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Density and localized states’ impact on amorphous carbon electron transport mechanisms

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This work discusses the electron transport mechanisms that we obtained as a function of the density of amorphous carbon (a-C) ultra-thin films. We calculated the density of states (total and projected), degree of electronic states’ localization, and transmission function using the density functional theory and nonequilibrium Green’s functions method. We generated 25 sample a-C structures using ab-initio molecular dynamics within the isothermal-isobaric ensemble. We identified three transport regimes as a function of the density, varying from semimetallic in low-density samples (<2.4 g/cm³) to thermally activated in high-density (≥2.9 g/cm³) tetrahedral a-C. The middle-range densities (2.4 g/cm³ ≤ ρ ≤ 2.9 g/cm³) are characterized by resonant tunneling and hopping transport. Our findings offer a different perspective from the tight-binding model proposed by Katkov and Bhattacharyya [J. Appl. Phys. 113, 183712 (2013)], and agree with experimental observations in low-dimensional carbon systems [see S. Bhattacharyya, Appl. Phys. Lett. 91, 21 (2007)]. Identifying transport regimes is crucial to the process of understanding and applying a-C thin film in electronic devices and electrode coating in biosensors. Published by AIP Publishing.

I. INTRODUCTION

Amorphous carbon (a-C) is a noncrystalline solid allotropic form of carbon with a combination of sp² and sp³ hybridized atoms. The ratio of bonds determines its density and different material classes and properties: from low-density sp²-rich non-graphitizing a-C, to high-density sp³-rich tetrahedral a-C (ta-C). The a-C thin films’ electrical properties make them attractive for electronic devices and as electrode coating in electrochemical biosensors.1–5 as shown by their inclusion in the Emerging Research Devices chapter of the 2014 report, International Technology Roadmap for Semiconductors.6 For this reason, understanding and predicting electron transport mechanisms and electrical properties is vital to advance a-C-based technology.

Bhattacharyya7 explained different transport mechanisms from experimental measurements of a-C, correlating the conduction with the density of states (DOS) at the Fermi level (E_F). Some recent models8–10 attempted to describe the transport phenomena of a-C at the tight-binding level of theory, considering regions of concentrated sp³ and sp² atoms as barriers and wells. These models need further refinement to include the random nature of the material’s atomic structure, low-dimensionality, and coupling with electrodes in very thin films, because they affect the transport mechanisms in low-dimensional structures, as observed experimentally.11,12

Although atomistic simulations have been used to model a-C, the lack of long-range order limits the system size possible to simulate. Parametric and/or semi-empirical approaches have been proposed to overcome this limitation,13 but at the expense of calculated properties’ accuracy. Atomistic ab-initio simulations based on density functional theory (DFT) remain the most accurate and affordable technique to calculate a-C properties from the electronic structure, for small- to medium-size systems.14,15 These simulations have focused on computational generation strategies, analysis of the atomic structure, and optical and mechanical properties.14–20 Combining DFT and nonequilibrium Green’s functions (DFT-NEGF) has proven useful for qualitatively describing and investigating electron transport in nano and molecular systems, materials interfaces, and functional electronic devices and sensors.21–24 Despite the importance of the theoretical models established for electrical and transport properties,25 to our knowledge, a-C transport properties have not been studied from first-principles calculations. Thus, here, we used DFT-NEGF to investigate the electron transport mechanisms in ultra-small a-C structures, analyzing the effects of density, dimensionality, and contact with metallic electrodes. We were able to explain different types of conductions observed experimentally,7,26,27 semimetallic in low-density a-C films, and thermally activated for sp³ rich a-C. Our calculations also support experimental evidence of resonant tunneling (RT) in ultra-thin films,11,12 and establish three regimes of dominant mechanisms, as a function of the density.

Recent work19 demonstrated that the compressive pressure is an important variable to generate a-C structures. In their work, the authors obtained the desired pressure by

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readjusting the cell size followed by relaxation until reaching the target conditions. We propose a generalization of this generation strategy using Born-Oppenheimer Molecular Dynamics (MD) within the isothermal-isobaric ensemble (NPT). With this generation strategy, we can increase the sampling size for statistical purposes.

The computational results presented here are important in the context of predictive modeling and engineering of materials with specific electrical properties, especially when considering the materials’ aimed industrial applications.

II. COMPUTATIONAL DETAILS AND METHODS

We carried out our calculations with the Kohn-Sham DFT method,28 using the Perdew-Burke-Ernzerhof (PBE) generalized gradient approximation (GGA) for the exchange-correlation functional,29 as implemented in the GPAW code.30 GPAW is a grid-based implementation of the projector-augmented wave (PAW) method.31 Unless otherwise specified, the calculations used a plain waves (PW) basis set with the cut-off set to 400 eV for the a-C generation and a double-ζ polarized (DZP) linear combination of atomic orbitals (LCAO) basis set for the transport calculations, which yields good convergence for a-C systems.32,33 The k-space integration was done over a (4 × 4 × 4) Monkhorst-Pack grid.34 The atomic positions were optimized following the Broyden-Fletcher-Goldfarb-Shanno (BFGS) algorithm35–38 until all the forces acting on the atoms were below 0.05 eV/Å.

A. Bulk a-C generation method

We built 25 orthogonal supercells of 64 randomly positioned carbon atoms. The dimensions of each supercell in the xy plane are multiples of the diamond unit cell (111) plane (5.8 Å × 5.0 Å), and the z component was adjusted so that the sum of the atoms’ masses divided by the volume of the cell yielded a target density. After a full relaxation, the system was equilibrated at 300 K (±10 K) using the Berendsen algorithm39 within the canonical ensemble (NVT dynamics). After equilibration at 300 K, the temperature was kept constant using the Berendsen algorithm within the frame of isothermal-isobaric (NPT) dynamics, with a target compressive pressure of 1.9 GPa, considered the bonding distance: an sp³ atom has three neighbors, while an sp³ atom has four neighbors. This definition and cutoff radius have been used in the literature, obtained from the analysis of the radial distribution function.15–17,19,33,42

After labeling each atomic site, it is easy to qualitatively assess the contributions of sp² and sp³ sites to the DOS by summing the contributions of all the sites in the system

\[ g_{sp²}(E) = \sum_{x=sp²} \sum_{i=2}^{sp³} g_{z,i}(E), \]

where \( g_{sp²}(E) \) is the contribution of sp² or sp³ sites to the DOS.

Finally, we define the local DOS (LDOS) as the sum of \( g_{z,i} \) over the valence states i of all the sites z in a particular region in space.

B. Transport system

The generated a-C samples were used to build a system for electron transport calculations that consists of a scattering region and left and right electrodes. The scattering region was built with a molecule (a-C sample) and two principal layers (PL), which in our case are slabs of 3 fcc (111) Au atomic layers (a more extensive discussion on PL is found elsewhere43–47). The Au/a-C interfaces in the scattering region were relaxed to prevent abrupt changes in the electrostatic potential. The interface was modeled with a reduced system consisting of an a-C slab, and built taking the atoms within a 5 Å surface measured from the topmost C atom. For the bottom atoms a similar slab was defined. The surface was passivated with hydrogen to compensate for dangling bonds.19 This slab and the PL were relaxed at six different distances from each other, between 1.5 and 3.5 Å, in a unit cell adding 10 Å vacuum up and down in the z direction to avoid orbital overlapping and periodicity effects. The atoms in the outermost 2 Å surface of the a-C slab and last atomic layer of the PL were fixed, so that the transport system could be reconstructed easily. Boundary conditions can cause an electrostatic potential gradient in the vacuum region (see the blue line in Fig. 1). The gradient creates an artificial electric field in the system, which may lead to nonphysical effects during the relaxation.48 To ensure vacuum neutrality we set periodic boundary conditions in the x and y directions, and fixed boundary conditions with dipole correction in the z direction (see the red line in Fig. 1).
The total energy of the six relaxed interfaces as a function of the Au/a-C distance was fixed to a quadratic polynomial. The minimum is assumed to be the optimum interface distance, which lies between 2.5 and 3 Å for all the generated samples, as Fig. 2 shows. The interfaces are finally relaxed using the calculated optimum distance shown in the top panel of Fig. 1, and the scattering region is reconstructed and connected to the electrodes shown in Figs. 3(a) and 3(b).

The transport properties were calculated by means of the retarded Green’s function,\(^49,50\) as implemented in the GPAW code\(^47\)

\[ G^r(E) = [ES - H - \Sigma'(E)]^{-1}, \]  

where \( E \) is the energy variable, \( \Sigma'(E) = \sum_i \Sigma'_i(E) \) is the sum of self-energies \( \Sigma'_i(E) \) (one for each electrode, in our case), and \( S \) and \( H \) are the overlap and Hamiltonian matrices of the scattering region, respectively (calculated self-consistently from the nonequilibrium electron density as detailed elsewhere\(^47\)). Using Equation (6), we can calculate the Transmission function as follows:

\[ T(E) = Tr[\Gamma'_1(E)G^r(E) \times \Gamma'_2(E)G^r(E)^\dagger], \]  

where \( \Gamma'_i(E) = i(\Sigma'_i(E) - \Sigma'_i(E)^\dagger) \) is the broadening matrix, the antihermitian part of the individual self-energies for each electrode. We can calculate the current as a function of energy, and therefore applied bias, by integrating Equation (7) in the Landauer Formula\(^49\)

\[ I = \frac{e}{h} \int T(E)(f_1(E) - f_2(E)) dE, \]  

FIG. 1. The top panel shows the fully relaxed atomic structure at the Au/a-C interface. Yellow, gray, and white spheres represent Au, C, and H atoms, respectively. The bottom panel shows the electrostatic potential with (see the red solid line) and without (blue dotted line) dipole correction of a 3.2 g/cm\(^3\) sample. The Fermi level is shown as a dashed horizontal line.

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FIG. 2. Normalized energy of the relaxed interfaces, calculated as a function of the Au/a-C distance for a 3.2 g/cm\(^3\) sample. We show the fixed quadratic curves and the equilibrium distance (dashed vertical lines).

FIG. 3. (a) Atomic structure of low-density (2.29 g/cm\(^3\)) and (b) high-density (3.20 g/cm\(^3\)) systems used in quantum transport calculations. Yellow and gray spheres represent Au and C atoms, respectively.
where \( q \) is the electron charge, \( h \) the Planck’s constant, and \( f_n(E) = f(E - \mu_n) \) is the Fermi distribution function of electron \( n \). We used an electron smearing of \( kT = 0.1 \text{ eV} \) and \( \mu_n = E_F + qV/2 \), with \( E_F \) being the equilibrium Fermi level of the Au electrodes, and \( V \) the applied voltage.

### III. RESULTS AND DISCUSSION

#### A. Bulk a-C generation

Fig. 4(a) shows the results of \( \rho \) versus \( sp^3 \) for the generated a-C structures. The NPT method yields structures closer to the experiment, compared with other theoretical approaches, such as Carr-Parrinello molecular dynamics (CPMD) by Marks et al. The slope fitted for CPMD is 70.23 (g/cm\(^3\))\(^{-1}\), while for our NPT calculation it is 58.51 (g/cm\(^3\))\(^{-1}\). The slope fitted from experimental data is 16.67%, while for the CPMD it is 46.56%. A bigger set would yield better precision and confidence to the statistical measures, but it is out of the scope of this work.

![Figure 4](image)

**FIG. 4.** (a) Calculated \( sp^3 \) fraction as a function of the sample density (red dotted line). For comparison, we include experimental results from Fallon et al. (black solid line), and computational CPMD results from Marks et al. (blue dashed line). (b) Pressure as a function of the \( sp^3 \) fraction (red dotted line). We also include experimental results from Fallon et al. (black solid line).

Figure 4(b) shows the pressure as a function of the \( sp^3 \) fraction for the same points in Fig. 4(a). The error of the fitted slope with respect to experiment is 6.03%. Samples with a pressure closer to the experimental values, accurately predict \( sp^3 \) fraction as a function of \( \rho \), and are therefore, used here for the electronic structure analysis and electron transport calculations.

#### B. Bulk a-C electronic structure

Amorphous semiconductors have a pseudogap, i.e., there are midgap states (traps) around the Fermi level. These are \( \pi \) states that lie between the \( \sigma \) gap and are localized due to the random potential around them (Anderson localization). The degree of electronic states’ localization can be determined using a participation ratio analysis. This method was introduced by Tritsaris et al. to the analysis of a-C structures, and it has the advantage of using the PDOS on each atom, instead of the single electron wave functions, which facilitates the localization calculation.

The degree of localization \( w(E) \) can be defined as follows:

\[
w(E) = \frac{\sum_{i=1}^{N} (N \times g_x(E)/g'(E) - 1)^2}{N(N-1)}.
\]

Here \( g_x(E) = \sum_{i=2,3} g_{2,3}(E) \) is the PDOS on site \( x \), \( g'(E) = \sum_{i=1}^{N} g_x(E) \) is the sum of all PDOS, and \( N \) is the number of atoms in the system \((N=64 \text{ atoms})\). For an extended state \( g_x(E) \rightarrow g'(E)/N \) and \( w(E) \rightarrow 0 \), while for a state strictly localized on atom \( x \), \( g_x(E) = g'(E) \) and \( w(E) = 1 \).

In Fig. 5, we show the DOS and \( w(E) \) for two of the samples with low \((2.29 \text{ g/cm}^3)\) and high \((3.2 \text{ g/cm}^3)\) densities. We see that \( w(E) \) increases with density, as others reported in related work. For the low-density sample, a continuum of states exists around \( E_F \) delimited by two gaps around \(-1.11 \text{ and } 1.86 \text{ eV} \) (averaged from three samples, all with \( \rho = 2.3 \text{ g/cm}^3 \), approximately), while for the high-density sample, states around \( E_F \) are distinguishable from each other and the gaps are located around \(-1.91 \text{ and } 1.80 \text{ eV} \) (averaged from two samples with \( \rho = 3.2 \text{ g/cm}^3 \)). There is also a small gap between the states around \( E_F \) that opens with density, but its particular value for each sample is distributed statistically, as the presence of the gap is not only determined exclusively by the amount of \( sp^3 \), but also by the spatial correlation between \( sp^2 \) sites. This behavior is characteristic of a pseudogap where localized “defect” states are present within the optical gap.

Contributions due to \( sp^3 \) and \( sp^2 \) sites to the DOS and \( w(E) \) are shown in Fig. 6 for a \( 3.20 \text{ g/cm}^3 \) a-C sample. Most of the localized electronic states are contributed by the \( sp^2 \) sites, while the localization due to \( sp^3 \) sites is located mainly at the conduction band’s edge, in good agreement with previous work. The PDOS for \( sp^2 \) sites is higher than for \( sp^3 \) sites within the pseudogap, which confirms that the localized electronic states that determine conduction are \( sp^2 \)-like. Because the \( sp^2/\!sp^3 \) fraction is directly related to the density, the latter will determine the transport mechanism of a-C.

Keeping in mind that \( \pi \) states around \( E_F \), contributed by \( sp^2 \) sites, will determine the transport mechanism for low
bias, and that these are highly localized, we fitted a weighted Lorentzian function to the $w(E)$ points within an energy window between $-1.5$ and $1.5$ eV around $E_F$, in order to quantify localization parameters that help us determine different transport regimes as a function $\rho$ and $sp^3$ (see Fig. 7)

$$w_{fit}(E) = A \frac{\Gamma}{(E - \bar{\omega})^2 + \Gamma^2}. \quad (10)$$

Here, $A$ is a fitting constant proportional to the peak magnitude of $w(E)$, $\Gamma$ is a parameter, which determines the width over the energy, and $\bar{\omega}$ determines the position of the peak, relative to the Fermi level. The magnitude of the peak of $w_{fit}(E)$, and its value at the Fermi level can then be calculated as $w_{fit}(\bar{\omega}) = A/\Gamma$ and $w_{fit}(E_F) = A\Gamma/(\bar{\omega}^2 + \Gamma^2)$ correspondingly.

In Fig. 8, we plotted both the $w_{fit}(\bar{\omega})$ and $w_{fit}(E_F)$ as a function of the density and $f_{sp^3}$. Note that $w_{fit}(\bar{\omega}) \approx w_{fit}(E_F)$, which indicates that the $sp^3$ localized states’ energies lie around $E_F$. The deviations from $E_F$ can be explained by the fluctuations of our calculated points from the experimental $\rho$ vs. $f_{sp^3}$ curve. These are significantly higher for densities around 2.8 g/cm$^3$, as Fig. 4(a) shows.

We approximated the trend of the data points to a Bezier polynomial. Three different trends can be differentiated with respect to the density and $f_{sp^3}$: (1) for low densities ($\rho \leq 2.55$ g/cm$^3$) and $f_{sp^3} \leq 44.94$, the localization increases linearly; (2) for middle-range densities ($2.55$ g/cm$^3 \leq \rho \leq 2.80$ g/cm$^3$) and (44.94 $\leq f_{sp^3} \leq 59.64$), the localization is constant; (3) for high densities ($\rho \geq 2.80$ g/cm$^3$) and $f_{sp^3} \geq 59.64$, the localization increments exponentially with respect to the density and $f_{sp^3}$. Essentially, we can expect these three regions, where the Fermi localization behaves differently to have different transport mechanisms. Next, we show how this is indeed the case.

### C. Electron transport in a-C films

After contacting with the electrodes, the system loses periodicity in the transport direction ($z$). Furthermore, the relaxation at the interfaces with the metallic electrodes changes the atoms’ positions and their bonding nature, becoming more $sp^3$, similar to what happens at a-C surfaces.\cite{19,20}

The variation in the atoms’ positions changes the electronic structure of the sample, see Fig. 9. The new contacted system loses the electronic states gap for both high- and low-density samples and the electronic states’ localization (calculated using the PDOS [see Equations (4) and (9)] on the carbon atoms within the scattering region) drops. The effect is stronger in samples with higher densities, in which localized $sp^2$ states are very confined in space due to the potential of dominant $sp^3$ sites. Confinement is lost when the interfaces reconfigure. We have shown that the high-density systems are shorter (in the $z$ direction) than the low-density ones for a constant number of atoms. In the case of the highest density studied (3.2 g/cm$^3$), the length of the unit cell in the transport direction is 13.83 Å, because we used 5 Å a-C slabs on each
side to build the interfaces, from which only 2 Å were fixed, 43% of the original bulk structure is free to reconfigure. This percentage explains the drop on the localization peak, stronger for high density than for low-density samples.

Figure 10 shows the electron transmission function together with the PDOS on the $sp^2$ and $sp^3$ sites, providing qualitative information on the contribution of each hybridization to the available electronic states in the scattering region. PDOS has been widely used in the literature to analyze the transport properties of molecular junctions and nanoscale systems. However, the strategy of PDOS for transport analysis is subject to debate. PDOS changes depending on the chosen basis set and the subspace on which the Hamiltonian is diagonalized. This is a challenge for the analysis of transport from the electronic structure.

We see that the shape of the transmission function follows the shape of the $sp^3$ PDOS for both low- and high-density samples, since it determines the pseudogap. However, the transmission function of the low-density sample contains resonances around the Fermi level, which are neither present in the $sp^2$ or the $sp^3$ PDOS. Resonances in the transmission function are usually associated to highly localized or quasi-bound (QB) states in the transport path. Such localized states lead to RT, which can be measured by changes in the $I-V$ curve’s slope and/or negative differential resistance. RT transport is characteristic of composite materials and semiconductor heterostructures; it has been studied...
On the contrary, we propose that conduction in high-density decreases due to the filling of exponentially distributed traps. We divided the system into six slices along the transport direction and determined the available states and localization from the PDOS and \( w(E) \) for each individual slice. We also calculated the density distribution along the transport directions and the contributions to the density due to \( sp^2 \) and \( sp^3 \) sites: \( \rho_{sp^2} \) and \( \rho_{sp^3} \), respectively. To calculate \( \rho \), \( \rho_{sp^2} \), and \( \rho_{sp^3} \) along the transport direction, we divided the system here into 50 slices. The points were smoothed using a Bezier curve with the calculated points as the polynomial parameters. In Fig. 11 we plotted \( w(E) \) and the density distribution for the same samples used in Fig. 10. These results show that the localization is directly related to the \( sp^3 \) distribution within the system. For the longer, low-density sample, a region exists between 10 Å and 12 Å where the \( sp^2 \) content rises, creating a potential barrier and causing the \( sp^2 \) electronic states to get localized, around 7 and 15 Å. We have only considered the features of \( \rho \) and \( w(E) \) in one dimension, but it is worth mentioning that clusters of \( sp^2 \) sites might be distributed in 3D, influencing the transport mechanism. From Fig. 11(a) we see the similarity of the a-C with a square quantum well, which was the model proposed by Katkov and Bhattacharyya. However, their model considered symmetric wells with sharp barriers in the electrostatic potential (associated with \( sp^3 \) clusters), which is not the case for realistic disordered structures as the ones we present. The influence of this geometric asymmetry can be seen in the difference in the energy position of the QB states, their energy width, and localization degree. Although the localized states at each side of the barrier are not at the same energy (due to the wells’ asymmetry), there is a considerable tail overlapping that is responsible for the resonance in \( T(E) \), thus, evidencing RT transport as reported by Bhattacharyya et al.1, 12

Our results can also be contrasted with the experimental results from Protopopova et al.61 They studied the thickness dependence of the conductivity in a-C ultra-thin films, and proposed that for thinner films (\( \leq 4 \) nm) the conduction decreases due to the filling of exponentially distributed traps. On the contrary, we propose that conduction in high-density ultra-thin films decreases due to the reduction in the number of charge traps (localized electronic states) induced by the atomic reconfiguration.

Computing \( T(E) \) self-consistently using the DFT-NEGF method for each of the samples is a computationally expensive task. Moreover, to calculate the \( I-V \) curves with Equation (8), one has to compute \( T(E) \) at each bias point. Therefore, we first calculated \( I-V \) characteristic curves using a reduced single-\( \zeta \) polarized (SZP) basis set, as shown in Fig. 12. As the density of the sample increases, there is a change in the shape of the \( I-V \) characteristic curve. The \( I-V \) curve in Fig. 12(a) is similar to a linear response with minor slope changes. We attribute these changes to weakly localized states, but the dominant transport mechanism is semimetallic through extended \( sp^2 \) states around \( E_F \). On the other hand, in Fig. 12(b) we see notorious changes in the slope of the \( I-V \) curves, and for the case of the 2.67 g/cm\(^3\) sample, negative differential resistance around 1.3 and 1.7 V. This is evidence of RT and variable range hopping (VRH) transport, induced by strongly localized \( sp^2 \) states around \( E_F \). Finally, in Fig. 12(c), we see linear characteristics for low biases. The slope changes as the voltage increases, which is a common feature in systems whose dominant transport mechanism is thermally activated.27,62

These results are only qualitative, since a SZP basis set is unable to represent accurately the energetics of a-C, especially in high-density samples, as shown in previous works.32,33 Therefore, we calculated the \( I-V \) curves for a set of samples with densities ranging from 2.23 to 3.20 g/cm\(^3\), using a DZP basis set, as shown in Fig. 13. The \( I-V \) curve for the sample with \( \rho = 2.46 \) g/cm\(^3\), shows fluctuations in the slope, characteristic of RT transport, which is consistent with the localized states observed in Fig. 11(a), while the \( I-V \) characteristic for the sample with \( \rho = 3.20 \) g/cm\(^3\) is linear, for activated transport and consistent with the results in Fig. 13(b).

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**FIG. 11.** The top panel shows the degree of localization \( w(E) \) in different spatial regions along the transport direction \( z \). The bottom panel shows the density \( \rho \) (black solid line), and the \( sp^2 \) and \( sp^3 \) contributions to the total density \( \rho_{sp^2} \) (blue dotted) and \( \rho_{sp^3} \) (red dashed), respectively, averaged in the \( xy \) plane along the transport direction \( z \). The system’s average density is shown as a horizontal black dashed line for reference. Two systems are shown: (a) \( \rho = 2.46 \) g/cm\(^3\) and (b) \( \rho = 3.20 \) g/cm\(^3\).

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**FIG. 12.** Current-voltage (\( I-V \)) characteristic curves of a-C thin films with different densities, calculated using DFT-NEGF self-consistent method and a SZP basis set. We identify three types of conduction: (a) semimetallic for the low-density samples, (b) RT and VRH for middle-range densities and (c) activated for the high-density samples.
These transitions from one conduction mechanism to the other have been observed previously by Chimowa, Churochkin, and Bhattacharyya\textsuperscript{63} for boron-doped nano-crystalline films. They showed that the transitions are heavily dependent on Nitrogen concentrations at low temperatures. However, the dominant conduction mechanism is also dependent on the temperature. They found that for concentrations of 10\% to 15\% Nitrogen, at temperatures below 118 K, the dominant conduction mechanism is VRH, while for concentrations of 20\% Nitrogen, activated conduction was evidenced in the range of 28 K to 89 K. These conclusions are applicable to our findings, since the role of Nitrogen doping is to increase the level of disorder, i.e., to increase the fraction of $sp^2$ sites in the film. A more rigorous computational study of the temperature dependence of these transitions is needed to establish the range of temperatures in which our model is valid.

IV. CONCLUSIONS AND FUTURE WORK

Using an algorithm based on the underlying physical process modeled with NPT ensemble dynamics, we generated 25 bulk a-C samples, which were used to build ultra-thin film systems. This allowed us to study the transport’s evolution, the localization as function of the $sp^3$ content, and density. We used Kohn-Sham DFT and the NEGF formalism to calculate the electron transmission and transport mechanisms. We identified three transport regimes, depending on the density. For very low densities ($\rho \leq 2.4 \text{ g/cm}^3$), the dominant transport mechanism is semimetallic. In high-density systems ($\rho \geq 2.9 \text{ g/cm}^3$), transport is driven by variable range and thermally activated hopping. For middle-range densities ($2.4 \text{ g/cm}^3 \leq \rho \leq 2.9 \text{ g/cm}^3$), RT and hopping (induced due to the appearance of localized states) dominate.

We demonstrated that contact with electrodes drops the localization of ultra-thin films’ electronic states, affecting the electron transport mechanisms. The contacts’ effect is especially strong for high-density films, and our findings are in agreement with experimental works. These findings can impact materials engineering significantly, providing essential information for designing future sensors and devices envisioned by the electronics industry.

While we were able to differentiate between density regions, we have yet to establish sharp thresholds. This requires further calculations and adopting more sophisticated techniques to obtain a quantitative model. We must also take into account electron-phonon interactions and phase-incoherent processes (which have not been considered), to improve the description of electron transport in a-C films. A different approach needs to be taken, either by considering a mean-field approach, similar to the recent studies by McIntosh et al.\textsuperscript{64} on superconductivity in boron-doped diamond, or by using the DFT-NEGF method, addressing electron-phonon interaction by means of perturbation theory at the level of self-consistent Born approximation.\textsuperscript{46}

SUPPLEMENTARY MATERIAL

See supplementary material for a more detailed view of the pseudogap and electronic gap and for more $I$–$V$ characteristic curves of individual samples.

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