be viewed as the "photoconductance". It implies that the absorption of "photons" from the "microwave field" determines completely the linear transport. *A posteriori*, this appears plausible. It is nevertheless surprising, that the same features are also found for a strongly correlated quantum system.

For g slightly less than 1

$$\Gamma_\Omega(0) \approx \frac{1}{R_t}\left(\frac{\Omega}{\omega_c}\right)^{2/g-2}\left[1 - J_0^2\left(\frac{eV_\Omega}{\Omega}\right)\right]. \tag{21}$$

The maxima of the conductance are roughly determined by the zeros of the square of the J_0 Bessel function. Thus mathematically, the oscillations originate from the absence of the $n=0$ channel in the DC-limit of (19) at zero temperature.

Physically, they are a signature of the Coulomb blockade phenomenon or pinning of a charge density wave at a barrier in a correlated 1D electron system [17]. The correlations due to repulsive electron-electron interactions lead to the nontrivial behavior. As the bias voltage is absent in the $n=0$ DC-channel, the respective contribution to the conductance drops to zero $\propto T^{2/g-2}$ for $g<1$ as $T \to 0$. The missing DC-channel has non-monotonous weight $J_0^2(eV_\Omega/\Omega)$. It becomes noticeable as the sum which describes the contribution of the open channels, for

$g \approx 1$ is $1 - J_0^2(eV_\Omega/\Omega)$. With decreasing g, the oscillations are reduced. At $g = 1/2$ the sum adds up to the monotonous form $(eV_\Omega/\Omega)^2/4$. Together with the prefactor, this gives a conductance which is independent of Ω. With increasing temperature for g near 1, the $n = 0$ channel is switched on, and compensates the oscillating part in (21), $J_0^2(eV_\Omega/\Omega)$, partially. The oscillations of the conductance are increasingly suppressed with increasing T. Eventually, they vanish, when the rate $\Gamma_0(0)$ becomes numerically comparable to $\Gamma_0(\Omega)$, i.e., when the temperature is of the order of Ω.

V. CONCLUSIONS

We presented results of AC-transport properties of the Luttinger model with finite-range interaction. First, we investigated the linear AC– conductance of a quantum wire. Its frequency behavior shows features related to the elementary excitations. They are due to the finite range of the interaction, and were used to define a quantum capacitance and inductance.

We found that the finite range of the interaction was also crucial for the occurrance of the Coulomb blockade in the current-voltage characteristic of a tunnel junction between two Luttinger liquids. In this case the spectrum of the elementary excitations determines completely the impedance of the system.

Finally we investigated the photo-induced tunneling of electrons with zero-range interaction. We showed that the photo-induced conductance is given by the linear AC-conductance. For low temperature and weak repulsive interaction the DC-photoconductance shows characteristic oscillations as a function of the driving frequency. They are signature of photo-induced charge density wave depinning.

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TRANSPORT D'ÉLECTRONS CORRÉLÉS À UNE DIMENSION

Nous rappelons plusieurs résultats sur le transport non-linéaire et dépendant de la fréquence d'électrons corrélés à une dimension. La conductance AC d'un liquide de Luttinger est également discutée. Pour des interactions à portée nulle, la conductance est semblable à celle d'un fil quantique idéal sans interaction, mais avec une limite à fréquence nulle renormalisée. Pour un potentiel Coulombien écranté, la conductance en fonction de la fréquence présente certaines caractéristiques dues aux excitations élementaires. Elles peuvent être modélisées par une capacité, une inductance et une résistance. Si la portée de l'interaction est finie on observe du blocage de Coulomb dans la conductance DC non-linéaire du liquide de Luttinger en présence d'une barrière tunnel. La capacité se comporte de la même façon que pour le système sans barrière tunnel. La réponse DC linéaire d'une barrière tunnel en fonction de la fréquence et de la tension appliquée possède des oscillations qui peuvent être interprétées comme la signature du décrochage induit par la fréquence d'une onde de densité de charge.