Signatures of Klein tunneling in disordered graphene p-n-p junctions

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We present a method for obtaining quantum transport properties in graphene that uniquely combines three crucial features: microscopic treatment of charge disorder, fully quantum mechanical analysis of transport, and the ability to model experimentally relevant system sizes. As a pertinent application we study the disorder dependence of Klein tunneling dominated transport in p-n-p junctions. Both the resistance and the Fano factor show broad resonance peaks due to the presence of quasi bound states. This feature is washed out by the disorder when the mean free path becomes of the order of the distance between the two p-n interfaces.

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A realistic modeling of quantum transport in graphene [1] has two main requirements: microscopic treatment of disorder and a fully quantum mechanical calculation of transport, taking into account the finite (but large with respect to the lattice constant) size of the samples. This is especially important close to the Dirac point [2], the charge neutrality point at which transport is dominated by evanescent modes and conventional analytical methods fail, and for interference effects. In this letter we report on the first such calculation in a study of the resistance oscillations in a disordered graphene p-n-p junction which incorporates both realistic Coulomb disorder and quantum mechanical transport.

The physics of the p-n-p junction is governed by “Klein tunneling”, an effect first discovered in the context of a relativistic particle tunneling through a potential barrier with a height comparable to its rest energy [3]. Since the occurrence of relativistic Klein tunneling is a crucial feature distinguishing graphene from ordinary electronic materials, a clear understanding of how to unambiguously observe Klein tunneling in realistic graphene samples, where disorder is unavoidably present, is important. In graphene, transmission through a potential barrier has a pronounced angular dependence, with perfect transmission for perpendicular incidence and a quick decrease of transmission probability for finite angle of incidence [4]. As a consequence, electrons can be confined within a single potential barrier, such as the p-n-p junction [5]. Resonant tunneling through those confined states then leads to pronounced oscillations in the resistance as a function of system parameters [5, 6, 7].

The existence of this phenomenon relies heavily on the presence of a well defined interface between the p and n regions. Early experiments failed to reproduce the predicted resistance oscillations [8], presumably because of too much disorder. More recent experiments that seek to minimize the effects of disorder have been more successful [9, 10, 11]. However a theoretical understanding of whether or to what extent these experiments are really observing Klein tunneling phenomena and the issue of the experimental conditions needed to see Klein tunneling unambiguously have remained open.

There exists strong evidence that disorder in current experiments is dominated by remote charged impurities [12, 13] invariably present in the graphene environment. Due to the long range character of the Coulomb potential of impurity and gate charges, all scattering potentials and charge densities in the graphene sheet vary slowly on the scale of the lattice constant. This fact leads to two simplifications that make an accurate modeling of current experimental setups [9, 10, 11] possible. First, the Thomas-Fermi-Dirac (TFD) self-consistent density functional method is sufficient to obtain the ground state carrier density in the presence of disorder [14, 15]. Second, intervalley coupling can be neglected for a smooth scattering potential, so that one can model transport using the single valley Dirac Hamiltonian. The transfer matrix method of Refs. 16 and 17 is ideal for this purpose and can be used for experimentally relevant system sizes, as opposed to tight binding models which are generally limited to small system sizes.

Our calculation thus consists of two steps. In the first step the ground state carrier density n(\mathbf{r}) and the corresponding screened Coulomb potential \(V_{sc}\) is obtained using the TFD approach [14]. In the second step the resistance and Fano factor, the ratio of shot noise power and electrical current, are calculated using \(V_{sc}\) as an input into the fully quantum mechanical transfer matrix approach [16, 17]. We first describe the two steps of our method in more detail and then turn to the results.

A schematic of the p-n-p junction under consideration is shown in Fig. 1a. A top gate of length \(L_{tg}\) is located a distance \(d_{tg}\) above the graphene layer — in an air-bridge setup the medium in between is air with dielectric constant \(\epsilon_1 = 1\). A back gate that controls the average carrier density is separated from the graphene by a SiO₂
substrate (dielectric constant $\epsilon_2 = 4$). The effective dielectric constant in this setup is $\epsilon = (\epsilon_1 + \epsilon_2)/2$. Together, the voltages on these two gates define the junction. Examples of potential profiles are shown in Fig. 1b. A number of charged impurities is trapped in the substrate and just below the graphene layer. We model this by a random distribution $C(\mathbf{r})$ of impurity charges at a fixed distance $d$ with properties

$$\langle C(\mathbf{r}) \rangle = 0; \quad \langle C(\mathbf{r}_1) C(\mathbf{r}_2) \rangle = n_{\text{imp}} \delta(\mathbf{r}_2 - \mathbf{r}_1);$$

with $n_{\text{imp}}$ the average 2D charge impurity density.

In the TFD approximation, the ground state carrier density $n(\mathbf{r})$ in the graphene layer is obtained by minimizing the energy functional

$$E[n] = \int d^2 \mathbf{r} n(\mathbf{r}) \left[ \frac{2}{3} \hbar v_F |\pi n(\mathbf{r})|^{1/2} + V_s(\mathbf{r}) \right] + E_{\text{xc}}[n]$$

with respect to $n(\mathbf{r})$. Here $v_F \approx 10^6$ m/s is the Fermi velocity for graphene,

$$V_s(\mathbf{r}) = \hbar v_F r_s \left[ V_d(\mathbf{r}) + V_{tg}(\mathbf{r}) + \frac{1}{2} \int d^2 \mathbf{r}' \frac{n(\mathbf{r}')}{|\mathbf{r} - \mathbf{r}'|} \right] - \hbar v_F \mu,$$

with $r_s \approx e^2/\hbar v_F \epsilon$, $V_d$ and $V_{tg}$ the potentials induced by the impurity density $C(\mathbf{r})$ and the top gate, respectively, and $\mu$ the chemical potential. $E_{\text{xc}}[n]$ is the exchange correlation energy (see [13] [15] for details) which we include even though it gives only minor quantitative corrections to the transport properties of the $p-n-p$ junction. The minimization is subject to the constraint

$$\int_{A'} n(\mathbf{r}) d^2 \mathbf{r} = n_{\text{tg}},$$

where $A'$ is the area of the sample away from the top gate. The constraint is enforced self-consistently by varying $\mu$. The electrostatic potential $V_{tg}$ is expressed in terms of the top gate charge density $n_{tg}$,

$$V_{tg}(\mathbf{r}) = \int d^2 \mathbf{r}' \frac{n_{tg}(\mathbf{r}')}{(|\mathbf{r} - \mathbf{r}'|^2 + d_{tg}^2)^{1/2}},$$

and is obtained self-consistently, for a fixed voltage difference between the top gate and graphene, $\varphi_{tg}$, by requiring that in the region below the top gate $\Delta n_{tg} = n_{tg}(\mathbf{r}) - n(\mathbf{r}) = C_{tg} \varphi_{tg}$, where $C_{tg}$ is the top gate capacitance.

The potential $V_s$ defines a scattering problem through the Dirac Hamiltonian

$$H = v_F \mathbf{p} \cdot \sigma + V_s(\mathbf{r})$$

with $\sigma = (\sigma_x, \sigma_y)$ the Pauli matrices. The Schrödinger equation $H \psi = 0$ generates a transfer matrix $M$, which relates the wavefunction at $x = 0$ to the one at $x = L$, $\psi_L = M \psi_0$. In order to numerically calculate $M$ we divide the interval $(0, L)$ into $N$ equal subintervals of length $\delta x = L/N$ and calculate the transfer matrix in each subinterval in the Born approximation. This gives

$$M = \prod_{n=1}^N e^{-\frac{i\delta x}{\hbar} \sigma_z} e^{-i u_n \sigma_x} e^{-\frac{i\delta x}{\hbar} \sigma_z},$$

where

$$u_n(y) = \frac{1}{\hbar v_F} \int_{(n-1)\delta x}^{n\delta x} dx V_s(x, y).$$

We then take the limit $N \rightarrow \infty$ in which the Born approximation becomes exact. From the transfer matrix $M$ we calculate the matrix $t$ of transmission amplitudes, which in turn gives us the two-terminal conductance, which in turn gives us the two-terminal conductance and Fano factor as

$$G = \frac{1}{4 \epsilon^2} \text{tr} \left[ (1 - t t^\dagger) t t^\dagger \right]/\text{tr} t t^\dagger.$$

We now turn to our results. We first consider the clean case $n_{\text{imp}} = 0$. By varying the top gate potential, the junction type can be varied from $p-n-p$ to $p-p-p$ (or $n-p-n$ to $n-n-n$ depending on the back gate voltage), as shown in Fig. 1b. The dependence of the resistance $R$ and Fano factor $F$ on the two gate voltages is shown in Fig. 2a-b. In the $p-p-p$ ($n-n-n$) region of the upper left (lower right) corner the resistance is small compared to the resistance of the $p-n-p$ ($n-p-n$) region of the upper right (lower left) corner. Pronounced oscillations in the resistance are seen in the $p-n-p$ and $n-p-n$ regions. The behavior of the Fano factor mostly follows that of the resistance.

The resistance oscillations can be understood as arising from resonances through quasi bound states inside the barrier created by the two $p-n$ interfaces [5] [6] [7]. The larger the transverse momentum $q_y$, the larger the incident angle for scattering off the barrier and the tunneling amplitude decreases. The broad oscillations arise from resonant tunneling of the few modes with smallest $q_y$. Larger $q_y$ give rise to very narrow resonance seen in...
The double peak structure of the Fano factor resonances reflects the increase of the transmission probability from zero to one and back to zero and the fact that the Fano factor, where the average was taken over 10^3 realizations, is much larger than the system size.

We do emphasize that for weak impurity density (from bottom to top n_{imp} = 0, 1, 2.5, 5, 10, and 15 × 10^{11} cm^{-2}). The disordered curves have been shifted for clarity, with dashed lines showing their zero.

(d) The top gate density dependence of the resistance at a narrow resonance for different values of the impurity density.

In contrast to the narrow resonances the resistance minima for the broader oscillations do not depend on boundary conditions and can be estimated by a semiclassical argument. One can think of the p-n-p junction as a Fabry-Perot etalon in which waves scattered off the two p-n interfaces interfere destructively if the WKB phase

\[ \theta_{\text{WKB}} = - \int_{x_1}^{x_2} V_{sc}(x', y) dx' \equiv 0 \pmod{\pi}, \]

with the turning points \( x_i \) defined by the condition \( V_{sc}(x_i, y) = 0, \ i = 1, 2 \). The positions of the resistance minima that follow from this argument are indicated by arrows in Fig. 2.

We now consider the effect of disorder. A single realization of the scattering potential \( V_{sc} \) in a disordered p-n-p junction is shown in Fig. 3a. For realistic impurity concentrations, the correlation length of \( V_{sc} \) is \( \sim 10 \) nm.

While one can still clearly make out the different \( p \) and \( n \) regions in the presence of disorder, the boundary between the two is no longer as sharp, leading to a weaker confinement of particles inside the barrier. The effects of Klein tunneling are thus expected to be suppressed.

Figures 3b-c give the transport properties as a function of top gate voltage for a single disorder realization and for different impurity densities. Mesoscopic fluctuations due to the disorder are superimposed on the oscillations from the Klein tunneling. As the impurity strength increases these fluctuations become stronger, eventually dominating the signal. The Fano factor is more sensitive to disorder than the resistance. The effect of disorder on a narrow resonance is shown in Fig. 3d.

In Fig. 4 we show the disorder averaged resistance and Fano factor, where the average was taken over 10^3 disorder realizations. We find that the clean limit broad quantum oscillations survive for impurity densities up to \( n_{\text{imp}} \sim 10^{12} \) cm^{-2} while the narrow resonances disappear at \( n_{\text{imp}} \sim 10^{11} \) cm^{-2}. We do emphasize that for weak enough disorder the narrow resonances will be present, and it is therefore conceivable that they may actually be observed at low enough temperatures, when the phase coherence length is much larger than the system size.

As in the case of the single disorder realization, the dependence of the Fano factor on the disorder strength is similar to that of the resistance. The dependence of the
resistance on the back gate voltage for a fixed impurity density is shown in Fig. 4. The main effect of the back gate is an overall reduction of the resistance and a shift of the top gate density at which the central region changes polarity.

One expects impurity scattering effects to start dominating the ballistic oscillations when the mean free path $l$ becomes comparable to the sample width $W$ (for the narrow resonances) or the length of the central $n$ region (for the broad oscillations). The mean free path $l$ can be estimated through its relation to the conductivity $\sigma$, $l = (\hbar/2e^2)(\sigma/k_F)$, where $k_F$ is the Fermi wavevector. Away from the Dirac point, $\sigma = (2e^2/h)(n)/n_{imp} f(r, x, d)$ \[20\], where $f(r, x, d)$ is a known function of $r_s$ and $d$ only \[21\]. Setting $\langle n \rangle = n_{bg}$ and using the impurity densities quoted above, we find that the narrow resonances vanish for $l \approx 400$ nm, which is the same order of magnitude as expected. The impurity density at which the broad oscillations disappear corresponds to a mean free path $l \approx 40$ nm. This is roughly the same as the length between the two $p$-$n$ interfaces, see Fig. 1 and consistent with the experimental results of Refs. \[9\] and \[10\].

In summary, we have presented a powerful theoretical method that is generally applicable to transport problems in realistic graphene samples where the disorder is dominated by charge impurities and transport properties need to be obtained fully quantum mechanically. We have applied this method to understand the effects of disorder on transport through $p$-$n$-$p$ junctions. The crossover from the ballistic transport governed by Klein tunneling, to the disordered diffusive transport is found to take place as the mean free path becomes of the order of the distance between the two $p$-$n$ interfaces consistent with recent experiments \[9\] \[10\].

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