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Dissertation

**MANIPULATING GRAPHENE'S LATTICE: PSEUDO VECTOR
POTENTIALS, ANOMALOUS FRICTION, TUNABLE THERMAL
CONDUCTIVITY, AND METASTABLE BAND GAP**

by

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ABSTRACT

Graphene, a single atomic sheet of graphite, exhibits a diverse range of unique properties. The electrons in intrinsic graphene behave like relativistic Dirac fermions, it has a record Young's modulus but extremely low bending rigidity, and suspended graphene has record thermal conductivity. These properties are made more intriguing because with a thickness of only one atomic layer, graphene is uncommonly affected by its environment. This allows for graphene's atomic lattice to be manipulated and its amazing properties to be altered. In this dissertation, we investigate the exciting new physics realized by manipulating graphene's lattice.

In an exotic coupling, strain causes alterations in graphene's electrical dispersion mathematically analogous to the vector potential associated with a magnetic field. Here we complete the standard description of the strain-induced vector potential with the explicit inclusion of the lattice deformations. Additionally, we propose several strain engineered devices which generate interesting pseudo magnetic fields.

Accurate strain engineering requires a complete understanding of the often bizarre interactions between a two dimensional material and its environment. We extract the pressure dependent sliding friction between the SiO₂ substrate and mono-, bi-, and tri-layer graphene using Raman spectroscopy of circular, graphene sealed microchambers

under variable external pressure. The sliding friction for trilayer graphene behaves normally, scaling with the applied load, whereas the friction for monolayer and bilayer graphene is anomalous, scaling with the inverse of the strain in the graphene.

Strain also modifies the phonon which is believed to provide the majority of graphene's record thermal conductivity, the out of plane acoustic phonon. Using similar techniques we demonstrate that graphene's thermal conductivity decreases with strain, providing additional insight into the record thermal conductivity of suspended graphene.

Finally, we can also manipulate the lattice at the smallest length scale by exciting optical phonons. It is predicted that a particular optical phonon can be used to continually push the system into an out of equilibrium state that has a transport band gap. We propose a measurement which uses a neon seeded helium beam to excite the optical phonon and electrical transport to measure the development of the band gap.

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Chapter 1

Unmodified band structure

The linear electronic dispersion of graphene was first published by Wallace (Wallace, 1947) 57 years before Geim, Novoselov and co-workers spurred graphene research forward with their method of mechanical exfoliation (Novoselov et al., 2004). Later, in 1984 Semenoff formalized the equivalence between the low energy electrons in graphene and Dirac-Weyl electrons (Semenoff, 1984). Remarkably, the nearest neighbor tight binding approach used by these authors has accurately described the majority of the low energy physics in graphene. This chapter will follow in the spirit of these derivations with additional emphasis provided to enhance the discussions of the strain induced pseudo vector potential and the phonon induced Kekulé transition.

1.1 Graphene's lattice and Brillouin zone

In its unperturbed state, the carbon atoms in the graphene lattice are arrayed in a hexagon as shown in Figure 1.1(a). Throughout the x axis will be taken in the zigzag direction as shown. Since a hexagonal lattice is not a Bravais lattice, the lattice is treated as a triangular Bravais lattice with a two atom bases. In figure 1.1(a) the A sub-lattice is colored orange and the B sub-lattice is colored blue. The lattice is

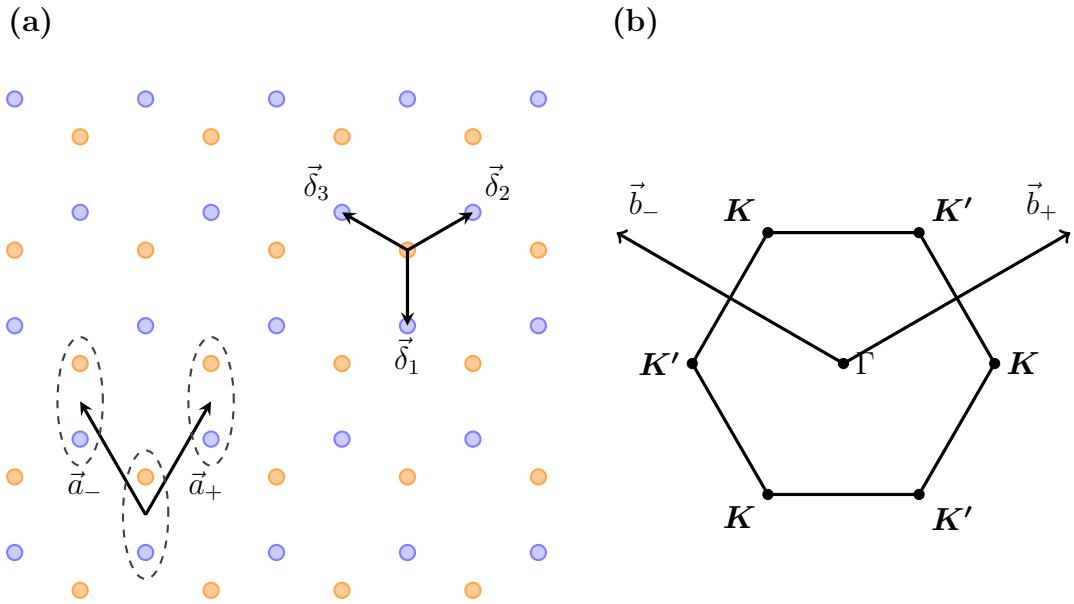


Figure 1.1: Geometry of intrinsic graphene. (a) Real space graphene lattice with the A sub-lattice in orange, the B sub-lattice in blue, nearest neighbor vectors ($\vec{\delta}_1$, $\vec{\delta}_2$, and $\vec{\delta}_3$) shown as arrows, and lattice vectors (\vec{a}_+ and \vec{a}_-) shown translating the two atom basis. (b) First Brillouin zone with the Γ point and equivalent \mathbf{K} and \mathbf{K}' points indicated.

created by arraying the two atoms basic using the lattice vectors

$$\vec{a}_+ = \frac{\sqrt{3}a}{2} \left(+\hat{x} + \sqrt{3} \hat{y} \right)$$

$$\vec{a}_- = \frac{\sqrt{3}a}{2} \left(-\hat{x} + \sqrt{3} \hat{y} \right)$$

where $a = 1.4 \text{ \AA}$ is the nearest neighbor distance. The three nearest neighbor vectors,

$$\vec{\delta}_1 = -a\hat{y}$$

$$\vec{\delta}_2 = \frac{a}{2} \left(+\sqrt{3} \hat{x} + \hat{y} \right)$$

$$\vec{\delta}_3 = \frac{a}{2} \left(-\sqrt{3} \hat{x} + \hat{y} \right) ,$$

connect each atom in the A sub-lattice to its three nearest neighbors in the B sub-lattice.

The primitive reciprocal lattice vectors are

$$b_+ = \frac{2\pi}{3} \left(+\sqrt{3} \hat{x} + \hat{y} \right)$$

$$b_- = \frac{2\pi}{3} \left(-\sqrt{3} \hat{x} + \hat{y} \right)$$

The BZ of graphene is shown in Figure 1.1(b). It is a hexagon rotated 30 degrees relative to the crystal axis. The Γ point is at the center of the Brillouin zone while the \mathbf{K} and \mathbf{K}' are at the corners. It should be noted that only 2 of the 6 corners of the hexagon are unique, the others can be connected by reciprocal lattice vectors.

The position of the \mathbf{K} and \mathbf{K}' points are given by:

$$\begin{aligned}\mathbf{K} = -\mathbf{K}' &= \frac{4\pi}{3\sqrt{3}a} \hat{x} \\ \mathbf{K} = -\mathbf{K}' &= \frac{2\pi}{3\sqrt{3}a} \left(-\sqrt{3} \hat{x} + \hat{y} \right) \\ \mathbf{K} = -\mathbf{K}' &= \frac{2\pi}{3\sqrt{3}a} \left(-\sqrt{3} \hat{x} - \hat{y} \right),\end{aligned}$$

with the first pair the easiest to work with.

In later sections this discussion will be expanded to take into account strain and phonons perturbing the graphene lattice. In both cases the changes in graphene's electronic dispersion originates from changes in the position of the atoms.

1.2 Electronic dispersion

Since the tight binding formalism is used universally in this work, this section will start with a brief motivation of its form. Afterward, the nearest neighbor tight binding formalism will be applied to graphene.

In second quantization, the total electronic energy in the system is written as:

$$H = \sum_{\vec{k},\sigma} c_{\vec{k},\sigma}^\dagger c_{\vec{k},\sigma} \epsilon_{\vec{k}}$$

where $\epsilon_{\vec{k}}$ is the energy of the electron with wave-vector \vec{k} and $c_{\vec{k},\sigma}^\dagger$ and $c_{\vec{k},\sigma}$ are the creation and annihilation operators for an electrons with wave-vector \vec{k} and spin σ . The product $c_{\vec{k},\sigma}^\dagger c_{\vec{k},\sigma}$ is the number operator, providing the number of electrons with the given spin and wave-vector. This is a very natural definition of the energy; the equation simply sums the energy one wave-vector at a time.

When the atomic wave functions of the atoms in the material do not overlap con-

siderably, the electrons are sufficiently localized to consider real space creation and annihilation operators which create or annihilate electrons at specific lattice points. The reciprocal space creation and annihilation operators can then be Fourier expanded about these points,

$$\begin{aligned} c_{\vec{k},\sigma}^\dagger &= \frac{1}{\sqrt{N}} \sum_{\vec{R}_i} e^{i\vec{k}\cdot\vec{R}_i} c_{i,\sigma}^\dagger \\ c_{\vec{k},\sigma} &= \frac{1}{\sqrt{N}} \sum_{\vec{R}_i} e^{-i\vec{k}\cdot\vec{R}_i} c_{i,\sigma} . \end{aligned}$$

It will be important to note that the phase term must include the position which corresponds to the position at which the operator works. Applying this Fourier transform to the systems Hamiltonian yields the tight binding Hamiltonian,

$$\begin{aligned} H &= - \sum_{\vec{R}_i, \vec{R}_j, \sigma} (-) \sum_{\vec{k}} \frac{1}{N} e^{i\vec{k}\cdot(\vec{R}_i - \vec{R}_j)} \epsilon_{\vec{k}} c_{i,\sigma}^\dagger c_{j,\sigma} \\ &= - \sum_{\langle i,j \rangle, \sigma} t_{i,j} c_{i,\sigma}^\dagger c_{j,\sigma} + \text{H.C.} , \end{aligned}$$

where the final sum is over $\langle i, j \rangle$ nearest neighbor pairs and $t_{i,j} = - \sum_{\vec{k}} \frac{1}{N} e^{i\vec{k}\cdot(\vec{R}_i - \vec{R}_j)} \epsilon_{\vec{k}}$ is the hoping. This seemingly complicated parameter has an easy interpretation when the final form of the Hamiltonian is considered. The hoping energy is the energy associated with removing an electron from atom j and putting it on atom i . It is usually measured empirically or calculated by matching the tight binding model to other more powerful methods such as density functional theory. As will be illustrate for graphene, this tight binding formalism provides a simple but powerful method of calculating band structures.

The physics relevant for this work is captured by this nearest neighbor tight binding model. For graphene, when an electron hops between nearest neighbor, it changes

sub-lattice. The nearest neighbor Hamiltonian reflects this:

$$H = -t_0 \sum_{\langle i,j \rangle} (a_i^\dagger b_j + \text{H.C.}), \quad (1.1)$$

where the hoping energy t_0 gives the energy it takes to remove an electron from the j th atom in the B sub-lattice with the B sub-lattice annihilation operator b_j and put an electron on its nearest neighbor, the i th atom in the A sub-lattice, using the A sub-lattice creation operator a_i . The sum is restricted to $\langle i,j \rangle$ nearest neighbor pairs. Finally the hoping from the A sub-lattice to the B sub-lattice is taken into account by the Hermitian conjugate, H.C..

The Hamiltonian is simplified by writing the creation and annihilation operators in Fourier space. There is some freedom in choosing the phase factors. They can either be chosen so that the operators are expanded around the position of each atom or instead they can be expanded around the position of the atomic basis occupied by the atom. Both approaches yield the same result if one is consistent (Bena and Montambaux, 2009). Here the operators are expanded about the atomic basis. This will make the correspondence with the Dirac-Weyl equation more clear. The expanded operators are

$$\begin{aligned} a_i^\dagger &= \frac{1}{\sqrt{N}} \sum_{\vec{k}} e^{i\vec{k} \cdot \vec{R}_i} a_{\vec{k}}^\dagger \\ b_j &= \frac{1}{\sqrt{N}} \sum_{\vec{k}'} e^{-i\vec{k}' \cdot \vec{R}_j} b_{\vec{k}'} \end{aligned} \quad (1.2)$$

where R_i and R_j are the position of the i -th and j -th atomic basis respectively. Since we are limited to nearest neighbor pairs the nearest neighbor is either in the same atomic basis or in one of the neighboring atomic basis. Thus, \vec{R}_j is restricted to

$\vec{R}_j \in \{\vec{R}_i, \vec{R}_i + \vec{a}_+, \vec{R}_i + \vec{a}_-\}$ and the difference $\vec{\Delta}_j = \vec{R}_j - \vec{R}_i$ is independent of i . In reciprocal space the Hamiltonian is simplified,

$$\begin{aligned} H &= -t_0 \frac{1}{N} \sum_{\vec{k}, \vec{k}'} \sum_i \left(e^{i(\vec{k}-\vec{k}') \cdot \vec{R}_i} e^{-i\vec{k}' \cdot \vec{\Delta}_j} a_{\vec{k}}^\dagger b_{\vec{k}'} + \text{H.C.} \right) \\ &= -t_0 \frac{1}{N} \sum_{\vec{k}, \vec{k}'} \left(N \delta_{\vec{k}, \vec{k}'} e^{-i\vec{k}' \cdot \vec{\Delta}_j} a_{\vec{k}}^\dagger b_{\vec{k}'} + \text{H.C.} \right) \\ &= -t_0 \sum_{\vec{k}} \left(\sum_j e^{-i\vec{k} \cdot \vec{\Delta}_j} a_{\vec{k}}^\dagger b_{\vec{k}} + \text{H.C.} \right), \end{aligned} \quad (1.3)$$

where $\delta_{\vec{k}, \vec{k}'}$ is the Kronecker delta function. In matrix notation this reads,

$$H = \sum_{\vec{k}} \begin{pmatrix} a_{\vec{k}}^\dagger & b_{\vec{k}}^\dagger \end{pmatrix} \begin{pmatrix} 0 & -t_0 \sum_j e^{-i\vec{k} \cdot \vec{\Delta}_j} \\ -t_0 \sum_j e^{-i\vec{k} \cdot \vec{\Delta}_j} & 0 \end{pmatrix} \begin{pmatrix} a_{\vec{k}} \\ b_{\vec{k}} \end{pmatrix}. \quad (1.4)$$

It is sometimes useful to express the sum as,

$$\sum_j e^{-i\vec{k} \cdot \vec{\Delta}_j} = 1 + 2e^{-i\frac{3}{2}k_y} \cos\left(\frac{\sqrt{3}}{2}ak_x\right).$$

The geometry of the lattice was the only input required to find this nearest neighbor tight binding Hamiltonian.

A straightforward calculation provides the electron dispersion,

$$E(\vec{k}) = \pm t_0 |h(\vec{k})| = t_0 \sqrt{1 + 4 \cos^2\left(\frac{\sqrt{3}}{2}ak_x\right) + 4 \cos\left(\frac{\sqrt{3}}{2}ak_x\right) \cos\left(\frac{3}{2}ak_y\right)},$$

which is plotted in figure 1.2. As shown, the two energy bands meet at points at the corners of the Brillouin zone. For pristine graphene there is one electron per carbon atom leaving the low energy band completely filled and the high energy band

empty. The Fermi energy can be shifted by charge transfer from contaminates or it can be purposely modified by adding adding or removing charges through capacitive coupling. The resulting shifts in Fermi energy are relatively small and thus the low energy excitations happen in a narrow energy window around these points. This just happens to be the energy window for which the nearest neighbor tight binding approach is most accurate. A higher order model is required to account for things such as trigonal warping which alter the dispersion at higher energies.

The two points of convergence between the high and low energy bands are referred to as the Dirac points because of the characteristic energy dispersion in their vicinity. This interesting low energy physics can be best captured by expanding the Hamiltonian in Equation 1.4 about these points. The wave-vectors are approximated as $\vec{k} = \mathbf{K} + \vec{q}$ and $\vec{k} = \mathbf{K}' + \vec{q}$ and for small qa the sum in Equation 1.4 is approximated as,

$$\begin{aligned}
\sum_j e^{-i\vec{k}\cdot\vec{\Delta}_j} &= \\
\mathbf{K} : &= \sum_j e^{-i(\mathbf{K}+\vec{q})\cdot\vec{\Delta}_j} \\
&\simeq \sum_j (1 - i\vec{q}\cdot\vec{\Delta}_j) e^{-i\mathbf{K}\cdot\vec{\Delta}_j} \\
&= -\frac{3}{2}a(q_x - iq_y) \\
\mathbf{K}' : &\simeq \sum_j (1 - i\vec{q}\cdot\vec{\Delta}_j) e^{-i\mathbf{K}'\cdot\vec{\Delta}_j} \\
&= -\frac{3}{2}a(-q_x - iq_y) . \tag{1.5}
\end{aligned}$$

These values are independent of which of the three identical \mathbf{K} or \mathbf{K}' points are

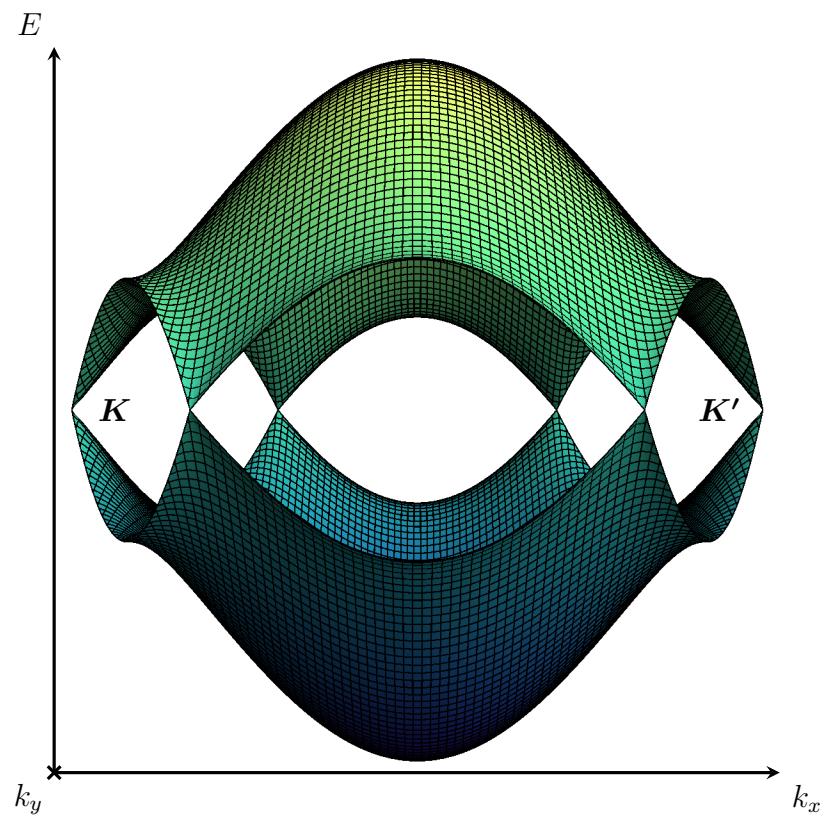


Figure 1.2: The electronic dispersion of intrinsic graphene in the first Brillouin zone calculated with a nearest neighbor tight binding model. The two energy bands meet at the \mathbf{K} and \mathbf{K}' points.

expanded about. The low energy Hamiltonian is thus,

$$\mathbf{K} : H \simeq \frac{3}{2}at_0 \sum_{\vec{q}} \begin{pmatrix} a_{\vec{q}}^\dagger & b_{\vec{q}}^\dagger \end{pmatrix} \begin{pmatrix} 0 & q_x - iq_y \\ q_x + iq_y & 0 \end{pmatrix} \begin{pmatrix} a_{\vec{q}} \\ b_{\vec{q}} \end{pmatrix} \quad (1.6)$$

$$= \frac{3}{2}at_0 \sum_{\vec{q}} \begin{pmatrix} a_{\vec{q}}^\dagger & b_{\vec{q}}^\dagger \end{pmatrix} \vec{q} \cdot \vec{\sigma} \begin{pmatrix} a_{\vec{q}} \\ b_{\vec{q}} \end{pmatrix} \quad (1.7)$$

$$\mathbf{K}' : H \simeq \frac{3}{2}at_0 \sum_{\vec{q}} \begin{pmatrix} b_{\vec{q}}^\dagger & a_{\vec{q}}^\dagger \end{pmatrix} \begin{pmatrix} 0 & -q_x + iq_y \\ -q_x - iq_y & 0 \end{pmatrix} \begin{pmatrix} b_{\vec{q}} \\ a_{\vec{q}} \end{pmatrix} \quad (1.8)$$

$$= \frac{3}{2}at_0 \sum_{\vec{q}} \begin{pmatrix} b_{\vec{q}}^\dagger & a_{\vec{q}}^\dagger \end{pmatrix} (-1) \vec{q} \cdot \vec{\sigma} \begin{pmatrix} b_{\vec{q}} \\ a_{\vec{q}} \end{pmatrix}, \quad (1.9)$$

$$(1.10)$$

or in the combined form:

$$H = \hbar v_f \sum_{\vec{q}, \vec{q}} \begin{pmatrix} a_{\vec{q}}^\dagger & b_{\vec{q}}^\dagger & b_{\vec{q}}^\dagger & a_{\vec{q}}^\dagger \end{pmatrix} \begin{pmatrix} 0 & q_x - iq_y & 0 & 0 \\ q_x + iq_y & 0 & 0 & 0 \\ 0 & 0 & 0 & -q_x + iq_y \\ 0 & 0 & -q_x - iq_y & 0 \end{pmatrix} \begin{pmatrix} a_{\vec{q}} \\ b_{\vec{q}} \\ b_{\vec{q}} \\ a_{\vec{q}} \end{pmatrix} \quad (1.11)$$

$$H = \sum_{\vec{q}, \vec{q}} \begin{pmatrix} a_{\vec{q}}^\dagger & b_{\vec{q}}^\dagger & b_{\vec{q}}^\dagger & a_{\vec{q}}^\dagger \end{pmatrix} \begin{pmatrix} v_f \vec{p} \cdot \vec{\sigma} & 0 \\ 0 & -v_f \vec{p} \cdot \vec{\sigma} \end{pmatrix} \begin{pmatrix} a_{\vec{q}} \\ b_{\vec{q}} \\ b_{\vec{q}} \\ a_{\vec{q}} \end{pmatrix} \quad (1.12)$$

with $\vec{\sigma} = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix} \hat{x} + \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix} \hat{y}$ is the vector of Pauli matrices and v_f is the Fermi velocity given by $\frac{3}{2} \frac{at}{\hbar} \sim .9 \times 10^6 \text{ m/s}$. To express the Hamiltonian near the \mathbf{K}' point in terms of Pauli matrices, the order of the raising and lower operators had to be switched. This subtle difference means that the first component of eigenvectors for electrons near \mathbf{K}' corresponds to the amplitude for an electron on the B sub-lattice, not the A sub-lattice. These approximate Hamiltonians will be central for the discussion of electron physics in manipulated graphene.

The eigenvalues near either \mathbf{K} or \mathbf{K}' are easy to solve for using the low energy Hamiltonians. At both points, the low energy electronic dispersion is identical,

$$E = \hbar v_f |\vec{q}| .$$

As shown in Figure 1.2, the high and low energy bands meet at points at the Dirac points. Away from this point, the energy scales linearly with momentum measured from the Dirac point. Reminiscent of the light cone, this canonical dispersion is characteristic of a massless particle.

From the low energy electronic dispersion the density of electronic states can be directly calculated for the two dimensional system. After taking account of the two full spin and two fold valley degeneracy, the density of states is

$$\rho(\epsilon) = \frac{2}{\pi} \left(\frac{\hbar v_f}{L} \right)^2 \epsilon , \quad (1.13)$$

where L^2 is the area of the graphene and ϵ the energy measured from the Dirac points. This is linear in energy due to the canonical energy dispersion. The number of states at a given energy is related to the circumference of the Dirac cone at this energy. Using

the density of states the chemical potential can be calculated in standard graphene transport experiments.

Graphene's two dimensional nature makes it simple to continuously modify the chemical potential. Rather than varying the concentration of dopants as must be done for most three dimensional systems, graphene only requires the application of a back gate voltage. A standard graphene gated device is shown in Figure 1.3. Changing the back gate voltage (V_{BG}) charges the graphene-backgate capacitor, introducing charges into the system. In general, the differential increase in the number of charges is $dQ = C(\mu)dV$. The capacitance is defined as:

$$C = \frac{dN}{d\mu} \quad (1.14)$$

where N is the number of electrons in the graphene and μ is the chemical potential. For intrinsic graphene the capacitance is continuous and the total number of charges is just $N = CV/e$. The gated graphene device shown in Figure 1.3 can be treated as a parallel plate capacitor with a plate separation of 285 nm and the dielectric constant of silicon dioxide. Matching the integrated density of states to the number of charges in the system relates the applied gate voltage to different measurements of the chemical potential in the system. Some useful relationships are:

$$n(cm^{-2}) \sim 7 \times 10^{10} V_{BG}(Volts)$$

$$\mu(meV) \sim 30\sqrt{V_{BG}(Volts)}$$

$$\mu(meV) \sim 1 \times 10^{-7}\sqrt{n(cm^{-2})}$$

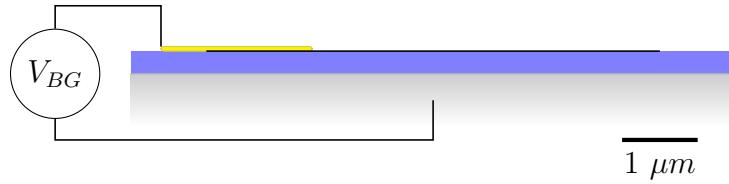


Figure 1.3: Side view of a back gated graphene device. The graphene is on top of 285 nm of thermal oxide grown on heavily doped Si.

1.3 Dirac-Weyl electrons

Graphene's linear electrical dispersion is peculiar. According to the usual expression for the effective mass of an electron in an electric field, $m^* = \frac{\hbar^2}{(\frac{d^2E}{dk^2})}$, (Kittel, 2005) the electrons in graphene have infinite mass. In a classical sense, this is understandable. The electron's group velocity, $d\omega/dk$ is constant; it can not be changed by applying a force. Classically, this corresponds to the case of an infinite mass which cannot be accelerated. An alternative, and more enlightening, interpretation of graphene's properties is as a massless, relativistic, spin 1/2, fermion. In this scenario the group velocity cannot be changed because the electrons are moving at the systems effective light speed. Hence, like a photon moving at the speed of light graphene's electrons cannot be accelerated.

In 1984 Semenoff demonstrated that there is an exact correspondence between the low energy Hamiltonian of graphene and the two dimensional Dirac-Weyl Hamiltonian (Semenoff, 1984) which governs massless, relativistic, spin 1/2, Fermions. In this section this correspondence will be briefly discussed.

The Hamiltonian of general, relativistic, spin 1/2, fermions is known as the Dirac equation. It is a matrix equation which is covariant, obeys the relativistic energy expression, and is first order in time. For the special case of massless particles, the

expression is simplified by a wisely chosen representation and is written as

$$H = \begin{pmatrix} c\vec{p} \cdot \vec{\sigma} & 0 \\ 0 & c\vec{p} \cdot \vec{\sigma} \end{pmatrix}$$

which, for particles in two spatial dimensions, is identical to Equation 1.12. It is clear now that the electrons in graphene behave like relativistic particles in a system where the speed of light is given by the fermi velocity.

As was the case in graphene, the Hamiltonian represents a decoupled system. In response to this, the wave function is written as the combination of two, two element spinors:

$$\psi = \begin{pmatrix} \chi_+ \\ \chi_- \end{pmatrix}$$

It can be shown that the spinners have well defined and opposite helicities. Since the energy of these massless particles is given by $\pm pc$, the Hamiltonian can be rewritten in a simpler form known as Weyl's equations,

$$(1 \mp \vec{\sigma} \cdot \hat{p})\chi_+ = 0$$

$$(1 \pm \vec{\sigma} \cdot \hat{p})\chi_- = 0 ,$$

where \hat{p} is the unit vector in the direction of the momentum. From this form it is clear that the spinors are eigenvectors of the helicity operator, $\hat{h} = \frac{1}{2}\vec{\sigma}\hat{p}$, with eigenvectors $\hat{h}\chi_+ = \pm 1/2$ and $\hat{h}\chi_- = \mp 1/2$. The χ_+ spinor is said to be right handed; for particles with positive energy the spin and momentum are in the same direction whereas for particles with negative energy they are in the opposite direction. The χ_- spinor then is left handed. Since the helicity operator commutes with the Hamiltonian, helicity

is conserved in this system. Gottfried and Yan provide a more detailed discussion of the Dirac equation and its consequences (Gottfried and Yan, 2003).

In graphene, the elements of the wavevector have a geometric interpretation. They correspond to the amplitude that the electron is at the A or B sublattice. Identifying these components with the relativistic components yields the pseudo spinners: $\psi_+ \equiv \begin{pmatrix} \psi_A^{\mathbf{K}} \\ \psi_B^{\mathbf{K}} \end{pmatrix}$ and $\psi_- \equiv \begin{pmatrix} \psi_B^{\mathbf{K}'} \\ \psi_A^{\mathbf{K}'} \end{pmatrix}$. It is a subtle point that the spinor for the electrons at the \mathbf{K}' point is in the opposite order than the \mathbf{K} point. Thus, the pseudo spin in this system refers to the sublattice. Further, the wavefunction for electrons at the \mathbf{K} point are right handed and the electrons at the \mathbf{K}' are left handed.

Since graphene's Hamiltonian is identical to that of massless, relativistic, spin 1/2 Fermions, the electrons should exhibit the same odd properties that have been predicted by high energy physicists. These unique properties include the Klein paradox (Young and Kim, 2009), Zitterbewegung (Castro Neto et al., 2009), and the anomalous quantum Hall effect (Novoselov et al., 2005; Zhang et al., 2005).

Chapter 2

Strain-induced vector potentials: Lattice-corrections and engineered pseudo magnetic fields

At the intersection of graphene's outstanding mechanical and electrical properties lies a peculiar and provocative coupling. Certain strain distributions in graphene can trick the electrons into acting as if they were in a magnetic field. Strain shifts the crystal momentum of the Dirac points much like the canonical momentum is shifted in the presence of a magnetic field. This exotic coupling is made more appealing by the two dimensional elastic nature of graphene. Not only does the coupling exist in theory, but the material also allows the physical realization of it.

A dazzling glimpse of the feasibility and potential of strain-engineered graphene (Pereira and Castro Neto, 2009; Guinea et al., 2010) has recently emerged with experiments reporting that certain strain profiles can induce Landau quantization and effective pseudomagnetic fields in excess of 300 T (Levy et al., 2010; Yan et al., 2012; Yeh et al., 2011). This strongly encourages the prospect of harnessing this unconventional interplay between electronic and mechanical properties to control electronic transport in graphene devices (Pereira and Castro Neto, 2009; Fogler et al., 2008).

This chapter starts by discussing the theory of the strain-induced vector potentials with an emphasis on the lattice-corrections first introduced by the author (Kitt et al., 2012; Kitt et al., 2013a). This is followed by an examination of the importance of

these lattice-corrections in different physical scenarios. Finally, methods of strain engineering graphene devices are examined with an emphasis on the over pressured hour glass shaped microchamber. This device cleverly takes advantage of plasmonics to enhance signals from high pseudo magnetic field regions.

2.1 Derivation of the pseudo vector potentials

By distorting the graphene lattice, strain causes the following (Pereira et al., 2009):

- (i) for any amount of strain, the Dirac points are displaced from the corners of the unstrained BZ and, furthermore, do not necessarily sit at the corners of the strained BZ; (ii) the gapless and conical nature of the energy dispersion remains robust, except when the deformation is so strong that the two inequivalent Dirac points merge in a Lifshitz transition (but that probably requires strains of the order 20%, where the tight-binding description is not reliable anymore); (iii) at any finite density the Fermi line is deformed from an isotropic circle to an elliptical shape, and two Fermi velocities can be defined along the principal directions (Pereira et al., 2009; Pereira et al., 2010; Choi et al., 2010). All these modifications are significant and happen concurrently.

From the theoretical as well as technical point of view, the effects of strain are frequently considered independently. One usually isolates the dominant effect for the physical observable of interest. For example, the strain-induced corrections to optical absorption arising from inter-band transitions are insensitive to the absolute position of the Dirac point in the BZ (i), but strongly depend on the velocity anisotropy (iii) (Pereira et al., 2010; Pellegrino et al., 2010). On the other hand, the local shift of the Dirac point (i) can hinder or completely suppress electronic propagation across regions of different strain states (Pereira and Castro Neto, 2009; Fogler et al., 2008). In a first approach the anisotropy (iii) is usually neglected in these situations(Fogler

et al., 2008; Pereira and Castro Neto, 2009).

When the strain-induced shift of the Dirac points (i) is considered independently of (ii) and (iii), it can be thought of as a pseudo vector potential (ichi Sasaki et al., 2005; Ando, 2006; Mañes, 2007; Castro Neto et al., 2009; Vozmediano et al., 2010). This can be done because of the peculiar form of the strain corrections to the electronic dispersion in graphene. Electrons in strained graphene are still governed by a Dirac equation, but one in which the strain modifications can be completely absorbed in the replacement $\hbar\vec{k} \rightarrow \hbar\vec{k} - e\vec{\mathcal{A}}$ where $\vec{\mathcal{A}}$ is the pseudo vector potential. This matches the conventional minimal coupling scenario used to describe electrons in a magnetic field. The strains response maps onto the response to a magnetic field, and thus, electrons in strained graphene can behave as if they were in a magnetic field. The pseudo vector potential is related to the shift in the Dirac point, $\Delta\vec{k}_D$, through $\Delta\vec{k}_D = -\frac{e}{\hbar}\vec{\mathcal{A}}$.

An omission in earlier work in the context of these pseudo vector potentials is the explicit consideration of the deformation of the lattice when computing the position of the new Dirac points. Here this effect is included and its importance in determining the pseudo vector potentials is shown. The inclusion of the lattice deformation yields new leading order terms in the strain-induced pseudo vector potential which are different at the three inequivalent Dirac points. The discussion is restricted to planar deformations, and hence ignore effects that might arise in the presence of curvature (Castro Neto et al., 2009; Vozmediano et al., 2010). A detailed derivation of the pseudo vector potential will be presented. It will begin with a geometric motivation of the strain-induced perturbations, continue by determining the strain dependencies of these perturbations, and finish with the derivation of the pseudo vector potentials.

2.1.1 Strain altered lattice

The key elements underlying strain-induced pseudo vector potentials are captured by generalizing the tight binding Hamiltonian discussed in Chapter 1 to the geometry of strained graphene. Figure 2.1 illustrates the changes in the lattice geometry due to strain by comparing the unstrained graphene lattice (top row) to the lattice under 20% uniaxial strain (bottom row). The large, 20% strain was chosen to better illustrate the deformation of the lattice and the BZ. Strain does not need to be this large; all of the effects discussed here are linear in strain. The method of calculating the distortion of the real and reciprocal space lattices is discussed in detail later in this chapter. Figure 2.1 will provide a qualitative geometric picture for how strain modifies the unstrained nearest neighbor tight binding approach discussed in Chapter 1.

The noticeable geometric distortion of both the nearest neighbor vectors and the lattice vectors corresponds to two distinct changes in the tight binding prescription. By making the lengths of the nearest neighbor vectors anisotropic, strain also makes the length dependent nearest neighbor hopping amplitudes anisotropic. This introduces a bond dependent hopping energy into the the unstrained real space Hamiltonian (Equation 1.1) (Hasegawa et al., 2006),

$$H = - \sum_{\langle i,j \rangle} \left(t_{i,j} a_i^\dagger b_j + t_{j,i} b_j^\dagger a_i \right) , \quad (2.1)$$

where $t_{i,j}$ is the bond specific hopping energy.

The second, and oft neglected alteration is a result of the distortion of the lattice vectors. Their alteration necessitates a change in the phase factors in the Fourier

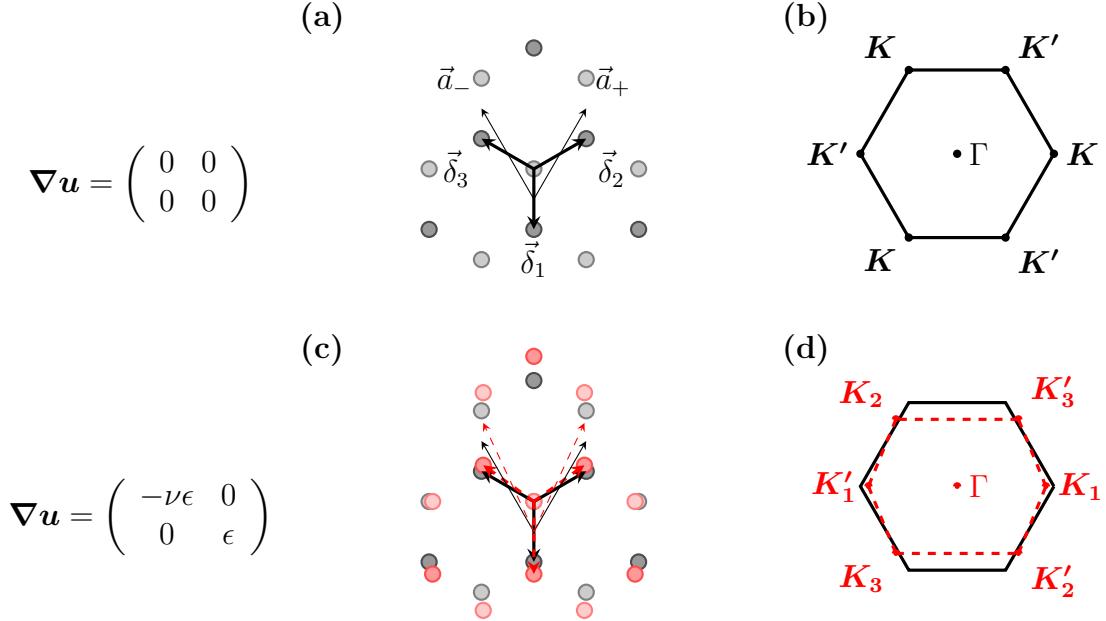


Figure 2.1: Geometry of strained graphene. (a) Unstrained graphene's real space lattice with labeled nearest neighbor vectors ($\vec{\delta}_i$), labeled lattice vectors, \vec{a}_+ and \vec{a}_- , and with light and dark dots representing the A and B sub-lattices, respectively. (b) The BZ of unstrained graphene with labeled high symmetry points. (c) The real space lattice for 20% uniaxial strain in the \hat{y} direction. Red dots represent the position of the strained atoms while the strained nearest neighbor vector, $\vec{\delta}'_i$, and the strained lattice vectors, \vec{a}'_i , are represented with red dashed lines. (d) The unstrained (black, solid) and 20% armchair uniaxial strained (red, dashed) BZ with the now inequivalent Dirac points labeled. $\nabla \mathbf{u}$ is the displacement gradient tensor describing the distortion for each row.

expansion of the creation and annihilation operators in Equation 1.2,

$$\begin{aligned} a_i^\dagger &= \frac{1}{\sqrt{N}} \sum_{\vec{k}} e^{i\vec{k} \cdot \vec{R}'_i} a_{\vec{k}}^\dagger \\ b_j &= \frac{1}{\sqrt{N}} \sum_{\vec{k}'} e^{-i\vec{k}' \cdot \vec{R}'_j} b_{\vec{k}'} . \end{aligned} \quad (2.2)$$

where \vec{R}'_i and \vec{R}'_j are the positions of the atoms in the strained A and B sub-lattices respectively. These two effects are not controllable independently in an actual physical system. Kitt *et al.* was the first to introduce the modification of the relative positions of the atoms (Kitt et al., 2012).

2.1.2 Strain altered lattice vectors

A necessary and non-trivial first step toward the inclusion of these modifications is the determination of how the lattice vectors and nearest neighbor vectors are altered by strain. This directly determines how the Fourier transforms in Equation 1.2 should be altered and is also useful for determining how the hoppings in Equation 1.1 are altered.

In general the strain is not uniform and the distortion of these vectors has a spatial dependence. In macroscopic continuum mechanics this distortion field is quantified by the elastic deformation field, $\vec{u}(\vec{r})$, which gives the position dependent deformation. Here we follow the Cauchy-Born rule which projects these macroscopic quantities onto the atomic lattice. In this approximation, the position of the i -th atom in the deformed configuration, \vec{R}'_i , is given with reference to the undeformed one, \vec{R}_i , in terms of the deformation field

$$\vec{R}'_i = \vec{R}_i + \vec{u}(\vec{R}_i) .$$

Although it is feasible that the A and B sublattices behave differently at the atomic scale but still yield the same macroscopic response, the Cauchy-Born rule provides a simple and reasonable starting point for the discussion of strain effects.

The strained lattice vector at the i -th lattice site, $\vec{a}'_{i,\pm}$, is then given approximately by

$$\begin{aligned}\vec{a}'_{i,\pm}(\vec{r}) &= \vec{R}'_j - \vec{R}'_i \\ &= (\vec{R}_j - \vec{R}_i) + \vec{u}(\vec{R}_j) - \vec{u}(\vec{R}_i) \\ &\simeq \vec{a}_{i,\pm} + (\vec{a}_{i,\pm} \cdot \vec{\nabla}) \vec{u}(\vec{R}_i) \\ &= (\mathbf{1} + \boldsymbol{\nabla} \mathbf{u}(\vec{R}_i)) \cdot \vec{a}_{i,\pm},\end{aligned}\quad (2.3)$$

where $\mathbf{1}$ is the two dimensional identity matrix and the dyadic product $\boldsymbol{\nabla} \mathbf{u}(\vec{R}_i)$ is known as the displacement gradient tensor,

$$\begin{aligned}[\boldsymbol{\nabla} \mathbf{u}]_{ij} = u_{i,j} &= \frac{u_{i,j} + u_{j,i}}{2} + \frac{u_{i,j} - u_{j,i}}{2} \\ &\equiv \tilde{\epsilon}_{ij} + \tilde{\omega}_{ij} \\ \rightarrow \quad \boldsymbol{\nabla} \mathbf{u} &= \tilde{\boldsymbol{\epsilon}} + \tilde{\boldsymbol{\omega}},\end{aligned}$$

where $\tilde{\boldsymbol{\omega}}$ is the rotation tensor and $\tilde{\boldsymbol{\epsilon}}$ is the *linear* strain tensor. It is only one part of the full (Lagrange) strain tensor given by $\boldsymbol{\epsilon} = \frac{1}{2}(\boldsymbol{\nabla} \mathbf{u} + \boldsymbol{\nabla} \mathbf{u}^\top + \boldsymbol{\nabla} \mathbf{u}^\top \boldsymbol{\nabla} \mathbf{u}) = \tilde{\boldsymbol{\epsilon}} + \frac{1}{2}(\boldsymbol{\nabla} \mathbf{u}^\top \boldsymbol{\nabla} \mathbf{u})$. It is important to stress that the often used approximating $\vec{a}'_{i,\pm}(\vec{r}) \simeq (\mathbf{1} + \tilde{\boldsymbol{\epsilon}}) \cdot \vec{a}'_{i,\pm}$ is only valid if the deformation does not involve local rotation ($\tilde{\boldsymbol{\omega}} = 0$). To simplify the notation, the position dependence has been left off of $\boldsymbol{\nabla} \mathbf{u}$. Similarly, the strain modified nearest neighbor vectors are $\vec{\delta}'_{i,j}(\vec{r}) = (\mathbf{1} + \boldsymbol{\nabla} \mathbf{u}) \cdot \vec{\delta}_j(\vec{r})$ (Kitt et al., 2013a) where i keeps track of the spatial dependence by indicating the unit cell and $j \in 1, 2, 3$ indicates the nearest neighbor vector.

2.1.3 Strain altered hopping energies

The distortion of the nearest neighbor vectors causes a bond specific alteration of the hopping energy. By comparing a tight binding treatment of strained graphene to density functional calculations of the same system, Ribeiro and coworkers showed that the dependence of the hopping energy on the distance between carbon atoms can be parameterized as

$$t(|\vec{\delta}_{i,j}|) = t_0 e^{-\beta(|\vec{\delta}_{i,j}|-a)/a} \approx t_0 - \beta t_0 \frac{|\vec{\delta}_{i,j}| - a}{a}, \quad (2.4)$$

with $\beta \approx 3$ (Ribeiro et al., 2009). An exponential decay is chosen because it correctly predicts both the nearest neighbor and next nearest neighbor hopping energies (Pereira et al., 2009). The fact that the hopping energy is mostly insensitive to the bond angle greatly simplifies the physics.

Having already established the form for the strained nearest neighbor vectors, their lengths can be easily computed for small strains

$$\begin{aligned} |\delta'_{i,j}|^2 &= (\vec{\delta}_j^\top + \vec{\delta}_j^\top \nabla \mathbf{u}^\top) \cdot (\vec{\delta}_j + \nabla \mathbf{u} \cdot \tilde{\delta}_j) \\ &= \vec{\delta}_j^\top \cdot \vec{\delta}_j + \vec{\delta}_j^\top \cdot (\nabla \mathbf{u} + \nabla \mathbf{u}^\top + \nabla \mathbf{u}^\top \nabla \mathbf{u}) \cdot \vec{\delta}_j \\ &\rightarrow |\delta'_{i,j}| \simeq a + \frac{1}{a} \vec{\delta}_j \cdot \epsilon \cdot \vec{\delta}_j. \end{aligned}$$

Here $\epsilon = \frac{1}{2}(\nabla \mathbf{u} + \nabla \mathbf{u}^\top + \nabla \mathbf{u}^\top \nabla \mathbf{u})$ is the full Lagrangian strain tensor. As was done for the displacement gradient tensor, the spatial dependence of the strain tensor is dropped from the notation for simplicity. As expected, the length of these vectors are dependent only on the the symmetric part of the displacement gradient tensor, the strain.

Thus, to first order the bond specific nearest neighbor hopping can be written as

$$t_{i,j} = t_0 + \delta t_{i,j} = t_0 - \beta t_0 \frac{1}{a^2} \vec{\delta}_j \cdot \boldsymbol{\epsilon} \cdot \vec{\delta}_j,$$

or

$$\begin{aligned} \delta t_{i,1} &= -\beta t_0 \epsilon_{yy} \\ \delta t_{i,2} &= -\beta t_0 \left(\frac{3}{4} \epsilon_{xx} + \frac{1}{4} \epsilon_{yy} + \frac{\sqrt{3}}{2} \epsilon_{xy} \right) \\ \delta t_{i,3} &= -\beta t_0 \left(\frac{3}{4} \epsilon_{xx} + \frac{1}{4} \epsilon_{yy} - \frac{\sqrt{3}}{2} \epsilon_{xy} \right). \end{aligned} \quad (2.5)$$

This fully defines the strain dependences of the modifications in Equations 2.1 and 2.2.

2.1.4 Strain altered Hamiltonian

All the pieces necessary to modify the nearest neighbor tight binding approach from Chapter 1 are now in place. The first modification to include the approximated hopping energy in Equation 2.1,

$$H = - \sum_{\langle i,j \rangle} \left((t_0 + \delta t_{i,j}) a_i^\dagger b_j + (t_0 + \delta t_{j,i}) b_j^\dagger a_i \right).$$

In situations such as the deformation of graphene around a sharp scanning tunneling microscopy tip, it is not necessarily true that $\delta t_{i,j} = \delta t_{j,i}$. In these extreme strain situations the continuum approximation can break down, the sub-lattice symmetry can be broken, and it is possible that $\delta t_{i,j} \neq \delta t_{j,i}$. However, for strains that vary slowly with respect to the lattice spacing, the continuum approximation holds and $\delta t_{i,j} \simeq \delta t_{j,i}$ (Sloan et al., 2013). This assumption will be made throughout. As a

result, the real space Hamiltonian in Equation 2.1 can be approximated to first order in strain as

$$H \simeq - \sum_{\langle i,j \rangle} (t_0 + \delta t_{i,j}) a_i^\dagger b_j + \text{H.C.},$$

with the strain dependence of $\delta t_{i,j}$ given by Equation 2.5.

Using the Fourier transforms of the creation and annihilation operators in Equation 2.2, the Hamiltonian is written in reciprocal space as

$$H = -\frac{1}{N} \sum_{\langle i,j \rangle} \sum_{\vec{k}, \vec{k}'} (t_0 + \delta t_{i,j}) e^{i(\vec{k}-\vec{k}') \cdot \vec{R}'_i} e^{-i\vec{k}' \cdot \vec{\Delta}'_{i,j}} a_{\vec{k}}^\dagger b_{\vec{k}'} + \text{H.C.}, \quad (2.6)$$

where $\vec{\Delta}'_{i,j} = \vec{R}'_j - \vec{R}'_i$. Comparing this Hamiltonian to the unstrained Hamiltonian in Equation 1.3 reiterates the effects of strain. The modifications of the hopping energy and the lattice positions are evident. Additionally, the \vec{R}'_i term is no longer the only i dependent term. Instead, the strain can vary throughout the lattice resulting in the i dependence of both $\delta t_{i,j}$ and $\vec{\Delta}'_{i,j}$. In Appendix B it is shown that the i dependences can be neglected to first order in small parameters as long as $\delta t_{i,j}$ and $\nabla \mathbf{u}$ do not have Fourier components with frequencies $\mathbf{K} - \mathbf{K}'$. Throughout this slowly varying approximation will be applied. The i dependence is then eliminated through the substitutions $\delta t_{i,j} \rightarrow \delta t_j = \langle \delta t_{i,j} \rangle$ and $\nabla \mathbf{u}_i \rightarrow \langle \nabla \mathbf{u} \rangle$ where the averages are over i .

After making the slowly varying approximation, the Hamiltonian simplifies to

$$H = - \sum_{\vec{k}} \sum_j (t_0 + \delta t_j) e^{-i\vec{k} \cdot \vec{\Delta}'_{i,j}} a_{\vec{k}}^\dagger b_{\vec{k}} + \text{H.C.}$$

By applying the low energy approximation and using the form for the strained vectors in Equation 2.3, the Hamiltonian can be approximated in terms of the three small

parameters: qa , ϵ , $\nabla \mathbf{u}$. Near \mathbf{K}_i the Hamiltonian becomes

$$\begin{aligned}
H &\simeq - \sum_{\vec{q}} \sum_j (t_0 + \delta t_j) e^{-i(\mathbf{K}_i + \vec{q}) \cdot (\mathbf{1} + \nabla \mathbf{u}) \cdot \vec{\Delta}_j} a_{\vec{k}}^\dagger b_{\vec{k}} + \text{H.C.} \\
&\simeq - \sum_{\vec{q}} \sum_j (t_0 + \delta t_j) e^{-i\mathbf{K} \cdot \vec{\Delta}_j} (1 - i\vec{q} \cdot \vec{\Delta}_j) (1 - i\mathbf{K}_i \cdot \nabla \mathbf{u} \cdot \vec{\Delta}_j) a_{\vec{k}}^\dagger b_{\vec{k}} + \text{H.C.} \\
&\simeq -t_0 \underbrace{\sum_{\vec{q}} \sum_j (1 - i\vec{q} \cdot \vec{\Delta}_j) e^{-i\mathbf{K} \cdot \vec{\Delta}_j} a_{\vec{k}}^\dagger b_{\vec{k}}}_{H_0} - \underbrace{\sum_{\vec{q}} \sum_j \delta t_j e^{-i\mathbf{K} \cdot \vec{\Delta}_j} a_{\vec{k}}^\dagger b_{\vec{k}}}_{H_{hopping}} \\
&\quad + \underbrace{i t_0 \mathbf{K}_i \cdot \nabla \mathbf{u} \sum_{\vec{q}} \sum_j \vec{\Delta}_j e^{-i\mathbf{K} \cdot \vec{\Delta}_j} a_{\vec{k}}^\dagger b_{\vec{k}}}_{H_{lattice}} + \text{H.C.},
\end{aligned}$$

with a similar form near \mathbf{K}' except with \mathbf{K} replaced with \mathbf{K}' . The full Hamiltonian nicely breaks up into three parts. The first part, H_0 , exactly matches the low energy Hamiltonian of unstrained graphene from Equations 1.5 and 1.3. Strain, then, acts to perturb the unstrained Hamiltonian through $H_{hopping}$ and $H_{lattice}$. Returning to the geometry in Figure 2.1, $H_{hopping}$ originates from the deformation of the nearest neighbor vectors and the corresponding bond specific hopping energy. $H_{lattice}$ is a result of the deformation of the lattice vectors. Before it was originally introduced by Kitt *et al.* (Kitt et al., 2012), the second term was neglected. The lattice was treated as if it were undeformed and strain only acted to alter the interactions between nearest neighbors. However, both of the perturbations are first order in small parameters; $H_{hopping}$ is $\mathcal{O}(\epsilon)$ while $H_{lattice}$ is $\mathcal{O}(\nabla \mathbf{u})$. Consequently, they contributes on equal footing and $H_{lattice}$ should be included.

2.1.5 Pseudo vector potential derivation

The perturbations $H_{hopping}$ and $H_{lattice}$ mimic the effects of a magnetic field. In the minimum coupling scenario the presence of a magnetic field is included using the

substitution $\vec{p} \rightarrow \vec{p} - e\vec{A}$ where \vec{A} is the vector potential related to the magnetic field through $\vec{B} = \nabla \times \vec{A}$. The perturbations $H_{hopping}$ and $H_{lattice}$ contribute a similar shift to graphene's crystal momentum: $\hbar\vec{k} \rightarrow \hbar\vec{k} - e\vec{A}$ where \vec{A} is the pseudo vector potential. In this formalism the crystal momentum shifts are referenced to the *undeformed* BZ. This identification can be made because like the matrix elements of H_0 , the perturbations are off diagonal and do not couple the \mathbf{K} and \mathbf{K}' points. Further, they provide only a strain dependent complex offset that is independent of the crystal momentum.

The form of H_0 shown in Equation 1.12 allows the isolation of the x and y components of \vec{A} from the real and imaginary parts of the perturbations

$$\begin{aligned}\mathcal{A}_{x,\mathbf{K}} &= -\vec{\mathcal{A}}_{x,\mathbf{K}'} = -\frac{1}{v_f e} \text{Re}[h_{hopping} + h_{lattice}] \\ \mathcal{A}_{y,\mathbf{K}} &= \vec{\mathcal{A}}_{y,\mathbf{K}'} = \frac{1}{v_f e} \text{Im}[h_{hopping} + h_{lattice}],\end{aligned}\quad (2.7)$$

where

$$\begin{aligned}h_{hopping} &= -\sum_j \delta t_j e^{-i\mathbf{K} \cdot \vec{\Delta}_j} = v_f e \frac{\phi_0}{2a} \left(-\frac{\beta}{2\pi} (\epsilon_{xx} - \epsilon_{yy}) \mp \frac{\beta}{\pi} \epsilon_{xy} \right) \\ h_{lattice} &= i t_0 \mathbf{K}_i^{(\prime)} \cdot \nabla \mathbf{u} \cdot \sum_j \vec{\Delta}_j e^{-i\mathbf{K}_i^{(\prime)} \cdot \vec{\Delta}_j} = \hbar v_f \mathbf{K}_i^{(\prime)} \cdot \nabla \mathbf{u} \cdot (\pm \hat{x} - i \hat{y}) .\end{aligned}\quad (2.8)$$

with the top sign in the \pm for \mathbf{K} the bottom sign for \mathbf{K}' , and $\mathbf{K}^{(\prime)}$ indicating the

appropriate \mathbf{K} or \mathbf{K}' . The resulting pseudo vector potentials are

$$\begin{aligned}\vec{\mathcal{A}}_{\mathbf{K}_1} = -\vec{\mathcal{A}}_{\mathbf{K}'_1} &= \frac{\phi_0}{2a} \begin{pmatrix} -\frac{4}{3\sqrt{3}}[\nabla \mathbf{u}]_{xx} \\ -\frac{4}{3\sqrt{3}}[\nabla \mathbf{u}]_{xy} \end{pmatrix} + \vec{A}_p \\ \vec{\mathcal{A}}_{\mathbf{K}_2} = -\vec{\mathcal{A}}_{\mathbf{K}'_2} &= \frac{\phi_0}{2a} \begin{pmatrix} \frac{2}{3\sqrt{3}}[\nabla \mathbf{u}]_{xx} - \frac{2}{3}[\nabla \mathbf{u}]_{yx} \\ -\frac{2}{3}[\nabla \mathbf{u}]_{yy} + \frac{2}{3\sqrt{3}}[\nabla \mathbf{u}]_{xy} \end{pmatrix} + \vec{A}_p \\ \vec{\mathcal{A}}_{\mathbf{K}_3} = -\vec{\mathcal{A}}_{\mathbf{K}'_3} &= \frac{\phi_0}{2a} \begin{pmatrix} \frac{2}{3\sqrt{3}}[\nabla \mathbf{u}]_{xx} + \frac{2}{3}[\nabla \mathbf{u}]_{yx} \\ \frac{2}{3}[\nabla \mathbf{u}]_{yy} + \frac{2}{3\sqrt{3}}[\nabla \mathbf{u}]_{xy} \end{pmatrix} + \vec{A}_p, \\ \text{with } \vec{\mathcal{A}}_p &= \frac{\phi_0}{2a} \begin{pmatrix} \frac{\beta}{2\pi}(\epsilon_{xx} - \epsilon_{yy}) \\ -\frac{\beta}{\pi}\epsilon_{xy} \end{pmatrix},\end{aligned}\tag{2.9}$$

with $\phi_0 = \frac{\hbar}{e}$ and the various \mathbf{K}_i points defined in Figure 2.1(d). The common term $\vec{\mathcal{A}}_p$ is proportional to the logarithmic derivative of the hopping, β , and arises from the hopping perturbations, δt_j , alone. It agrees with past derivations (Castro Neto et al., 2009; Vozmediano et al., 2010). The additional terms are the corrections due to lattice deformations completing the derivation of the strain-induced pseudo vector potential.

2.2 Pseudo vector potential discussion

The pseudo vector potential in Equation 2.9 was found using an analogy between the minimal substitution formalism and the shift in crystal momentum caused by strain. Even though this analogy is very strong, there are several discrepancies. First, unlike a magnetic field, a spatial deformation can not break time reversal symmetry. This limitation manifests itself in Equation 2.9 through the sign difference between pseudo vector potentials at time reversed points: $\vec{\mathcal{A}}_{\mathbf{K}_i} = -\vec{\mathcal{A}}_{\mathbf{K}'_i}$. Second, strictly speaking

the pseudo vector potential does not have gauge freedom. It is true that there are a class of pseudo vector potentials which all have the same curl. However, each member of this class corresponds to a measurably different physical realization of the system. Each member has a different strain distribution and correspondingly different strain-induced shifts in the momentum. A real vector potential is determined only by its curl whereas a pseudo vector potential has a concrete interpretation beyond its curl. Finally, in the minimal substitution formalism the gauge independent momentum is replaced with the gauge dependent canonical momentum. This does not occur for strain-induced pseudo vector potentials. The definition of the crystal momentum is not change by strain. To recognize these differences, the strain-induced momentum shifts are referred to as pseudo vector potentials.

It should be noted that the pseudo vector potential depends on the crystallographic orientation relative to the strain. Uniaxial strain of the same macroscopic value but at a different crystallographic orientation will result in different pseudo vector potentials. The form for the pseudo vector potentials in Equations 2.9 assume that the tensors are written with the x axis aligned with a zig zag direction. They can be rewritten for different orientations by rotating the tensors so that their x axis is aligned with the zig zag direction, calculating the pseudo vector potential using Equations 2.9, and then rotate the resulting pseudo vector potential back into the original reference frame.

Calculations preceding Kitt *et al.* (Kitt et al., 2012) considered only the term arising from hopping alterations, \vec{A}_p in Equation 2.9. Since $\beta \approx 3$, the deformation to the lattice contributes equally to hopping alterations. The terms are not only the same order in strain, but they have similar numerical coefficients. It is also worth emphasizing that taking explicit consideration of the lattice deformations leads to a vector potential that is different for all the corners of the BZ. This is of course

expected because under an arbitrary deformation the equivalence among the three \mathbf{K} and \mathbf{K}' points is lost.

When the lattice terms are neglected, the pseudo vector potential does not properly account for the shift in the Dirac points. Consider, for example, the seemingly trivial case of tensile isotropic strain. The traditional form of the pseudovector potential, \vec{A}_p , predicts that there should be no shift in the Dirac points ($\epsilon_{xx} = \epsilon_{yy}$ and $\epsilon_{xy} = 0$ in equation 2.9). Physically however, the BZ must shrink isotropically under tensile isotropic strain, and, by symmetry, the Dirac points should follow the corners of the shrinking BZ. The pseudo vector potential should correspond to a shift of each Dirac point toward the Γ point. Further, the direction of this shift is different for \mathbf{K}_1 than it is for \mathbf{K}_2 then it is for \mathbf{K}_3 . This illustrate two points: Preceding calculations were incomplete and the pseudo vector potential is not necessarily the same for each \mathbf{K} point.

In Figure 2.2, the reciprocal space shifts of the Dirac points predicted by the traditional and the corrected forms of the pseudovector potential are compared. The contours represent the dispersion of strained graphene calculated using a nearest neighbor tight binding model which accounts for both the strain-induced changes in hopping amplitudes and the lattice deformation (Pereira et al., 2009). The calculation of the shape of the strained BZ is described in Appendix A. For isotropic tensile strain, the lattice-corrected pseudo vector potential in Equation 2.9 properly predicts the displacement of each Dirac point due to strain. The less trivial cases of uniaxial or shear strain are also shown in Figure 2.2. The differences between the red (traditional) and orange (corrected) arrows make it clear that the lattice corrections are needed to determine the absolute position of the Dirac points in reciprocal space.

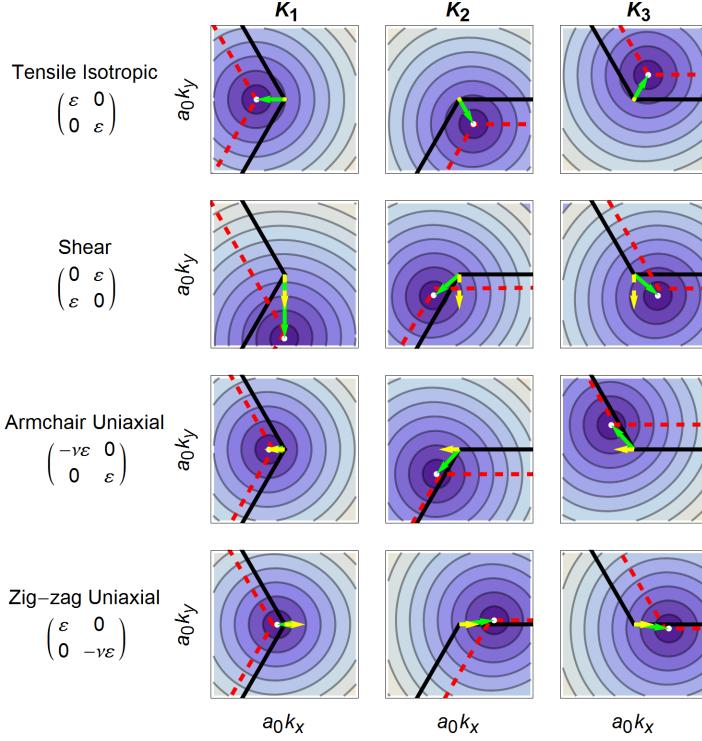


Figure 2.2: A geometric depiction of the need for the corrections to the strain-induced pseudo vector potentials. Contours of the band structure of graphene under tensile isotropic strain, shear strain, armchair uniaxial strain, and zig-zag uniaxial strain (rows), for $\epsilon = 1\%$ near the three \mathbf{K} points (columns) are overlaid with the BZ of unstrained (solid, black) and strained (dashed, red) graphene. Vectors mark the displacement of the Dirac points predicted by the traditional (\vec{A}_p , dashed, yellow arrow) and the corrected (\vec{A}_{K_i} , solid, green arrow) form of the pseudo vector potential from Equations 2.9, with the white dots marking the positions of the Dirac points for strained graphene. The yellow vectors appears as a dot for isotropic strain because the traditional form of the vector potential does not predict a shift in the Dirac points. Each plot is square with an area of 0.12^2 .

2.3 Pseudo Magnetic Fields

Having discussed the strain-induced analog of a real vector potential, the so called pseudo vector potential, it is natural to extend the analogy to pseudo magnetic fields. Although Landau level quantization is usually associated with an out of plane magnetic field, it turns out that this field is not required in graphene. In fact, Landau level quantization requires only the criterion used in deriving the pseudo vector potential: That there is a momentum shift which can be expressed as $\vec{p} - e\vec{A}$ where \vec{p} is the momentum operator which obeys the canonical commutation relationship with the position operator x , $[x, p_x] = i\hbar$ (Goerbig, 2011). It does not require a specific origin of \vec{A} . Thus, strain alone can trick the electrons into quantizing into Landau levels as if they were in a magnetic field given by $\vec{B} = \vec{\nabla} \times \vec{A}$. This effect has recently been observed using scanning tunneling microscopy of accidentally strained graphene (Levy et al., 2010; Yan et al., 2012; Yeh et al., 2011) and of engineered graphene analogs (Gomes et al., 2012).

This phenomenology is very powerful but has several subtleties. First, it should be reiterated that despite the fact that the electrons act as if there were a real magnetic field, there is not one. Strain does not somehow generate a magnetic field in a region near the graphene. Second, strain can not break time reversal symmetry. This symmetry is preserved via the relationship $\vec{A}_{\mathbf{K}_i} = -\vec{A}_{\mathbf{K}'_i}$ in Equation 2.9 which causes the pseudo magnetic field to point in opposite direction at time reversed \mathbf{K}_i points. An additional subtlety is the lack of spatial dependence in the pseudo vector potential derived here. If these calculations are treated as a local computation in the vicinity of position \vec{r} then the spatial dependence of ϵ and $\nabla \mathbf{u}$ provide the spatial dependence for $\vec{A}(\vec{r})$ as long as this pseudomagnetic field is relatively constant on the scale of the magnetic length (Castro Neto et al., 2009).

The discussion of the pseudo magnetic field in the following sections begins with a discussion of the unimportance of the lattice corrections when calculating pseudo magnetic fields. Next, a pressurized triangular graphene sealed microchamber (Guinea et al., 2010) is used as an example to illustrate how pseudo magnetic fields might be engineered and measured. This is followed by the prediction of a particularly interesting device, the hourglass shaped graphene sealed microchamber. This device should have large localized pseudo magnetic fields which can be accessed using plasmonic field enhancement. Finally, the importance of proper continuum modeling will be demonstrated by considering circular graphene sealed microchambers.

2.3.1 Contribution of lattice corrections to the pseudo magnetic field

Although the lattice corrections in Equation 2.9 are finite and, in general, have a position dependence, it happens that the lattice corrections to the pseudo vector potential do not effect the pseudo magnetic field. This has been pointed out recently by de Juan *et al.* (de Juan et al., 2013). Using Equations 2.7 and 2.8 the pseudo vector potential from lattice corrections can be recast as

$$\vec{\mathcal{A}} = -\frac{\hbar}{e} \mathbf{K}_i^{(')} \cdot \nabla \mathbf{u} = -\frac{\hbar}{e} \vec{\nabla} \left(\mathbf{K}_i^{(')} \cdot \vec{u} \right) ,$$

where the order was changed using $K_j u_j \nabla_i = \nabla_i K_j u_j$. Since the above is a total derivative, it cannot contribute to the pseudomagnetic field because $\nabla \times \nabla \phi \equiv 0$. Thus, the only term in Equation 2.9 which contributes to the pseudo magnetic field is $\vec{\mathcal{A}}_p$. As a result, even though the pseudo vector potential is different at the three \mathbf{K} points, the pseudo magnetic field is the same at each \mathbf{K} point and opposite in sign at the \mathbf{K}' points.

The lattice corrections are still needed to correctly describe the shift in the po-

sitions of the Dirac points due to strain relative to a global frame of reference. This includes the momentum-sensitive electronic tunneling to/from strained graphene from/to another system, probe, or contact (Fogler et al., 2008).

2.3.2 Pressurized graphene sealed microchambers: Pseudo magnetic field test bed

The potential of continuously replicating physics which usually only occurs at large magnetic fields is tantalizing. Experimental measurements of pseudo magnetic fields (Levy et al., 2010; Yan et al., 2012; Yeh et al., 2011; Gomes et al., 2012) have validated the theory and energized the idea of strain engineering. However, to date no strain device has been predicted, fabricated, and then shown to have the desired properties. The challenge now is to design and build devices which use this new science to enable unique device functionality.

This process of strain engineering pseudo magnetic field devices is non trivial from almost any stand point. There are no established techniques for generating spatially varying strains in atomically thin materials. In the past, graphene has deliberate been put under spatially uniform strain using strain transferred from an abutting polymer (Yu et al., 2008; Ni et al., 2008; Tsoukleri et al., 2009; Huang et al., 2009; Mohiuddin et al., 2009; Frank et al., 2010; Yoon et al., 2011), exfoliating graphene onto a piezoelectric substrate (Ding et al., 2010; Jie et al., 2013), or through the application of hydrostatic pressure (Proctor et al., 2009; Clark et al., 2013). Generating specifically designed spatially varying strains, however, is a whole other ball game. Since there are no predefined techniques for the device design, there are also no constraints provided to theorists to dictate what systems might be experimentally reasonable. Identifying and specifying new methods of generating designed spatially varying strain is thus a major priority. Beyond this challenge is the theoretical dif-

ficulty of the inverse problem. The desired property, the pseudo magnetic field, is a function of the curl of the strain which is itself a function of the surrounding environment through continuum mechanics. To complicate these non trivial mathematics is a general lack of understanding in how an atomically thin material, like graphene, responds to its surroundings. All in all, there are several challenges which must be overcome before the Quantum Hall effect might be measured at room temperature with no magnetic field.

One early, experimentally accessible design which is predicted to yield fairly uniform pseudo magnetic fields is a pressurized triangular graphene sealed microchamber (Guinea et al., 2010). To make the Landau levels as narrow and well defined as possible Guinea *et al.* proposed a variety of designs for nearly uniform pseudo magnetic fields. They noticed that strain profiles with an overall trigonal symmetry tend to generate smooth effective pseudomagnetic fields. The pressurized graphene sealed microchamber represents the most experimentally realizable of their designs. In this system equilateral triangle microchambers are etched into the underlying substrate and graphene is put across the top, sealing the gas inside. When external pressure is applied the graphene is pushed into the microchamber generating a spatially varying strain which yields fairly uniform pseudo magnetic fields.

Shown in Figure 2.3 is the spatial distribution of the pseudo magnetic field for a graphene sealed equilateral triangle microchamber with 50 nm legs oriented along the zigzag direction pressurized to 14 MPa. The pictured pseudo magnetic field is for the \mathbf{K} points, the pseudo magnetic field for \mathbf{K}' points would have the same magnitude but the opposite sign. At this pressure the graphene sheet is under less than 0.26% strain. The pseudo magnetic fields are relatively large because the strains very quickly over the relatively small length scale of the device.

To generate this figure, the strain fields calculated using finite element analysis

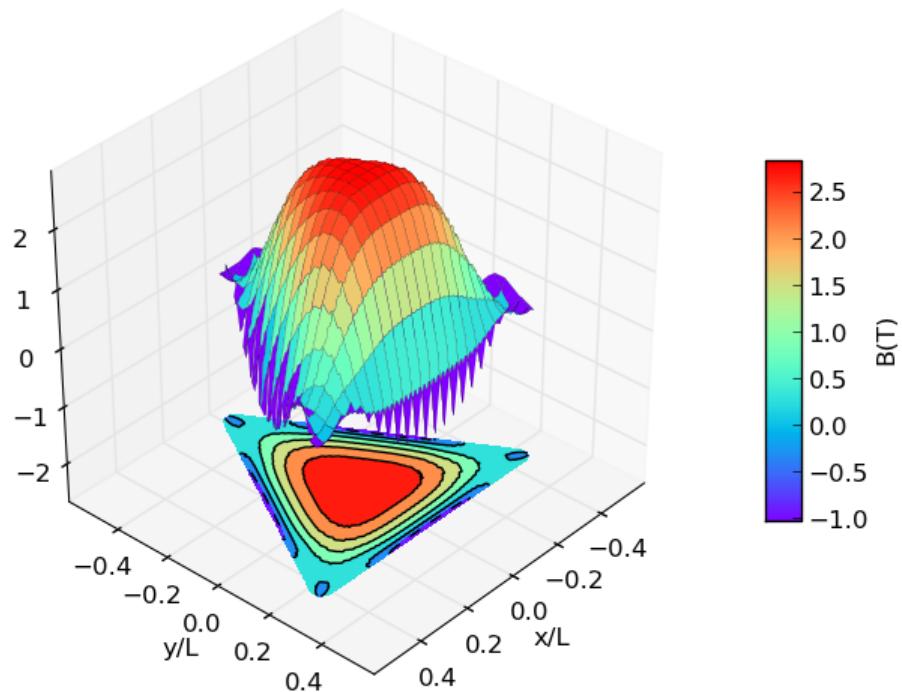


Figure 2.3: Pseudomagnetic field generated by pressurizing a triangular graphene sealed microchamber. A surface plot of the predicted spatial distribution of the pseudo magnetic field is shown above a contour plot with the same color scale. The equilateral triangular microchamber with edge lengths of 50 nm and a 2 nm fillets is pictured under 14 MPa of pressure. The graphene zigzag crystallographic orientation is in the \hat{x} direction.

were used to calculate the spatial distribution of the pseudo vector potential from \vec{A}_p in Equation 2.9. Although \vec{A}_p was derived for the special case of in plane strains, it applies to out of plane strains as well. This is because it is a function of the full Lagrangian strain tensor which properly references the deformation back to the undeformed coordinates. Finally, the curl of \vec{A}_p is taken numerically to get the pseudo magnetic field.

Finite element analysis was performed using Comsol Multiphysics with a two-dimensional thin plate model including geometric non-linearity. The edges were fixed and the pressure was applied using a face load. Graphene's Young's Modulus of 1 TPa and thickness of 3.5 Å(Lee et al., 2008) were used along with the Poisson ratio of graphite of 0.165 (Blakslee et al., 1970). To make the triangles more realistic 2 nm radius fillets were included on the corners to smooth the sharp boundaries. The surface was meshed with triangles with a maximum element size of 1 nm and strain fields were evaluated in the mid-plane of the plate.

The importance of crystallographic orientation is illustrated by this triangular graphene sealed microchamber. If the graphene is fixed while the microchamber is rotated underneath it, the pseudo magnetic field changes drastically. Figure 2.4 compares the two extreme cases: The Zigzag edge along the legs of the triangle and the armchair edge along the legs of the triangle. This 30 degree rotation ruins the large uniform pseudo magnetic field. Thus, when fabricating these devices care need to be taken to orient the graphene in the appropriate direction. One possible method to accomplish this is placing the graphene on an elastic polymer as an intermediate step. This would allow the determination of the crystallographic orientation using the polarization dependence of the Raman strain response (Huang et al., 2009). The graphene on polymer could be incorporated pick and place transfer technique which often use an intermediate polymer layer (Dean et al., 2010; Zomer et al., 2011). Since

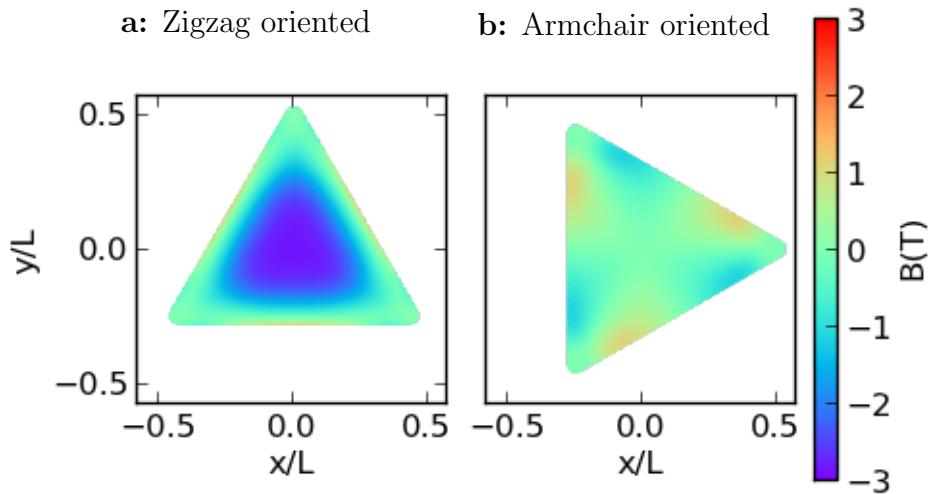


Figure 2.4: The effects of crystallographic orientation on the pseudo magnetic field for equilateral triangle graphene sealed microchambers with 50 nm legs and 2 nm fillets under 14 MPa of applied pressure. The zig zag crystallographic orientation is along the \hat{x} direction for both plots and both plots are referenced to the same color bar. In (a) and (b) the underlying microchamber is rotated such that the legs of the triangle are along the zigzag and armchair directions respectively.

different orientations can result in such different physics one should consider all crystal orientations when theoretically modeling a device.

Graphene sealed microchambers define a class of strain engineered devices which could provide a test bed for strain engineering. Even though graphene sealed microchambers do not possess the well defined electrical current path needed for electronic devices, they should still be useful as a strain engineering test bed. They represent an experimentally realizable method of generating measurable pseudo magnetic fields. Since similar devices have been made in the past, it is believable that

they can be fabricated. In fact, micron sized graphene sealed microchambers are used in the study of how graphene slides presented in Chapter 3. Additionally the Landau levels should be measurable using optical probes or local probes such as scanning tunneling spectroscopy. In particular, due to electron phonon coupling the phonon energy measured by Raman spectroscopy is renormalized when this energy matches specific Landau level transitions (Goerbig, 2011). These magneto phonon resonances have been observed on multilayer epitaxial graphene (Faugeras et al., 2009), on single layer graphene like regions on the surface of graphite crystals (Faugeras et al., 2011), and for CVD graphene on SiO₂ (Kim et al., 2013). Thus, it is easy to envision an experiment where a graphene sealed microchamber is pressurized while the Raman spectrum is measured *in situ*. Monitoring these energy fluctuations as the pseudo magnetic field is tuned by changing the applied pressure would provide an experimental test of the predicted pseudo magnetic fields.

In summary, pressurized graphene sealed microchambers represent a unique test bed for studying pseudo magnetic field effects. The strain distributions can be easily treated theoretically using finite element analysis, the predicted devices are experimentally realizable, and the resulting pseudo magnetic fields could be measured with optical techniques such as Raman spectroscopy.

2.3.3 Large, localized pseudo magnetic fields in plasmonically enhanced regions

Microchambers with different shapes were simulated using the same technique as for equilateral triangles. The predicted pseudo magnetic field distributions for simple convex shapes including circles, squares, rectangles, hexagons, and acute triangles were not compelling. Concave devices, however, proved much more interesting. When the graphene bends around a tip, the strain fields change rapidly causing large but

localized pseudo magnetic fields. Usually such a localized effect would be difficult to measure optically. However, the tip has a second function. It can enhance an optical field making the local effects measurable. If the optical field enhancement region and the region of large pseudo magnetic field were to overlap, Raman spectroscopy could be used to directly probe the high field regions.

An hourglass shaped device yields a near perfect agreement between the location of the plasmonic field enhancement and the regions of large pseudo magnetic fields. The hourglass device and plasmonic response is illustrated in Figure 2.5. The field enhancement, modeled by Arif Çetin, was modeled using a 3-dimensional finite difference time domain technique. Devices would consist of a periodic, two dimensional array of the unit cell shown in (a) and (b). In this way, any graphene sheet would cover multiple devices, increasing the signal. The unwanted signal from the supported graphene between hourglasses is minimized by destructive interference of the incident and emitted light upon reflection off the gold substrate. Device geometry, detailed in (a) and (b), was chosen so that the plasmon resonance wavelength, shown in (c), was in an experimentally accessible region. The dielectric function of the metallic layers are obtained from Palik (Palik, 1985). The predicted reflectivity in (c) was found using a power reflection monitor located $1 \mu m$ from the structure. In (d) the field enhancement of the device for light polarized in the x direction is shown. It is calculated using a field monitor located at the top surface of the top gold layer. The Raman enhancement goes as the square of the intensity enhancement since the intensity of the incident light and the intensity of the inelastically scattered light are both enhanced. The resulting Raman enhancement is greater than 5×10^5 in the region extending roughly 2 nm away from the tips. As a result, the signal measured from these tiny regions should be 3 orders of magnitude stronger than the signal from the rest of the hourglass.

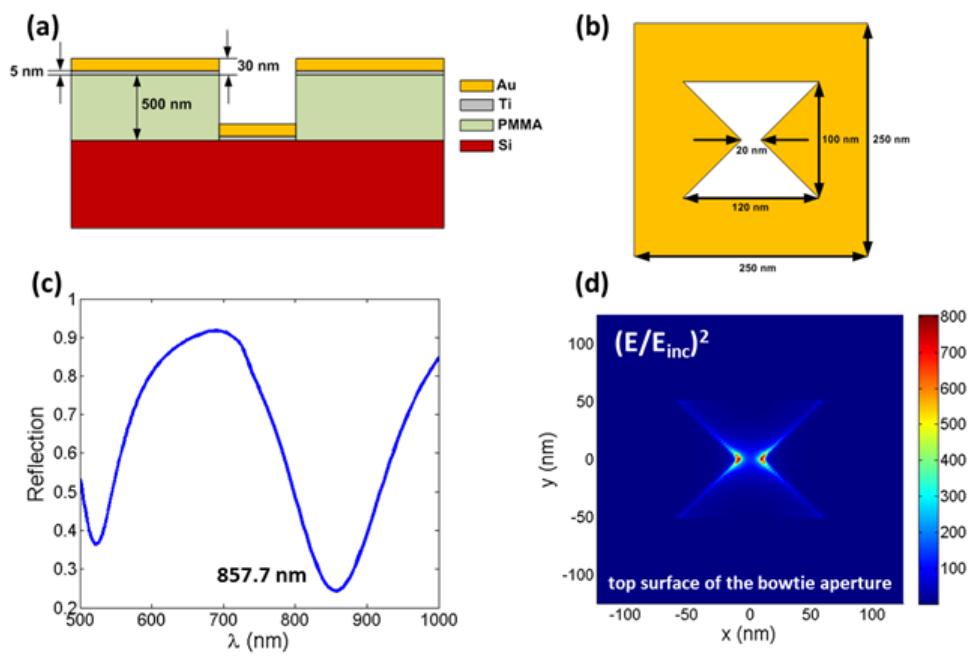


Figure 2.5: The design and plasmonic response of the hourglass device. The drawings in (a), the side view, and (b), the top view, detail the geometry of one unit cell of the plasmonic device. In (c) the predicted spectral reflectivity of the device is plotted. The intensity enhancement at the top surface of the device is plotted in (d).

The pseudo magnetic field distribution resulting from sealing the microchambers shown in Figure 2.5 with graphene and applying 14 MPa of pressure is shown in Figure 2.6. The pseudo magnetic fields reach the impressive value of 60 Tesla near the tip. This small region of large pseudo magnetic field is made experimentally accessible by the plasmonic field enhancement. When compared, the region of high pseudo magnetic field and the regions of large field enhancement are almost indistinguishable. This should allow for the measurement of pressure tunable, very large pseudo magnetic fields using Raman spectroscopy.

2.3.4 The necessity of proper continuum modeling

The first step toward modeling pseudo magnetic fields is the difficult task of determining the expected strains. This is complicated for several reasons. First, the large deflection of thin plates is often a difficult nonlinear problem. Second, many well accepted approximate continuum models are not valid for pseudo magnetic fields. Historically, most works sought approximations that accurately determined aspects of the deflection profile, sacrificing the accuracy of the strain profile in the process. Models that do not correctly predict the strain distribution can not be used to predict pseudo magnetic fields. Finally, the form of the interactions between atomically thin sheets and their surrounding environment is not yet completely understood. As a result, calculating strains can be a tricky task.

Pressurized circular graphene sealed microchambers illustrate the pitfalls of improper continuum modeling. The symmetry of the system makes analytic solutions for the strain possible. In fact, there are several different models of increasing accuracy which can be compared. The pseudo magnetic field for each is calculated from the curl of the vector potential which contributes to the pseudo magnetic field, \vec{A}_p in

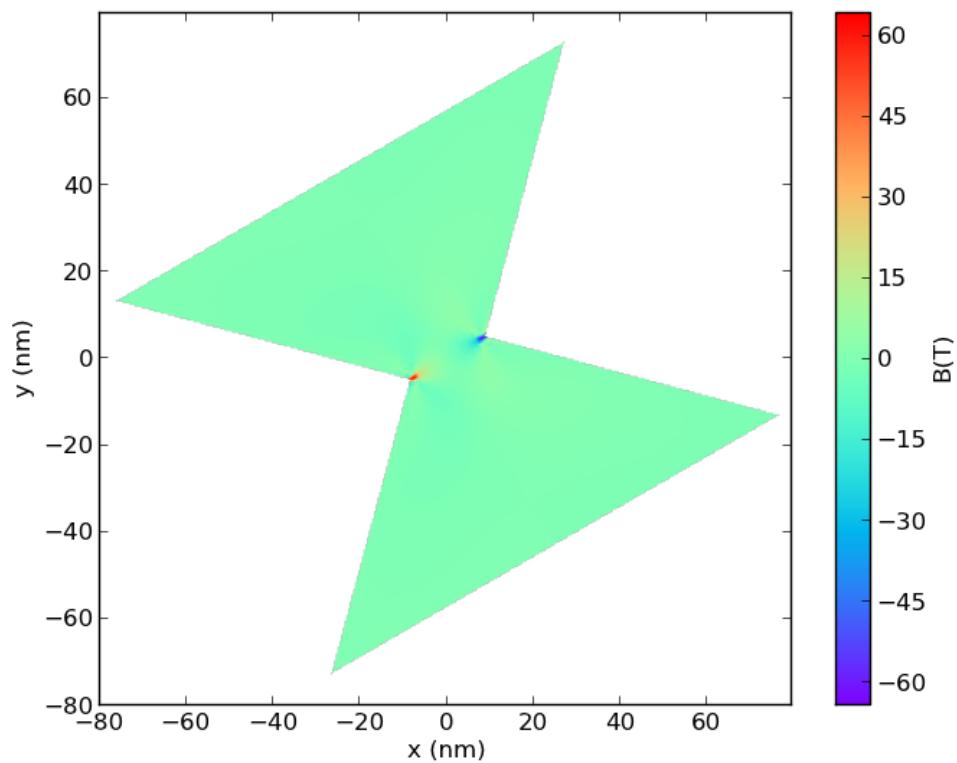


Figure 2.6: Large pseudo magnetic fields near the corners of a pressurized hourglass graphene sealed microchamber are visible in the image plot of the predicted pseudo magnetic fields. The microchamber is formed by the union of two right triangles. This leaves two 90 degree tips separated by 20 nm. The device is pressurized to 14 MPa of pressure. The graphene zigzag crystallographic orientation is in the \hat{x} direction.

Equation 2.9,

$$\vec{\mathcal{A}}_{p,cyl} = \frac{\phi_0}{2a} \frac{\beta}{2\pi} \begin{pmatrix} (\epsilon_{\rho\rho} - \epsilon_{\phi,\phi})\cos(3\phi) - 2\epsilon_{\rho,\phi}\sin(3\phi) \\ -(\epsilon_{\rho\rho} - \epsilon_{\phi,\phi})\sin(3\phi) - 2\epsilon_{\rho,\phi}\cos(3\phi) \end{pmatrix},$$

where $\vec{\mathcal{A}}_{p,cyl}$ is in cylindrical coordinates.

The spatial distributions of pseudo magnetic fields predicted using three different continuum models are shown in Figure 2.7. For each continuum model a surface plot of the pseudo magnetic field is plotted above a contour plot. The boundary of each plot is positioned at the point where the strains, and thus the pseudo magnetic fields, go to zero. In Figures 2.7a and 2.7b the plots terminates at the edge of the circular microchamber whereas in 2.7c the plot extends further because in this case the strain is allowed to spread to the supported graphene. From left to right the strain distributions used as input become more accurate. In each model the microchamber has a radius of 100 nm and is put under 70 MPa of pressure.

The strain distribution used in Figure 2.7a is based on an approximate solution useful for interpreting bulge tests. In this popular approximation the large deflection lateral displacement is assumed to have the same form as that of the small deflection limit (Timoshenko and Woinowsky-Krieger, 1959). This simple model provides everything needed to determine the Young's modulus of a thin film based on the pressure induced deflection at the center: A relationship between pressure, Young's modulus, and the center deflection. However, this approximation does not accurately describe the strain profiles and so is not applicable for studying pseudo magnetic fields. Nonetheless, Kyung-Joong Kim *et al.* use a similar model to predict exotic physical phenomena in pressurized circular graphene sealed microchambers (Kim et al., 2011). As shown in Figure 2.7a, when this model is naively used three distinct global max-

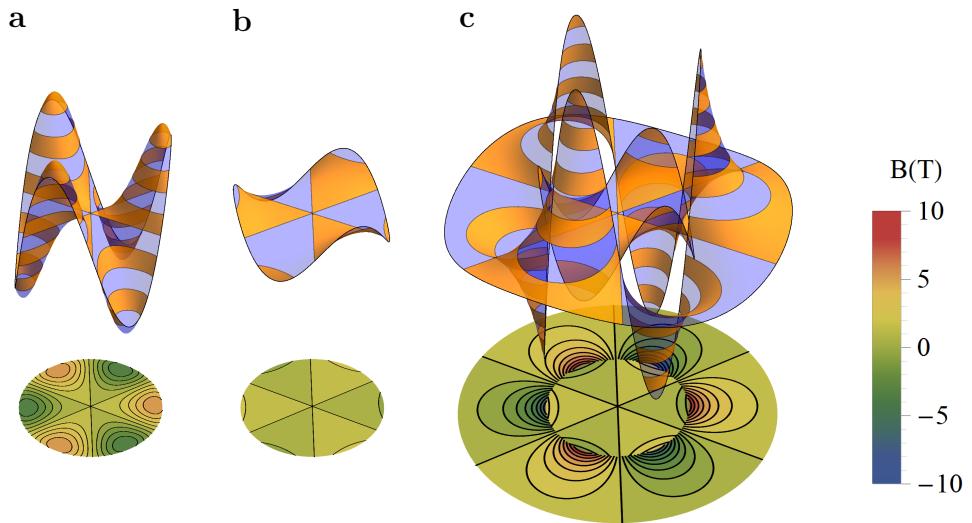


Figure 2.7: Different Pseudo magnetic field distributions resulting from three different models of a 100 nm radius circular graphene sealed microchamber pressurized to 70 MPa. The spatial distribution of the pseudo magnetic field for each model is illustrated using a surface plot with 1 Tesla bands plotted above a contour plot referenced to the color bar on the right. The z scale and the in plane scale are the same for each plot. For each model the pseudo magnetic field at the center is zero. In (a) the strains are determined using a small strain continuation model (Timoshenko and Woinowsky-Krieger, 1959), in (b) Hencky's model (Hencky, 1915) is used, and in (c) the extended Hencky model derived in Chapter 3 is used with a sliding friction of 31.5 MPa. The radial extent of (c) is larger because the extended Hencky model allows the strain to be distributed to the supported graphene.

imum and three distinct global minimum are visible with pseudo magnetic fields of nearly 5 Tesla. The six fold symmetry is a reflection of the lattice symmetry.

In Figure 2.7b Hencky's model is used to predict the pseudo magnetic field. This model is an exact series solution for the large deflection of a circular plate under a uniform vertical load with fixed boundaries. Although more complicated, this model should accurately describe the strains in the system as long as the boundaries are held fixed. The resulting pseudo magnetic fields are noticeably reduced, barely exceeding 1 Tesla. In this case the peaks occur at the boundary of the microchamber.

Finally, in 2.7c the pseudo magnetic fields are calculated using the extended Hencky model that is derived and experimentally confirmed in Chapter 3. In this model the fixed boundary conditions at the edge of the microchamber are relaxed. The supported graphene is allowed to slide against a resistive frictional force. The frictional force for this small of a device at this large of a pressure have not yet been measured experimentally. To get a qualitative idea of the changes resulting from sliding, the dimensionless frictional force found in Chapter 3 by comparison to atomistic modeling of a 6 nm radius microchamber was used. The resulting strain distribution of the suspended graphene is similar to the one predicted by the Hencky model. The strain distribution in the supported graphene, however, is much different. Since the derivative of the strain is not continuous at the edge of the microchamber, neither is the pseudo magnetic field. In fact, since the slope of the radial strain changes sign, so does the psuedo magnetic field. As friction decreases the distributed strain the pseudo magnetic field of the supported graphene decays to zero from its value of almost 9 Tesla at the edge of microchamber.

The example of cylindrical graphene sealed microchambers clearly shows that the validity of predicted pseudo magnetic fields is directly tied to the validity of the strain distribution predicted by the continuum model. One needs to be careful to ensure

that the approximations made in solving for the large deflection of the plate do not make the resulting strain distributions inaccurate. Further, one has to be careful to accurately account for how the graphene interacts with its environment.

2.4 Conclusion

The pseudo vector potential was derived in full and new, \mathbf{K}_i point dependent corrections to the pseudo vector potential were found. Even though these correction do not contribute to the pseudo magnetic field, they are necessary to complete the analogy between strain-induced momentum shifts and a real vector potential. Additionally, these shifts are required to accurately determine the position of the Dirac points and should be taken into account when considering electrical transport over a strain boundary.

Sealed microchambers were discussed as a test bed for predicting and measuring pseudo magnetic fields. It was shown that an hour glass microchamber should generate very large pseudo magnetic fields while also enabling their measurement through plasmonic enhancement. Finally, the need for accurate continuum modeling was demonstrated using cylindrical microchambers as an example.

Chapter 3

How graphene slides: Measurement and theory of strain-dependent frictional forces between graphene and SiO₂

Graphene is an amazing mechanical system with extreme elasticity (Lee et al., 2008), ultrastrong adhesion (Koenig et al., 2011), and impermeability to gases (Bunch et al., 2008). As a pure two dimensional material, graphene's interactions with its supporting substrate are unique. Amontons' first law states that macroscopic friction is proportional to the applied load, justified by arguing that increasing the load increases the microscopic contact area between two surfaces (Krim, 1996). Graphene, however, because of its ultrastrong adhesion (Koenig et al., 2011) and low bending rigidity requires no load to achieve nearly perfect conformation to the nanoscale topography of its substrate, especially the commonly used SiO₂ (Stolyarova et al., 2007; Lui et al., 2009; Cullen et al., 2010). Hence, the friction between graphene and SiO₂ might be expected to exhibit an atypical load dependence. The ability to control the thickness of few layer graphene (FLG) at an atomic scale makes it an excellent model system to study the role of thickness and load on friction, which has not previously been quantified or elucidated in detail. To date most tribological studies of FLG and graphitic materials have measured the interaction between graphene and a scanning probe tip using frictional force microscopy (Dienwiebel et al., 2004; Deng et al., 2012; Lee et al., 2010; Li et al., 2010b; Filleyer et al., 2009; Filleyer and Bennewitz, 2010;

Zhang et al., 2012). These nanoscale measurements have shown interesting effects such as superlubricity in graphite (Dienwiebel et al., 2004), negative frictional coefficient for chemically modified graphite (Deng et al., 2012), and increasing friction with decreasing FLG thickness (Lee et al., 2010; Li et al., 2010b; Filleter et al., 2009; Filleter and Bennewitz, 2010). Both the negative frictional coefficient and the increasing friction with decreasing thickness have been attributed to the puckering of graphene about the scanning probe tip (Lee et al., 2010; Li et al., 2010b; Deng et al., 2012).

Based on work by the author (Kitt et al., 2013b), this chapter discusses direct measurements of the intrinsic sliding of graphene over a SiO₂ substrate at the macroscopic device scale. Both the load and atomic layer dependence of sliding friction, or the substrate's resistance to graphene sliding are extracted from the data. Uniquely, the graphene-substrate interaction is isolated by avoiding the introduction of a scanning probe tip. Instead, the system is studied by using variable gas pressure applied to an FLG sealed microchamber as shown in Figure 3.1. The pressure acts as a tunable load, simultaneously pressing the supported graphene into the substrate while forcing the suspended FLG into the microchamber. *In situ* Raman measurements, which can easily measure FLG extensions of 1 nm over 1 μm , show that an annulus of the supported FLG reproducibly slides toward the center of the microchamber. By analyzing the strain response with a newly derived extension of the continuum Hencky model, the load-dependent sliding frictions for mono-, bi-, and tri-layer graphene is extracted. The layer dependence exhibits a crossover between bilayer and trilayer; the trilayer sliding friction obeys Amontons' first law, whereas the monolayer and bilayer sliding friction uniquely scales with the inverse of the strain in the graphene. These interesting results are attributed to the interplay between adhesion, in-plane strain and bending rigidity in this two dimensional tribological system.

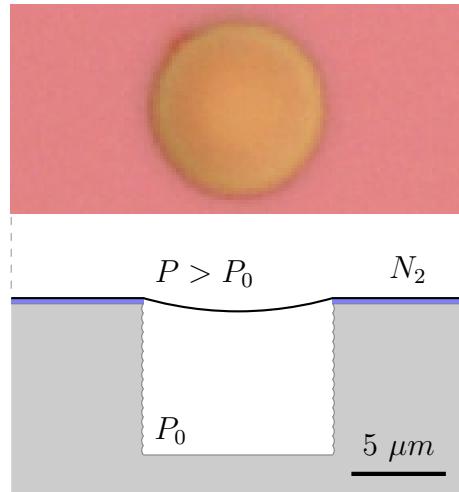


Figure 3.1: Top: An optical image of a trilayer graphene-sealed microchamber. Bottom: Device cross-section schematic showing the microchamber etched $8 \mu m$ into the underlying Si substrate and the supported graphene atop the 300 nm of thermal oxide. Pictured to scale is the largest pressure induced deflection of the graphene achieved in any of the analyzed experiments.

A firm understanding of graphene's sliding friction is necessary for a variety of exciting graphene devices such as flexible bistable displays (Bonaccorso et al., 2010), graphene electro-mechanical switches (Milaninia et al., 2009), high quality factor graphene mechanical resonators (Kim and Park, 2009; Bunch et al., 2007; Chen et al., 2009; Barton et al., 2011), and strain engineered devices (Pereira and Castro Neto, 2009) which take advantage of strain induced vector-potentials and pseudo magnetic fields (Castro Neto et al., 2009; Guinea et al., 2010; Kitt et al., 2012; Kitt et al., 2013a) as discussed in Chapter 2. Specifically, as demonstrated in Section 2.3.4, you can not predict pseudo magnetic fields without understanding the unique mechanics of how a floppy two dimensional system interacts with a rigid three dimensional body.

3.1 Raman G band strain response

Micro-Raman spectroscopy is a powerful tool to measure strain distributions in graphene. The Raman G band measures the zone center, in-plane optical phonons that are degenerate at zero strain (Tuinstra and Koenig, 1970; Ferrari et al., 2006). In the absence of shear strain, the G band shifts according to (Huang et al., 2009)

$$\Delta\omega_G = -\omega_0\gamma(\epsilon_r + \epsilon_t) \pm \frac{1}{2}\beta(\epsilon_r - \epsilon_t) , \quad (3.1)$$

where ϵ_r and ϵ_t are the strain in the radial and tangential directions, γ is the Grüneisen parameter and β is the shear deformation potential which details the amount of splitting between the G^+ and G^- bands.

Light scattered by the G^+ and G^- bands has orthogonal linear polarizations (Huang et al., 2009). To isolate the individual peaks spectra are often measured using linearly polarized light. Here, however, spectra are measured with circularly polarized light unless otherwise noted. This allows more data to be acquired because the G^+ and G^- bands are measured simultaneously. Figure 3.2 shows experimental verification that the spectra taken using circular polarized light matches the sum of the G^+ and G^- bands. The G^+ and G^- bands were measured in the usual way by setting the incident excitation along the radial strain direction and measuring the emitted spectra as a function of emission polarization. By summing over the spectra taken with emitted polarization varied from -90 to 90 degrees in 20 degree steps, a spectra with equal contributions from G^+ and G^- is created. As shown in Figure 3.2, these derived spectra match the single spectra taken with circularly polarized light. This verifies that spectra measured with circularly polarized represent both the G^+ and G^- bands.

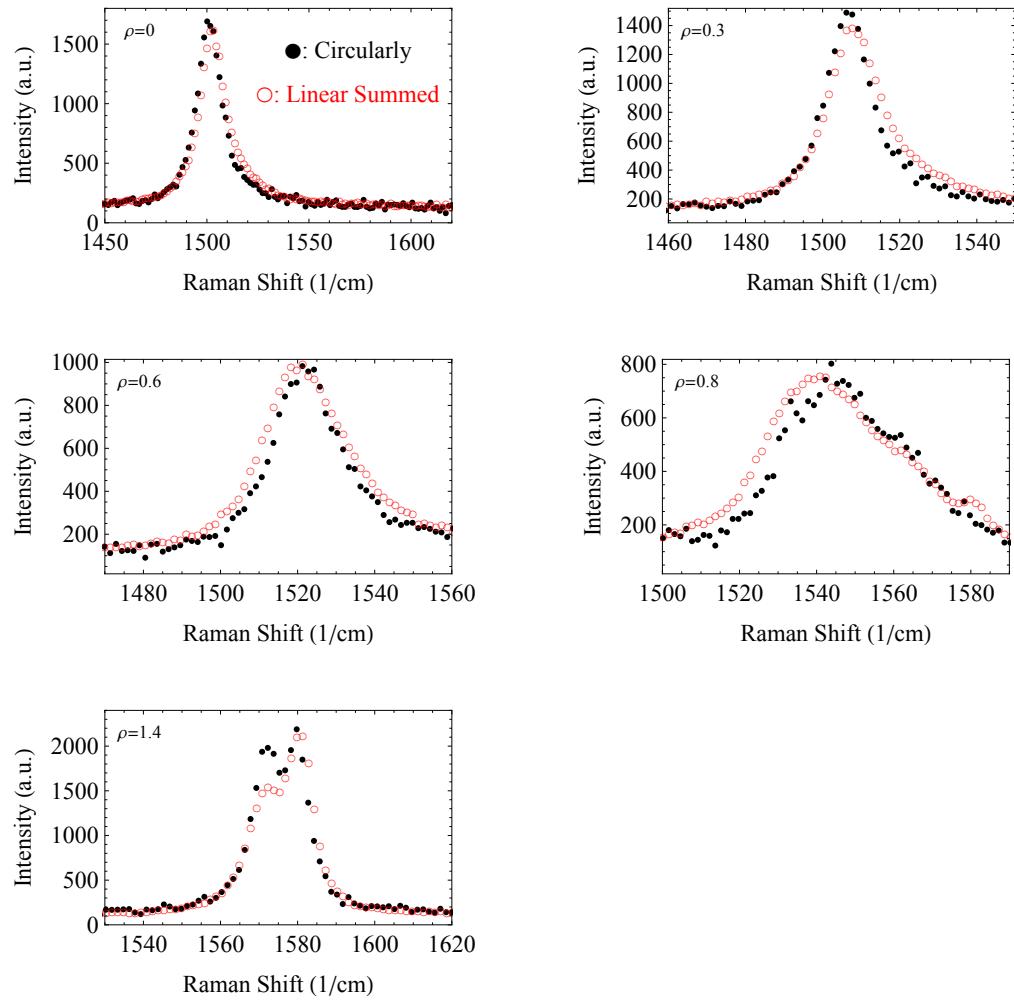


Figure 3.2: Comparison of Raman spectra measured using circularly and linearly polarized light. Measurements were taken at the labeled values of $\rho = r/R$ on a $R=5\text{ }\mu\text{m}$ graphene sealed microchamber under 0.80 MPa of applied absolute pressure. The black dots are spectrum measured using circularly polarized light. The red dots are the sum of spectra taken with outgoing polarizations varied between -90 to 90 degrees in 20 degree steps while the excitation is held in the \hat{x} direction. Spectra are scaled to match intensities.

The local Raman response is measured inside an optically accessible pressure chamber with a focused laser beam while variable pressures up to 0.80 MPa are used to push the FLG into the microchamber. Raman spectra are excited using the 514 nm line of an argon ion laser and collected using a Renishaw spectrometer with an 1800 lines/mm grating and a 63X, .7 NA, cover slip corrected objective. The laser power in the pressure chamber was kept below 0.5 mW to avoid sample heating. Optical access into the pressure chamber is through a 1 mm BK7 window. The beam waist of the focused laser was measured to be $0.81 \pm .01 \mu\text{m}$ by scanning a gold pad under the laser as shown in Figure 3.3.

3.2 Experimental Design

A cross-section of one of the microchambers sealed with mechanically exfoliated FLG is depicted in Figure 3.1. Microchambers are fabricated using standard optical lithographic methods to define holes ranging from 1 to 5 μm in radii, reactive ion etching to etch through the 300 nm thermal oxide layer, and deep reactive ion etching to etch roughly 8 μm into the underlying silicon layer. Before the graphene is mechanically exfoliated to seal the microchambers, the substrate is oxygen plasma ashed for 10 minutes at 300 sccm and 500 Watts to ensure the substrate is cleaned of any residues. The number of graphene layers sealing each device was determined using Raman spectroscopy (Ferrari et al., 2006) and optical contrast (Blake et al., 2007; Casiraghi et al., 2007).

The large microchamber depth of $\sim 8 \mu\text{m}$ is 10 times the largest FLG deflection of 700 nm. This allows for the changes in internal pressure as the microchamber shrinks with pressure to be ignored. It also enables longer measurement times because of slower leak-out rates through the silicon substrate. To eliminate surface residues, the

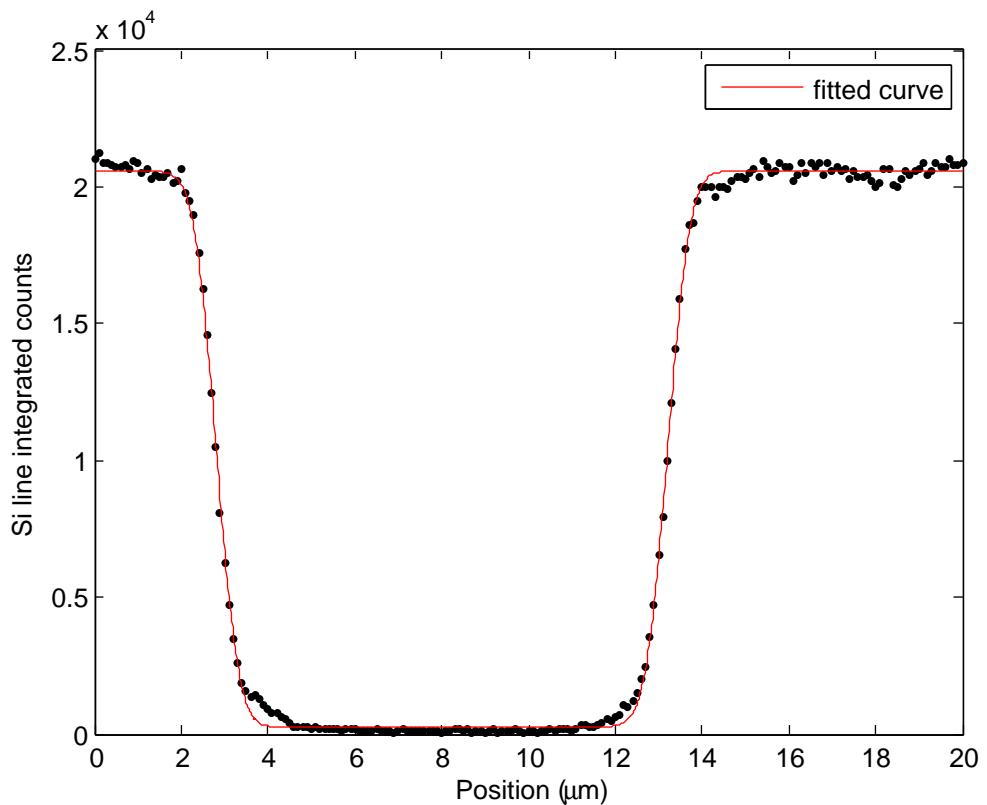


Figure 3.3: Determination of the beam waist of the focused laser in the pressure chamber by scanning a sample with a gold pad under the focused beam while measuring the Si Raman line intensity. As the sharp edge of the gold pad moves under the laser, the Si line is blocked giving a measure of the beam size. Fitting the signal to an error function gives a beam waist of $0.81 \pm .01 \mu\text{m}$.

substrate was oxygen plasma ashed before FLG exfoliation. It is important to note that different surface treatments may yield different sliding frictions providing a new degree of freedom in device engineering.

Two complementary Raman measurements are performed *in situ* to fully characterize the strain distributions. First, as the absolute applied pressure is varied between atmospheric pressure (0.10 MPa) and 0.80 MPa, Raman spectra at the center of the microchamber are recorded. Also, at selected pressures, Raman G band line scans with 0.5 μm point spacing are taken across the microchamber. The former is used to determine the pressure trapped inside the microchamber while the latter is used to determine the spatial distribution of the strain. In conjunction with low force ($\approx 1 \text{ nN}$) contact mode atomic force microscopy (AFM), the Raman spectra taken at the center of the microchamber reveal interesting ambient pressure behavior exhibited by the graphene sealed microchambers.

3.2.1 Pressure trapped in microchambers

The ambient pressure behavior of the measured devices is split roughly evenly between two representative cases. For half of the devices the pressure inside of the microchamber was greater than one atmosphere, as shown in the left side (a) of Figure 3.4. Here, both the AFM and Raman measurements indicate that the pressure inside of the microchamber is greater than ambient. The topographic image shows the graphene bulging out and the Raman G band response demonstrates that a non-zero gauge pressure is necessary for the G band to reach its greatest, or least strained, value. By fitting the Raman data with the $p^{2/3}$ form predicted by the standard Hencky model for the strain at the center of the hole (blue dashed line), the internal pressure and unstrained G band position are found.

In the second half of the devices, the pressure inside the microchamber was an

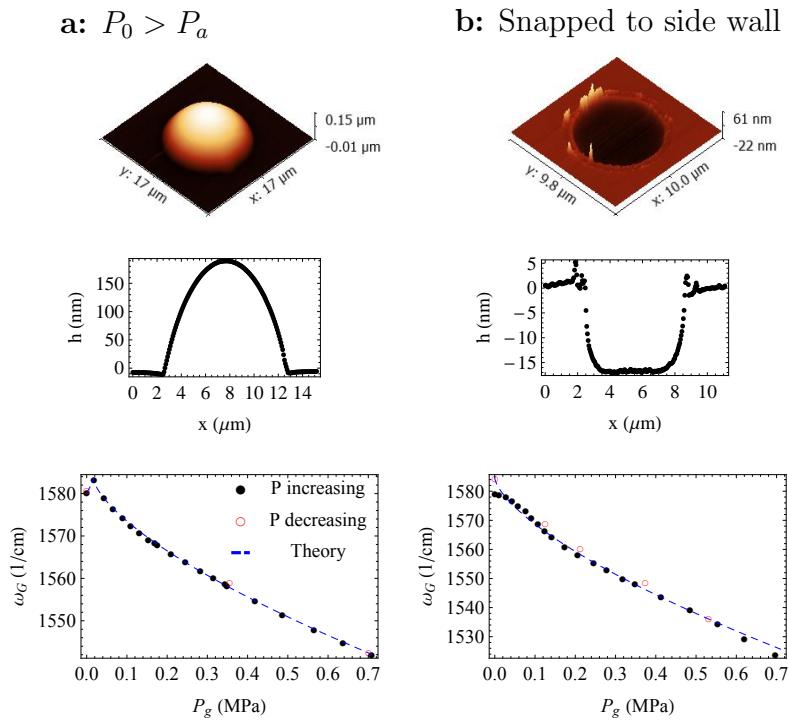


Figure 3.4: Characteristic ambient pressure behavior of FLG sealed microchambers. (Top) Low contact force (≈ 1 nN) AFM taken at ambient pressure with line cuts above (bottom) the center frequency of single Lorentzian fit spectra taken at the center of the hole as a function of the applied gauge pressure for (a) (left hand side) a $5 \mu\text{m}$ radius trilayer sealed microchamber and (b) (right hand side) a $3 \mu\text{m}$ radius monolayer sealed microchamber.

atmosphere, as shown on the right side (b) of Figure 3.4. The AFM topography shows that the graphene is flat across the aperture indicating the internal pressure matches the ambient pressure. The topography also shows that graphene is stuck to the walls of the microchamber for an AFM determined distance of roughly 7 nm. Additionally, the Raman G band response in these devices shows interesting small strain behavior with increasing pressure which seems to be a result of graphene peeling off the side walls as it is pressurized. At atmospheric pressure the G band is already red shifted due to pre-strain. Increasing the gauge pressure to ~ 0.1 MPa causes a slower than expected downshift in the G band position. However, when the gauge pressure is increased beyond 0.1 MPa, the shift rate converges to the expected trend, $P^{2/3}$, and remains on this trend when the pressure is reduced back to atmospheric pressure. This is consistent with the FLG unsticking from the walls as pressure is applied and remaining unstuck when the pressure returns to ambient pressure. Values for friction are extracted from data taken at 0.17 MPa and above, well beyond the pressure domain where side wall sticking effects are observed. For devices that behave in this way, the internal pressure is taken to be an atmosphere and the unstrained G band position is set as the G band position upon completing the pressure cycle.

Other groups have observed similar snap-to-sidewall behavior (Lee et al., 2008; Bunch et al., 2008). Bunch and Dunn noted that the AFM measured distance over which FLG snaps to sidewalls may be overestimated (Bunch and Dunn, 2012). This position is supported by the comparison of the strain measured with Raman and the strain estimated from topography. For a 3 μm hole, a change in radius of 7 nm should generate .2% strain whereas a Raman shift of -5 1/cm indicates a strain of .08%. Additionally, side wall sticking may explain the small Young's Modulus measured by Lee *et al.* using a similar measurement over the a 0.1 MPa pressure range (Lee et al., 2012).

3.3 Qualitative results

Raman line scans over pressurized microchambers show that the supported graphene around the microchamber has slid inward toward the center. Figure 3.5 shows the center frequency of single Lorentzian fits to the G band as a function of position across a $6 \mu\text{m}$ diameter monolayer covered graphene sealed microchamber with applied absolute pressures of 0.45 and 0.80 MPa during three separate pressure cycles from atmospheric to 0.80 MPa. As expected, as the suspended graphene is pushed down into the microchamber, the G band red shifts or softens from its unstrained value. Unexpectedly, the G band of the supported graphene *outside* the edge of the microchamber also shows softening, and thus significant strain. The observed softening decreases with the distance from the edge of the microchamber until the G band returns to its unstrained energy. This strain is real; the G band red shift cannot be attributed to the averaging over the finite spot size of the beam because the measured downshifts persist much further from the edge of the microchamber than the $0.83 \mu\text{m}$ beam waist. As the applied pressure increases, more strain is distributed outside of the microchamber causing both a larger redshift and a larger region over which the strain is distributed. The strain distributed outside of the microchamber's edge is a clear indicator that the graphene is not rigidly fixed to the substrate outside of the microchamber. Instead of a line force acting at the circumference of the microchamber to fix the graphene at the edge, there must be a distributed sliding frictional force, f , acting between the graphene and the substrate.

This behavior is reproducible, stable, and azimuthally symmetric. The four 0.45 MPa line scans in Figure 3.5 include one line scan in the x direction for each of the first two pressure cycles and a line scan from both the x and y direction during the third pressure cycle. The five 0.80 MPa line scans include one line scan in the x direction

from the first pressure cycle, two sequential line scans in the x direction which took 35 minutes each from the second pressure cycle, and a line scan from both the x and y direction during the third pressure cycle. Other than the development of a dimple at the center of the microchamber, the spectra and G band shifts are nearly identical. This dimple is the result of laser deposition of dirt at the center of microchamber due to tens of hours of high pressure resolution, single point measurements. An SEM image of the schmutz dimple is shown in the inset. This dirt seems to stabilize the graphene underneath, reducing the strain in its vicinity. Similar behavior is observed in each of the eight measured microchambers which have radii between 1.2 and 5 μm and are covered with from one to three layers.

To determine the nature of the strain outside of the microchamber, Raman spectra 2 μm outside the edge of a 5 μm radius monolayer covered microchamber were analyzed in detail using linearly polarized light. As shown in Figure 3.6, there are two discrete peaks in the spectra at 1570.9 1/cm and 1581.3 1/cm which are tuned on and off by rotating the analyzer on the outgoing light. The polarization dependence of the integrated peak areas is consistent with the orthogonality of the G^+ and G^- peaks (Huang et al., 2009). The peak positions indicate a tensile radial strain of 0.6% and a compressive tangential strain of -0.3% at this location. The compressive tangential strain is expected: When an annulus of the supported FLG is pulled inward, its circumference shrinks and, if the adhesion energy between FLG and its substrate is large enough to suppress out-of-plane wrinkling, this shrinkage causes compressive tangential strain.

In the Raman and AFM experiments there is no evidence of the FLG wrinkling to relieve its compressive strain. Shown in Figure 3.7 is a partial Raman map of this 5 μm radius monolayer covered microchamber at 0.80 MPa of applied pressure. To generate the Raman maps Raman spectra were taken at each spatial location

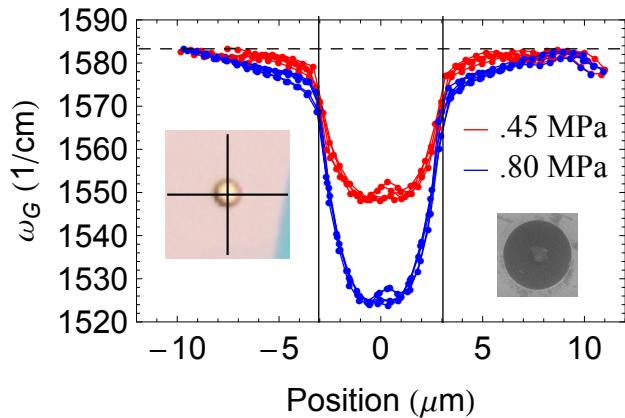


Figure 3.5: Qualitative observations of graphene sliding. The frequency shift of the Raman G band as a function of position for four line scans taken at 0.45 MPa and five line scans taken at 0.80 MPa from 3 separate pressure cycles scanned across a single 6 μm diameter monolayer sealed microchamber. Each point represents the position of the center of a single Lorentzian fit to the Raman spectra at that position. Solid vertical black lines are positioned at the edges of the microchamber and the dashed horizontal line indicates the zero strain position of the G band. Data points are separated by 0.5 μm ; the focused beam has a waste of 0.81 μm . Inset left is an optical image of the microchamber with line scan direction indicated. The bottom right inset is an SEM image of the device taken after all of the data was acquired.

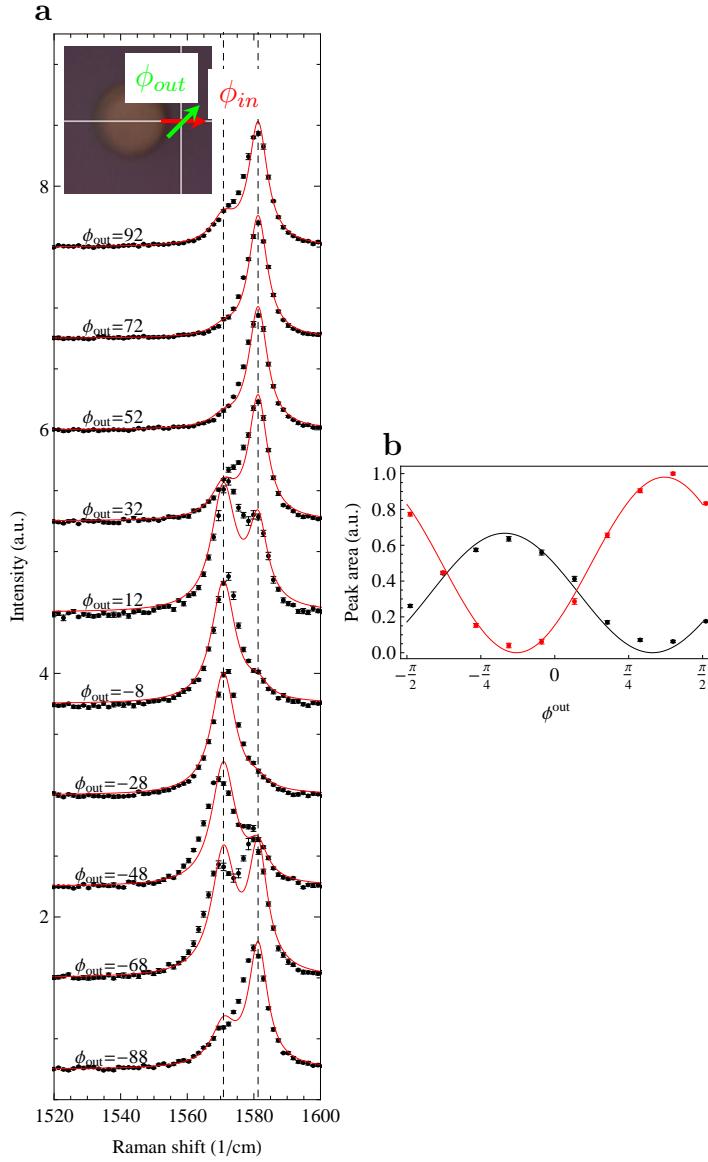


Figure 3.6: (a) Linearly polarized Raman spectra of supported graphene taken 2 μm outside a 5 μm radius graphene sealed microchamber pressurized to 0.80 MPa absolute pressure with the incident light polarized in the \hat{x} direction ($\phi = 0$) and the outgoing light linearly analyzed at the labeled angle. (b) The areas under the 1570.9 1/cm and the 1581.3 1/cm peak as a function of outgoing analyzer angle in black and red, respectively fit with π out-of-phase sine squared functions.

on the map. Each spectra was fit with Lorentzians of energies ω^+ and ω^- and the best fit energies were plotted as a function of the spatial location where the spectra was taken. This two Lorentzian fit matches the characteristic two peak shape of the supported graphene shown in Figure 3.2 for $\rho = 1.4$. Both maps in Figure 3.7 were scaled to emphasize the shifts in the supported graphene peak positions. As a result, the more highly strained suspended graphene appears black in the maps. The ω^- Raman map, (a), shows very high radial symmetry while the ω^+ map, (b), has slightly reduced symmetry. At the highest point on the microchamber, there is a variation of around four wavenumbers in the ω^+ map which could be due to small variations in local adhesion or could be due to nonuniform doping. The degree of radial symmetry exhibited by both of these Raman maps is a strong indicator that the supported graphene is not wrinkling. Since these Raman maps are for the microchamber with the largest diameter and thinnest thickness at the highest applied pressure, the compressive strains are the highest of any of the microchambers measured. Thus, wrinkling in all other measured microchambers which have lower compressive strain can be precluded.

3.4 Continuum model of strain distributions

In this section, a continuum model to extract the sliding friction, f , from the Raman determined strain distributions is developed. In 1915, Hencky proposed a continuum model for the non-linear pressure induced deflection of a thin circular plate with fixed boundary conditions (Hencky, 1915; Fichter, 1997). This model has been successfully used to describe a variety of systems including inflatable membrane mirrors (Meinel and Meinel, 2000), electrostatic actuators for micro gas pumps (Zhang and Wang, 2011), and the topography of FLG bulging from sealed microchambers (Koenig et al.,

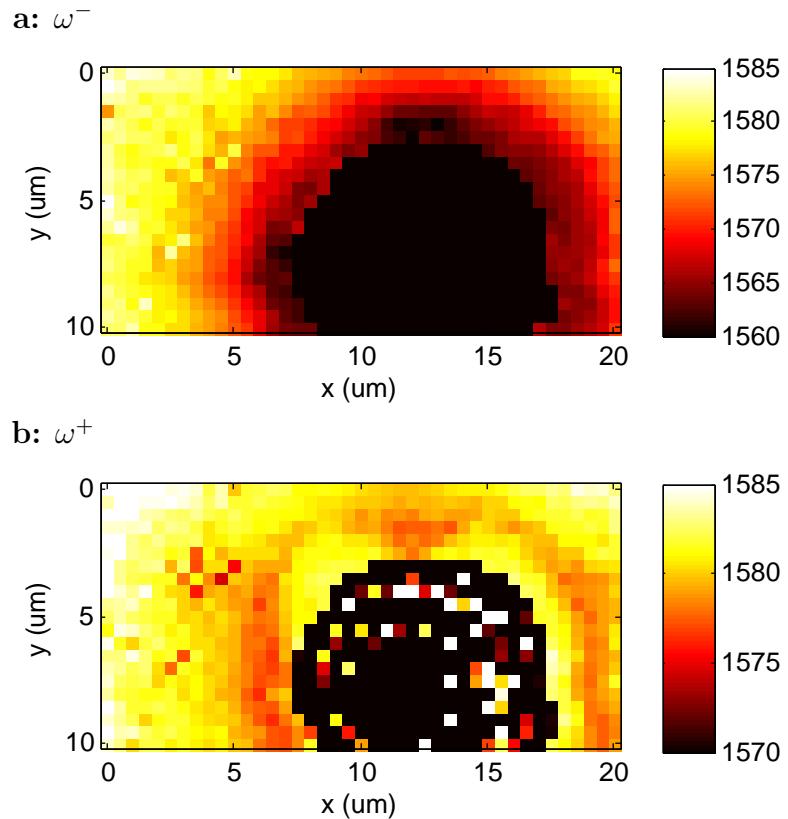


Figure 3.7: Raman map of pressurized graphene sealed microchamber. (a) The ω^+ peak and (b) ω^- peak for a 5 μm radius graphene sealed microchamber pressurized to 0.80 MPa. Color scale are set to emphasize the peak shifts experienced by the supported graphene. The speckles in (b) appear because the suspended graphene spectra are not fit well by two Lorentzians. Both maps show a high degree of radial symmetry inconsistent with the formation of wrinkles.

2011). However, the fixed boundary conditions assumed by this model preclude its application to the strain distributions observed here. The fixed boundary conditions can be relaxed by matching the radial and tangential stresses inside the hole, derived from Hencky's model before the application of boundary conditions, to the radial and tangential stresses of the supported material outside of the hole found by including a sliding friction, f , acting against the radial displacement. The stresses and, using Hooke's law, the strains, are then fully determined as a function of $\frac{\Delta P^2 E_{2D}}{f^3 R}$ where R is the radius of the microchamber measured by AFM, ΔP is the differential pressure, and E_{2D} is the 2D Young's modulus of FLG taken to be $n \times 340 \text{ N/m}$ where n is the number of layers (Lee et al., 2008; Koenig et al., 2011).

Figure 3.8 compares the radial and tangential strains from the standard Hencky solution, the extended Hencky model derived here, and an atomistic, molecular dynamics model. The solid lines of the extended model demonstrate the desired features; strain is distributed *outside* of the hole with compressive tangential and tensile radial strain. The strain distribution depends on the friction as expected: At constant pressure and radius, a greater sliding friction holds the graphene more firmly to the substrate surrounding the microchamber, and thereby increases ϵ_c , the strain in the center of the microchamber, while also reducing ρ_0 , the largest radial distance that the strain acts outside of the hole. The extended model, with $f=520 \text{ MPa}$, is in good agreement with the dots in Figure 3.8 that are the results of an atomistic molecular dynamics simulation of a 6 nm radius microchamber under 500 MPa of pressure performed using the open source simulation package LAMMPS (Lammps, 2012; Plimpton, 1995) developed at Sandia National Labs.

This work represents the first time that Hencky's model has been generalized to allow strain to be distributed outside of the microchamber's edge. The derivations provides a framework for including the sliding of graphene over a rigid substrate in a

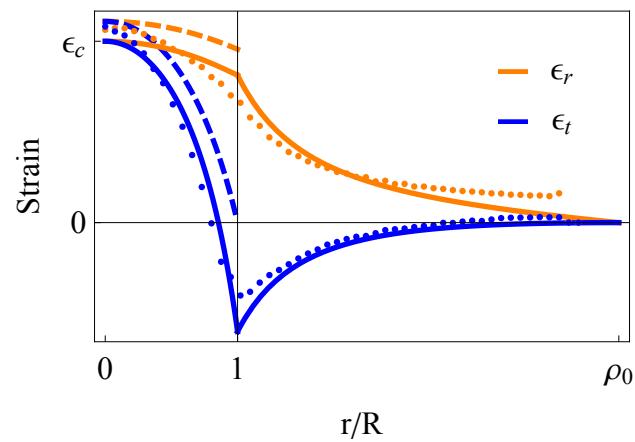


Figure 3.8: Theoretical strains in FLG sealed microchambers. Comparisons of the radial and tangential strains predicted by Hencky's model (dashed), the extended Hencky model that includes strain outside the microchamber for a sliding friction of $f=520$ MPa (solid), and atomistic simulations of a 6 nm radius microchamber with 500 MPa of applied pressure (dots). The extended Hencky model used to extract friction agrees very well with the atomistic model.

continuum model. This should be useful in applications such as the strain engineering of pseudo magnetic fields. For those interested, a full derivation of the extended Hencky model as well as a detailed description of the atomistic modeling is included in the next sections.

3.4.1 Detailed derivation of the extended Hencky model

To derive the extended Hencky model the general forms of the stress inside and outside of the hole will be derived, and then the boundary condition will be determined and applied. The resulting solution will be a function of the dimensionless loading parameter, $q = \frac{\Delta PR}{E_{2D}}$, and the dimensionless friction coefficient $F = Rf/E_{2D}$ in the combination $q^{2/3}/F$. This means that large dimensionless loading, $q^{2/3}$, causes the same effects as small friction, F ; pushing hard has the same effect as sliding easily. A simple analytic relationship does not exist between $q^{2/3}/F$ and the two parameters which define the shape of the strain profile: the distance over which the strain spreads outside of the hole, ρ_0 , and the coefficient of stress at the center of the hole, b_0 . However, these can be easily solved for analytically.

Throughout this section graphene is treated in the continuum limit as a thin plate: A solid which is much thinner in one dimension than the other two. Assuming that there is no shear, the stress-strain relationship for the thin plate is the same inside and outside the hole (Landau and Lifshitz, 1986)

$$\epsilon_r = \sigma_r - \nu\sigma_t \quad (3.2)$$

$$\epsilon_t = \sigma_t - \nu\sigma_r , \quad (3.3)$$

where ν is the Poisson Ratio and σ_r and σ_t are the dimensionless radial and tangential stresses given by the stress divided by the effective 3D Young's modulus ($E_{3D} = E_{2D}/t$

where t is the thickness).

Both the strain-displacement and the governing equations, however, depend on the region of interest. For the suspended material, the plate flexes down into the hole as pressure is applied. The resulting strain-displacement relationships are

$$\epsilon_r^i = \frac{dU}{d\rho} + \frac{1}{2} \left(\frac{dW}{d\rho} \right)^2 \quad (3.4)$$

$$\epsilon_t^i = \frac{U}{\rho}, \quad (3.5)$$

where U is the dimensionless radial deflection, W is the dimensionless out of plane deflection, and ρ is the dimensionless radius, all of which are made dimensionless by dividing by the radius of the hole. The governing equations inside the hole are calculated by balancing the forces on a radial element. The stresses and pressure acting on such an element are shown in Figure 3.9. Using the area of the radial element ($rdrd\theta$) to convert the pressure to a force, and the cross sectional area to convert the strains into forces results in the two governing equations

$$\sigma_t^i = \frac{d}{d\rho}(\rho\sigma_r^i) \quad (3.6)$$

$$\sigma_r^i \frac{dW}{d\rho} = -\rho \frac{q}{2}. \quad (3.7)$$

the first from radial equilibrium and the second from lateral equilibrium. These six equations can be combined to form a single differential equation for σ_r

$$\frac{1}{8}\rho q^2 + (\sigma_r^i)^2 \frac{d}{d\rho}[\sigma_r^i + \frac{d}{d\rho}(\rho\sigma_r^i)] = 0. \quad (3.8)$$

This is most easily found by using Equation 3.6 and then Equation 3.7 in the combination of Equations: [(3.4) \rightarrow (3.2)] $- \frac{d}{d\rho}\rho$ [(3.5) \rightarrow (3.3)]. Once this equation is used

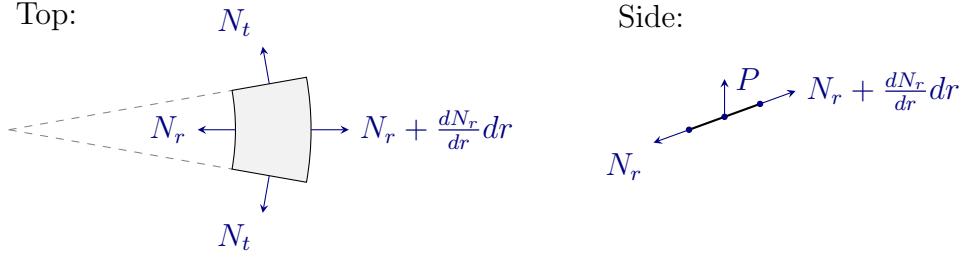


Figure 3.9: Force diagram of the suspended section ($\rho < 1$) of a pressurized graphene sealed microchamber. Depicted are the stresses, $N_r = E_{3D}\sigma_r$ and $N_t = E_{3D}\sigma_t$, and pressure, P , which act on a radial area element as viewed from the top and side.

to solve for σ_r , Equations 3.6, 3.2, and 3.3 allow the easy determination of σ_t , ϵ_r , and ϵ_t . Since there is no analytic solution to this differential equation, so, following Hencky's lead, a series expansion of the radial stress even in powers of ρ to match symmetry is assumed

$$\sigma_r^i = \frac{1}{4}q^{2/3} \sum_{n=0}^{\infty} b_{2n}\rho^{2n} . \quad (3.9)$$

When used in eqn. 3.8 all of the higher order coefficients are determined in terms of one free parameter, b_0 , the stress coefficient at the center of the hole. To this point the derivation has not deviated from Hencky's original work. However, instead of continuing forward to determine the value of b_0 by requiring that there is no radial deflection at the edge of the hole, b_0 will be left as a free parameter to match to the strain outside of the hole, which is now derive.

Outside of the hole the plate is constrained to move in the x-y plane, radial symmetry is preserved, and friction is present. The stress strain relationships (eqns. 3.2 and 3.3) are for a thin plate with no shear, and thus, apply equally well outside the hole as inside. The strain displacement relationships are based on the geometry of a general radially symmetric differential element. The only change that needs to

be made to these is the elimination of the out of plane motion,

$$\epsilon_r^o = \frac{dU}{d\rho} \quad (3.10)$$

$$\epsilon_t^o = \frac{U}{\rho}. \quad (3.11)$$

For sufficiently negative ϵ_t , the plate is expected to buckle out-of-plane, forming wrinkles. However, in the experimental regime there is no evidence for these effects. The forces acting on the plate are different inside and outside the hole, and thus, so are the governing equations. The stresses and forces acting on the graphene outside of the hole are shown in Figure 3.10. The interaction between the plate and its underlying substrate is included via the frictional force per unit area, f , which is pointed to oppose the radial stresses outside of the hole. If there were a component of friction acting against the tangential strain, it would go as $d\theta^2$ rendering it negligible. One power of $d\theta$ from the resultant and the other from the area element. Care needs to be taken in this mathematical treatment because this friction term does not turn off when the stress goes to zero. Instead, it works in positive feedback amplifying the stress. As a result, the stresses are only physical until they decay to zero at a position ρ_0 . The governing equation outside the hole is then

$$\sigma_t^o = \frac{d}{d\rho}(\rho\sigma_r^o) + F\rho. \quad (3.12)$$

Again these equations can be combined to form a single differential equation for σ_r . The easiest way to do this is by using Equation 3.12 twice in the combination of Equations: $[(3.10) \rightarrow (3.2)] - \frac{d}{d\rho}\rho[(3.11) \rightarrow (3.3)]$. The result is

$$\frac{d}{d\rho}[\sigma_r^o + \frac{d}{d\rho}(\rho\sigma_r^o)] = -(2 + \nu)F. \quad (3.13)$$

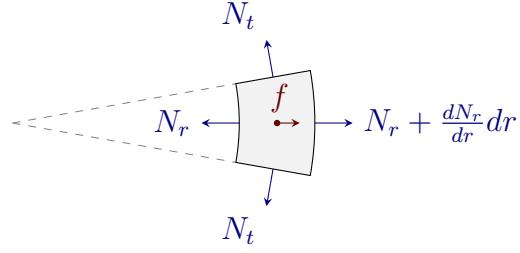


Figure 3.10: Force diagram of the supported section ($\rho > 1$) of a pressurized graphene sealed microchamber. Depicted are the stresses, $N_r = E_{3D}\sigma_r$ and $N_t = E_{3D}\sigma_t$, and friction, f , which act on a radial area element.

Unlike the situation for $\rho < 1$, this equation can be solved exactly. The general result for the radial and tangential stresses are

$$\sigma_r^o = (2 + \nu)F\left(-\frac{\rho}{3} - \frac{c_2}{2\rho^2}\right) + c_1 \quad (3.14)$$

$$\sigma_t^o = (2 + \nu)F\left(-\rho\frac{1 + 2\nu}{3(2 + \nu)} + \frac{c_2}{2\rho^2}\right) + c_1 , \quad (3.15)$$

where the tangential stress was solved for using Equation 3.12. To reduce the number of free coefficients, the radial and tangential strains are both forced to come to zero at the position ρ_0 . This natural restriction imposes a limit on the region where the theory gives physical results. For $\rho > \rho_0$, the mathematical friction persists even though there is no stress, yielding unphysical results. For these regions, the stresses are taken to be zero. Applying this restriction gives $c_2 = \rho_0^3 \frac{\nu-1}{3(2+\nu)}$ and $c_1 = \rho_0 \frac{1+\nu}{2} F$. This defines in closed form the strains for $\rho > 1$ as a function of one free parameter, ρ_0 .

Next, the boundary condition for the stresses at the edge of the hole will be determined so that the stresses can be matched across the boundary. The stress is

related to the force per unit area through a divergence operation

$$F_i = \frac{\partial N_{ik}}{\partial x_k} , \quad (3.16)$$

where F_i is the applied force per unit volume in the i direction, N_{ik} is the strain tensor, and repeated indices are summed over. Integrating the equations and using the divergence theorem gives

$$\int_V F_i \, dV = \oint_S \sigma_{ik} \, df_k . \quad (3.17)$$

Hence, the total force acting on a volume entity is given by the surface integral of the stress. A symmeterized volume element that spans the edge of the hole is shown in Figure 3.11. In the limit that the width of the volume element, ϵ , goes to zero the contributions from the surface forces, f and P , also go to zero and so must the surface integral of the stresses. This would not be the case if these were 1D edge forces. The undrawn tangential stresses also go to zero since the cross section that they act on, ϵt where t is the thickness, goes to zero as well. Thus, the remaining stresses must be equal,

$$N_r^i = N_r^o . \quad (3.18)$$

The boundary condition on the tangential stress is found by the requirement that the radial displacement, U , must be continuous. If it were not, the material on one side of the discontinuity would separate from the material on the other side leaving a gap. If U is continuous, so must ϵ_t be by eqns. 3.5 and 3.11. Finally, if ϵ_t and N_r are continuous, so must N_t be by equation 3.3. This proves that the radial and tangential stresses must be continuous over the edge of the hole.

All of the necessary relationships have now been derived. The series solution for

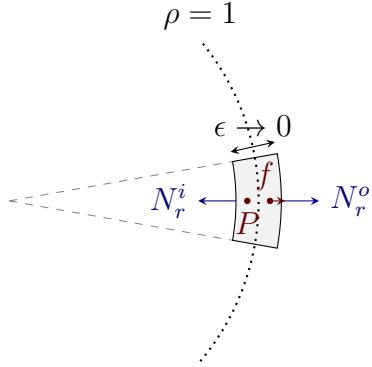


Figure 3.11: The boundary condition at the edge of the microchamber is determined by the stresses and forces acting on a thin, symmetrized volume element which spans the edge of the hole. The tangential strains inside and outside are ignored because their contribution goes to zero as the width of the element, ϵ goes to zero. P is the pressure which acts out of plane.

the stresses for $\rho < 1$ as a function of the dimensionless loading, q , are known with b_0 , the stress at the center, a free parameter. Similarly, the closed form solution for the stresses for $\rho > 1$ as a function of the dimensionless friction, F , have been found, with ρ_0 , the furthest distance the stress spreads, as a free parameter. The derived boundary conditions at $\rho = 1$ can now be applied

$$\begin{aligned}\sigma_r^i(\rho = 1) &= \sigma_r^o(\rho = 1) \\ \sigma_t^i(\rho = 1) &= \sigma_t^o(\rho = 1).\end{aligned}\quad (3.19)$$

This sets the value of ρ_0 and b_0 in terms of $q^{2/3}/F$. Regrettably, the series form of the solution inside of the hole makes the presentation of an exact expression for ρ_0 and b_0 impossible. Numerical solutions for the values of ρ_0 and b_0 are presented in Figures 3.12 for graphite's Poisson ratio of 0.165 (Blakslee et al., 1970). As the

friction increases or the dimensionless loading coefficient decreases, $q^{2/3}/F$ decreases, and ρ_0 decreases toward unity, indicating that no strain distribution outside of the hole. At the same time b_0 increases toward the limiting value of 1.66, the result for the original Hencky model which assumes perfectly fixed boundary conditions at the edge of the hole.

This completes the derivation of the extended Hencky model. The shape of the strain distributions shown in Figure 3.8 has been determined along with the pressure, friction, radius, and layer thickness dependencies shown in Figure 3.12.

3.4.2 Detailed description of the Atomistic model

In this section the atomistic model used to verify the extended Hencky continuum model is discussed. Classical molecular dynamics simulations (MD) were performed using open source simulation package LAMMPS (Lammps, 2012; Plimpton, 1995) developed at Sandia National Labs. The model used argon gas to compress a graphene monolayer that lies atop an amorphous silicon dioxide substrate, as illustrated in Figure 3.13. The substrate had dimensions 46 x 46 x 3 nm with a 6 nm radius hole in the center. A circular graphene monolayer with radius 22 nm was placed on top of the hole in the substrate and argon atoms were randomly distributed within the simulation box after relaxation of the graphene-substrate system. There were 562,110 atoms in total, with the simulation run in parallel for maximum computational efficiency.

The covalent carbon bonds were modeled using the AIREBO (Stuart et al., 2000) potential, which has been shown to have good accuracy in describing hydrocarbon systems (Qi et al., 2010; Zhao and Aluru, 2010). The Tersoff (Tersoff, 1988) potential was utilized for the Si-Si, Si-O and O-O interactions, while a Lenard-Jones potential was used for all other interactions with a cutoff distance of 8 Å and the corresponding interaction parameters chosen as follows: $\epsilon_{Ar-Ar}=0.0104$ eV, $\sigma_{Ar-Ar}=3.405$

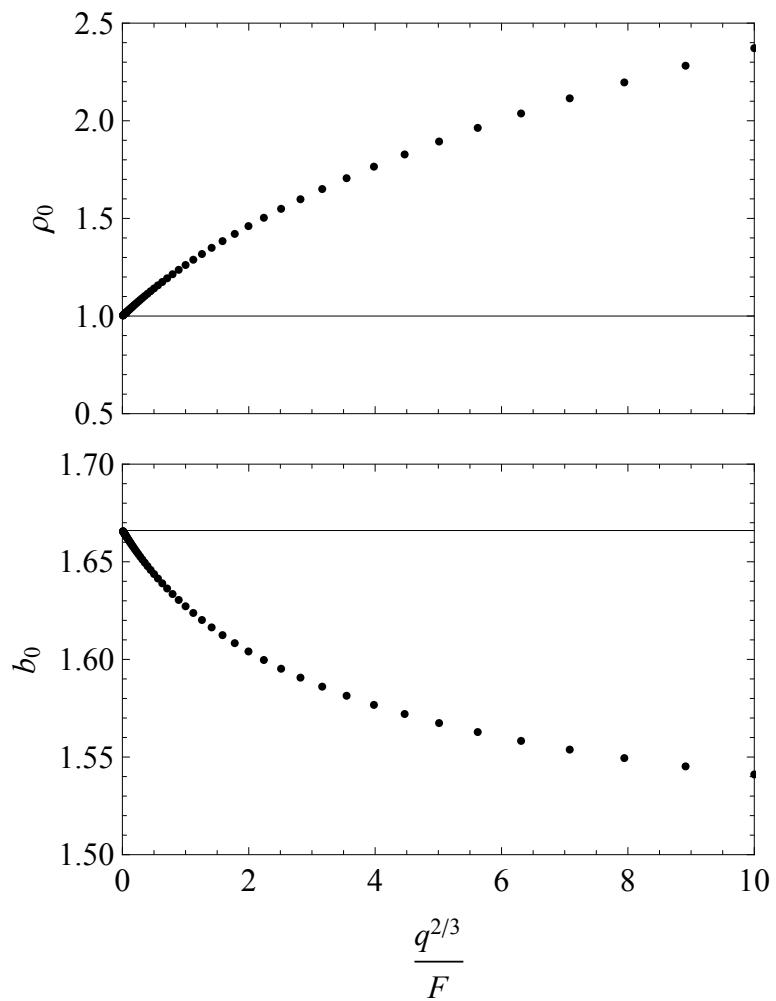


Figure 3.12: Numerical solutions for the values of ρ_0 and b_0 , the furthest normalized radii that the strain reaches and the coefficient for the strain at the center of the hole, respectively. Note the large friction convergence to the values of the original Hencky model, $\rho_0 = 1$ and $b_0 = 1.66$.

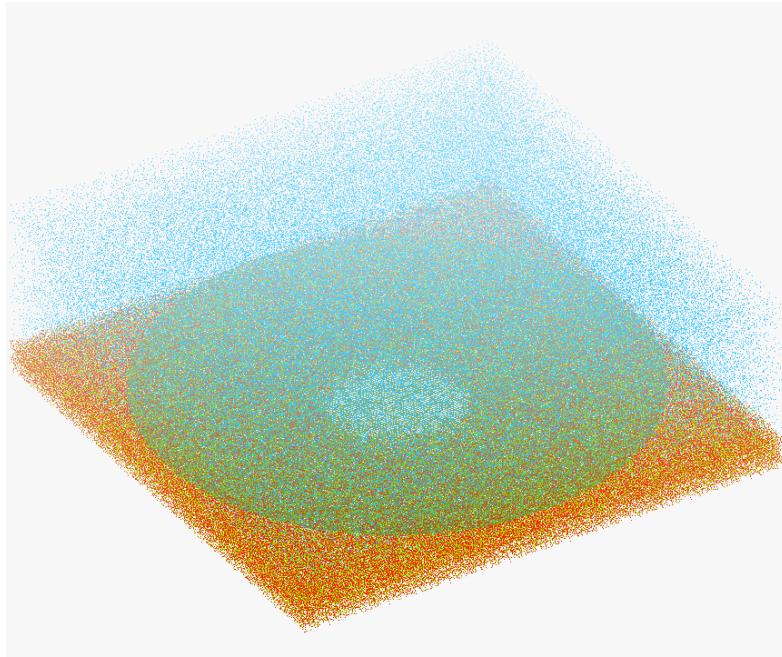


Figure 3.13: Schematic diagram of the atomistic simulation of the pressurized graphene sealed microchamber including (from bottom to top) amorphous SiO₂ substrate (red and orange atoms), circular graphene monolayer (brown atoms) and argon gas (cyan atoms). Visualization was performed using Visual Molecular Dynamics (Humphrey et al., 1996).

\AA (Rytkonen et al., 1998); $\epsilon_{Ar-C}=0.0123 \text{ eV}$, $\sigma_{Ar-C}=3.573 \text{ \AA}$ (Tuzun et al., 1996); $\epsilon_{Ar-Si}=0.0028 \text{ eV}$, $\sigma_{Ar-Si}=3.778 \text{ \AA}$ (Li et al., 2010a); $\epsilon_{Ar-O}=0.0058 \text{ eV}$, $\sigma_{Ar-O}=3.3075 \text{ \AA}$ (Everitt and Skinner, 1999); $\epsilon_{Si-C}=0.008909 \text{ eV}$, $\sigma_{Si-C}=3.326 \text{ \AA}$ (Ong and Pop, 2010); $\epsilon_{O-C}=0.003442 \text{ eV}$, $\sigma_{O-C}=3.001 \text{ \AA}$ (Ong and Pop, 2010). The Ar-Si and Ar-O interaction parameters were obtained using the standard Lorentz-Berthelot mixing rule.

The system was first relaxed at room temperature (300K) with a pressure of .03 MPa, at which point the gas density was slowly increased by reducing the volume of gas to increase the pressure. When the desired deflection of graphene was reached,

the gas density was then kept fixed, and the system was allowed to relax at that pressure for 10 ps. The final equilibrium configuration of graphene was computed by averaging over all of the atomic positions of graphene during the relaxation period with constant gas density. The substrate was assumed to be rigid and fixed during the simulation. As a boundary condition on the graphene, the atoms in the outer 1 nm annulus of the supported graphene were fixed during the simulation matching the experimental observation that the supported graphene only slides in a local region around the hole, remaining fixed at large radii. A canonical ensemble (NVT) at room temperature (300K) was used for the entire simulation and reflection boundary conditions were used to ensure that gas atoms remained inside the simulation box.

To calculate the strain in the deformed graphene sheet, the atomistic displacement field was first found (Zimmerman et al., 2009) by calculating the difference between the reference (initial) configuration and the final, deformed configuration. The strain tensor from continuum mechanics was then calculated as

$$\epsilon_{ij} = \frac{1}{2} \left(\frac{\partial u_i}{\partial X_j} + \frac{\partial u_j}{\partial X_i} + \frac{\partial u_k \partial u_k}{\partial X_i \partial X_j} \right), \quad (3.20)$$

where ϵ_{ij} are the components of the strain, u is the displacement, and X denotes the position of a point in the reference configuration.

The displacement field \mathbf{U} for each atom was computed by linearly interpolating, using standard finite element method shape functions for triangular elements (Hughes, 1987). The displacements of its three nearest neighbors is calculated as $\mathbf{U} = \mathbf{N} \bullet \mathbf{u}_N$, where \mathbf{N} are the finite element shape functions and \mathbf{u}_N are the displacements for each atom. From the displacement field \mathbf{U} , the strain was calculated by taking derivative of \mathbf{U} as $\epsilon = \mathbf{B} \bullet \mathbf{u}_N$, where $\mathbf{B} = \frac{\partial \mathbf{N}}{\partial \mathbf{X}}$. The atomistic data plotted in the paper was obtained from the MD simulation results by the method above, and was expressed in

polar coordinates.

3.5 Fitting Raman spectra to the continuum model

In conjunction with the measured Raman spectra, the extended continuum model can be used to determine the sliding friction, the Grüneisen parameter, and the shear deformation potential. A simple comparison of the strains predicted by the extended Hencky model to strains found by directly inverting the positions of the G^+ and G^- peaks is not possible because of the finite size of the focused laser beam. Instead, spectra are predicted by integrating against the system point spread function. A set of spectra predicted for each point in the line scan can then be compared to the measured line scan spectra. The best fit is found by predicting spectra for a series of different fitting parameters and choosing the one that best matches the measured spectra. This is done quantitatively by choosing the spectra with the smallest χ^2 . A detailed description of the fitting algorithm used as well as the determination of fitting error is presented in Appendix C.

Unlike the sliding friction, the Grüneisen parameter and shear deformation potential should be the same for every line scan. As such, they were included as fitting parameters in only the two lines scans which best defined β based on the splitting of the supported graphene's G band just outside the edge of the microchamber: the $\sim 5\mu m$ radius monolayer and the $\sim 5\mu m$ radius trilayer at 0.80 MPa of applied pressure. Figure 3.14 shows reduced dimensionality plots of the χ^2 per degree of freedom space. As described in Appendix C, each data point in this type of plot represents the best case for the given value of the plotted fitting parameters. The other, not plotted, fitting parameters (β and F for the γ plot, and γ and F for the β plot) are chosen to minimize χ^2 at each data point. The sharp minimum and the agreement between

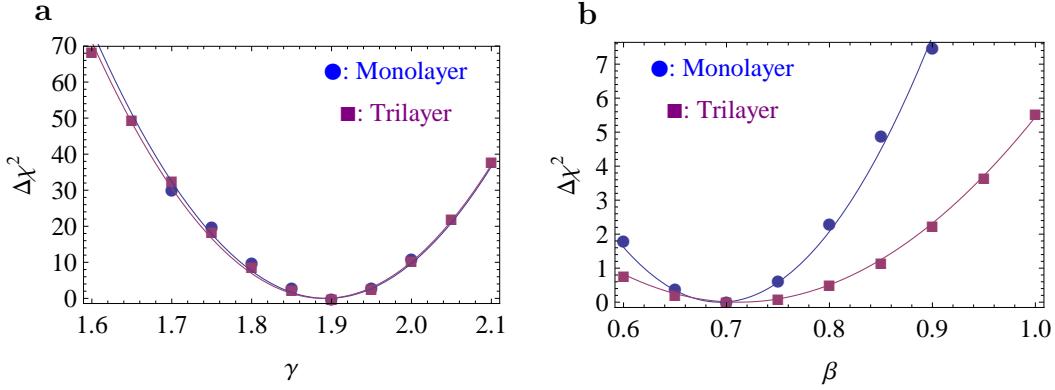


Figure 3.14: Determination of the Grüneisen parameter and shear deformation potential based on global fits to a monolayer covered (blue, circles) and a trilayer covered (purple, squares) $\sim 5 \mu m$ radius sealed microchamber at 0.80 MPa of applied absolute pressure. All three fitting parameters (F , γ , and β) were included in the fits. In (a) and (b) the reduced dimensionality plots of the deviation of χ^2 per degree of freedom from the minimum are plotted for the Grüneisen parameter and shear deformation potential respectively. The parabolic fits show sharp minimum.

monolayer and trilayer indicates a good determination of γ and β . Using an increase in χ^2 per degree of freedom of 0.25 to define uncertainties gives $\gamma_{mono} = 1.89 \pm .02$, $\gamma_{tri} = 1.89 \pm .02$, $\beta_{mono} = 0.69 \pm .04$, and $\beta_{tri} = 0.71 \pm .06$. The averages of the monolayer and the trilayer, $\gamma = 1.89 \pm 0.01$ and $\beta = 0.70 \pm 0.04$, were treated as known material parameters when fitting the remaining twenty Raman line scans.

Table 3.1 compares the measurements of the Grüneisen parameter and shear deformation potentials to previous measurements. The substrate on which the parameters were measured in also included. Our measured γ is commensurate with most of the other measured values and agrees particularly well with the *ab initio* calculations of Cheng *et al.* (Cheng et al., 2011). On the other hand, the measured shear de-

	γ	β
This work	1.89	.70
SiO ₂ depression (Metzger et al., 2010)	2.4	
On PDMS (Huang et al., 2009)	.69	.38
On SU8 (Mohiuddin et al., 2009)	1.99	.99
Embedded (Frank et al., 2010)	2.01	1.01
On Acrylic (Yoon et al., 2011)	2.2	.93
Bubble (Zabel et al., 2012)	1.8	
<i>Ab initio</i> (Thomsen et al., 2002)	2.0	0.66
<i>Ab initio</i> (Cheng et al., 2011)	1.86	.96

Table 3.1: Summary of the Grüneisen parameter, γ , and shear deformation potential, β , as measured on different substrates.

formation potential is lower than most other measurements. Buckling out-of-plane cannot explain this result since the mono and trilayers would buckle differently for a microchamber with the same pressure and radius. These are the first measurements of γ and β for which the sliding of FLG over its substrate was included.

The extended Hencky model agrees extremely well with the measured spectra. Figure 3.15 shows a global data fit for a $\sim 5 \mu\text{m}$ radius monolayer covered graphene sealed microchamber at 0.80 MPa of applied pressure. The spectra and fits from each position along the line scan are stacked vertically in the direction of the line scan. The extended continuum model successfully fits the softening and splitting of the G band of the supported graphene and successfully predicts the downshift and sharpening of the G band of the suspended graphene as the center of the microchamber is approached. In comparison, without the theoretical extension the standard Hencky model would fail to reproduce the supported graphene spectra.

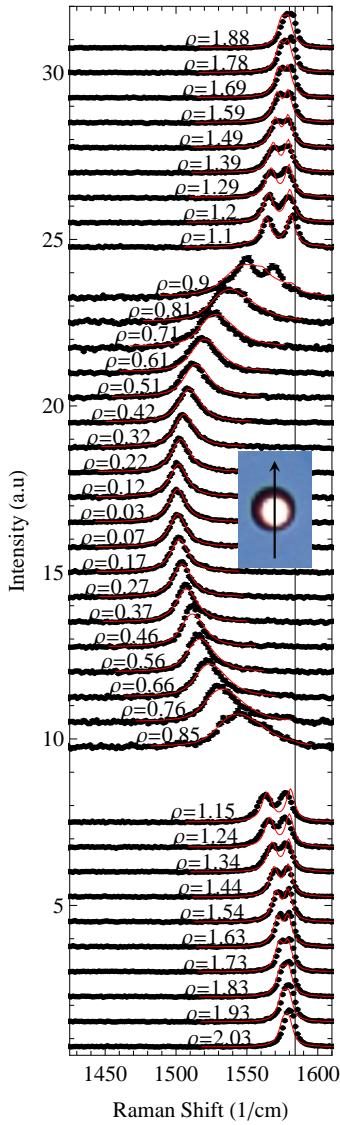


Figure 3.15: Raman spectra from a line scan over a $\sim 5 \mu\text{m}$ radius monolayer graphene sealed microchamber with 0.80 MPa of applied pressure analyzed with the extended Hencky model (red lines). The spectra taken along the path shown in the inset are arrayed vertically with spectra taken too close to the edge of the microchamber omitted (see text). The black vertical line is positioned at the supported graphene's unstrained G band energy.

3.6 Measured frictional dependencies

The sliding friction extracted for eight microchambers with radii between 1.2 and 5 μm and with applied absolute pressures from 0.10 to 0.80 MPa exhibits fundamentally different behavior for trilayer graphene than for monolayer and bilayer. In Figure 3.16a (left) the friction is plotted as a function of absolute applied pressure. The data for trilayer graphene (black dots) shows a linearly dependent sliding friction in accordance with Amontons' law with a coefficient of friction of 0.11 ± 0.01 . For reference, Teflon on Teflon has $\mu = 0.04$ while clean steel on clean steel has $\mu = 0.6$ (Resnick et al., 2002). A zoom in of the trilayer data is shown in Figure 3.17. The sliding friction for monolayer and bilayer graphene behave much differently. They decrease generally with applied pressure and the wide scatter of the points for different radii and layer number clearly indicate that the sliding friction is dependent on the geometry of the microchamber. Our theoretical analysis shows that the radial strain at the edge due to the pressure pushing the graphene into the microchamber has the same radii and layer number dependence as the friction, so the sliding friction is replotted as a function of the radial strain in Figure 3.16b (right). The monolayer and bilayer data for all different radii microchambers now *collapse* to a single curve versus radial strain, well described by $1/\epsilon_{r,\text{edge}}$ behavior (dashed line). The best fit to the strain dependence is $0.002\text{MPa}/\epsilon_{r,\text{edge}}$.

The gross difference in behavior between trilayer on one hand, and mono-and bilayer on the other, illustrates the two roles of the applied pressure. The pressure load pushes the graphene more firmly onto the substrate so that sliding friction should increase (Amontons' law), yielding a positive coefficient of friction. This is the case for the trilayer graphene. On the other hand, as the pressure pushes the graphene into the microchamber, it creates a radial tension in the supported graphene outside

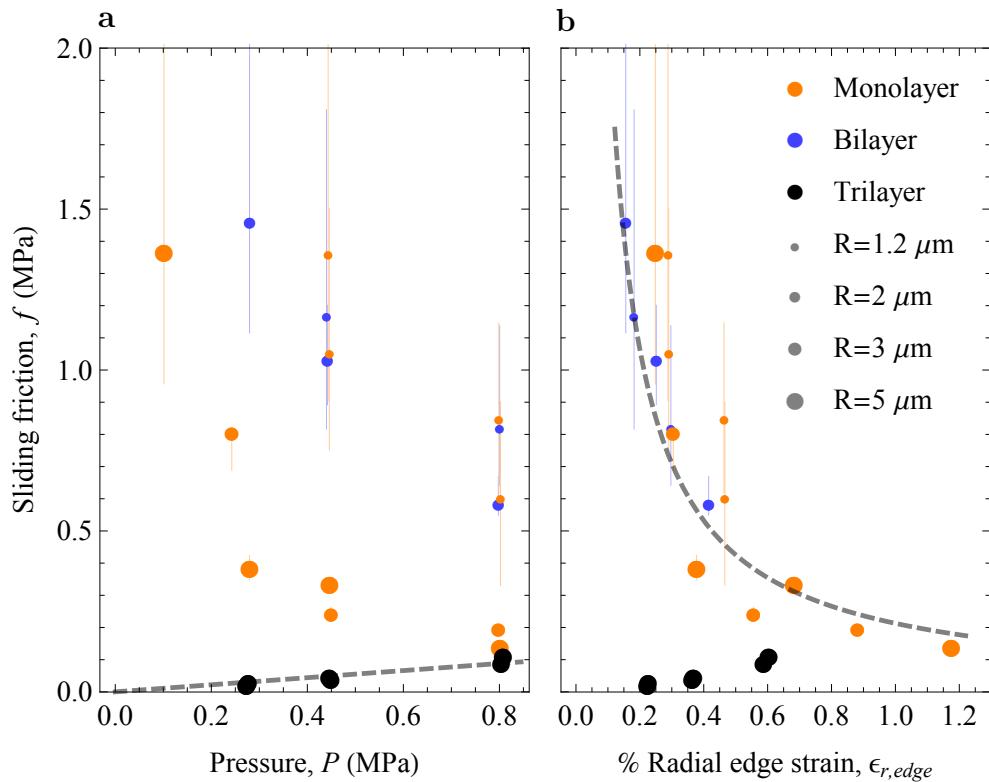


Figure 3.16: The dependencies of sliding friction for FLG extracted by analyzing Raman line scans with the extended Hencky model. In panel (a) (left) friction is plotted as a function of absolute applied pressure where in panel (b) (left) it is plotted as a function of the radial strain at the edge of the microchamber. The size of each data point represents the radius of the FLG-sealed microchamber corresponding to that point. The error bars are given by an increase in global fit χ^2 per degree of freedom of 0.25. The sliding friction of trilayer graphene depends linearly on the absolute applied pressure in agreement with Amontons' law. The sliding friction of monolayer and bilayer graphene, however, go as the inverse of the radial strain at the edge of the microchamber. Gray dashed lines are the best fits to these trends.

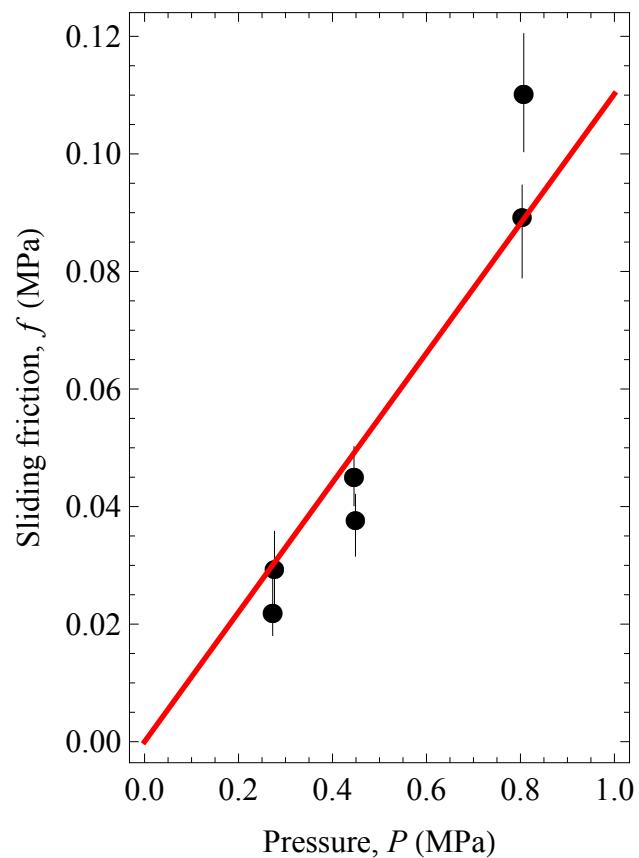


Figure 3.17: The sliding friction for trilayer graphene as a function of the absolute applied pressure extracted from Raman lines scans over two $5\ \mu\text{m}$ radii trilayer sealed microchambers. The data is fit with Amonton's law, $f = \mu P_{abs}$, giving a coefficient of friction between trilayer graphene and SiO_2 of $\mu = 0.11 \pm 0.01$.

the microchamber. The data collapse in Figure 3.16b for monolayer and bilayer graphene demonstrate that the pressure dependence of the sliding friction is not due to loading the supported graphene but is instead due to the graphene being pulled and stretched by the applied pressure. This is the only mechanism that would depend on the geometric parameters of the microchamber while also being consistent with the data. It is not surprising that the sliding friction for mono- and bilayer graphene is dependent on the strain and not the load because thin graphene conforms nearly perfectly to a SiO₂ substrate (Stolyarova et al., 2007; Lui et al., 2009; Cullen et al., 2010). Increasing the load cannot further increase the contact area, but increasing the radial strain beyond the edge of the microchamber may act to smooth out the graphene sheet, decreasing the contact between the graphene and the substrate, and thus decreasing the sliding friction. The bending rigidity, which goes as thickness cubed, of trilayer graphene must be high enough to counteract the adhesion energy, causing lower conformation and allowing for a traditional pressure and load response. The existence of a bilayer to trilayer crossover in FLG-substrate interactions is also observed in GPa range pressure measurements of silicon dioxide supported graphene (Proctor et al., 2009; Nicolle et al., 2011). Nicolle *et al.* observed a decrease in the pressure response of the G band between bilayer and trilayer graphene which was attributed to the transition from biaxial compression mediated by the substrate to hydrostatic compression mediated by the pressure transmitting medium (Nicolle et al., 2011).

3.7 Summary

In summary, it has been demonstrated that few layer graphene slides along the substrate when pulled. Furthermore, using a newly developed extension of the continuum

Hencky model, the sliding friction as a function of the number of atomic layers and the load was extracted. Trilayer graphene shows a typical load response whereas the sliding friction for monolayer and bilayer graphene goes as the inverse of strain. The data collapse of the friction for mono- and bilayer graphene when plotted versus strain is strong experimental evidence for a reduction in surface conformation when graphene is pulled as the fundamental origin of the negative coefficient of friction. These results will be important for the design of strain engineered devices (Pereira and Castro Neto, 2009), while the sliding of a flexible surface along a bulk object should be of fundamental, tribological interest. Finally, the method used in generalizing Hencky's solution should be useful for including distributed strains in other continuum models for use in designing strain-engineered graphene devices and in understanding other, few-layer material systems.

Chapter 4

The thermal conductivity of strained graphene

Graphene's very high thermal conductivity has been linked to out of plane acoustic (ZA) phonons. As a result of their quadratic dispersion, these phonons are believed to have two conflicting contributions to the thermal conductivity. Strain tunes the character of these phonons providing a method to study their influence on thermal transport.

Lindsay and coworkers have argued that the ZA phonons contribute more than 70% of the very high thermal conductivity of suspended graphene (Seol et al., 2010). The quadratic dispersion gives a high density of states throughout the BZ and the in plane reflection symmetry provides a selection rule that limits the scattering phase space (Lindsay et al., 2010). ZA phonon dominated thermal transport is consistent with observations that the thermal conductivity of supported graphene (Seol et al., 2010; Cai et al., 2010) is a factor of four smaller than suspended graphene (Balandin et al., 2008; Cai et al., 2010; Chen et al., 2011; Lee et al., 2011) and that encased graphene is a factor of two smaller yet (Jang et al., 2010). The ZA phonons either leak into the surrounding media, are scattered by it, or are damped by it.

On the other hand the ZA phonons are believed to limit the potentially divergent thermal conductivity of the in plane phonons. In the absence of ZA phonons Klemens showed that the linear in plane acoustic phonons should provide a logarithmically

divergent thermal conductivity in two dimensions. This divergence follows from the frequency dependence of the terms which make up the thermal conductivity, $\kappa = \frac{1}{3} \int Cv_G l(\omega) \rho(\omega) d\omega$, where κ is the thermal conductivity, C is the specific heat, v_G is the group velocity of the acoustic phonons, $l(\omega)$ is the scattering length, and $\rho(\omega)$ is the density of states. For a two dimensional phonon gas with linear dispersion $\rho(\omega) \propto \omega$ while for anharmonic scattering between linear acoustic phonons $l(\omega) \propto 1/\omega^2$. Hence, in the absence of extrinsic scatterers, the integrand scales as $1/\omega$ and κ diverges logarithmically. The logarithmic divergence in ω should translate to a logarithmic divergence in device size (Klemens, 2001). To date this has not been observed (Chen et al., 2011) and Pereira and coworkers believe that the ZA phonons are the cause. Very close to the center of the BZ there is a divergent population of zero group velocity ZA phonons which can only contribute to the thermal conductivity as scatterers. It is predicted that they alter the frequency dependence of $l(\omega)$ and remove the divergent contribution of the in plane phonons (Pereira and Donadio, 2013). Thus, the ZA phonons are the main contributor to graphene's thermal conductivity only because the low energy ZA phonons suppress the divergent thermal conductivity of the in plane phonons.

Here we use strain to tune the thermal conductivity of

4.1 Past measurements of graphene's thermal conductivity

The first measurement of graphene's thermal conductivity was on suspended graphene and was reported to have the highest thermal conductivity ever measured (Balandin et al., 2008). Shortly thereafter a more modest value with a thermal conductivity roughly one six of the original value was measured for supported graphene (Seol et al., 2010). This discrepancy has been the source of heavy study in the years since.

The original measurements were performed on suspended graphene samples using an optical technique. This creative technique uses a laser excitation to heat the sample and the temperature induced energy shift of the Raman scattered light to measure temperature. In this way a simple Raman measurement serves a dual role as both a heat source and a thermometer. When used in conjunction with a heat transfer model, this is enough to determine the thermal conductivity. This technique has advantages and disadvantages. Uncertainties in the optical absorption and the laser spot size effect the measured thermal conductivity, good results depend on an accurate thermal transport model, and these models are limited by their inability to include ballistic thermal transport. Nonetheless, the required samples are very simple and the measurement is fairly standard. This simplicity has resulted in a variety of publications using this technique to study different aspects of thermal transport in suspended graphene (Balandin et al., 2008; Faugeras et al., 2010; Cai et al., 2010; Ghosh et al., 2010; Lee et al., 2011; Chen et al., 2011; Chen et al., 2012).

The measurement which yielded the more modest values of thermal conductivity were direct measurements of supported graphene. In these more traditional measurements resistive heaters are used to create a temperature gradient and solid state thermometers are used to measure the temperature. The thermal conductivity of graphene is then determined directly from the thermal resistance of the graphene. Although certainly less error prone, there has been only one other measurement of the thermal conductivity of graphene using this technique (Jang et al., 2010). This is a result of the more difficult measurement and much more difficult sample fabrication.

Figure 4.1 summarizes all of the reported room temperature thermal conductivity measurements of single layer graphene. The thermal conductivity decrease from a value of roughly 2000 W/m-K for suspended graphene to a value close to 500 W/m-K for graphene supported on one side, to a value of less than 160 W/m-K for graphene

encased in SiO_2 (Jang et al., 2010). The general observation is that graphene's thermal conductivity is negatively effected by the presence of a supporting or encasing bulk material. This difference in thermal conductivity is not an artifact of the measurement techniques, a nanoscale thermal conductivity measurement also observed an increase in thermal conductivity for suspended few layer graphene (Pumarol et al., 2012).

The Raman measurement technique has been advanced in a series of publications. Faugeras et al was the first to use graphene suspended over a circular aperture. This circular symmetry of this geometry simplifies the heat transfer model. Additionally, they used the Raman Stokes to anti-Stokes ratio to measure the temperature and validated the previously used temperature induced phonon energy shift (Faugeras et al., 2010). Cai et al advanced the thermal transport model by including the thermal transport to the substrate instead of treating supported graphene as a heat sync. Additionally, they brought up concerns about ballistic thermal transport (Cai et al., 2010). Chen et al showed that if not accounted for, the thermal conductivity to the gas could inflate the measured value of graphene's thermal conductivity (Chen et al., 2011).

A variety of potentially interesting physics surrounding the thermal transport of few layer graphene has been studied. The thickness dependence of the thermal transport has also generated some interest. Supported graphene's thermal conductivity was shown to converge upward toward graphite's thermal conductivity as the number of layers increased above 34 layers (Sadeghi et al., 2013). On the other hand the thermal conductivity of suspended graphene was shown to converge downward toward graphite's thermal conductivity as the number of layers increased to four layers (Ghosh et al., 2010). Ballistic phonon transport has also been of interest. Since phonons with longer mean free path are able to thermalize and participate in thermal transport for larger samples a size dependence has been expected and observed for

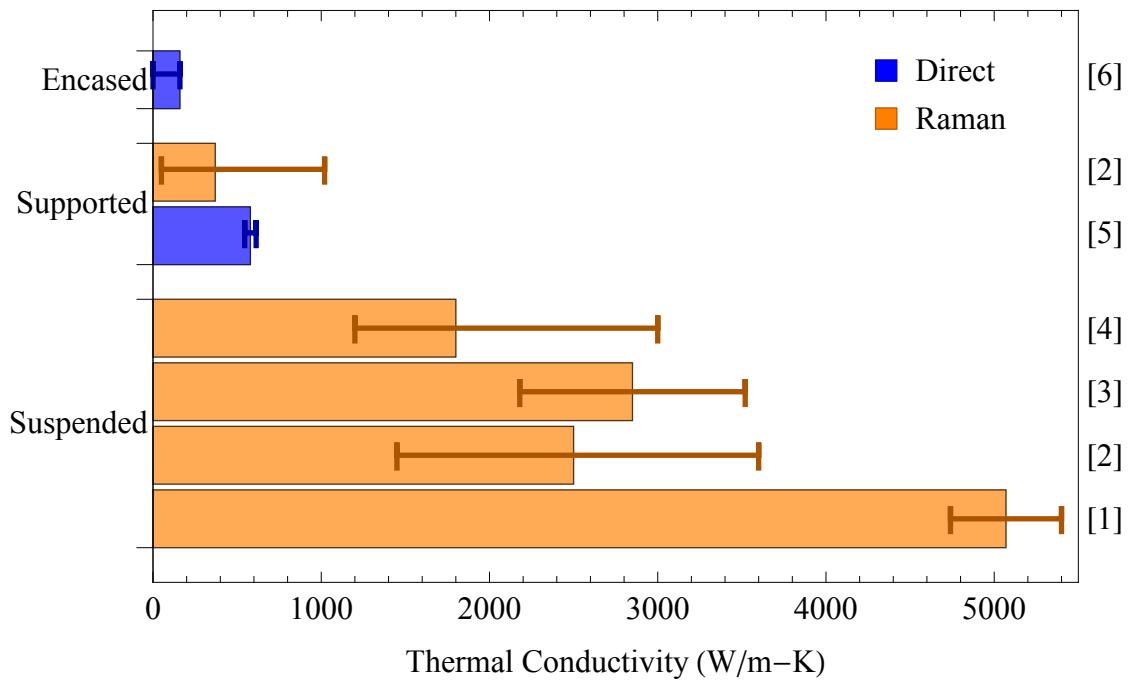


Figure 4.1: Summary of the reported values of the room temperature thermal conductivity of monolayer graphene grouped by the environment of the graphene: Suspended, supported on a bulk substrate, and encased in amorphous SiO_2 above and below. The color code indicates whether the measurements were performed using the Raman technique or a direct technique. The number on the right indicates the source: [1] is (Balandin et al., 2008), [2] is (Cai et al., 2010), [3] is (Chen et al., 2011), [4] is (Lee et al., 2011), [5] is (Seol et al., 2010), [6] is (Jang et al., 2010). Error bars are taken from the literature.

trilayer graphene (Wang et al., 2011) but not for monolayer CVD graphene (Chen et al., 2011). Other studies have concentrated on the effects of scattering on the thermal properties of graphene (Wang et al., 2011; Pettes et al., 2011; Jang et al., 2013; Chen et al., 2012).

4.2 Heat transport model

An accurate heat transport model is necessary for the determination of the thermal conductivity using the Raman measurement technique. In this section the heat transfer model first proposed by Cai et al is summarized. It predicts the temperature profile which develops in a circular graphene sealed microchamber due to central laser heating (Cai et al., 2010). The temperature profile is then used to predict the anti-Stokes to Stokes ratio following Herman et al (Herman, 2011). Finally, by predicting how the anti-Stokes to Stokes ratio depends on unknown parameters, the importance of the various experimental unknowns is investigated.

The heat transfer model is simplified by the concentric circular symmetry of the laser heating source and the circular microchamber. The solution is determined by solving the steady state heat equation,

$$\vec{\nabla} \cdot (\kappa(r) \vec{\nabla} T(r)) + \dot{Q} = 0 ,$$

where $T(r)$ is the temperature distribution, κ is the thermal conductivity, and \dot{Q} is the generated heat flux per unit time. The general solution is found by matching the solution for inside, $r < R$, and outside, $r > R$, the edge of the microchamber with boundary conditions at $r = R$.

Isolating the suspended and supported graphene, the generated heat flux includes the laser heating, the energy lost to the gas, and the energy lost to the supporting

substrate. The laser heating is described by,

$$\dot{Q}_L = \frac{\alpha P}{t} \frac{1}{2\pi\sigma^2} e^{-r^2/\sigma^2}, \quad (4.1)$$

where \dot{Q}_L is the energy from the incident laser, P is the laser power at the graphene surface, α is the factor of this power which is absorbed by the graphene, t is the thickness of the graphene, σ is half the beam waste measured to be $0.81 \mu\text{m}$ as shown in Section 3.1, and $\frac{1}{2\pi\sigma^2}$ is the normalization factor. The energy lost to the substrate and gas is modeled using Newton's law of cooling,

$$\begin{aligned}\dot{Q}_S &= -g_S \frac{1}{t} (T(r) - T_0) \\ \dot{Q}_g &= -g_g \frac{1}{t} (T(r) - T_0),\end{aligned}$$

where \dot{Q}_S and \dot{Q}_g are the energy lost to the substrate and to the gas respectively, g_S and g_g are the heat transfer coefficients per unit area, and T_0 is ambient temperature. In the case of the suspended graphene, heat can be lost to both the gas above and below the graphene. This will be differentiated using g_g^\uparrow and g_g^\downarrow .

When applying Newton's law of cooling for the heat conduction to the silicon dioxide we implicitly assume that the temperature of the silicon dioxide remains unchanged. To test this assumption the thermal oxide is treated as a cylindrical thermal wire connected to a heat sync on one side and with 0.03 mW of power (matching a 1.5 mW laser with 2.3 % absorption) heating the other side. The cylinders width was taken to be 200 nm to match the width of the supported graphene that is expected to have elevated temperatures. An inner radius of $5 \mu\text{m}$, a thickness of 300 nm, and a thermal conductivity of 1.4 W/m-K () gives a temperature rise of 1.5 K in the thermal oxide. In the physical system the transport in the silicon dioxide would flare

outward resulting in a larger conduction cross section. Thus, this 1.5 K temperature rise should be treated as an upper bound. This validates the assumption that the thermal oxide is not significantly heated in our measurements.

Assuming a piecewise uniform thermal conductivity,

$$\kappa(\vec{r}) = \begin{cases} \kappa_{SS} & , r < R \\ \kappa_{SP} & , r \geq R \end{cases}$$

eliminates the spatial dependence of κ in the two regions, simplifying the heat equation in Equation 4.2 while still allowing for the disparate suspended and supported thermal conductivities observed in the literature. While needed to simplify the problem, this simple form for the thermal conductivity may not be entirely accurate. A strain dependent thermal conductivity would inherit the spatial dependence of the strain field. The measurement described here can not provide the spatial information which would inform this more detailed model. Instead the discussion will be limited to the effective suspended thermal conductivity of the suspended graphene κ_{SS} . This factor is a metric which represents the departure of the system from the unstrained case.

Including the generated heat fluxes and the piecewise thermal conductivity in the heat equation, Equation 4.2, results in two differential equations,

$$\begin{cases} \kappa_{SS} \nabla^2 T + \dot{Q}_L + \dot{Q}_g^\uparrow + \dot{Q}_g^\downarrow = 0 & , \text{for } r < R \\ \kappa_{SP} \nabla^2 T + \dot{Q}_S + \dot{Q}_g^\uparrow = 0 & , \text{for } r \geq R \end{cases}$$

In the dimensionless variables

$$\theta = \frac{1}{T_a} (T(r) - T_a)$$

$$\rho = \frac{r}{\sqrt{2}\sigma}$$

the differential equations become

$$\begin{cases} \nabla^2 \theta + \beta e^{-\rho^2} - \gamma \theta = 0 & , \text{ for } \rho < \frac{R}{\sqrt{2}\sigma} \\ \nabla^2 \theta - \Gamma \theta = 0 & , \text{ for } \rho \geq \frac{R}{\sqrt{2}\sigma} \end{cases},$$

where

$$\beta = \frac{1}{\pi T_0} \frac{\alpha P}{t \kappa_{SS}}$$

$$\gamma = 2\sigma^2 \frac{g_g^\uparrow + g_g^\downarrow}{t \kappa_{SS}}$$

$$\Gamma = 2\sigma^2 \frac{g_g^\uparrow + g_S}{t \kappa_{SP}}.$$

The boundary conditions,

$$\theta(\infty) = 0 \quad (4.2)$$

$$|\theta(0)| < \infty \quad (4.3)$$

$$\theta \left(\frac{R}{\sqrt{2}\sigma} \right)^- = \theta \left(\frac{R}{\sqrt{2}\sigma} \right)^+ \quad (4.4)$$

$$-\kappa_{SS} \frac{d}{d\rho} \theta^- \Big)_{r=\frac{R}{\sqrt{2}\sigma}} = -\kappa_{SP} \frac{d}{d\rho} \theta^+ \Big)_{r=\frac{R}{\sqrt{2}\sigma}}, \quad (4.5)$$

are written by assuming that the temperature profiles are finite, decay asymptotically to ambient temperature, are continuous at the edge of the microchamber, and have

continuous heat flux at the edge of the microchamber respectively. The heat flux, $\vec{\phi}$, is found using Fourier's law of heat conduction, $\vec{\phi} = -\kappa \vec{\nabla} T(\vec{r})$.

The heat equation for $r \geq R$ has a relatively simple solution. It can be recast in the form of the modified Bessel's equation and after applying the boundary condition in Equation 4.2 it has the solution,

$$\theta(\rho) = c_2 K_0(\sqrt{\gamma}\rho), \text{ for } \rho \geq \frac{R}{\sqrt{2}\sigma}, \quad (4.6)$$

where c_2 is a integration constant and $K_0(\sqrt{\gamma}\rho)$ is the modified Bessel function of the second kind.

The solution for the suspended graphene is complicated by the incident laser. Using the substitution $x = \sqrt{\gamma}\rho$ the heat equation can again be cast in the form of a modified Bessel function, but this time it is inhomogeneous

$$\rho^2 \frac{d^2\theta}{dx^2} + x \frac{d\theta}{dx} - x^2\theta = -\frac{\beta}{\gamma} x^2 e^{-x^2/\gamma}.$$

This differential equation has the general solution

$$\theta(\rho) = c_3 I_0(\sqrt{\gamma}\rho) + c_4 K_0(\sqrt{\gamma}\rho) + \theta_P(\sqrt{\gamma}\rho), \text{ for } \rho \geq \frac{R}{\sqrt{2}\sigma},$$

where $I_0(\sqrt{\gamma}\rho)$ is the modified Bessel function of the first kind and $\theta_P(\rho)$ is the particular solution to the inhomogeneous differential equation. Using the variation of parameters technique, the particular solution can be written as an integral function,

$$\begin{aligned} \theta_P(x) = & \frac{\beta}{\gamma} \left\{ I_0(x) \int_0^x \frac{K_0(x') e^{-x'^2/\gamma}}{-I_0(x')K_1(x') - K_0(x')I_1(x')} dx' \right. \\ & \left. - K_0(x) \int_0^x \frac{I_0(x') e^{-x'^2/\gamma}}{-I_0(x')K_1(x') - K_0(x')I_1(x')} dx' \right\}. \end{aligned}$$

Since the particular solution is finite at the origin, the boundary condition in Equation 4.3 requires that $c_4 = 0$ leaving two integration constants.

These two remaining integration constants are fully determined by the boundary conditions in Equations 4.4 and 4.5. This results in two equations with two unknowns,

$$\begin{aligned} c_2 K_0(\sqrt{\Gamma/2}R/\sigma) &= c_3 I_0(\sqrt{\gamma/2}R/\sigma) + \theta_P(\sqrt{\Gamma/2}R/\sigma) \\ c_2 \kappa_{SP} \sqrt{\Gamma} K_1(\sqrt{\Gamma/2}R/\sigma) &= -\kappa_{SS} \sqrt{\gamma} \left\{ c_3 I_1(\sqrt{\gamma/2}) \right. \\ &\quad + I_1(x) \frac{\beta}{\gamma} \int_0^x \frac{K_0(x') e^{-x'^2/\gamma}}{-I_0(x') K_1(x') - K_0(x') I_1(x')} dx' \\ &\quad \left. - K_1(x) \frac{\beta}{\gamma} \int_0^x \frac{I_0(x') e^{-x'^2/\gamma}}{-I_0(x') K_1(x') - K_0(x') I_1(x')} dx' \right\}_{x=\sqrt{\frac{\gamma R}{2\sigma}}}, \end{aligned}$$

which can be solved for numerically for given κ_{SS} , κ_{SP} , g_g , and g_S .

An example predicted temperature distribution is shown in figure 4.2. As one might expect, the temperature rise is the greatest where the laser is heating, near the center of the microchamber. At the edge of the microchamber the slope is discontinuous due to the change in thermal conductivity. Where the thermal conductivity is lower it takes a larger temperature gradient to conduct the same heat energy. While the qualitative features of the temperature distribution are simple, the quantitative dependence on the free parameters are non trivial. The functional dependencies of κ_{SS} , κ_{SP} , g_G , and g_S are made indecipherable by the form of the particular solution. The exception is the absorbed power which acts as a linear scaling factor. Doubling the power simply scales the temperature distribution by a factor of two.

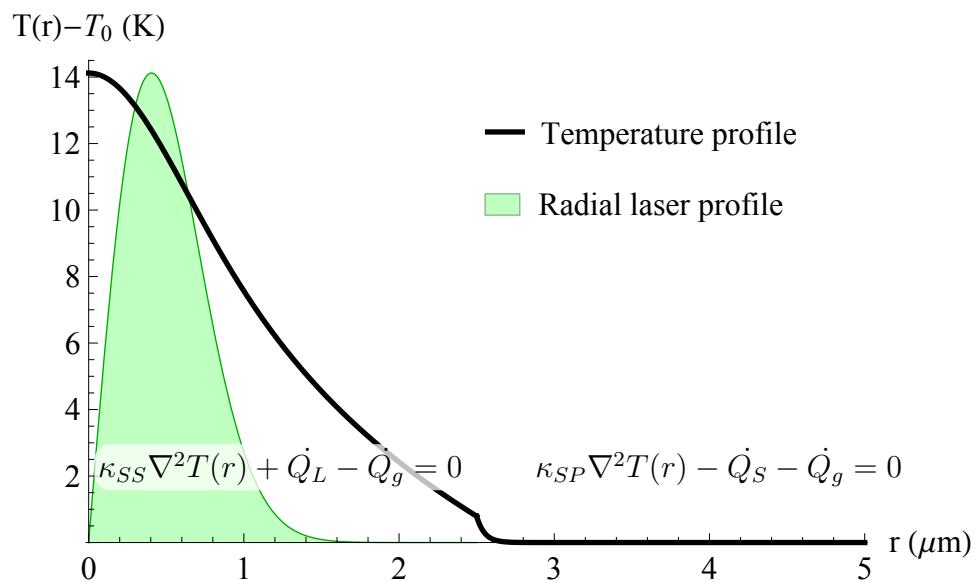


Figure 4.2: Theoretical temperature distribution in a graphene sealed microchamber heated by a centered laser. The black curve shows the temperature increase for a $5 \mu\text{m}$ diameter microchamber assuming that $\kappa_{SS} = 2000 \text{ W/m} - \text{K}$, $\kappa_{SP} = 370 \text{ W/m} - \text{K}$, $g_G = .03 \text{ MW/m}^2 - \text{K}$, and $g_S = 50 \text{ MW/m}^2 - \text{K}$. The radial profile of the 1.5 mW laser heat source with a $0.81 \mu\text{m}$ waste is overlaid in green. Also shown are the differential equations which govern the heat transfer in the two regions.

Chapter 5

Kekulé

$$u^{A,B}(\vec{r}_{A,B}, t) = c e^{i\vec{r}_{A,B} \cdot \vec{K}} \quad (5.1)$$

5.1 Theory

In this section the electronic band structure of the Kekulé lattice is calculated using a tight binding model for the expanded unit cell. This calculation is done in two parts. First, the effects of the change in lattice periodicity are discussed including a description of the zone folded electrical dispersion. Second, it is shown that by perturbing the zone folded Hamiltonian the phonons gap the electronic dispersion.

5.1.1 Kekulé geometry

The Kekulé distortion causes an expansion of the unit cell, a reduction of the BZ, and a modification of the primitive lattice vectors and reciprocal lattice vectors. When referencing the new geometry back to the intrinsic geometry, the notation previously developed in Chapter 1 is used. The expanded periodicity of the Kekulé lattice is determined by the periodicity of the distortion in Equation 5.1. This term repeats

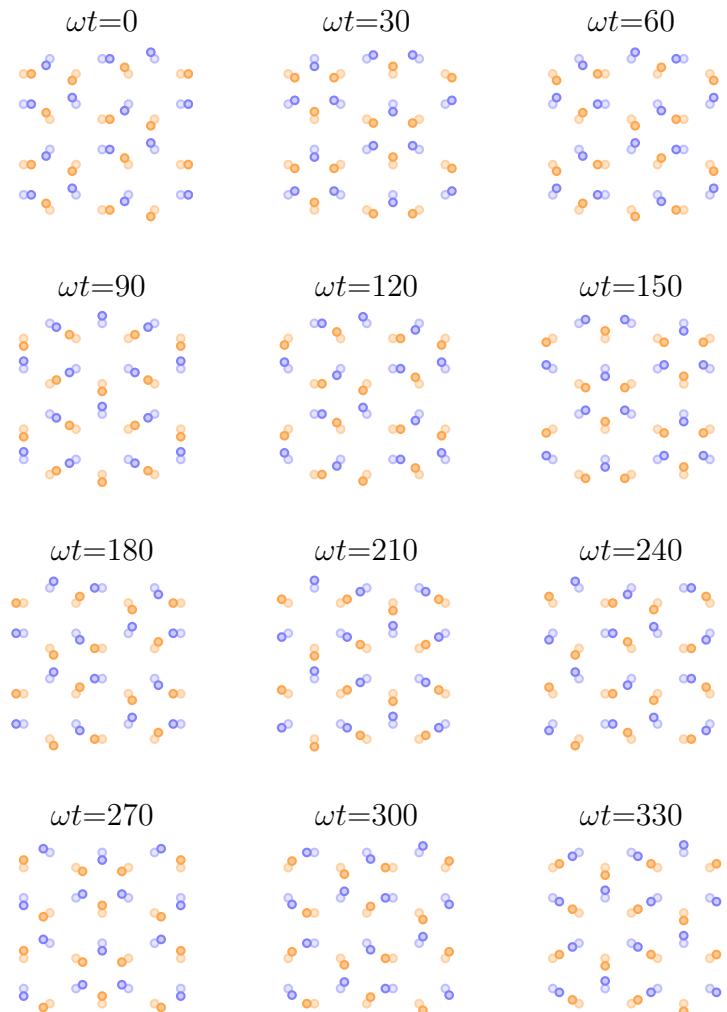


Figure 5.1

whenever

$$2\pi = \vec{r}_{A,B} \cdot \mathbf{K} = (m\vec{a}_+ + n\vec{a}_-) \cdot (\vec{b}_+ - \vec{b}_-)/3 = \frac{2\pi}{3}(m-n) ,$$

resulting in the enlarged unit cell shown in Figure 5.2.

The primitive lattice vectors of the Kekulé lattice,

$$\begin{aligned}\vec{A}_+ &= 2\vec{a}_+ - \vec{a}_- = \frac{3a}{2}(+\sqrt{3}, 1) \\ \vec{A}_- &= 2\vec{a}_- - \vec{a}_+ = \frac{3a}{2}(-\sqrt{3}, 1) ,\end{aligned}$$

represent a triangular lattice rotated by 90 degrees and expanded by a factor of $\sqrt{3}$ relative to the intrinsic lattice. This tripling of the area of the unit cell is accompanied by a corresponding decrease in the area of the BZ as shown in Figure 5.2. The primitive reciprocal lattice vectors,

$$\begin{aligned}\vec{B}_+ &= \frac{1}{3}(2\vec{b}_+ + \vec{b}_-) = \frac{2\pi}{3\sqrt{3}a}(+1, \sqrt{3}) \\ \vec{B}_- &= \frac{1}{3}(2\vec{b}_- + \vec{b}_+) = \frac{2\pi}{3\sqrt{3}a}(-1, \sqrt{3}) ,\end{aligned}$$

generate a hexagonal BZ rotated 90 degrees and shrunken by a factor of three relative to the intrinsic lattice. Unlike for intrinsic graphene, the interesting physics occurs near the center of the BZ. For completeness, three equivalent corners of the BZ are positioned at

$$\frac{2\pi}{9a}(-\sqrt{3}, -1), \quad \frac{2\pi}{9a}(+\sqrt{3}, -1), \quad \text{and} \quad \frac{2\pi}{9a}(0, 2) .$$

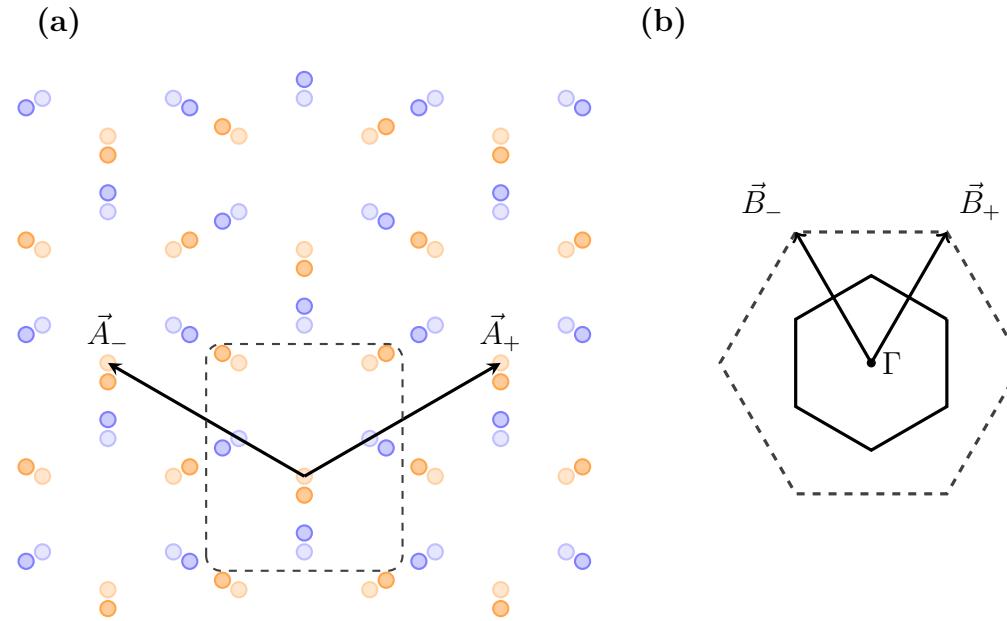


Figure 5.2: The real space (a) and reciprocal space (b) geometry of the Kekulé lattice. Figure (a) shows a snapshot of the atomic positions with the A sub-lattice shown as orange dots and the B sub-lattice shown as blue dots. For reference, the intrinsic graphene lattice is shown using faded dots. The dashed rectangle outlines the time independent unit cell and the labeled arrows represent the primitive lattice vectors. In (b) the hexagon with dashed lines indicates the BZ of intrinsic graphene while the hexagon with solid lines indicates the shrunken BZ of the Kekulé lattice. The primitive reciprocal lattice vectors are labeled.

5.1.2 Zone folding

When the BZ is reduced in size the two band dispersion of intrinsic graphene becomes a six band dispersion. The energy bands outside the new BZ are folded into the new BZ by translation by a reciprocal lattice vectors. Since the size of BZ is reduced by a factor there are two zone folding schemes resulting in six energy bands.

The zone folding schemes shown in Figure 5.3 describe the two distinct ways the symmetry reduced area of the new BZ can be mapped onto via translations of reciprocal lattice vectors. The rest of the BZ can be constructed using symmetry operations on the symmetry reduced area. In both zone folding schemes, the Dirac point is translated from the corner of the old BZ to the zone center. Since this will still correspond to the charge neutrality point, this is the point in the new BZ we will be most interested in. It is also worth noting that right edge of the symmetry reduced area in zone fold 1 shares its right edge with the symmetry reduced area already inside the new BZ. Also, the top edge of the symmetry reduced area is shared between zone fold 1 and zone fold 2.

The hierarchy of the folded energy bands away from degeneracies can be determined by comparing the zone folded areas to the electronic dispersion of intrinsic graphene shown in Figure 1.2. Just like for intrinsic graphene the the energy bands should be symmetric about the charge neutrality point. The new BZ occupies the basin in the intrinsic graphene dispersion. Thus, the unfolded area gives the lowest and highest energy bands. Zone fold 1 extends further into the basin of intrinsic graphene's dispersion and so gives the second lowest and second highest energy bands. Finally, the zone fold 2 gives the two middle energy bands.

The zone folded electronic dispersion can either be calculated by folding the dispersion calculated in Chapter 1 or by performing a tight binding calculation using

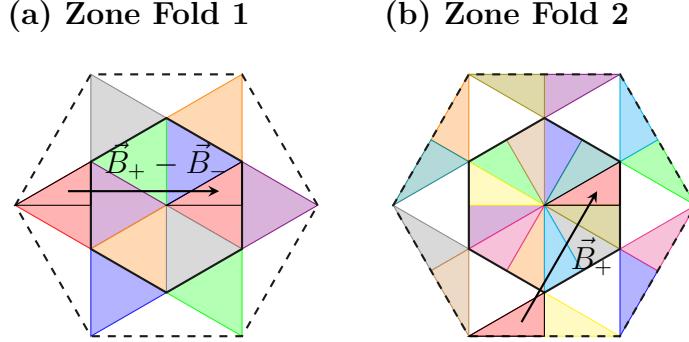


Figure 5.3: The zone foldings introduced by the Kekulé distortion. The outer hexagon with dashed lines is the BZ of intrinsic graphene and the inner hexagon with solid lines is the new BZ of the Kekulé lattice. Specifically shown, for both foldings the symmetry reduced area represented by the black outlined triangle is shifted by the labeled reciprocal lattice vector.

the expanded unit cell. Although it is more involved, the tight binding calculation will be done here because it will be needed later. Each of the six atoms in the unit cell must have its own raising and lowering operator. To simplify this bookkeeping the operators will be referenced back to the three two atom bases which make up the six atom Kekulé basis. These two atom bases are shown in Figure 5.4. The operators $a_{i,l}$ and $b_{i,l}$ are then the lowering operators for the three A sub-lattice atoms and the three B sub-lattice atoms in the l th Kekulé basis respectively. The index i runs over the three two atom bases. In this notation the real space nearest neighbor tight binding Hamiltonian is given by

$$\begin{aligned}
 H = & - \sum_l (ta_{1,l}^\dagger b_{1,l} + ta_{1,l}^\dagger b_{2,l} + ta_{1,l}^\dagger b_{3,l} \\
 & + ta_{2,l}^\dagger b_{1,l\uparrow} + ta_{2,l}^\dagger b_{2,l} + ta_{2,l}^\dagger b_{3,l\rightarrow} \\
 & + ta_{3,l}^\dagger b_{1,l\uparrow} + ta_{3,l}^\dagger b_{2,l\leftarrow} + ta_{3,l}^\dagger b_{3,l} + \text{H.C.}) . \tag{5.2}
 \end{aligned}$$

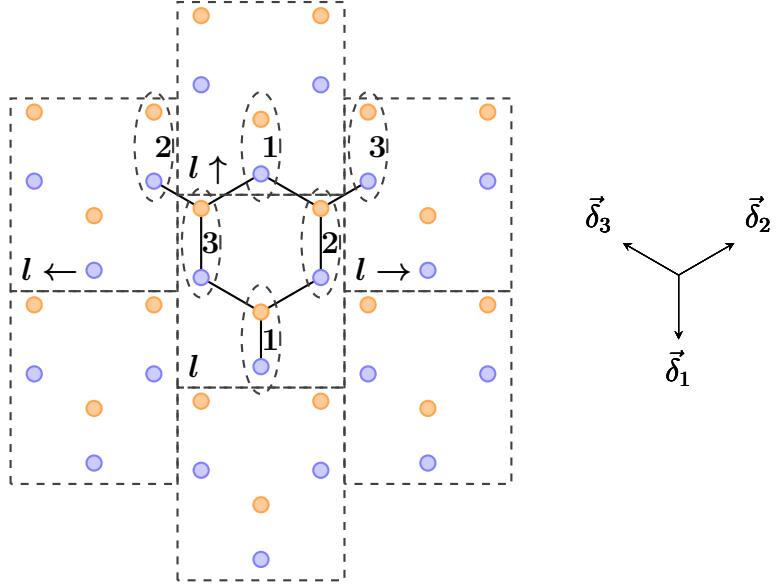


Figure 5.4: A diagram of the hoppings included in the Hamiltonian. Hoppings connect atoms originally from the A sub-lattice (orange) to atoms originally from the B sub-lattice (blue). They can pass between the labeled intrinsic unit cells (dashed ellipses) and they can also pass between the labeled extended unit cells (dashed rectangles). For reference the directions of the nearest neighbor vectors are included.

Later the hopping energies t will be made bond dependent to account for the altered nearest neighbor distances. In the meantime the bond independent hopping energy of intrinsic graphene, t_0 , will be used.

Similarly to intrinsic graphene, the individual terms in the sum can be simplified by writing the operators in Fourier space,

$$a_{m,l}^\dagger = \frac{1}{\sqrt{N}} \sum_{\vec{k}} e^{i\vec{k}\cdot\vec{R}_l} a_{m,\vec{k}}^\dagger , \quad (5.3)$$

where we are expanding about the positions of the Kekulé unit cells. The individual

terms are then

$$\begin{aligned} -t_0 \sum_l a_{m,l}^\dagger b_{m',l'} &= -\frac{t_0}{N} \sum_l \sum_{\vec{k}, \vec{k}'} a_{m,\vec{k}}^\dagger b_{m',\vec{k}'} e^{i\vec{R}_l \cdot (\vec{k} - \vec{k}')} e^{i(\vec{R}_l - \vec{R}'_l) \cdot \vec{k}'} \\ &= -t_0 \sum_{\vec{k}} a_{m,\vec{k}}^\dagger b_{m',\vec{k}} \underbrace{e^{i(\vec{R}_l - \vec{R}'_l) \cdot \vec{k}'}}_{s_{l-l'}} , \end{aligned}$$

which is only dependent on the l independent distance $l - l' \in \{0, \leftarrow, \uparrow, \rightarrow\}$ between the Kekulé unit cells that are being hopped between.

The Hamiltonian can then be expressed in matrix form as

$$H_0 = -t_0 \sum_{\vec{k}} \psi^\dagger \begin{pmatrix} 0 & 0 & 0 & s_0 & s_0 & s_0 \\ 0 & 0 & 0 & s_\uparrow & s_0 & s_\rightarrow \\ 0 & 0 & 0 & s_\uparrow & s_\leftarrow & s_0 \\ s_0^* & s_\uparrow^* & s_\uparrow^* & 0 & 0 & 0 \\ s_0^* & s_0^* & s_\leftarrow^* & 0 & 0 & 0 \\ s_0^* & s_\rightarrow^* & s_0^* & 0 & 0 & 0 \end{pmatrix} \psi , \quad (5.4)$$

where $\psi^\dagger = (a_1^\dagger, a_2^\dagger, a_3^\dagger, b_1^\dagger, b_2^\dagger, b_3^\dagger)$. As expected, the six by six Hamiltonian will provide six energy levels.

The zone folded electronic dispersion is shown in Figure 5.5. To best show the shapes of the bands the dispersion is plotted both over the full Kekulé BZ and also over only the symmetry reduced area. The six energy bands are clearly visible with the highest and lowest energy bands appearing as caps. As expected the Dirac point has been shifted to zone center where the four lower energy bands converge to touch at a single point. In agreement with the zone folding schemes the highest energy band is degenerate with the second highest energy band on the BZ border. Also, the second and third highest energy bands are degenerate along the lines connecting the Γ

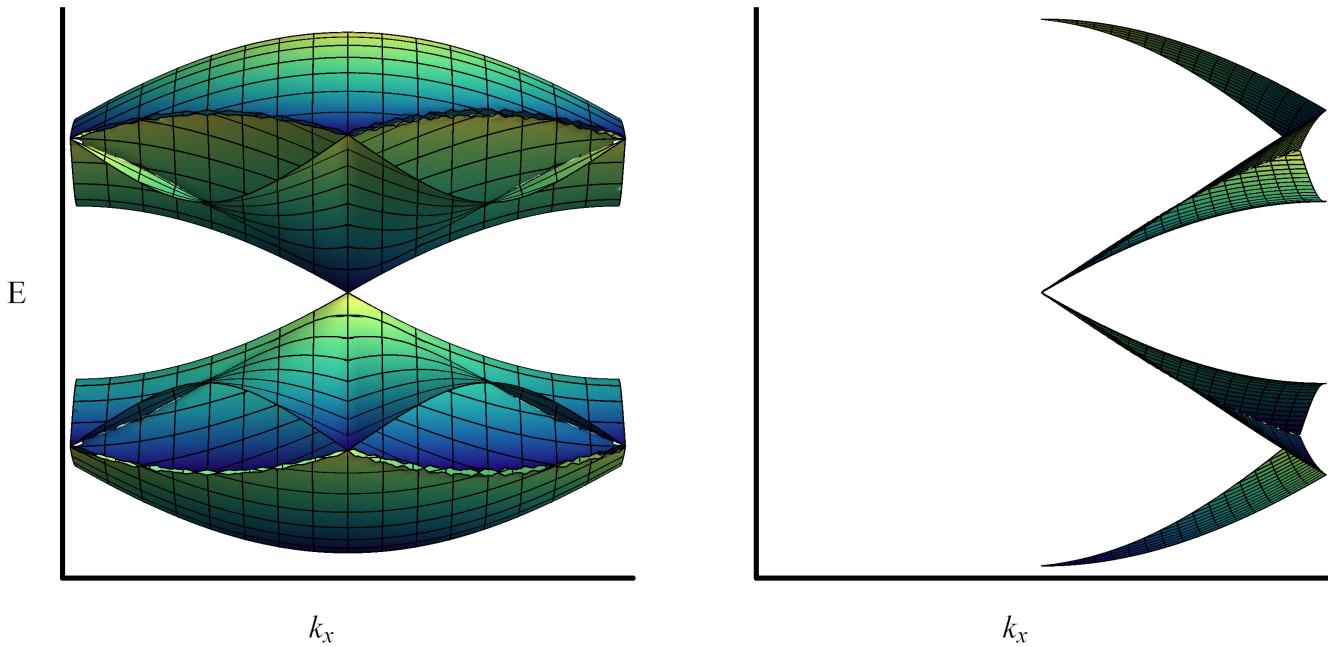


Figure 5.5: Surface plots of the folded electronic dispersion of the Kekulé lattice including all six energy bands. In the left plot the surfaces are plotted for the full Kekulé BZ whereas in the right plot only the symmetry reduced area is plotted.

point to the BZ border. In summary, the electronic dispersion calculated with a tight binding model of the expanded unit cell agrees with our zone folding predictions.

5.1.3 Altered hoppings

Coherent excitation of the iTO phonon mode does more than just enlarge the unit cell, it also modifies the hopping energies. This is similar to the case of strained graphene where the altered bond lengths caused altered hopping energies and generated new physics. In this case, however, the bond lengths vary with a much higher spatial frequency and the slowly varying approximation described in Appendix B is not applicable. In fact, the atomic displacements in Equation 5.1 will lead to hopping

alterations with a spatial frequency of $\mathbf{G} = \mathbf{K} - \mathbf{K}'$. As mentioned in Appendix B, a distortion with this frequency is expected to couple the inequivalent \mathbf{K} and \mathbf{K}' points. This coupling will generate a band gap at the Γ point of the zone folded dispersion.

Before determining the hopping alterations the bond length changes must be found. Although it was helpful in finding the lengths of the strained nearest neighbor vectors in Section 2.1.2, the Cauchy-Born rule cannot be used here. The iTO phonon causes the atoms in the A and B sub-lattices to rotate in opposite directions, an effect which can not be captured in the Cauchy-Born frame work. Instead, the bond lengths must be calculated directly from the displacements. Iadecola *et al.* showed that the bond length alterations generate hopping alterations,

$$\delta t_{m,j} = \frac{1}{3} \Delta(t) e^{i\mathbf{K}\cdot\vec{\delta}_j} e^{i\mathbf{G}\cdot\vec{r}_m} + \text{c.c.}$$

with $\Delta(t) = -i3\beta t_0 \frac{c^*}{a} e^{i\omega t}$, (5.5)

which have a spatial frequency component of \mathbf{G} that couples the Dirac points (Iadecola et al., 2013). Here \vec{r}_m is the position of the A sub-lattice atom involved in the hopping. Index m indicates which of the three old unit cells the atom is in. The B sub-lattice atom is specified through $\vec{\delta}_j$, the unperturbed nearest neighbor vector which connects the A sub-lattice atom to the B sub-lattice atom. Figure 5.4 summarizes these indices. For completeness, the calculation of Equation 5.5 is included in Appendix D.

5.1.4 Tight binding of the expanded Kekulé lattice

The Kekulé mode causes the hopping energies in Equation 5.2 to be bond and time specific. Taking $t = t_0 + \delta t_{m,j}$ with $\delta t_{m,j}$ defined in equation 5.5 breaks Equation 5.2 into two pieces. The first, corresponding to t_0 is just the zone folding Hamiltonian in

Equation 5.4. The second is the additional perturbation which opens the band gap.

Similar to the individual terms in H_0 , each term in H' can be simplified by writing the operators in Fourier using Equation 5.3

$$\begin{aligned} -\sum_l \delta t_{m,j} a_{m,l}^\dagger b_{m',l'} &= -\frac{1}{N} \sum_l \sum_{\vec{k}, \vec{k}'} \delta t_{m,j} a_{m,\vec{k}}^\dagger b_{m',\vec{k}'} e^{i\vec{R}_l \cdot (\vec{k} - \vec{k}')} e^{i(\vec{R}_l - \vec{R}'_l) \cdot \vec{k}'} \\ &= \frac{1}{3} \sum_{\vec{k}} a_{m,\vec{k}}^\dagger b_{m',\vec{k}} \underbrace{\Delta(t) \left\{ e^{i\mathbf{K} \cdot \vec{\delta}_j} e^{i\mathbf{G} \cdot \vec{r}_m} + \text{c.c.} \right\}}_{g_{m,j,l-l'}} e^{i(\vec{R}_l - \vec{R}'_l) \cdot \vec{k}} . \end{aligned}$$

In addition to the l independent distance between involved Kekulé unit cells, each term depends on m which indicates the old unit cell which the A sub-lattice atom occupies, and j which indicates which $\vec{\delta}_j$ the hopping is along. The associated vectors are $\vec{R}_l - \vec{R}'_l \in \{0, \vec{A}_+, \vec{A}_+ + \vec{A}_-, \vec{A}_-\}$ and $\vec{r}_m \in 0, \vec{a}_+, \vec{a}_-$.

Figure 5.4 is useful in determining the elements of perturbed Hamiltonian.

$$H' = -\frac{1}{3} \sum_{\vec{k}} \psi^\dagger \begin{pmatrix} 0 & 0 & 0 & g_{1,1,0} & g_{1,2,0} & g_{1,3,0} \\ 0 & 0 & 0 & g_{2,3,\uparrow} & g_{2,1,0} & g_{2,2,\rightarrow} \\ 0 & 0 & 0 & g_{3,2,\uparrow} & g_{3,3,\leftarrow} & g_{3,1,0} \\ g_{1,1,0}^* & g_{2,3,\uparrow}^* & g_{3,2,\uparrow}^* & 0 & 0 & 0 \\ g_{1,2,0}^* & g_{2,1,0}^* & g_{3,3,\leftarrow}^* & 0 & 0 & 0 \\ g_{1,3,0}^* & g_{2,2,\rightarrow}^* & g_{3,1,0}^* & 0 & 0 & 0 \end{pmatrix} \psi ,$$

5.1.5 Results and discussion

5.2 Experimental design

5.2.1 Phonon excitation

5.2.2 Band gap measurement

Appendix A

The first Brillouin zone of strained graphene

In this appendix an approximate analytic expression for the positions of the corners of the BZ of deformed graphene is presented. The BZ can then be constructed by connecting these points. The strain dependence is found by applying a general method for determining the positions of the BZ corners for close to hexagonal lattices:

1. The lattice vectors are used to determine the reciprocal lattice vectors.
2. The combination of reciprocal lattice vectors which give the important points in reciprocal space are determined.
3. The conditions for Bragg refraction will be used to determine the corners of the BZ.

After establishing this general methodology, the explicit strain dependence can easily be determined to first order.

The lattice vectors \vec{a}_+ and \vec{a}_- determine the reciprocal lattice vectors through the relationship $\vec{a}_i \cdot \vec{b}_j = 2\pi\delta_{ij}$, where \vec{b}_j are the two reciprocal lattice vectors, i and j are $\in \{+, -\}$, and δ_{ij} is the Kronecker delta function (Kittel, 2005). In two dimensions

this can be cast as a matrix relationship,

$$\begin{pmatrix} a_{\pm x} & a_{\pm y} \\ a_{\mp x} & a_{\mp y} \end{pmatrix} \begin{pmatrix} b_{\pm x} \\ b_{\pm y} \end{pmatrix} = \begin{pmatrix} 2\pi \\ 0 \end{pmatrix}, \quad (\text{A.1})$$

which can be easily solved by inverting the matrix. Having determined the form for the reciprocal lattice vectors, the next step is to determine the boundaries of the BZ.

The traditional method of determining the BZ does not lend itself to simple algorithmic implementation. In this method, one draws the perpendicular bisector of each reciprocal lattice vector given by $\vec{G} = n\vec{b}_+ + m\vec{b}_-$ where m and n are integers. The most inner polygon formed by the perpendicular bisectors is then the BZ (Kittel, 2005). The first step to simplify this method is to restrict the number of perpendicular bisectors which are considered. The construction of the BZ using the minimum number of reciprocal lattice vectors is shown for a perfect hexagonal lattice in Figure A.1. This six combinations of reciprocal lattice vectors needed to construct the BZ is relatively robust to distortions of the hexagonal lattice. They determine the BZ for strains as large as 20% armchair uniaxial, 20% armchair uniaxial, or 20% shear strain. The reciprocal lattice vectors themselves are altered by strain but the combinations used to construct the BZ are the same. This was confirmed by comparing the BZ predicted by the method presented here with that calculated using a geometric construction. The phrase “close to hexagonal” lattices is used to refer to those lattice for which the minimal set of reciprocal lattice vectors is given by the combinations shown in Figure A.1. For this discussion, the strained graphene lattice is then a close to hexagonal lattice.

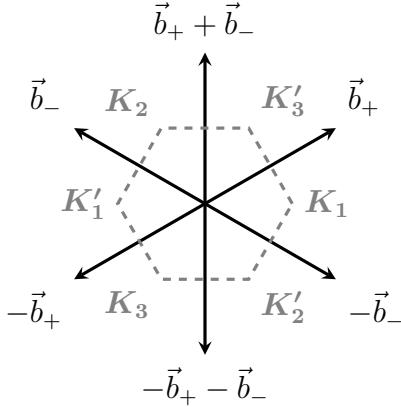


Figure A.1: The construction of the BZ for a hexagonal lattice. The edges of the BZ (dashed, gray) are made up of the perpendicular bisectors of the labeled reciprocal lattice vectors (black, arrows). Close to hexagonal lattices are lattices for which the same set of reciprocal lattice vectors define the BZ.

Having restricted the reciprocal lattice vectors, the corners of the BZ can be constructed from the condition for Bragg reflection. This condition,

$$\vec{k} \cdot \left(\frac{1}{2} \vec{G} \right) = \left(\frac{1}{2} \vec{G} \right)^2 ,$$

defines the wave vectors, \vec{k} , which make up the perpendicular bisector of the reciprocal lattice vector (Kittel, 2005). The corners of the BZ can be found by noticing that if the wave vector is on the perpendicular bisector sequential reciprocal lattice vectors in Figure A.1, it is a corner of the BZ. For example, the corner K_1 is a perpendicular bisector of both $-\vec{b}_-$ and \vec{b}_+ . Thus, the corners of the BZ can be calculated analytically by using the matrix identity,

$$\begin{pmatrix} G_{1,x} & G_{1,y} \\ G_{2,x} & G_{2,y} \end{pmatrix} \begin{pmatrix} k_x \\ k_y \end{pmatrix} = \begin{pmatrix} \frac{1}{2} G_1^2 \\ \frac{1}{2} G_2^2 \end{pmatrix} \quad (\text{A.2})$$

where \vec{G}_1 and \vec{G}_2 are sequential reciprocal lattice vectors from Figure A.1. This matrix identity can be inverted to determine the wave-vector at the corner of the BZ. This completes the general methodology for determining the BZ based on the lattice vectors for close to hexagonal lattices. In summary, the reciprocal lattice vectors are calculated from the lattice vectors using Equation A.1 and then the corners of the BZ can be found using Equation A.2 for the reciprocal lattice vector combination shown in Figure A.1.

The final step is to calculate the positions of the BZ corners using the form of the strained lattice vectors in Equation 2.3. The first order approximation found using Wolfram Mathematica version 9.0 is

$$\begin{aligned}\mathbf{K}_1 = -\mathbf{K}'_1 &\simeq \frac{4\pi}{3\sqrt{3}a} \begin{pmatrix} 1 \\ 0 \end{pmatrix} + \frac{4\pi}{3\sqrt{3}a} \begin{pmatrix} -\frac{1}{2}u_{xx} - \frac{1}{2}u_{yy} \\ -\frac{1}{2}u_{yx} - \frac{3}{2}u_{xy} \end{pmatrix} \\ \mathbf{K}_2 = -\mathbf{K}'_2 &\simeq \frac{4\pi}{3\sqrt{3}a} \begin{pmatrix} -\frac{1}{2} \\ \frac{\sqrt{3}}{2} \end{pmatrix} + \frac{4\pi}{3\sqrt{3}a} \begin{pmatrix} u_{xx} - \frac{1}{2}u_{yy} - \frac{\sqrt{3}}{2}u_{yx} \\ -\frac{\sqrt{3}}{2}u_{yy} - \frac{1}{2}u_{yx} \end{pmatrix} \\ \mathbf{K}_3 = -\mathbf{K}'_3 &\simeq \frac{4\pi}{3\sqrt{3}a} \begin{pmatrix} -\frac{1}{2} \\ -\frac{\sqrt{3}}{2} \end{pmatrix} + \frac{4\pi}{3\sqrt{3}a} \begin{pmatrix} u_{xx} - \frac{1}{2}u_{yy} + \frac{\sqrt{3}}{2}u_{yx} \\ \frac{\sqrt{3}}{2}u_{yy} - \frac{1}{2}u_{yx} \end{pmatrix}.\end{aligned}$$

The shifts in the corners of the BZ are opposite for time reversal pairs respecting time reversal symmetry. The positions of the corners are determined by the terms of the displacement gradient tensor and not the strain tensor. These approximation were successfully tested against the geometric construction. They are used in Chapter 2 to visualize the distortion of Reciprocal space which accompanies the deformation of the real space lattice. In particular, Figure 2.2 shows the shifts in the corners of the BZ for several deformation geometries.

Appendix B

Slowly varying approximation

The pseudo vector potential interpretation of strain developed in Chapter 2 is an approximation to a full rigorous treatment of a non-uniform strain field. In a strict treatment of non-uniform strain the i dependencies of Equation 2.6 cannot be neglected and the transformation into Fourier space is obstructed. Without elimination of the spatial dependence, the Hamiltonian of the full system must be painstakingly solved. In this case, pseudo magnetic field phenomena would be captured by a Landau level quantized density of states. Although rigorous, this method is more arduous and also less conceptually pleasing. In the slowly varying approximation the i dependencies in Equation 2.6 are eliminated. In this approximation the pseudo vector potential formalism is valid, providing a qualitative framework for understanding the observed effects. In this Appendix this approximation will be developed and the limits of validity will be discussed.

Remembering that $\nabla \mathbf{u}$ is i dependent, the i dependent terms in Equation 2.6 are

$$\begin{aligned}
H_i &= \sum_i (t_0 + \delta t_{i,j}) e^{i(\vec{k} - \vec{k}') \cdot \vec{R}'_i} e^{-i\vec{k}' \cdot \vec{\Delta}'_{i,j}} \\
&= \sum_i (t_0 + \delta t_{i,j}) e^{i(\vec{k} - \vec{k}') \cdot \vec{R}'_i} e^{-i\vec{k}' \cdot (\mathbf{1} + \nabla \mathbf{u}) \vec{\Delta}_j} \\
&\simeq e^{-i\vec{k}' \cdot \vec{\Delta}_j} \sum_i (t_0 + \delta t_{i,j}) e^{i(\vec{k} - \vec{k}') \cdot \vec{R}'_i} \left(1 - i\vec{k}' \cdot \nabla \mathbf{u} \cdot \vec{\Delta}_j \right) \\
&\simeq e^{-i\vec{k}' \cdot \vec{\Delta}_j} \left\{ t_0 \sum_i e^{i(\vec{k} - \vec{k}') \cdot \vec{R}'_i} - it_0 \vec{k}' \cdot \sum_i (\nabla \mathbf{u} e^{i(\vec{k} - \vec{k}') \cdot \vec{R}'_i}) \cdot \vec{\Delta}_j + \sum_i \delta t_{i,j} e^{i(\vec{k} - \vec{k}') \cdot \vec{R}'_i} \right\} \\
&= e^{-i\vec{k}' \cdot \vec{\Delta}_j} \left\{ Nt_0 \delta_{\vec{k}, \vec{k}'} - it_0 \vec{k}' \cdot \widetilde{\nabla \mathbf{u}}_{\vec{k}-\vec{k}'} \cdot \vec{\Delta}_j + \widetilde{\delta t}_{\vec{k}-\vec{k}'} \right\},
\end{aligned}$$

where $\widetilde{\nabla \mathbf{u}}_{\vec{k}-\vec{k}'}$ and $\widetilde{\delta t}_{\vec{k}-\vec{k}'}$ are the Fourier transforms of $\delta t_{i,j}$ and $\nabla \mathbf{u}$ respectively. Only terms first order in products of the small quantities $\nabla \mathbf{u}$ and $\delta t_{i,j}$ were kept. All of the i dependence has been absorbed by the Fourier transforms.

Only specific Fourier components yield relevant $\widetilde{\nabla \mathbf{u}}_{\vec{k}-\vec{k}'}$ and $\widetilde{\delta t}_{\vec{k}-\vec{k}'}$ when working in the low energy regime. The wave-vectors will again be approximated as $\vec{k} = \mathbf{K} + \vec{q}$ and $\vec{k} = \mathbf{K}' + \vec{q}$ with the additional small parameter qa , giving

$$\begin{aligned}
\widetilde{\delta t}_{\vec{k}-\vec{k}'} &= \sum_i \delta t_{i,j} e^{i(\mathbf{K}'') + \vec{q} - \mathbf{K}' - \vec{q}' \cdot \vec{R}'_i} \\
&\simeq \sum_i \delta t_{i,j} e^{i(\mathbf{K}'') - \mathbf{K}'') \cdot \vec{R}'_i} (1 + (\vec{q} - \vec{q}') \cdot \vec{R}'_i) \\
&= \sum_i \delta t_{i,j} e^{i(\mathbf{K}'') - \mathbf{K}'') \cdot \vec{R}'_i},
\end{aligned}$$

where \mathbf{K}'' refers to either \mathbf{K} or \mathbf{K}' depending on the wave-vector. The low energy approximation for $\widetilde{\nabla \mathbf{u}}_{\vec{k}-\vec{k}'}$ is similar. Thus, to first order in the small parameters the only relevant Fourier components are for $\vec{k} - \vec{k}' \in \{0, \mathbf{K} - \mathbf{K}', \mathbf{K}' - \mathbf{K}\}$. Interestingly, the high frequency components could act to couple the \mathbf{K} and \mathbf{K}' points. However,

here we apply the slowly varying approximation for which we eliminate the high frequency components and limit $\vec{k} - \vec{k}' \rightarrow 0$, yielding

$$\begin{aligned}\tilde{\delta t}_{\vec{k}-\vec{k}'} &\simeq N\delta_{\vec{k},\vec{k}'}(\langle \delta t_{i,j} \rangle) \\ \widetilde{\nabla \mathbf{u}}_{\vec{k}-\vec{k}'} &\simeq N\delta_{\vec{k},\vec{k}'}(\langle \nabla \mathbf{u} \rangle).\end{aligned}$$

where $\langle \delta t_{i,j} \rangle = \delta t_j$ and $\langle \nabla \mathbf{u} \rangle$ are the average value over i of $\delta t_{i,j}$ and $\nabla \mathbf{u}$ respectively.

Thus, in the slowly varying approximation the i dependence of Equation 2.6 becomes

$$H_i = e^{-i\vec{k}' \cdot \vec{\Delta}_j} \left\{ N\delta_{\vec{k},\vec{k}'} \left(t_0 + \delta t_j - it_0 \vec{k}' \cdot \nabla \mathbf{u} \cdot \vec{\Delta}_j \right) \right\}$$

which gives the same result as the simple substitutions $\delta t_{i,j} \rightarrow \delta t_j = \langle \delta t_{i,j} \rangle$ and $\nabla \mathbf{u}_i \rightarrow \langle \nabla \mathbf{u} \rangle$ in Equation 2.6. This approximation required low energies, $qa \ll 1$, small strains, $\delta t_j \ll t_0$ and $\nabla \mathbf{u} \ll 1$, and slowly varying strain fields, $\tilde{\delta t}_{\mathbf{K}-\mathbf{K}'} = \widetilde{\nabla \mathbf{u}}_{\mathbf{K}-\mathbf{K}'} = 0$. It eliminates the spatial dependence in the Hamiltonian of strained graphene allowing the identification of the pseudo vector potential.

Appendix C

Global fitting algorithm details

This Appendix is devoted to the method used to extract the fitting parameters from the Raman line scans discussed in Chapter 3. Throughout, fitting parameters refers to the non-linear parameters of interest, namely the dimensionless friction, F , the Grüneisen parameter, γ , and the shear deformation potential, β . The determination of these values is complicated; there is no out of the box technique which can be blindly applied. Additionally it is computationally expensive. The fits including all three fitting parameters took two days to run on 7 processors. Accordingly, this Appendix will endeavor to give enough detail to easily reproduce the fitting algorithm while providing tricks to maximize efficiency in Wolfram Mathematica 8.0 along the way.

Since the functional dependencies of the fitted parameters are very complicated and non-linear, there is no applicable out of the box fitting algorithm. Instead, a brute force algorithm is used. A range of fitting parameters sets, (F, γ, β) , are iterated through and for each set the global reduced χ^2 ,

$$\tilde{\chi}^2 = \frac{1}{D.O.F.} \sum_i \left(\frac{model_i - data_i}{\sigma_i} \right)^2 ,$$

is calculated as a goodness of fit metric. Here, $D.O.F.$ is the degrees of freedom in

the fit equal to the number of fit data points less the number of parameters fit to the data, $model_i - data_i$ is the difference between the modeled and measured values at point i , and σ_i is the measured uncertainty in $data_i$. Most importantly, i includes the data points from *all* of the spectra in the line scan. Instead of fitting each spectra individually, the full line scan is taken into account for a global determination of the fitting parameters. The best fit value is then found by looking for the fitting parameters which yield the smallest $\tilde{\chi}^2$.

The algorithm used to find the best fit is detailed in the flowchart in Figure C.1. Broadly speaking, the algorithm works by looping across a range of fitting parameters and checking how accurately the fitting parameters represent the global line scan spectra. Individual steps are described in the proceeding. Steps with red outlines in the flowchart will be described in the most detail as they are the most difficult.

Load spectra, known parameters

The first step is to load in the measured spectra and the independently determined parameters. The majority of the many free parameters are measured independently. The applied pressure is measured using a digital pressure gauge and recorded as a function of time. Often, the pressure regulator will allow the pressure to decrease by several tenths of a PSI during a measurement. In this case, the average of the pressure over the measurement window is used. The determination of the pressure trapped inside the microchamber as well as the unstrained G band position for the suspended FLG is detailed in Section 3.2.1. The G band width for the suspended graphene is taken from the center of the atmospheric pressure line scan. For supported FLG, the width and position of the unstrained G band is taken from the spectra at the largest radial distance measured during a line scan at either 0.17 MPa or atmospheric

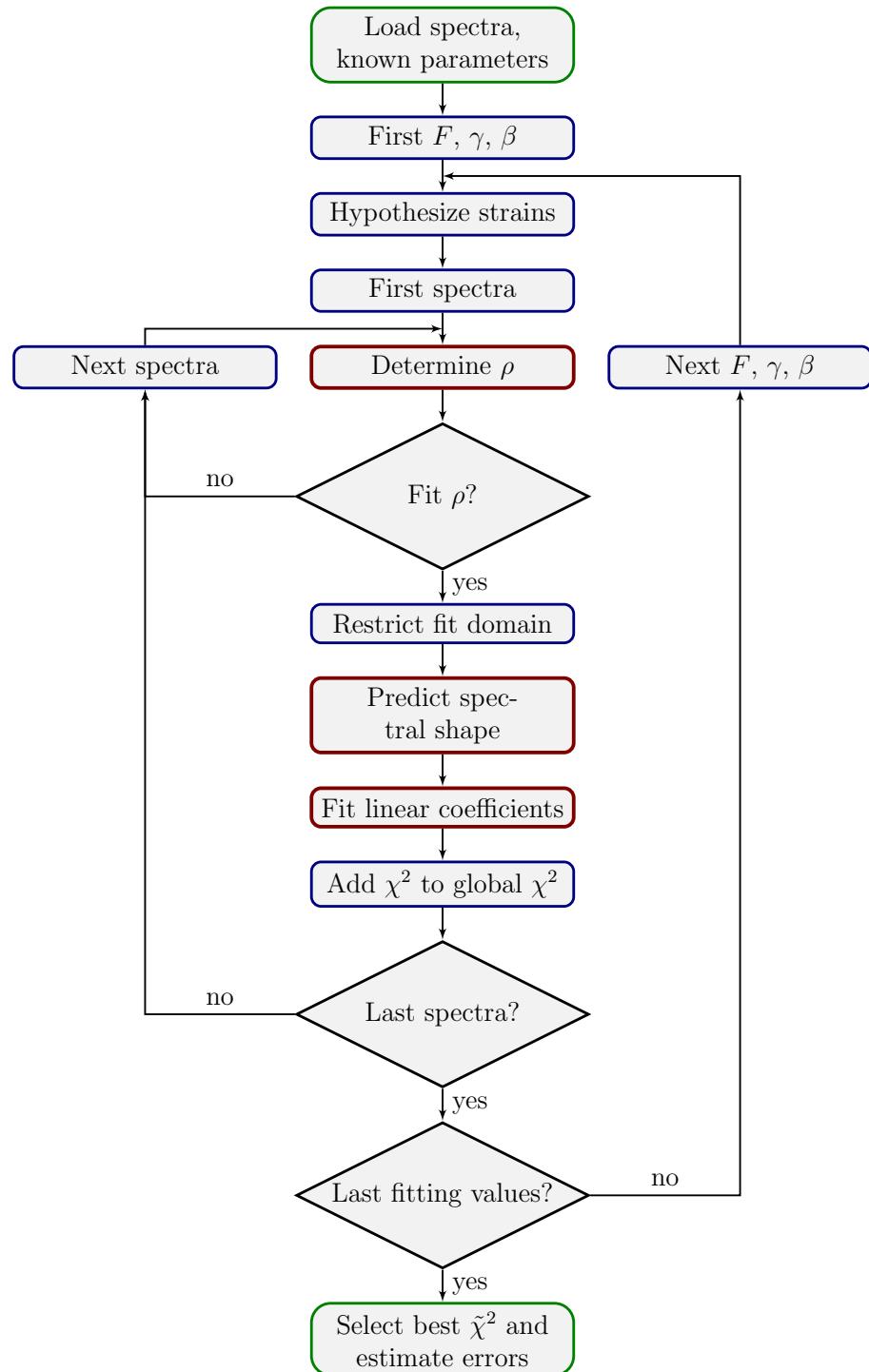


Figure C.1: Flow chart of global fitting algorithm to determine F , γ , β from line scans over pressurized graphene sealed microchambers.

pressure. The number of graphene layers is found using Raman spectroscopy (Ferrari et al., 2006) and optical contrast (Blake et al., 2007; Casiraghi et al., 2007). The measurement of the laser spot size is described in Section 3.1 Microchamber radius is determined from ambient pressure AFM which is analyzed depending on the devices ambient pressure behavior. For those devices that bulge up, the radius is found by fitting an AFM section cut which spans the microchamber center with the lowest order approximation to the Hencky model: A parabola. For those devices which stick to the side walls, the radius is measured manually based on the two dimensional topography. This leaves only the friction, Grüneisen parameter, shear deformation potential, signal intensity, and signal background as unknowns.

First F, γ, β

The outer loop cycles through the range of sampled fitting parameters. A balance must be struck when defining the range of fitting parameters to be tested. Too few fitting parameters increases the chances of missing the global minimum while too many makes the computational load too high. Supplementing this fitting algorithm with a manual fitting technique can build an intuition for the system while also decreasing the chances the global minimum is not found. By guessing individual fitting parameters instead of looping through a range of values, the user can observe how the fitting parameters effect the predicted spectra and $\tilde{\chi}^2$. This intuition gained should assist the user in determining a reasonable range of fitting parameters.

Hypothesize strains

Having selected one set of fitting parameters, the guessed dimensionless friction, F , is used with the known dimensionless loading, q , to generate a hypothetical strain

distribution. This amounts to solving Equations 3.19 for the parameters which determine the shape of the strain distribution. Mathematica's Solve routine can be effectively used for this. When including terms up to tenth order in the strain expansion a solution is found in seconds. The hypothetical strain distribution is defined by the single real solution. The next task is to determine how well the measure spectra agree with this strain distribution.

First spectra

The inner loop cycles through the spectra in the line scan. Although the fitting is global in scope, individual spectra must still be isolated to fit linear coefficients.

Determine ρ

Even though the relative distance the sample is moved between spectra is set experimentally, the absolute position of the individual spectra measured from the center of the microchamber is not necessarily fully determined. This uncertainty could arise due to a slight offset in the beam position relative to the targeted position, from sample drift, or from targeting uncertainty. It can be best overcome using a retroactive determination of the central spectra us single Lorentzian fits to the data. As shown in Figure 3.5 the single Lorentzian fit to the line scan spectra comes to a symmetric minimum. This minimum occurs at point closest to the center of the microchamber where the strains, and thus the down-shifts, are the greatest. Its position can be accurately determined by fitting a similar symmetric function to the single Lorentzian fits such as a Gaussian or a parabola. The dimensionless radius, $\rho = r/R$, of each spectra is then determined by the position of the center of the microchamber, the spacing between spectra, and the radius of the microchamber.

Fit ρ ?

Not every spectra in the line scan is included in the fit. For microchambers with radii larger than $1.5 \mu\text{m}$, the global $\tilde{\chi}^2$ includes every spectra except for those located less than half the beam waist away from the edge of the microchamber. This is done to avoid the inclusion of the fitting parameters that would need to be added to account for the different optical interference conditions for suspended and supported graphene.

When fitting the smaller radius microchambers ($R < 1.5 \mu\text{m}$), the spectra from the suspended graphene are ignored. Figure C.2 shows an example of the resulting best fit. Even though they were not included in the fit, the spectra predicted by the extended Hencky model matches the major feature of the suspended spectra. Spectra from the suspended region ($\rho < 1$) exhibit a high energy shoulder that is not predicted by the extended Hencky model. This feature is observed in each of the three measured graphene sealed microchambers with radius less than $1.5 \mu\text{m}$ but not for the five larger radius microchambers. This feature is believed to come from airy rings in the focused laser beam. These rings of intensity would contribute signal from lower strain regions giving a distinct higher energy signal. For larger radius microchambers, the laser spot is smaller relative to the microchamber radius and so it samples a more uniform distribution of strains, making an extra contribution from the airy rings unimportant. Complications of fitting these features are avoided by only including the spectra from the supported graphene when fitting.

Restrict fit domain

The numerical integration over the point spread function is the most numerically intensive step in the algorithm and it is linear in the number of frequency points sampled. To hasten this process, many of the measured data points can be eliminated.

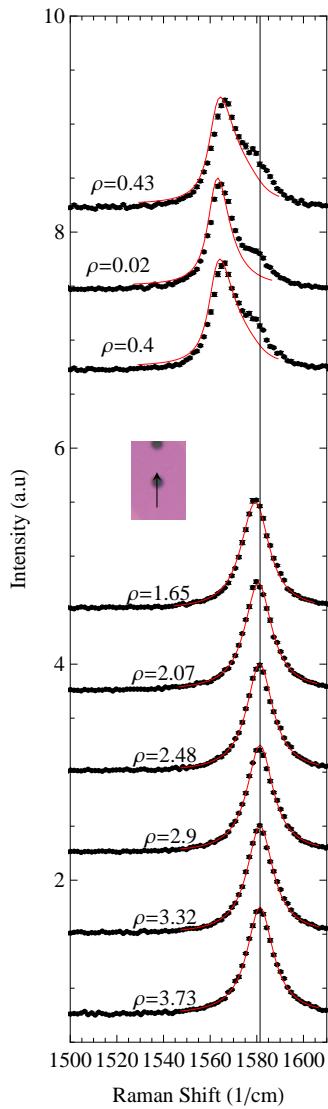


Figure C.2: An example globally fit line scan with radii less than 1.5 microns ($R \sim 1.2 \mu m$, bilayer, 0.8 MPa). Spectra taken along the path shown in the inset are arrayed vertically with those spectra within half a beam waste of the microchamber edge omitted. The black vertical line indicates the supported graphene's unstrained G band energy. Plotted in red are global fitting results which include only supported spectra as described in the text.

They are flat background which adds no value to the fit. Since the energy of the features vary by as much as 100 1/cm, a static fitting domain is not ideal. Instead the expendable data points are eliminated using a position dependent fitting domain. The domain boundaries are calculated using the hypothetical energies of the G^- (ω_G^-) and G^+ (ω_G^+) peaks at the position the spectra is taken given the values of γ and β . The fitting domain is taken as $[\omega_G^- - 4\Gamma_{in}, \omega_G^+ + 3\Gamma_{in}]$ where Γ_{in} is the width of the G band for the suspended graphene. The asymmetry of the fitting domain is chosen to better match the fitted features. In this way computation time is decrease without jeopardizing the quality of the fit.

Predict spectral shape

Measured spectra do not represent the signal from a single point. In fact, as shown in Section 3.1 the focused laser excitation has a Gaussian profile with a beam waist of $0.81 \pm .01 \mu\text{m}$. The need to account for the finite excitation spot size is illustrated in Figure C.3 which shows the intensity envelope of the excitation beam overlaid on the strain profile. Even for this large radii microchamber, each spectra represents a continuum of strain states. This section explains how the shape for the current values of the fitting parameters is predicted by integrating over the excitation point spread function. In the next section, this spectral shape will be scaled to best fit the data.

It should be noted that when the suspended graphene is pushed into the microchamber, it is also pushed slightly out of the focal plane. This effect is minimal. The confocal length of a 514 nm laser with a $0.81 \mu\text{m}$ beam waste is $4 \mu\text{m}$, yielding a beam expansion of $< 3\%$ for the $< 1 \mu\text{m}$ graphene deflections encountered in the experiments. This was considered small enough to ignore.

The finite spot size is accounted for by integrating over the spectra contributed by

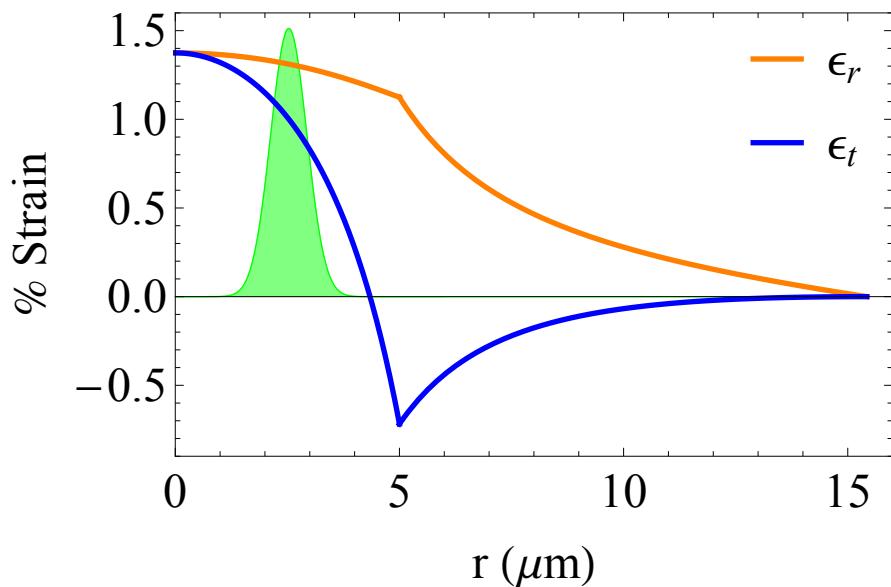


Figure C.3: Laser excitation profile overlaid on the best fit strain distribution for the $\sim 5 \mu\text{m}$ radius monolayer graphene sealed microchamber with 0.80 MPa of applied pressure.

each point weighted by the excitation point spread function. Each point contributes a spectra with two Lorentzian peaks centered at

$$\begin{aligned}\omega^+(\rho) &= \omega_0(\rho) - \omega_0(\rho)\gamma(\epsilon_r(\rho) + \epsilon_t(\rho)) + \frac{1}{2}\beta(\epsilon_r(\rho) - \epsilon_t(\rho)) \\ \omega^-(&\rho) = \omega_0(\rho) - \omega_0(\rho)\gamma(\epsilon_r(\rho) + \epsilon_t(\rho)) - \frac{1}{2}\beta(\epsilon_r(\rho) - \epsilon_t(\rho)) .\end{aligned}$$

The resulting spatial dependence of the Raman spectra is

$$g(\rho, \omega) = \frac{\left(\frac{1}{2}\Gamma(\rho)\right)^2}{(\omega - \omega^+(\rho))^2 + \left(\frac{1}{2}\Gamma(\rho)\right)^2} + \frac{\left(\frac{1}{2}\Gamma(\rho)\right)^2}{(\omega - \omega^-(\rho))^2 + \left(\frac{1}{2}\Gamma(\rho)\right)^2} ,$$

A position dependence is given to $\omega_0(\rho)$ and $\Gamma(\rho)$ because the suspended graphene is in general less doped than supported graphene. This is taken into account by treating $\omega_0(\rho)$ and $\Gamma(\rho)$ as step functions which change value at the edge of the microchamber.

To match the circular symmetry of the strain distribution, the excitation point spread function should be expressed as an envelope function in radial coordinates with the origin at the center of the microchamber. Accounting for the point spread function, the shape of the measured spectra is given by

$$\begin{aligned}f(\omega, \rho_\star) &= \int \int dx dy e^{-\frac{(x-x_\star)^2+y^2}{2\sigma^2}} g(\sqrt{x^2+y^2}/R, \omega) \\ &= e^{-\frac{r_\star^2}{2\sigma^2}} \int dr r e^{-\frac{r^2}{2\sigma^2}} g(r/R, \omega) \int d\theta e^{\frac{rr_\star \cos(\theta)}{\sigma^2}} \\ &= 2\pi R^2 e^{-\rho_\star^2/2\bar{\sigma}^2} \int d\rho g(\rho, \omega) \rho e^{-\rho^2/2\bar{\sigma}^2} I_0\left(\frac{\rho\rho_\star}{\bar{\sigma}^2}\right) \\ &\propto \frac{1}{Norm(\rho_\star)} \int d\rho g(\rho, \omega) P(\rho, \rho_\star) ,\end{aligned}$$

where x_\star , r_\star , and ρ_\star are the position of the laser in Cartesian, cylindrical, and normalized cylindrical coordinates respectively, σ , is equal to half the beam waste, $\bar{\sigma} = \sigma/R$,

$I_0(x)$ is the modified Bessel function of the first kind, and $P(\rho, \rho_\star) = \rho e^{-\rho^2/2\bar{\sigma}^2} I_0\left(\frac{\rho\rho_\star}{\bar{\sigma}^2}\right)$ is the envelope function. A normalization constant, $Norm(\rho_\star)$, which is close to the maximum value of the envelope function was included to keep the widely varying values of the envelope function in a computationally reasonable range. The maximum of the envelope function is well approximated as

$$Norm(\rho_\star) = \begin{cases} .03 & , \rho_\star < .1 \\ P(\rho_\star, \rho_\star) & , \rho_\star \geq .1 \end{cases}$$

and is independent of ω . Thus, it should be calculated once per spectra as an overhead step.

The integral in $f(\omega, \rho_\star)$ can not be solved analytically. The strains fields alone are too complex, let alone the additional complexity added by the Lorentzian functions and the envelope function. Instead, $f(\omega, \rho_\star)$ must be integrated numerically at each frequency. Since the form of the strain distribution are different in the domains $[0, 1]$, $[1, \rho_0]$, $[\rho_0, \inf]$ it might seem simplest to treat $f(\omega, \rho_\star)$ as three distinct numeric integrals. Alternatively, one numerical integral could be used with piecewise functions. The fastest technique in Wolfram Mathematica 8.0, however, is to minimize the number of numeric integrals without introducing piecewise functions. The integral over the second domain can be scaled into the first domain using the change of variables

$\rho' = \frac{\rho-1}{\rho_0-1}$ giving

$$f(\omega, \rho_*) = \int_0^1 d\rho \left\{ \underbrace{g(\rho, \omega) P(\rho, \rho_*)}_{\text{inside}} + \underbrace{(\rho_0 - 1) g(\rho(\rho_0 - 1) + 1, \omega) P(\rho(\rho_0 - 1) + 1, \rho_*)}_{\text{outside}} \right\} + g(\infty, \omega) \underbrace{\int_{\rho_0}^{\infty} d\rho P(\rho, \rho_*)}_{\text{int3}(\rho_*)} .$$

Similar to $\text{Norm}(\rho_*)$, $\text{int3}(\rho_*)$ is independent of ω and thus needs to be calculated only once for each ρ_* as an overhead step. In this way three numeric integrals have been combined into two numeric integrals with one of which only calculated once per ρ_* .

Finally, it should be noted that since the numeric integration is the slowest step it should only be calculated once for each ω in the restrict data domain. Everything should then reference back to these values.

Fit linear coefficients

Having established the predicted shape of the spectra given the current values of the fitting parameters, the modeled spectra needs only to be scaled to match the measured spectra for a comparison to be made. The scaled modeled spectra is

$$\text{model}(\omega, \rho_0) = A f(\omega, \rho_*) + b ,$$

where A and b are linear coefficients which correspond to the signal amplitude and the background levels respectively. These will be independently fit to the measured

spectra on a spectra by spectra basis.

Fitting these linear coefficients is much simpler than the non-linear fitting parameters. Press *et al.* provide a simple matrix method to determine these coefficients (Press et al., 2007). It is based on the direct minimization of χ^2 by requiring that $\frac{\partial\chi^2}{\partial A} = \frac{\partial\chi^2}{\partial b} = 0$. This yields the system of equations

$$\begin{pmatrix} \sum_j \frac{data_j f(\omega_j, \rho_*)}{\sigma_j^2} \\ \sum_j \frac{data_j}{\sigma_j^2} \end{pmatrix} = \begin{pmatrix} \sum_j \frac{f(\omega_j, \rho_*)^2}{\sigma_j^2} & \sum_j \frac{f(\omega_j, \rho_*)}{\sigma_j^2} \\ \sum_j \frac{f(\omega_j, \rho_*)}{\sigma_j^2} & \sum_j \frac{1}{\sigma_j^2} \end{pmatrix} \begin{pmatrix} A \\ b \end{pmatrix},$$

where j runs over the data points in the restricted domain of the current spectra. Thus, fitting for the linear coefficients requires only the inversion of a 2 by 2 matrix.

More linear coefficients could be included to account for effects like the different signal intensity for suspended and supported graphene or to allow for different signal intensities from the G^+ and G^- peaks. However, there is no way to restrict the above procedure to give only positive values for the linear coefficients. Extra coefficients will often lead to non physical, negative coefficients that make the interpretation of the fit metrics more complicated. Two linear coefficients are enough to capture the complexity of the system while maintaining the simplicity of the fit.

Add χ^2 to global χ^2

Having calculated the best possible fit of the hypothetical spectra to the measured data, the goodness of fit can be quantified using the χ^2 metric. The calculated, not yet reduced by dividing the degrees of freedom, χ^2 of this spectra should be added to the global χ^2 which includes the χ^2 of each spectra in the data set for the current fitting parameters. In the off chance that the current fitting parameters are the best fit, it is also useful to save the linear coefficients so that they will not need to be

redetermined later.

Select best $\tilde{\chi}^2$ and estimate errors

After each spectra in the line scan is considered, the global χ^2 is reduced by dividing by the degrees of freedom. The number of degrees of freedom are equal to the number of fit data points, which was reduced when the fitting domain was restricted, minus the number of fitted parameters, including the 3 fitting parameters and the two linear coefficients per spectra. This global $\tilde{\chi}^2$ should be saved so that the $\tilde{\chi}^2$ space can be visualized and the minimum found. It is useful to keep track of the best $\tilde{\chi}^2$ and the corresponding linear coefficients so that the linear coefficients won't need to be refit afterward.

The sampled fitting parameter with the lowest $\tilde{\chi}^2$ best represent the data. However, these values are not necessarily the best fit. To ensure that these fitting parameters represent a global minimum and not just a local minimum the $\tilde{\chi}^2$ space should be visualized. If only the dimensionless friction was varied, this is done by plotting $\tilde{\chi}^2$ as a function of F . A global minimum indicates a best fit. If all three fitting parameters were varied the four dimensional $\tilde{\chi}^2$ space is more difficult to visualize. One useful way to reduce the dimensionality is to choose the unploted fitting parameters such that they minimized the $\tilde{\chi}^2$ for the plotted value. In this way, the plotted values represent the best case scenario. Figure C.4 shows a reduced dimensionality $\tilde{\chi}^2$ plot for a three fitting parameter fit. An obvious minimum in the $\tilde{\chi}^2$ near $f=0.14$ MPa is visible.

The departure from the minimum of $\tilde{\chi}^2$ can be used to estimate uncertainties in the fitting parameters. A very steep $\tilde{\chi}^2$ space corresponds to tightly defined fitting parameters. When viewing the $\tilde{\chi}^2$ space in reduced dimensionality, one needs to be

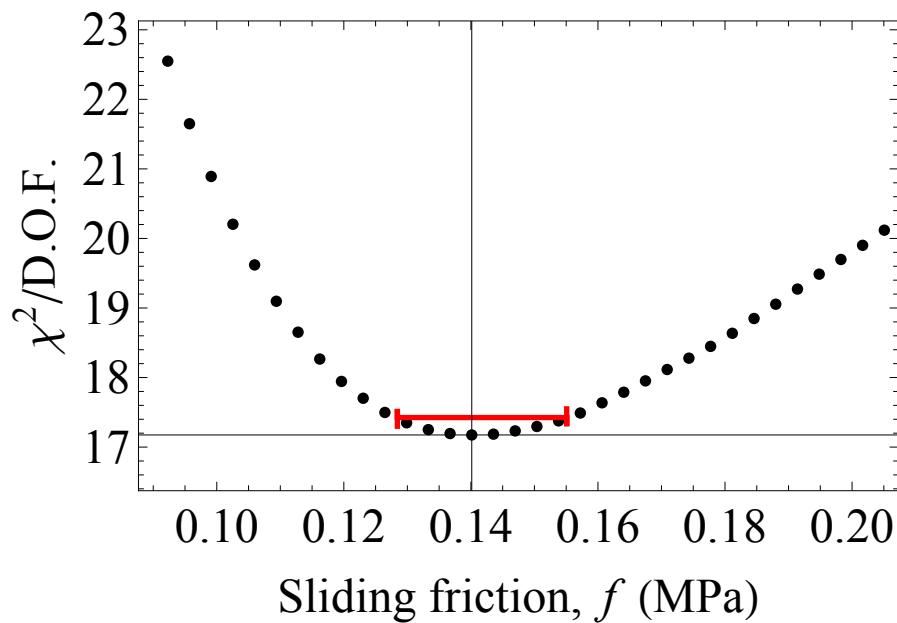


Figure C.4: $\tilde{\chi}^2$ as a function of the dimensionless friction for the $\sim 5 \mu\text{m}$ radius monolayer graphene sealed microchamber with 0.80 MPa of applied pressure. For each value of F , the γ and β pair which minimized $\tilde{\chi}^2$ was chosen. The black cross hair sits at the best fit and the red, side ways error bar indicates the uncertainty in the fit based on an increase in $\tilde{\chi}^2$ of 0.25.

careful that the dimensionality is reduced properly. Taking an arbitrary slice in $\tilde{\chi}^2$ space can yield artificially low uncertainties. Errors in the fitted parameters found by using the increase in χ^2 away from its minimum value (Press et al., 2007) are better than one part in one hundred, much smaller than we can claim to have achieved in our experiment. This discrepancy is due to an underestimation of our uncertainties which include only photon counting and ignore effects due to inhomogeneous doping, sample drift, and laser assisted deposition of dirt on the FLG. To better illustrate the relative uncertainties amongst different fitted friction values we use a 0.25 increase of χ^2 per degree of freedom above the best fit value to define confidence intervals. Using this method, the uncertainty in the fitted value in Figure C.4 is indicated by the red horizontal error bar.

Appendix D

Hopping energies in the Kekulé geometry

In this appendix the hopping energies between nearest neighbors in the Kekulé distorted lattice are calculated. This calculation closely follows the work of Chamon *et al.* (Chamon et al., 2013) in which they cleverly find an expression for the nearest neighbor bond length by working in complex coordinates where vectors are expressed as $a = \vec{a} \cdot \hat{x} + i\vec{a} \cdot \hat{y}$. They then use the change in bond length to determine the change in hopping energy using Equation 2.4

The displacements generated by the iTO phonon at the \mathbf{K} point (Equation 5.1) can be rewritten in complex notation as

$$\vec{u}_A(\vec{r}_A, t) = \begin{pmatrix} Re[c e^{i\vec{r}_A \cdot \mathbf{K}} e^{-i\omega t}] \\ Im[c e^{i\vec{r}_A \cdot \mathbf{K}} e^{-i\omega t}] \end{pmatrix} \rightarrow u_A(\vec{r}_A, t) = c e^{i\vec{r}_A \cdot \mathbf{K}} e^{-i\omega t}$$

$$\vec{u}_B(\vec{r}_B, t) = \begin{pmatrix} Re[c e^{i\vec{r}_B \cdot \mathbf{K}} e^{-i\omega t}] \\ -Im[c e^{i\vec{r}_B \cdot \mathbf{K}} e^{-i\omega t}] \end{pmatrix} \rightarrow u_B(\vec{r}_B, t) = c^* e^{-i\vec{r}_B \cdot \mathbf{K}} e^{i\omega t} .$$

The nearest neighbor vector in the Kekulé lattice connecting an A sub-lattice atom originally at position \vec{r} and a B sub-lattice atom originally at position $\vec{r} + \vec{\delta}_j$ is

$\vec{\delta}'_j(\vec{r}) = \vec{\delta}_j + \vec{u}_B(\vec{r} + \vec{\delta}_j) - \vec{u}_A(\vec{r})$. The magnitude of this vector is approximated as

$$\begin{aligned} |\vec{\delta}'_j(\vec{r})| &= \sqrt{\delta'_j(\vec{r}) \delta'_j(\vec{r})^*} \\ &\approx \sqrt{a^2 + \left\{ \delta_j^* \left(u_B(\vec{r} + \vec{\delta}_j) - u_A(\vec{r}) \right) + \text{c.c.} \right\}} \\ &\approx a + \frac{1}{2} \frac{1}{a} \left\{ \delta_j^* \left(u_B(\vec{r} + \vec{\delta}_j) - u_A(\vec{r}) \right) + \text{c.c.} \right\}. \end{aligned}$$

This can be simplified by using $\vec{\delta}_j = -iaz_j$ where $z_j = e^{i\mathbf{K}\cdot\vec{\delta}_j}$ and $z_j^2 = z_j^*$. The relative extension of the nearest neighbor vector is then

$$\begin{aligned} \frac{|\vec{\delta}'_j(\vec{r})| - a}{a} &= \frac{1}{2} \frac{1}{a^2} (iaz_j^*) \left(c^* e^{-i\vec{r}\cdot\mathbf{K}} z_j^* e^{i\omega t} - c e^{i\vec{r}\cdot\mathbf{K}} e^{-i\omega t} \right) \\ &\quad + \frac{1}{2} \frac{1}{a^2} (-iaz_j) \left(c e^{i\vec{r}\cdot\mathbf{K}} z_j e^{-i\omega t} - c^* e^{-i\vec{r}\cdot\mathbf{K}} e^{i\omega t} \right) \\ &= i \frac{1}{2a} \left\{ (z_j + (z_j^*)^2) c^* e^{-i\vec{r}\cdot\mathbf{K}} e^{i\omega t} - (z_j^* + z_j^2) c e^{i\vec{r}\cdot\mathbf{K}} e^{-i\omega t} \right\} \\ &= -i \frac{c}{a} z_j^* e^{i\vec{r}\cdot\mathbf{K}} e^{-i\omega t} + \text{c.c.} \end{aligned}$$

The use of the complex relationship between the δ_j and the z_j was paramount in deriving such a simple relationship.

The term $e^{i\vec{r}\cdot\mathbf{K}}$ can be rewritten to have the spatial frequency which connects the inequivalent \mathbf{K} points, $\mathbf{G} = \mathbf{K} - \mathbf{K}'$. This can be shown by taking advantage of the fact that \vec{r} is a lattice vector of the old lattice and that

$$\mathbf{K} - (\vec{b}_+ - \vec{b}_-) = \mathbf{K} - 3\mathbf{K} = -2\mathbf{K} = -(\mathbf{K} - \mathbf{K}') = -\mathbf{G}.$$

Thus, $e^{i\vec{r}\cdot\mathbf{K}} = e^{i\vec{r}\cdot(\mathbf{K}-\vec{b}_+-\vec{b}_-)} = e^{-i\vec{r}\cdot\mathbf{G}}$. In addition to making it clear that the Kekulé mode connects the \mathbf{K} points this also supports the simplification of \vec{r} . The spatial frequency \mathbf{G} is a reciprocal lattice vector of the Kekulé lattice equal to $\vec{B}_+ - \vec{B}_-$.

Thus, when specifying $\vec{r} = \vec{R}_l + \vec{r}_m$ there is no need to include \vec{R}_l . Instead we only specify \vec{r}_m , the position in the Kekulé unit cell of the A sub-lattice atom involved in the hopping.

The change in hopping energy resulting from the change in bond length can then be found using Equation 2.4. The resulting Kekulé altered hopping energies are given by

$$\delta t_{m,j} = \frac{1}{3} \Delta(t) e^{i\mathbf{K}\cdot\vec{\delta}_j} e^{i\mathbf{G}\cdot\vec{r}_m} + \text{c.c.}$$

with $\Delta(t) = -i3\beta t_0 \frac{c^*}{a} e^{i\omega t}$.

The interesting effects of these modulated hoppings are discussed in detail in Chapter 5.

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