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Quantum corrections to conductivity in graphene with vacancies

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ABSTRACT

In this work, different regions of a graphene device were exposed to a 30 keV helium ion beam creating a series of alternating strips of vacancy-type defects and pristine graphene. From magnetoconductance measurements as function of temperature, density of carriers and density of strips we show that the electron-electron interaction is important to explain the logarithmic quantum corrections to the Drude conductivity in graphene with vacancies. It is known that vacancies in graphene behave as local magnetic moments that interact with the conduction electrons and leads to a logarithmic correction to the conductance through the Kondo effect. However, our work shows that it is necessary to account for the non-homogeneity of the sample to avoid misinterpretations about the Kondo physics due the difficulties in separating the electron-electron interaction from the Kondo effect.

1. Introduction

The unique properties of graphene such as transparence, high mobility, biocompatibility and high mechanical strength have put this material in evidence over the last decade as a promising material for technological applications [1]. However, one of the main challenges for the modern electronic industry is the controlled fabrication of high-performance devices based on 2D materials such as graphene. Some of the required processes for device processing unavoidably introduce defects that have a deleterious effect in some properties of graphene such as conductivity and carrier mobility. It is therefore of great importance to understand how defects affect the electronic transport in order to improve the quality of graphene-based devices.

In particular, recent magnetotransport studies of graphene devices intentionally modified by vacancies created in the lattice showed a logarithmic behavior of the conductivity with temperature [2–4]. This behavior was attributed to quantum corrections in the classical Drude transport model that must be taken into account in disordered systems. In these systems, both the well-known weak localization (WL) [5–8] and electron-electron interaction (EEI) [9–13] lead to a logarithmic temperature dependence of the conductivity. Other works have shown that in graphene this dependence can also be due to interactions of the conducting electron with local magnetic moments produced by vacancies – an effect known as Kondo effect [3,14,15].

In this work, we report on a study of the temperature dependence of the conductivity of a graphene device in which disorder was created by bombardment with a helium ion beam, which introduces mainly vacancies in the lattice [16,17]. Different density of defects were introduced in separate regions of a graphene device and the conductivity of each region was measured for varying temperatures and gate voltages. Our results show that if non-homogeneity is not taken into account, the weak localization and electron-electron interactions on their own cannot explain the quantum correction to the conductivity dependence with temperature, giving a wrong result that could be misinterpreted as a Kondo effect.

2. Sample preparation

Our device, shown in Fig. 1 (a), consists of a graphene monolayer deposited on top of a heavily-doped Si- substrate coated with a 300 nm-thick SiO_2 and processed as a Hall bar using conventional nanofabrication techniques [4]. A series of alternating bands of pristine graphene and defect strips were created by scanning a 30 keV focused helium ion beam perpendicularly to the main channel of the device. Each defect strip is 5 nm wide and has a density of defects $n_d \sim 10^{13}$ cm^{-12}. The strips are uniformly distributed over three regions – D1, D2 and D3 – located between the contacts of the Hall bar. Each region is 1 μm wide and the distances between strips in regions D1, D2 and D3 are 40 nm, 20 nm and 10 nm, respectively. This alternating bands of pristine graphene and graphene with defects introduces a controlled inhomogeneity

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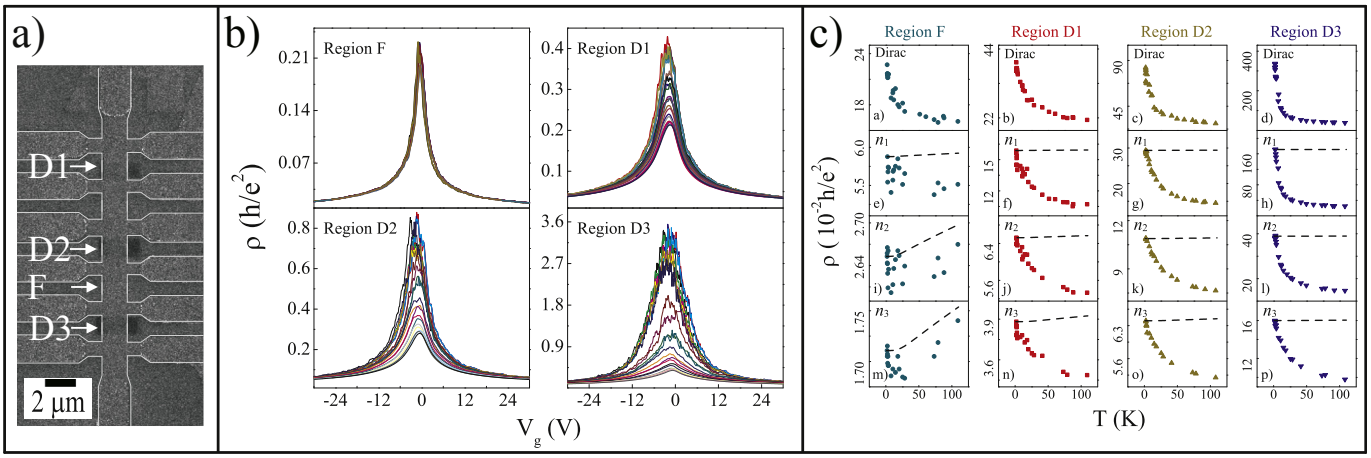

Fig. 1. a) Helium-ion microscopy image of the graphene device processed as a Hall bar. The darker regions (D1, D2, D3) have been exposed to different doses of helium ions. b) Resistivity as a function of gate voltage for each device region (from bottom to top, 1.8 up to 110 K) c) Resistivity as a function of temperature for each device region and for different carrier densities $n_1 = (0.4 \pm 0.1) \times 10^{12}\,\mathrm{cm}^{-2}$, $n_2 = (1.2 \pm 0.1) \times 10^{12}\,\mathrm{cm}^{-2}$ and $n_3 = (2.3 \pm 0.1) \times 10^{12}\,\mathrm{cm}^{-2}$, for $B = 0$ T. The dashed lines are the calculated contributions to the resistivity due to electron-acoustic phonon scattering and are shifted by the value of the resistivity measured at the lowest temperature $T \sim 1.8$ K.

within the modified regions. The control region F was not exposed to the He ions.

3. Results and discussion

Fig. 1 (b) shows measurements of the resistivity ρ as function of gate voltage for regions F, D1, D2 and D3, at zero magnetic field. While in pristine graphene, region F, the resistivity depends weakly on temperature, in modified graphene, regions D1, D2 and D3, it shows an increasing temperature dependence. Fig. 1 (c) shows the temperature dependence of the resistivity, for all measured regions, at different carrier densities n. These temperature dependences cannot be explained by the classical Drude model for resistivity. In order to identify the quantum corrections to this model, we have calculated and added each one in a step-by-step manner, and compared the results with our measurements.

As a first attempt, one could attribute this change in resistivity with temperature to electron-acoustic phonon scattering [11,18,19] which is given by,

$$\rho_{e-ph}(T) = \frac{1}{\pi}\frac{D_a^2 E_F^2}{\rho_s v_f \hbar^3 v_F^3}\int_0^{\pi}\frac{2v_f\hbar k_F}{Tk_B}\cos^2\frac{\theta}{2}\sin^4\frac{\theta}{2}\sinh^{-2}\left(\frac{2v_f\hbar k_F}{k_B T}\sin\frac{\theta}{2}\right)d\theta, \quad (1)$$

where θ is the scattering angle, $v_f \approx 2 \times 10^4\,\mathrm{m\,s^{-1}}$ is the sound speed, $\rho_s \approx 7.6 \times 10^{-7}\,\mathrm{kg\,m^{-2}}$ is the density, and $D_a \approx 18$ eV is the deformation potential of graphene [18,20]. Equation (1) was calculated numerically and plotted as dashed lines in Fig. 1 (c). In each plot, we have added to ρ_{e-ph} the value of the resistivity measured at the lowest temperature. These results show that electron-phonon scattering is significant only for region F (pristine graphene), at higher carrier densities. However, as the density of defects increases, the effect of electron-phonon interactions becomes increasingly negligible, even at the highest carrier density.

The results in Fig. 1 (c) show that electron-phonon scattering cannot be the only mechanism to explain the temperature dependence of the conductivity of the sample. Then, we have calculated an electron-phonon scattering-independent conductivity $\sigma_c = (\rho - \rho_{e-ph})^{-1}$ by subtracting from the measured resistivity the contribution of the electron-phonon scattering at each temperature [11]. This conductivity can be written as a sum of the Drude conductivity σ_0 with a quantum correction $\sigma_c^*(T)$, *i.e.*,

$$\sigma_c(T) = \sigma_0 + \sigma_c^*(T). \quad (2)$$

In order to study the temperature dependence of quantum correction, we calculate the $\Delta\sigma_c(T) = \sigma_c(T) - \sigma_c(T_0)$, where $T_0 \sim 7$ K is the lowest temperature at which we observe a Mott regime of transport [4]. As σ_0 is temperature independent, we can write

$$\Delta\sigma_c(T) = \Delta\sigma_c^*(T), \quad (3)$$

which includes all the quantum corrections at temperature T except the phonon contribution, which has already been removed. The plots of $\Delta\sigma_C(T)$ in Fig. 2 show that a logarithmic temperature dependence is still present.

In the next step, we analyze if the WL effect, which also determines a logarithmic temperature dependence to the conductivity, can account for the remaining quantum correction to our data. The contribution σ_{WL}^* of WL to the conductivity is given by Refs. [5–8].

$$\sigma_{WL}^*(B=0,T) = \frac{e^2}{\pi h}\left(2\ln\left(\frac{\tau_\phi(T)}{2\tau_p}\right) - \ln\left(1+\frac{2\tau_\phi(T)}{\tau_{iv}}\right) - 2\ln\left(1+\frac{\tau_\phi(T)}{\tau_{iv}} + \frac{\tau_\phi(T)}{\tau_*}\right)\right), \quad (4)$$

where T_ϕ, T_{iv} and T_* are, respectively, the phase-coherence time, the intervaley and intravaley scattering times. We have obtained these characteristic times (see Ref. [4]) at different temperatures and carrier

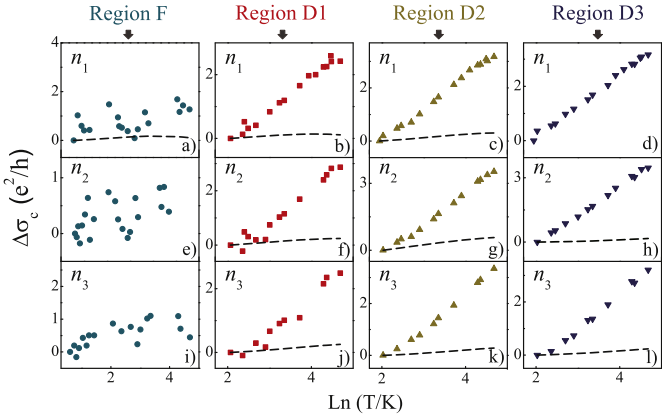

Fig. 2. Plots of the electron-phonon scattering independent conductivity (dots) minus its value at $T_0 = 7$ K. The dashed lines are the corrections to the conductivity due to the weak-localization.