We report on electron transport measurements in high-quality carbon nanotube devices with a total transmission of about 1/2. At liquid helium temperature the linear conductance oscillates with moderate amplitude as a function of the gate voltage around an average value of the conductance quantum. Upon decreasing temperature, we observe an intriguing fourfold increase in the period of the oscillations accompanied by an enhancement in their amplitude. While the high-temperature oscillations are suggestive of charging effects in an open interacting quantum dot, the low-temperature transport characteristics is reminiscent of single-particle Fabry-Pérot interference in a carbon nanotube waveguide. A similar crossover is observed in the low-temperature differential conductance by tuning the source-drain voltage. We reconcile these observations by attributing the four-fold increase at low energies to the interplay of interactions and quantum fluctuations, leading to a correlated Fabry-Pérot regime.
Electron interactions and quantum interference are central in mesoscopic devices. The former are due to the electronic charge and give rise to many-body effects; the latter emerges due to the wave-like properties of an electron. Resonant ballistic devices with a few conduction modes and moderate coupling to electrodes are sensitive to both of these electronic properties. On the one hand, quantum interference between electron waves backscattered at the boundaries between the mesoscopic system and the metallic electrodes gives rise to resonant features in the transmission, analogous to the light transmission in an optical Fabry-Pérot cavity [1]. On the other hand, if the electron spends enough time in the mesoscopic device before being transmitted, Coulomb repulsion can also become important giving rise to Coulomb blockade and single-charge tunneling effects [2]. Despite considerable efforts, the interplay between electron interactions and quantum interference remains poorly understood from both an experimental and a theoretical point of view, due to the many-body character of the problem. This is the topic of the present Letter.

Carbon nanotubes (CNTs) are an ideal system to study both electron correlations and quantum interference [3]. In fact, various many-body effects including Coulomb blockade [4, 5], Luttinger liquid behavior [6–9], Wigner phases [10–13], and Kondo physics [14–24] as well as Fabry-Pérot oscillations resulting from electron interference [25–27] have been observed in CNTs. It is possible to switch from interaction- to interference-governed transport regimes by tuning the tunnel couplings at the interface between the nanotube mesoscopic region and the electrodes, \( \Gamma_S \) and \( \Gamma_D \) for the source (S) and drain (D) electrodes. Which transport regime is dominant crucially depends on how large the tunneling broadening \( \hbar \Gamma = \hbar (\Gamma_S + \Gamma_D) \) is compared to other energy scales, in particular to the charging energy \( E_C \), being the electrostatic cost to add another (charged) electron to the CNT [3]. In the so-called quantum dot limit when \( \hbar \Gamma \ll E_C \), tunneling events in and out of the CNT are rare and Coulomb charging effects are dominant. They give rise to Coulomb blockade phenomena and single-electron tunneling in the regime \( k_B T \sim E_C \), when no clear hierarchy of energy scales exists.

An experimental hallmark of both interaction- and interference-dominated transport is the modulation of the conductance when sweeping the electrochemical potential, that is, by varying the gate voltage \( V_g \). In the incoherent tunneling regime, the alternance of single-electron tunneling and Coulomb blockade physics results in finite conductance peaks with a period in \( V_g \) of the order of \( e/C_g \) [2], where \(-e\) is the (negative) electron charge and \( C_g \) is the capacitance between the nanotube and the gate electrode, see Fig. 1(a). In contrast, in the interference-dominated regime the conductance modulation of the Fabry-Pérot oscillations arises from the electron wave phase accumulated

![Figure 1](image-url)
during a round trip along the nanotube. The presence of valley and spin degrees of freedom gives rise in CNT interferometers to oscillations of period $\Delta V_g = 4e/C_g$ [25].

In this work, we improve the device quality to an unprecedented level. In the open quantum dot configuration, we observe an intriguing change of the conductance oscillation period from $e/C_g$ to $4e/C_g$ by decreasing the temperature and source-drain voltage. This indicates a crossover from charging dominated to quantum interference dominated resonant tunneling. As discussed in more detail below, we interpret the Fabry-Pérot oscillations within a full many-body framework: above crossover, coherent single-electron tunneling in an interacting finite-length nanowire is the dominant transport mechanism; as the temperature is lowered below crossover, spin and charge fluctuations become increasingly important and finally lead to a Fabry-Pérot-like interference pattern.

Experimental results.- We use a new method to fabricate nanotube devices with improved transport measurement quality. We grow nanotubes by chemical vapor deposition on prepatterned electrodes [28]. The nanotube is suspended between two metal electrodes, see Fig. 1(a). We clean the nanotube in the dilution fridge at base temperature by applying a high constant source-drain voltage $V_{sd}$ for a few minutes (see Sec. I of the Supplemental Material). This current-annealing step cleans the nanotube surface from contamination molecules adsorbed when the device is in contact with air. The energy gap of the two nanotubes discussed in this work is on the order of 10 meV. The length of the two suspended nanotubes inferred by scanning electron microscopy (SEM) is about 1.5 µm.

Figure 1(b) shows the modulation of the differential conductance $G_{\text{diff}}$ of device I as a function of $V_g$ in the hole-side regime at 15 mK after annealing. Rapid conductance oscillations are superposed on slow modulations. Since the conductance remains always large, that is above $e^2/h$, we attribute the rapid oscillation to the Fabry-Pérot interference with period in gate voltage being $\Delta V_g = 4e/C_g$. The slow modulation may be caused by the Sagnac interference [26, 27], the additional backscattering due to a few residual adatoms on the CNT, the symmetry breaking of the electronic wave function by the planar contacts of the device, or any combination of these (for further discussion see Sec. I and IIA of the Supplemental Material). A crossover to a regime dominated by the charging effects in an open interacting quantum dot is observed upon increasing temperature. Specifically, by sweeping the temperature from 15 mK to 8 K the amplitude of the oscillations gets smaller. Further, the oscillation period gets four times lower, changing from $4e/C_g$ at 15 mK to $e/C_g$ at 8 K, see Figs. 2(a) and (c-e). The period in $V_g$ is calibrated in units of

![Figure 2](image-url)
$e/C_g$ using the measurements in the electron-side regime, where regular Coulomb oscillations are observed at 8 K, as shown in Fig. 2(b). The same behavior is observed in device II, Figs. 3(a) and (b). The $4e/C_g$ oscillations vanish above 4 K in both devices, whereas the $e/C_g$ oscillation amplitude is suppressed to almost zero below $\sim$ 1 K in device I and below $\sim$ 0.1 K in device II, see Figs. 2(f) and 3(b).

Our interpretation of a temperature-induced crossover between two seemingly distinct transport regimes is confirmed by measured maps of the differential conductance as a function of source-drain and gate voltages at $T=15$ mK and $T=8$ K, as shown in Fig. 4(a) and (d), respectively. The low-temperature data feature the regular chess-board-like Fabry-Pérot interference pattern [25], while the high-temperature data show smeared Coulomb diamonds. Such measurements further allow us to extract important energy scales for our device. The characteristic bias $V_{sd}^*$ indicated by the arrow in Fig. 4(a) yields a single-particle excitation energy $\Delta E = eV_{sd}^* \approx 1.7$ meV. This value is consistent with what is expected from a nanotube with length $L \approx 1.5 \mu$m. Assuming the linear dispersion $\varepsilon(k) = h v_F k$ with longitudinal quantization $k_n = n\pi/L$ and the Fermi velocity $v_F = 10^6$ m/s, yields $\Delta E = \varepsilon(k_{n+1}) - \varepsilon(k_n) = h v_F \pi/L \approx 1.4$ meV. The charging energy is estimated from the charge stability diagram measurements at 8 K, Fig. 4(d); from the Coulomb diamond, indicated by the dashed lines, a charging energy $E_C \approx 3.6$ meV is extracted. Further, we estimate $\hbar \Gamma \approx E_C$ because of the strong smearing of the diamonds in Fig. 4(d) and the weak conductance modulation at 8 K in Fig. 2(a) [29]. The energy hierarchy in our experiment is thus $E_C \approx \hbar \Gamma \gg \Delta E \gg k_B T$.

There is also a change of the conductance oscillation period upon sweeping the source-drain voltage bias at 15 mK. The period changes from $4e/C_g$ at zero bias to $e/C_g$ at high bias, see Figs. 4(a,b). The source-drain bias dependence of the amplitude of both oscillations is non-monotonous as seen in Fig. 4(c), in contrast to the temperature dependence in Fig. 2(f). **Discussion.** - The high-temperature measurement of the charging effect in an open quantum dot indicates electron correlation. When reducing temperature, the associated $e/C_g$ conductance oscillations disappear smoothly to give rise to the $4e/C_g$ oscillations. This smooth crossover occurring by lowering temperature suggests that the Fabry-Pérot-like oscillations also occur in a regime where electrons are correlated. As shown in the Sec. II of the Supplemental Material, neither a single-particle description of Fabry-Pérot interference nor the charging physics alone can account for the observed temperature induced crossover. Further, due to the high aspect ratio of nanotubes, our devices are one-dimensional in nature. Hence, we expect that such Fabry-Pérot oscillations should be well captured within a

![Figure 3](image-url) Figure 3. Temperature-induced crossover for device II. (a) Conductance traces for a series of different temperatures. (b) Temperature dependence of the FFT amplitude associated with the $4e/C_g$ period oscillations and the $e/C_g$ period oscillations.
Luttinger liquid framework for interacting one-dimensional electrons. This interpretation in terms of correlated Fabry-

Figure 4. From Fabry-Pérot patterns to blurred Coulomb diamonds in device I. (a) Map of the differential conductance as a function of $V_{sd}$ and $V_g$ at 15 mK. From the position of the arrow the single-particle excitation energy is extracted. (b) Differential conductance traces for a series of different source-drain voltages at 15 mK. (c) Source-drain voltage dependence of the FFT amplitude associated with the $4e/C_g$ and the $e/C_g$ period oscillations at 15 mK. The curves are obtained by doing a FFT of the $G_{diff}(V_g)$ trace for each $V_{sd}$ value. (d) Map of the differential conductance as a function of $V_{sd}$ and $V_g$ at 8 K. The dashed lines highlight the contours of the Coulomb diamonds.

Pérot oscillations is in line with the predictions in Refs. [30–33] of Fabry-Pérot interference in a Luttinger liquid. In particular, Peça et al. [30] studied the zero-temperature differential conductance of a CNT in the weak backscattering limit. They found that the simple single-particle approach and the many-body description result in qualitatively indistinguishable interference patterns when interactions in the CNT are treated within a Luttinger liquid picture. In the weak backscattering regime however, increasing temperature smoothens the conductance oscillations in the Luttinger liquid, but does not change the oscillation period [32]. This hints at the fact that a perturbative calculation in the backscattering strength does not account for the experimental observations. In the opposite weak tunneling limit, Coulomb blockade effects in a finite CNT quantum dot are also fully captured in a Luttinger liquid description of the low-energy physics [34]. Also in this case, however, decreasing the temperature does not change the period of the conductance oscillations, suggesting that our experiments also require a treatment non-perturbative in the tunneling coupling. Alicea et al. [31] bridged the gap between the weak backscattering limit and the weak tunneling limit. They showed that the average charge accumulated in an interacting CNT wire, modeled as a Luttinger liquid, oscillates in gate voltage – when increasing the tunneling coupling, Coulomb oscillations associated to single-electron filling collapse into Fabry-Pérot oscillations with a fourfold increase of the period due to the four-electron shell filling. However, the work [31] focused on the zero-temperature limit and did not investigate the crossover measured by
The temperature-induced crossover measured in our work has some connections to the predicted conductance enhancement due to Kondo-like exchange correlations in combination with spin-charge separation in an open one-channel quantum dot [35]. The conductance is predicted to oscillate in gate voltage and to exhibit a universal behavior below an appropriate characteristic temperature, which plays the role of the Kondo temperature in an open system. Our data indeed show similarities but also differences with the SU(4) Kondo effect in carbon nanotubes in the weak tunneling regime $\hbar \Gamma \ll E_C$ [22, 36]. In such case, the tunneling coupling is low enough compared to the charging energy to allow full localization of the charge within the dot, but it is large enough compared to the Kondo energy to enable spin correlations. This results in a crossover from charging effects at high temperature to the increased conductance of Kondo resonances at zero temperature with a fourfold enhancement of the oscillation period in the case of SU(4) Kondo [3, 17, 36]. In contrast to our experiment though, in the SU(4) Kondo effect the conductance alternates between large values close to $4e^2/h$ at the peak and almost zero at the minimum [22, 36]. In our devices, the tunneling coupling is large, $\hbar \Gamma \simeq E_C$. The charge is no longer strongly localized within the dot. Both spin and charge fluctuations are expected, which we attribute to be at the origin of the crossover from charging oscillations to fluctuations-dominated Fabry-Pérot oscillations observed when lowering temperature and source-drain voltage.

**Conclusion.**- Our work provides a comprehensible phenomenology of transport in nanotubes when both interference and interaction are involved. We measure a fourfold change of the oscillation period of $G_{\text{diff}}(V_g)$ upon varying the temperature and source-drain voltage. These findings hint to temperature and voltage induced crossover from weak tunneling to weak backscattering in a Luttinger liquid. However, a quantitative description of our experiment still constitutes a theoretical challenge. It will be interesting to measure shot noise [32, 37–39] and the backaction of the electro-mechanical coupling [40, 41] to further characterize these correlated Fabry-Pérot oscillations.

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Crossover from Fabry-Pérot to charging oscillations in correlated carbon nanotubes
Supplemental Material

EXPERIMENTAL SECTION

High-quality nanotubes obtained by current annealing

We grow nanotubes by chemical vapor deposition on prepatterned electrodes using the technique described in Ref. [28]. The nanotube is suspended between two metal electrodes Fig. 5(a). We clean the nanotube in the dilution fridge at base temperature by applying a high constant source-drain voltage $V_{sd}$ for a few minutes. The highest applied value of $V_{sd}$ is usually chosen by ramping up the bias until the point when the current starts to decrease, see Fig. 5(b). This current-annealing step cleans the nanotube surface from contaminations. This procedure allows us to adsorb helium monolayers uniformly along nanotubes, indicating that the nanotube is essentially free of adsorbate contamination [42]. Figures 5(c-g) show the modulation of the differential conductance $G_{diff}$ of device I as a function of $V_{sd}$ in the hole-side regime at 15 mK before and after annealing, respectively. The current annealing results in regular conductance modulation.

![Image](image_url)

**Figure 5.** Current annealing and low-temperature transport characteristics. (a) Three-terminal device with a suspended CNT contacted to source (S), drain (D), and gate (G) electrodes. (b) Current-voltage characteristic of device I at $T=15$ mK. The arrow indicates when the current starts to decrease while increasing $V_{sd}$. The highest voltage used for current annealing is usually around this value. (c-g) Gate voltage dependence of the conductance $G_{diff}(V_{sd})$ of device I at $T=15$ mK measured before current annealing and after different current annealing steps. The measurements in d-g have been carried out in a second cool-down, while all the other presented data of device I have been recorded in the first cool-down. An oscillating voltage with amplitude smaller than $k_B T/e$ is applied to measure the differential conductance.

In the annealed sample rapid conductance oscillations are superposed on slow modulations, see Fig. 5(d). Since the conductance remains always large, we attribute the rapid oscillation to the Fabry-Pérot interference with period in gate voltage being $\Delta V_g = 4e/C_g$. The first interpretation of slow modulation coming to mind is the so-called Sagnac
interference, due to the gradual change of the Fermi velocity when sweeping $V_g$ \cite{26, 27}, caused by the trigonal warping. In the dispersion of non-interacting electrons trigonal warping manifests at energies further than $\sim 200$ meV away from the charge neutrality point, while the range of single-particle energies scanned in our experiment is of the order of $\sim 56$ meV (estimated from $\sim 40$ peaks visible in Fig. 1(b) of the main text, separated by $\Delta E \simeq 1.4$ meV). Unless the interactions bring the trigonal warping effects closer to the charge neutrality point, an alternative explanation of the slow modulation is needed. One possibility is the beating caused by the presence of a symmetry breaking mechanism which introduces additional valley mixing and/or another characteristic length scale into the system (see the discussion of Fig. 9). The pattern of the secondary interference is completely changed each time that we do a current-annealing of the device, see Fig. 5(d,e). We attribute this modification either to the atomic rearrangement of the platinum electrodes in the region near the nanotube, so that the intervalley backscattering rate at the contacts changes \cite{26}, or to the changed position of residual adatoms near the contacts.

Electron transport properties

The energy gap of the two nanotubes discussed in this work is on the order of 10 meV. The size of the energy gap can be obtained by recording the dependence of the resistance on $V_g$ at different temperatures \cite{3}, see Fig. 6. The order of magnitude of the band gap $E_G$ is obtained from the temperature at which the resistance in the gap gets high, $E_G \sim k_B T$.

In Fig. 7 is shown a selection of $G_{\text{diff}}(V_g)$ traces of device I at different temperatures. We select the $V_g$ ranges for which data are presented in the main text.

THEORETICAL CALCULATION OF TRANSPORT

Because of the lack of clear energy scales separation, i.e. $U \simeq \Gamma \gg k_B T$, the theoretical description reproducing the results of the experiments is very challenging; $U = E_C$ stands for the characteristic strength of the Coulomb interaction between the electrons in the system. We can however provide theoretical support for our interpretation of the data as the interplay of correlations and interference effects by showing that neither of these mechanisms alone can explain the observed evolution of conductance with temperature. On one hand, we show in Sec. II A results for the Fabry-Pérot interference with $\Gamma \gg k_B T$ and $U = 0$. While such single-particle interference can explain the experimental results at 15 mK, it cannot reproduce the fourfold decrease in the oscillation period with increasing temperature. On the other hand, we analyze in Sec. II B the electronic transport across an interacting multilevel quantum dot with four-fold degenerate energy levels and level spacing $\Delta E$. We use a so-called coherent sequential tunneling approximation, which yields correct results for non-interacting ($U = 0$) and Coulomb blocked ($U \gg k_B T > \Gamma$) systems, but also in the regime $U > \Gamma \gtrsim k_B T$. For the parameters in the latter regime, $\Gamma = 0.5U$ and $k_B T \simeq 0.1U$, the theory yields a Coulomb oscillation behavior similar to that observed in the experimental data at 4 K, with its alternating higher
Figure 7. Series of $G_{\text{diff}}(V_g)$ traces at different temperatures of device I. We select the $V_g$ ranges for which data are presented in the main text.

and lower peak pairs. However, lowering the temperature again does not introduce any change in periodicity. An essential ingredient, the Kondo-like correlation, is missing from the theory.

**Single particle Fabry-Pérot interference**

In this section we shortly recall a single-particle approach to Fabry-Pérot interference and its prediction for a CNT-based electron waveguide. This approach is justified for devices with transparent contacts, when the electron transport through the system is usually too fast to show signatures of charging effects. Then the conductance assumes overall a high value; further, low-amplitude periodic oscillations in the conductance arise from constructive and destructive interference of the electronic trajectories shuttling between the two leads [25]. Besides the primary Fabry-Pérot interference, a slow oscillation of the average conductance due to Sagnac interference [26, 27] arises when the velocities of left- and right-moving electrons do not match in magnitude.

In the analytical approach the Fabry-Pérot interference is described through the different reflection and transmission coefficients of the two modes at the left and right interface, $t_{L/R}, r_{R/L}$, respectively. (Since all calculations presented here are at zero bias, instead of $S/D$ from the main text we use the convention of $L/R$ as in Fig. 8(a).) In the absence of mixing of the two intervalley channels (orange processes in Fig. 8(b)) the formula for the overall transmission is given by

$$T(V_g) = \sum_{j=a,b} \frac{2|t_{L_j}|^2|t_{R_j}|^2}{1 + |r_{L_j}|^2|r_{R_j}|^2 - 2|r_{L_j}||r_{R_j}|\cos(\phi_{j,k}(V_g))},$$

where $j$ labels the two independent channels for interference marked in Fig. 8(b) by green arrows, and $\phi_{j,k}(V_g) = (|k_{j,i}(V_g)| + |k_{j,r}(V_g)|)L$ is the phase accumulated by the electron after traversing the nanotube once back and forth,
i.e. once on a left-moving branch of the dispersion with momentum \( k_{j,l}(V_g) \) and once on the right-moving branch with the dispersion \( k_{j,r}(V_g) \). The momentum is related to the gate voltage through the dispersion relation \( \varepsilon(k_{j,r/l}) = \alpha eV_g \), where \( \alpha \) is the lever arm. The interference pattern in the transmission arises due to the \( \cos(\phi_{j,l}(V_g)) \) term.

Reproducing the experimental transmission curves requires the knowledge of the reflection and transmission coefficients \( t_{L/R}, r_{L/R} \), yielding four different parameters to adjust. Further, the simple formula \( 1 \) cannot account for the beating observed in the experiment due to combined intravalley and intervalley scattering \([60]\). Hence we turn to a numerical calculation of transmission, using a single particle Green’s functions approach\([1]\) with just the tunnel couplings \( \Gamma_L \) and \( \Gamma_R \) to the left and right lead, respectively.

We chose for the numerical simulation a \((20,5)\) nanotube with the diameter \( d = 1.8 \) nm and length \( L = 1.04 \mu m \), comparable with the experimental parameters. The leads are assumed to be wide band, since the experimental conductance is very high near the band gap\([43]\). The system is sketched in Fig. 8(a). The CNTs band structure in the Dirac regime is shown in Fig. 8(b), and the transmission (i.e. the zero temperature linear conductance) in Fig. 8(c). It has been obtained with the Landauer-Büttiker formula in the Fisher-Lee form\([1]\)

\[
T(E) = \text{Tr}[\Gamma_L G^R(E)\Gamma_R G^A(E)], \quad \text{with} \quad \Gamma_{L/R} = \Gamma_{L/R} \mathbb{1}_c, \tag{2}
\]

where \( \mathbb{1}_c \) is a diagonal matrix with 1 at the entries corresponding to atoms in contact with the leads and 0 elsewhere. The current is given by

\[
I(V_b) = \frac{2e}{h} \int_{-\infty}^{\infty} d\varepsilon \left[ f_L(\varepsilon) - f_R(\varepsilon) \right] T(\varepsilon), \tag{3}
\]

where \( f_{L/R}(\varepsilon) = [1 + \exp((\varepsilon - \mu_{L/R})/(k_B T))]^{-1} \) are the Fermi distribution functions of the leads. The lead chemical potentials are given by \( \mu_L = \mu_0 + \eta V_b \), \( \mu_R = \mu_0 + (\eta - 1) V_b \), where \( \mu_0 = E_F \) is the common Fermi energy of the whole system at zero bias; \( V_b \) is the bias voltage with a possibly asymmetric drop across the nanotube, with the asymmetry encoded in the factor \( \eta \in [0,1] \). In the absence of spin-orbit coupling we assume the two spin channels to be independent and the spin degeneracy is accounted for by the prefactor \( 2 \). Eq. \( 3 \) immediately yields the differential conductance \( G_{\text{diff}} = dI/dV_b \). The linear conductance follows in the limit of vanishing bias, and it has the usual form

\[
G = 2e^2/h \int_{-\infty}^{\infty} d\varepsilon \left( -\frac{\partial f(\varepsilon)}{\partial \varepsilon} \right) \bigg|_{V_b=0} T(\varepsilon). \tag{4}
\]

We set the zero of the energy at the charge neutrality point of the nanotube. The CNT Fermi energy is then determined by the gate voltage, \( E_F = e\alpha V_g \). For \( T \approx 0 \) the derivative of the Fermi function can be approximated by the Dirac \( \delta \) and the linear conductance simplifies even further to

\[
G_{T=0} = \frac{2e^2}{h} T(E_F). \tag{5}
\]

In our setup the linear conductance at \( T = 0 \) is plotted as the orange lines in the Fig. 8(c), while the conductance at \( T = 8 \) K (red line) is evaluated through the Eq. \( 4 \). The Sagnac interference due to the trigonal warping begins to be visible below the energy of \(-0.2\) eV.

While the results in Fig. 8(c) are obtained for a perfect lattice, the breaking of CNT’s symmetries may induce another way to mix the two interference channels. Two such scenarios are illustrated in Fig. 9. The rotational symmetry may be broken by different tunneling into the suspended part of the CNT from the top and bottom (in contact with the leads) atoms. In a CNT of the zigzag class this results in mixing the valleys and introducing a modulation of the Fabry-Pérot interference. This is shown in Fig. 9(a),(b) for a \((12,9)\) CNT, with the weaker tunneling at the top of the CNT modelled through increased on-site potential of the contact atoms. In Fig. 9(b) the potential configuration at the right lead is reversed with respect to the left lead (physically this would correspond to a CNT which is twisted by half a turn between the left and right lead).

The rotational (and translational) symmetry could also be broken by the presence of adatoms in the CNT lattice. The conductance shown in Fig. 9(c) has been calculated assuming the presence of an adatom, at the distance of \( \sim 36 \) nm from the left contact, modelled by adding to the Hamiltonian a local on-site energy of 24 eV. The presence of another scattering center and the tiny length scale associated with the adatom-contact distance causes a large scale modulation of the Fabry-Perot interference in the momentum space.

In both cases the resulting modification of the Fabry-Pérot interference reproduces some of the features of the experimental data in Fig. 1 of the main text and in Fig. 5, hinting that both may be occurring in the experiment.
Because the Fabry-Pérot interference relies on phase coherence, raising the temperature destroys the oscillation through decoherence, leaving only the slow modulation of the conductance, see Figs. 8 and 9. Hence, higher temperature clearly does not introduce the four-time faster oscillations seen in the experiment. This suggests that the low temperature experimental result cannot be simply interpreted in terms of Fabry-Pérot interference of non-interacting electrons. What we observe in the experiment is rather the interference of quasi-particle excitations of an interacting system.

Figure 8. Single-particle interference. (a) Sketch of the calculated setup. The central system with length $L_c = 1.04 \mu m$ is contacted to wide band leads by the couplings $\Gamma_L, \Gamma_R$. (b) Low energy dispersion of a (20,5) CNT. The interference channels with higher (a) and lower momentum (b) are marked by the green arrows. Since this nanotube belongs to the armchair class, the two channels are not independent and can be scattered into each other (this intra-valley scattering is marked by orange arrows). (c) Zero-bias conductance of a (20,5) CNT with the length of 1.04 $\mu m$, comparable to the one in experiment. The orange line is the zero temperature conductance and displays the fast Fabry-Pérot oscillations. The red line shows the conductance at $T = 8K$; no oscillations are discernible close to the band gap (see inset), and only the slow Sagnac oscillation can be seen far from the band edge.

**Transport with interactions: coherent sequential tunneling for the four-fold degenerate Anderson model**

The single-particle spectrum of a finite CNT is organized into subsets of nearly fourfold-degenerate energy levels, with each quadruplet corresponding to one quantized longitudinal mode. Our starting point is thus the Hamiltonian of a 4-fold degenerate Anderson model, corresponding to one such quadruplet. It has the form $H = H_d + H_T + H_R + H_L$, where $H_T = H_{TL} + H_{TR}$ describes the tunneling coupling of the dot (d) to left (L) and right (R) electrodes. The latter are described as an ensemble of non-interacting electrons and captured by the terms $H_L$ and $H_R$. Finally, the dot Hamiltonian has the form

$$H_d = \sum_j \varepsilon_d n_j + U \sum_{j<k} n_j n_k + \sum_j \alpha e V_g n_j =: \varepsilon_d \sum_j n_j + U \sum_{j<k} n_j n_k,$$

where the indices run over the quantum numbers of each of the four degenerate states. Further, $\varepsilon_d$ is the single-particle energy, $V_g$ the gate potential, and $\alpha$ is the lever arm of the quantum dot. In a carbon nanotube quantum dot the four-fold degeneracy arises from the presence of both valley and spin, but here we will number the degrees of freedom generally by $j = 1, 2, 3, 4$. The Coulomb interaction is denoted by $U$ and it corresponds to the charging energy $E_C$ in the main text. In order to recover the other longitudinal modes of the CNT, we will later extend this Hamiltonian to a sum of such 4-fold degenerate levels, separated by an energy $\Delta E$ which we shall take, following the experiment, to be $\Delta E \approx U/2$. 


The energies of the many-body states with $N = 0,...4$ electrons are $E(N) = N\bar{\epsilon}_d + N(N - 1)U/2$. The chemical potential for each occupation $N$ is then

$$
\mu(N) = E(N) - E(N - 1) = \bar{\epsilon}_d + (N - 1)U,
$$

$N = 1,...,4$. \hfill (7)

In the following we shall use the equation of motion technique (EOM) originally proposed in Ref. [44] for the spinful Anderson model to evaluate the retarded single particle Green’s functions $\tilde{G}_R(i,\varepsilon)$. Their knowledge will give us first indications for the current through the four-fold degenerate interacting Anderson model. In fact with $\nu(i,\varepsilon) = -2\text{Im} \tilde{G}_R(i,\varepsilon)$ being the spectral function of level $i$, the current follows from the Meir and Wingreen formula [45]

$$
I = \frac{e}{\hbar} \sum_{i=1}^{4} \int_{-\infty}^{\infty} \frac{d\varepsilon}{\Gamma_{Li}} \frac{\Gamma_{Li}}{\Gamma_{Li} + \Gamma_{Ri}} \nu(i,\varepsilon) [f_L(\varepsilon) - f_R(\varepsilon)].
$$

$\hfill (8)$

The coupling asymmetry parameter for the lead $\alpha$ and level $i$ is given by $\gamma_{\alpha i} = \Gamma_{\alpha i}/\Gamma_i$, with $\Gamma_i = \sum_{\alpha=L,R} \Gamma_{\alpha i}$. The parameter range of interest for the experiment, $U \simeq \Gamma \gg k_B T$, is highly non-trivial and in practice not accessible within the truncation schemes proposed in Ref. [44]. However, the EOM methods enables one to get the exact current in the non-interacting case; further, it well describes the tunneling dynamics in the coherent tunneling regime $U \simeq \Gamma \geq k_B T$, as discussed below.

**Atomic limit**

For a 4-fold isolated system with four single particle states, i.e., $H = H_d$, the equation of motion procedure closes after four iterations, yielding the exact set of coupled equations

$$
(\varepsilon - \mu(1) + i\eta) \tilde{G}_R(i,\varepsilon) = 1 + U \tilde{D}_R(i,\varepsilon),
$$

$\hfill (9a)$
(\varepsilon - \mu(2 + \imath \eta)) \bar{F}^R(i, \varepsilon) = \sum_{j \neq i} \langle n_j \rangle + U \tilde{F}^R(i, \varepsilon), \quad (9b)

(\varepsilon - \mu(3 + \imath \eta)) \bar{F}^R(i, \varepsilon) = \sum_{p \neq j, i \neq i} \langle n_p n_j \rangle + U \tilde{H}^R(i, \varepsilon), \quad (9c)

(\varepsilon - \mu(4 + \imath \eta)) \tilde{H}^R(i, \varepsilon) = \sum_{l \neq p, j \neq j, i \neq i} \sum_{n_p n_j} \langle n_l n_p n_j \rangle, \quad (9d)

with \( \eta = 0^+ \) a small infinitesimal. The tilded Green’s functions in the energy domain are the Fourier transforms of the time-dependent Green’s functions

\[
G^R(i, t) = -\frac{i}{\hbar} \theta(t) \langle \{c_i(t), c_i^\dagger\} \rangle, \quad (10a)
\]

\[
D^R(i, t) = -\frac{i}{\hbar} \theta(t) \sum_{j \neq i} \langle \{n_j c_i(t), c_i^\dagger\} \rangle, \quad (10b)
\]

\[
F^R(i, t) = -\frac{i}{\hbar} \theta(t) \sum_{j \neq i} \sum_{p \neq i, j} \langle \{n_j n_p c_i(t), c_i^\dagger\} \rangle, \quad (10c)
\]

\[
H^R(i, t) = -\frac{i}{\hbar} \theta(t) \sum_{j \neq i} \sum_{p \neq j, i} \sum_{m \neq p, j, i} \langle \{n_m n_p n_j c_i(t), c_i^\dagger\} \rangle. \quad (10d)
\]

Each of the four Green’s functions describes adding an electron to the level \( i \) if either the dot is empty \((G^R(i, t))\), or already hosts one \((D^R)\), two \((F^R)\) or three \((H^R)\) particles. Solving this set of coupled equations yields the single particle Green’s function \( \tilde{G}^R(i, \varepsilon) \), which can be conveniently expressed in the form

\[
\tilde{G}^R(i, \varepsilon) = \sum_{n=1}^{4} \frac{a_n(i)}{\varepsilon - \mu(n) + \imath \eta},
\]

with the coefficients \( a_n \) obeying the sum rule \( \sum_n a_n = 1 \). Let us introduce the occupation numbers

\[
\bar{N}_{1\Sigma} := \sum_{j \neq i} \langle n_j \rangle,
\]

\[
\bar{N}_{2\Sigma} := \sum_{j \neq i} \sum_{p \neq j, i} \langle n_j n_p \rangle,
\]

\[
\bar{N}_{3\Sigma} := \sum_{j \neq i} \sum_{p \neq j, i} \sum_{l \neq p, j, i} \langle n_j n_p n_l \rangle.
\]

Then in terms of such occupations the coefficients \( a_n(i) \) are given by

\[
a_1(i) = 1 - \bar{N}_{1\Sigma}(i) + \frac{\bar{N}_{2\Sigma}(i)}{2} - \frac{\bar{N}_{3\Sigma}(i)}{6}, \quad (13a)
\]

\[
a_2(i) = \bar{N}_{1\Sigma}(i) - \bar{N}_{2\Sigma}(i) + \frac{\bar{N}_{3\Sigma}(i)}{2}, \quad (13b)
\]

\[
a_3(i) = \bar{N}_{2\Sigma}(i) - \bar{N}_{3\Sigma}(i), \quad a_4(i) = \frac{\bar{N}_{3\Sigma}(i)}{6}. \quad (13c)
\]

In equilibrium it is possible to evaluate the expectation values \( \bar{N}_{n\Sigma}(i) \) using the Lehmann representation \([46]\). One finds

\[
\langle n_i \rangle = \int \frac{d\varepsilon}{2\pi} (-2 \Im \tilde{G}^R(i, \varepsilon)) f(\varepsilon), \quad (14)
\]
where $f(\varepsilon) = [1 + \exp((\varepsilon - \mu_0)/k_B T)]^{-1}$. Note that since we are now working with interacting particles, we replaced $E_F$ with the reference chemical potential $\mu_0$. Using the expression of the $G^R(i, \varepsilon)$ from Eq. (11), we find
\[
\langle n_i \rangle = \int \frac{d\varepsilon}{2\pi} \nu(i, \varepsilon) f(\varepsilon) = \sum_{n=1}^{4} a_n(i) \int d\varepsilon f(\varepsilon) \delta(\varepsilon - \mu(n)) = \sum_{n=1}^{4} a_n(i) f(\mu(n)).
\]

Similar relations hold for the higher Green’s functions. Introducing the shorthand notation $f(\mu(n)) =: f_n$, we find
\[
\sum_{j \neq i} \langle n_j n_i \rangle = \int \frac{d\varepsilon}{2\pi} (-2 \text{Im} \tilde{D}^R(i, \varepsilon)) f(\varepsilon) = a_2(i) f_2 + 2a_3(i) f_3 + 3a_4(i) f_4,
\]
\[
\sum_{p \neq j, i} \sum_{n_i} \langle n_p n_j n_i \rangle = \int \frac{d\varepsilon}{2\pi} (-2 \text{Im} \tilde{F}^R(i, \varepsilon)) f(\varepsilon) = 2a_3(i) f_3 + 6a_4(i) f_4,
\]
\[
\sum_{m \neq p, j, i} \sum_{n_p} \sum_{n_j} \langle n_m n_p n_j n_i \rangle = \int \frac{d\varepsilon}{2\pi} (-2 \text{Im} \tilde{H}^R(i, \varepsilon)) f(\varepsilon) = 6a_4(i) f_4.
\]

For a degenerate model the single particle occupation $\tilde{N}_1 := \langle n_i \rangle$ is independent of the index $i$. Likewise for the double and triple occupations $\tilde{N}_2 := \langle n_j n_i \rangle$ and $\tilde{N}_3 := \langle n_k n_k n_k \rangle$. This leads to the final result
\[
a_1(V_g) = 1 - [3\tilde{N}_1(V_g) - 3\tilde{N}_2(V_g) + \tilde{N}_3(V_g)],
\]
\[
a_2(V_g) = 3\tilde{N}_1(V_g) - 6\tilde{N}_2(V_g) + 3\tilde{N}_3(V_g),
\]
\[
a_3(V_g) = 3(\tilde{N}_2(V_g) - \tilde{N}_3(V_g)),
\]
\[
a_4(V_g) = \tilde{N}_3(V_g)
\]

together with
\[
\tilde{N}_1(V_g) = f_1 \left\{ 1 + 3(f_1 - f_2) - \frac{3f_2(f_1 - 2f_2 + f_3)}{1 + 2f_2 - 2f_3 - d(V_g)} + \frac{f_2f_3(f_1 - 3f_2 + 3f_3 - f_4)}{(1 + f_3 - f_4)(1 + 2f_2 - 2f_3 - d(V_g))} \right\}^{-1},
\]
\[
\tilde{N}_2(V_g) = \tilde{N}_1(V_g) \frac{f_2}{1 + 2f_2 - 2f_3 - d(V_g)},
\]
\[
\tilde{N}_3(V_g) = \tilde{N}_2(V_g) \frac{f_3}{1 + f_3 - f_4},
\]
\[
d(V_g) = \frac{f_3(f_2 - 2f_3 + f_4)}{1 + f_3 + f_4}.
\]

**Coherent sequential tunneling approximation**

When considering the influence of the coupling $H_T$ to external leads, the set of equations for the single particle Green’s function does not close anymore. This requires truncation and approximation schemes to properly account for the interplay of interactions and tunneling. We assume that the quantum numbers are conserved by the tunneling, i.e., $H_{T\alpha} = \sum_{i,k} t_{\alpha k,i} c_{\alpha k,i}^\dagger c_{\alpha k,i} + \text{h.c.}$, with $\alpha = L, R$. Further, $c_{\alpha k,i}^\dagger$ creates an electron in the dot and lead, respectively. The quantity $t_{\alpha k,i}$ describes the tunneling between the lead state with its continuous degree of freedom $k$ and the quantum number $i$. The dispersion of the states with quantum numbers $k,i$ in the lead $\alpha$ is given by $\varepsilon_{\alpha k,i}$. The most crude approximation, which is exact for a noninteracting Anderson model ($U = 0$) as well as in the atomic limit ($\Gamma \to 0^+$), amounts to truncating the hierarchy of equations for the higher order Green’s function $D^R$, $F^R$ and $H^R$ by neglecting some level non-conserving terms (spin-flip terms in the simpler spin-degenerate Anderson model) [46]. In this way the coupling to the leads enters only through a self-energy $\Sigma^R$, independent of $U$ and $T$, and defined by
\[
\Sigma^R(i, \varepsilon) = \sum_{\alpha k} \frac{|t_{\alpha k,i}|^2}{\varepsilon - \varepsilon_{\alpha k,i}}, \quad \alpha = L, R.
\]
In this approximation one finds

\[ (\varepsilon - \mu(1) + \Sigma^R(i,\varepsilon)) \tilde{G}^R(i,\varepsilon) = 1 + U\tilde{D}^R(i,\varepsilon), \]

\[ (\varepsilon - \mu(2) + \Sigma^R(i,\varepsilon)) \tilde{D}^R(i,\varepsilon) = \sum_{j \neq i} \langle n_j \rangle + U\tilde{F}^R(i,\varepsilon), \]

\[ (\varepsilon - \mu(3) + \Sigma^R(i,\varepsilon)) \tilde{F}^R(i,\varepsilon) = \sum_{p \neq j, i} \sum \langle n_p n_j \rangle + U\tilde{H}^R(i,\varepsilon), \]

\[ (\varepsilon - \mu(4) + \Sigma^R(i,\varepsilon)) \tilde{H}^R(i,\varepsilon) = \sum_{l \neq p, j, i} \sum \sum \langle n_l n_p n_j \rangle. \]

In the wide-band limit one finds \( \Sigma^R(i,\varepsilon) = -i(\Gamma_L + \Gamma_R)/2 = -i\Gamma/2 \). Hence, comparing with the results from the atomic limit, we obtain within this simple scheme that the leads induce a temperature independent broadening \( \Gamma \).

The Green’s function then read

\[ \tilde{G}^R(i,\varepsilon) = \frac{4}{\pi} \frac{a_n}{\varepsilon - \mu(n) + i\Gamma/2}, \]

with the coefficients \( a_n \) defined as in the atomic limit through Eqs. (17). However, due to the Lorentzian broadening of the Green’s functions, cf. Eqs. (20) and (21), the functions \( f_n \) yielding the coefficients \( N_n \) in Eqs. (18) should be replaced by \( F_n := F(\mu(n)) \), where

\[ F(\mu(n)) = \int \frac{de}{2\pi} f(\varepsilon)(-2) \text{Im} \left( \frac{1}{\varepsilon - \mu(n) + i\Gamma/2} \right) = \frac{1}{2} \frac{1}{\pi} \text{Im} \Psi \left( \frac{1}{2} + i\frac{\mu(n) - i\Gamma/2 - \mu_0}{2\pi k_B T} \right), \]

where \( \Psi(x) \) is the digamma function. The conductance within this Lorentzian scheme is shown in Fig. 10 for various values of the ratio \( \Gamma/U \) and varying temperatures. Similar to the single-particle interference discussed in the previous section, also in this case the conductance is only moderately dependent on temperature. In particular, a stronger increase of the conductance in the central valley by decreasing temperature, similar to the experimental observations, is not seen (the curves for \( k_B T/U = 0.01 \) and \( k_B T/U = 0.1 \) are essentially identical). This feature is well known from the studies of the spinful Anderson model within the EOM approach. A temperature dependent self-energy requires accounting for some of the neglected spin-flip contributions [44, 47]. However, an extension which recovers the unitary Kondo limit reached at low temperatures is already very intricate for the spinful case [47], and becomes intractable for the four-fold degenerate Anderson model. This generalisation is beyond the scope of this work.

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Transport through a multilevel Anderson model in the coherent sequential tunneling approximation. Left column: transport through an Anderson quantum dot with a 4-fold (spin and valley) degenerate single-particle energy level. With increasing broadening $\Gamma$ (approaching the non-interacting limit for $\Gamma/U = 2.5$) the four peaks merge into one, but temperature affects the conductance only quantitatively. Right column: conductance through a series of 4-fold degenerate shells with inter-shell spacing $\Delta E = 0.5U$ and $k_B T/U = 0.1$. The central pattern is reminiscent of the experimental zero-bias trace at $T = 4$ K shown in the inset of Fig. 7c; neighboring shells are enhancing the conductance maxima, but the structure of two higher and two lower peaks remains visible. In other words, an enhancement of the central valley similar to what is seen in the experiment is not captured by the coherent approximation.

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