

Title

Subtitle

by

Mikkel Metzsch Jensen

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Title

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Mikkel Metzsch Jensen

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Abstract

Abstract.

Acknowledgments

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List of symbols?

Maybe add list of symbols and where they are used like Trømborg.

Chapter 1

Background Theory and Method

Small introtext to motivate this chapter. What am I going to go over here.

1.1 Tribology - friction

Friction is a part of the wider field tribology which includes the study of friction, wear and lubrication between two surfaces in relative motion [1][p. 1]. In this thesis we will only concern ourselves with friction in cases without wear and without the use of lubrication. Dry friction (definition?)

Something about different scales of friction for the following sections.

1.1.1 Friction on a macroscopic scale - macroscale theories

1.1.1.1 Amontons' law

The work of Leonardo da Vinci (1452–1519), Guillaume Amontons (1663-705) and Charles de Coulomb (1736-1806) all contributed to what is commonly known as Amontons' law. In order to start and keep a solid block moving against a solid surface we must overcome certain frictional forces F_{extfric} . The static friction force F_s corresponds to the minimum tangential force required to initiate the sliding while the kinetic friction force F_k corresponds to the tangential force needed to sustain a steady sliding speed from that point on. This empirical law states that this frictional force, tangential to the sliding direction, is entirely independent of contact area and sliding velocity (at ordinary sliding velocities). Instead, it relies only on the normal force F_N acting perpendicular to the surface and the material specific friction coefficient μ as

$$F_{\text{fric}} = \mu F_N.$$

The friction coefficient is typically different for the cases of static μ_s and kinetic μ_k friction, usually with values lower than one and $\mu_s \geq \mu_k$ in all cases. [1][p. 6].

Statistical mechanics gives an simple overall explanation for this: A system at equilibrium has its kinetic energy uniformly distributed among all its degrees of freedom. When a macroscale object is sliding in a given direction it is clearly not in equilibrium since one of its degrees of freedom carries considerable more kinetic energy. Thus, the system will have a tendency to transfer the kinetic energy to the remaining degrees of freedom as heat while slowing down as a result. Hence friction is really the tendency of systems toward equilibrium energy equipartitioning among many interacting degrees of freedom. When friction occurs under an external force energy is reversed into the system in the form of frictional heat which will dissipate. [2]

Although Amontons' law is a common base for understanding friction on a macroscopic it has its limitations. For instance, it turns out that static friction is not time independent, but depends on the so-called contact history with increasing friction as the logarithm of time of stationary contact [3]. For the kinetic friction the independency of sliding velocity disappears at low velocities as thermal effects becomes important and for high velocities due to inertial effects. [1][pp. 5-6]. Additionally, this simple macroscale perspective disguises (better word) the underlying mechanism behind the friction phenomena arising from a more complicated contact interface. This includes taking the microscopic roughness into account together with surface chemistry. This more complex

perspective introduces new (...) as real contact area, contact stresses, surface adhesion which makes frictional properties dependent on sliding speed, temperature and environment in general [4]. In addition, the tribological systems take place across a broad range of time and length scales, from geological stratum layers involved in earthquakes [4] to microscopic atomistic processes, as in the gliding motion of nanocluster of a nanomotor [2].

Friction phenomena take place across a broad range of time and length scales, from microscopic atomistic processes, as in the gliding motion of a nanocluster or a nanomotor [1]

The conclusion is that the friction coefficient is not an intrinsic physical property [5].

1.1.2 Microscopic scale

Going beyond a macroscopic perspective to the microscopic scale we realise that most surfaces is in fact rough. The contact between two surfaces consist of numerous smaller contact point, so-called asperities, for which the friction between two such opposing surfaces involves interlocking of those asperities. Additionally, the adhesion between surfaces is taken to be a dominant factor, for which the chemical state of the surfaces has to be considered.

1.1.2.1 Surface roughness - Asperity theories

Asperity theories is based on the observation that microscopic rough surfaces with contacting asperities, each with a contact area of A_{asp} , will have a true contact area $\sum A_{\text{asp}}$ which is found to be much smaller than the apperent macroscopic area A_{macro} . The friction force was shown to be proportional (extra source on this) to this true contact area as

$$F_{\text{fric}} = \tau \sum A_{\text{asp}},$$

where τ is an effective shear strength of the contacting bodies. This is still compatible with Amontons' law for the cases where the real contact area depndt linearly on the applied normal force F_N .



Figure 1.1: Example figure shwoing asperity deformation as the load is incread from top to bottom. Source: [https://en.wikipedia.org/wiki/Asperity_\(materials_science\)](https://en.wikipedia.org/wiki/Asperity_(materials_science)) Put in reference list? Find better picture?

Thus, many studies have focused on the individual single asperity contact to reveal the relationship between the contact area and F_N (13-15 from [6]). By assuming perfectly smooth asperities with radii of curvature from nanometers to micrometres continuum mechanics can be used to predict the deformation of asperities as normal force is applied. A model for non-adhesive contact between homogenous, isotropic, linear elastic spheres was first developed by Hertz (17 [6]), which predicted $A_{\text{asp}} \propto F_N^{2/3}$. Later adhesion effects were included in a number of subsequent models, including Maugis-Dugdale theory (18 from [6]), which also predicts a sublinear relationship

between A_{asp} and f_N , leading to a similar sublinear relationship for $F_{\text{fric}}(F_N)$. (Address that this deviates with the macroscale).

What about multi asperity theory?

However, even though the successes of continuum mechanics there is no reason to believe that it will be capable of reproducing tribological behaviour at the nanometre length scale where the discreteness of atoms often has a direct effect on physical properties [5].

This approach enables the bottom-up derivation of the linear scaling laws of macroscopic friction with size, and their transition to the sublinear ones for incommensurate nanosized contacts. We can now understand that such transition takes place when the contact roughness becomes large compared to the range of interfacial interactions [162] [2].

1.1.3 Nanotribology - Atomic scale

On the smallest possible scale, atomic scale, the surfaces does not have structural asperities. Instead, atomic level friction is being modelled as a consequence of the rough potential of the atomic landscape and real contact area defined in terms of atomic bonding.

1.1.3.1 Frenkel-Kontorova (FK)

The FK (before FKT) model builds upon the time-honored PT model describing a point-like tip sliding over a space-periodic crystalline surface. The FK model extents this description to that of two 1D crystalline surfaces sliding on top of each other. The standard FK model consists of a 1D chain of N classical particles, representing atoms, interacting via a harmonic forces and moving in a sinusoidal potential as sketched in figure 1.2. The hamiltonian is

$$H = \sum_{i=1}^N \left[\frac{p_i^2}{2m} + \frac{1}{2}K(x_{i+1} - x_i - a_c)^2 + \frac{1}{2}U_0 \cos\left(\frac{2\pi x_i}{a_b}\right) \right], \quad (1.1)$$

where the first term $p_i^2/2m$ represents the kinetic energy with momentum p and mass m . The next term describes the harmonic interaction with elastic constant K , nearest neighbour distance $\Delta x = x_{i+1} - x_i$ with a corresponding equilibrium distance a_c . The final term represents the periodic substrate potential with amplitude U_0 and period a_b .

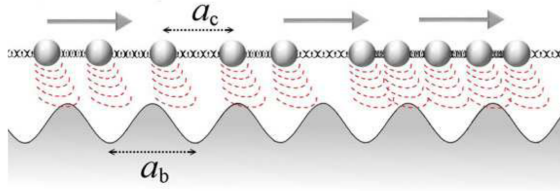


Figure 1. A sketch of the FK model, showing the two competing lengths: the average interparticle spacing and the lattice periodicity of the substrate.

Figure 1.2: Temporary figure from [2]

By applying an external force to the chain one can use the FK model to investigate static friction. The sliding properties are entirely governed by its topological excitations referred to as so-called kinks and anti-kinks

Commensurability We can describe the interactions of the model in a commensurate or incommensurate case. Let us begin with the commensurate case where the spacing of the atoms matches perfectly the substrate potential periodicity, i.e. $a_c = a_b$. The ground state (GS) is the configuration where each atom fits in one of the M minima in the substrate potential such that the coverage, $\Theta = N/M = a_b/a_c$, equals 1. By adding an atom we would excite the system with a so-called kink excitation (two atoms share the same potential well) while subtracting an atom results in an antikink excitation (one potential well becomes atomless). In order

to reach a local minimum the kink (antikink) will expand in space over a finite length such that the chain is imposed a local compression (expansion). When applying a tangential force to the chain it is much easier for a kink to move along the chain than it is for atoms since the activation energy ϵ_{PN} for a kink displacement is systematically smaller (often much smaller) than the potential barrier U_0 . Thus the motion of kinks (antikinks), i.e. the displacement of extra atoms or vacancies, is representing the fundamental mechanism for mass transport. In the ideal zero temperature case with additively (right term?) increase in force all atoms would be put into an accelerating motion as soon as the energy NU_0 is present. However, in reality any thermal excitation would excite the system with the creation of kink-antikink pairs that would travel down the chain. For a chain of finite length these often occur at the end of the chain running in opposite direction. As the kinks travel by the atoms is advanced by one atom spacing a_b along the force direction. This cascade of kink-antikink excitations is shown in figure 1.3

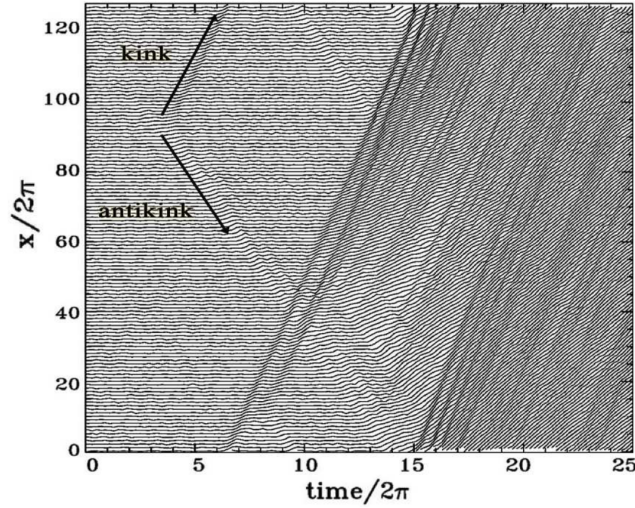


Figure 2. Time dependence of the atomic trajectories for the fully matched ($\theta = 1$) FK model at the (low-temperature) onset of depinning. Motion starts with the nucleation of a kink-antikink pair. The kink and the antikink depart in opposite directions cross the periodic boundary conditions, and collide quasielastically. A second kink-antikink pair forms in the wake of the initial kink. Further kink-antikink pairs are generated, with an avalanche-like increase of the kink-antikink concentration, eventually leading to a sliding state. Adapted from Ref. [2], Copyright (1997) by The American Physical Society.

Figure 1.3: Temporary figure from [2]

For the 2D case where an island is deposited on a surface, in our case the graphene sheet on the Si substrate, we generally expect the sliding to be initiated by kink-antikink pairs at the boundary.

For the case of incommensurability, i.e. $\theta = a_b/a_c$ is irrational, the GS is characterized by a sort of ‘staircase’ deformation. That is, the chain will exhibit regular periods of regions where the chain is slightly compressed (expanded) to match the substrate potential, separated by kinks (antikinks), where the increased stress is eventually realised through a localized expansion (compression) as illustrated in figure 1.4.



Figure 1.4: Temporary figure from [urlhttp://www.iop.kiev.ua/~obraun/myreprints/surveyfk.pdf](http://www.iop.kiev.ua/~obraun/myreprints/surveyfk.pdf) p. 14. Incommensurable case ($\theta = ?$) where atoms sit slightly closer than otherwise dictated by the substrate potential for which this regularly results in a kink here seen as the presence of two atoms close together in one of the potential wells.

The incommensurable FK model contains a critical elastic constant K_c , such that for $K > K_c$ (infinite chain, right?) the static friction F_s drops to zero (article says low-velocity kinetic friction as well which I do not get for no temperature?), making the chain able to slide at no energy cost. This can be explained by the fact that the displacement accounting in the incommensurable case will yield one atom climbing up a corrugation (introduce this word earlier) for every atom climbing down on. It turns out that this will exactly balance the energy making it non-resistant to sliding. Generally, incommensurability guarantees that the total energy (for $T = 0$) is independent of the relative position to the potential. However, when sliding freely a single atom will eventually occupy a maximum of the potential. When increasing the potential magnitude U_0 of softening the chain stiffness, lowering K , (the probability for a particle to occupy that position drops from a finite value to exactly zero, how can this work for zero temp?) this is no longer possible, which is described by the Aubry transition. The phase transition is marked by the critical elastic constant $K = K_c$ where the chain goes from free sliding to a pinned state with a nonzero static friction. K_c is a discontinuous function of the ratio θ , due to the reliance on irrational numbers for incommensurability (? get plot of this?). The minimal value $K_c \simeq 1.0291926$ in units $[2U_0(\pi/a_b)^2]$ is achieved for the golden-mean ratio $\theta = (1 + \sqrt{5})/2$ (compare to my model?). The pinning is provided despite translational invariance by the inaccessibility to move past the energy barrier which acts as a dynamical constraint. The Aubry transition can be investigated as a θ -order phase transition for which power laws can be defined for the order parameter (beyond the scope of this thesis).

The phenomenon of non-pinned configurations is named *superlubricity* in tribological context. Despite the misleading name this refers to the case where the static friction is zero while the kinetic friction is only reduced but nonzero.

Experimentally superlubricity has been studied for a graphene flake (single layer) sliding over a graphite surface (multiple layers) where the friction force was extremely weak (50 pN) for orientations between flake and substrate yielding incommensurability.

When considering the nonzero temperature case for the FK model one would have to add a thermostat such as the Langevin thermostat (see theory section ?). Thermal fluctuations can then overcome pinning effects even in fully commensurate cases. This also gives rise to a damping force proportional to sliding velocity as a means of modelling the dissipative force accounting.

By applying a finite driving force it is known that a pinned configuration will go through several first-order dynamical phase transitions as the system transfers from a pinned to a sliding state.

1. Properties depending on $\Theta = N/M = a_b/a_c$
2. Elastic constant K
3. Kinks anti kinks travelling
4. Importance of commensurability between lattice and potential
5. Aubry transition
6. Pinning and unpinning, stick slip
7. superlubricity (no stick slip, but still kinetic friction)
8. Adding a Langevin thermostat on top of this introduces temperature.
9. Phase transition?

1.1.3.2 FK extension: Frenkel-Kontorova-Tomlinson (FKT)

Several extensions has been provided for the FK model with modifications of the interactions or system dimensionality. Anharmonicity of the chain

1.1.3.3 Other stuff

At nanoscales things get a bit more unclear. SFM (explain) experiments have reported (copy sources 5, 6, 21 from [6]) where $F_f \propto F_N$ or even with these quantities being nearly independent of each other.

Physically relevant quantities, including the average friction force, the slider and the lubricant mean velocities, several correlation functions, and the heat flow can be evaluated numerically by carrying out suitable averages over the model dynamics of a sliding interface, as long as it is followed for a sufficiently long time. The modeling of friction must first of all address correctly ordinary equilibrium and near-equilibrium phenomena, where the fluctuation-dissipation theorem (Sec. 2) governs the smooth conversion of mechanical energy into heat, but most importantly it must also deal with inherently nonlinear dissipative phenomena such as instabilities, stick-slip, and all kinds of hysteretic response to external driving forces, characteristic of non-equilibrium dynamics. [2]

In several works by J. Fineberg's group [2–4] the transition from sticking to sliding is characterized by slip fronts propagating along the interface. [7][p. 2].

1.1.4 Multi scale models?

[2] p. 24.

1.1.5 Summary of expected frictional properties

1. Friction should decrease by increasing temperature.
2. We expect stick slip motion
3. What about dependence on normal force?
4. Dependence on contact area?
5. Dependence on speed?

- Different friction models on macro-and microscopic scale

Smooth kinetic friction generally increases with speed (velocity strengthening), but sometimes decreases with increasing speed in certain intervals [2].

the smallest force needed to set a slider in motion – is also dependent on the simulation time (a longer wait may lead to depinning when a short wait might not), and generally dependent on system size, often increasing with sub-linear scaling with the slider's contact area. To address this kind of behavior in MD simulations, it is often necessary to resort to scaling arguments in order to extrapolate the large-area static friction from small-size MD simulations [131, 140] [2].

In AFM and SFA experiments, stickslip and its associated characteristically high friction and mechanical hysteresis tend to transition into smooth sliding when the speed exceeds $\sim 1 \mu\text{m/s}$; in contrast, in MD modeling the same transition is observed in the $\sim 1 \text{ m/s}$ region. This 6-order-of-magnitude discrepancy in speed between experiments and simulations is well known and has been largely discussed [141–144] in connection with the effective mass distributions and spring-force constants, that are vastly different, and highly simplified in simulations. [2]

Concerning stick-slip friction, another problem is that, unlike simulations, real experiments contain mesoscale or macroscale component intrinsically involved in the mechanical instabilities of which stick-slip consists. Here the comforting observation is that stick-slip is nearly independent of speed, so that so long as a simulation is long enough to realize a sufficient number of slip events, the results may already be good enough [148] [2].

A serious aspect of stick-slip friction which MD simulation is unable to attack is ageing. The slip is a fast event, well described by MD, but sticking is a long waiting time, during which the frictional contact settles very

slowly. The longer the sticking time, the larger the static friction force necessary to cause the slip. Typically experiments show a logarithmic increase of static friction with time [150] [2].

Rate and state friction approaches, widely used in geophysics [151], describe phenomenologically frictional ageing, but a quantitative microscopic description is still lacking. Mechanisms invoked to account for contact ageing include chemical strengthening at the interface in nanoscale systems [152], and plastic creep phenomena in macroscopic systems [153]. [2].

See “Selected Results of MD Simulations” in [2] p. 24.

1.1.6 Graphene friction

Theory of friction experiment involving graphene.

Because of this frictional reduction, many studies indicate graphene as the thinnest solid-state lubricant and anti-wear coating [104–106]. (Current trends in the physics of nanoscale friction)

Accurate FFM measurements on few-layer graphene systems show that friction decreases by increasing graphene thickness from a single layer up to 4-5 layers, and then it approaches graphite values [97, 99, 101, 107, 108]. (Current trends in the physics of nanoscale friction)

1.2 Molecular Dynamics

Read [2][p. 18]

A promising compromise could possibly be provided by the so-called reactive potentials [120–122], capable of describing some chemical reactions, including interface wear with satisfactory computational efficiency in large-scale atomic simulations, compared to semi-empirical and first-principles approaches. [2]

- MD simulation (classical or ab initio)
- Basics of classical MD simulations: Integration and stuff
- Ab initio simulation (quantum mechanics, solving schrödinger)

1.2.1 Potentials

The choices of potentials used in the MD simulation is mainly based on the on [8] which have a somewhat similar MD friction simulation, the difference being that they impose a Si-tip on the graphene sheet supported by a Si-substrate where we impose drag the whole sheet upon the substrate. Nonetheless this serves as a good anchor for the methodology of the setup. The covalent bonds of C-C in graphene and Si-Si in the substrate is described by the Tersoff and Stillinger–Weber potentials, respectively. A typical 12-6 Lennard–Jones potential is used to describe the van der Waals adhesive interaction between graphene and the substrate.

1.2.1.1 General formulation of potentials...?

On a general note we can generalize the n-body potential as the expansion in orders of participating atoms as

$$E = \sum_i V_1(\mathbf{r}_i) + \sum_{\substack{i,j \\ i < j}} V_2(\mathbf{r}_i, \mathbf{r}_j) + \sum_{\substack{i,j,k \\ i < j < k}} V_3(\mathbf{r}_i, \mathbf{r}_j, \mathbf{r}_k) + \dots$$

where \mathbf{r}_n is the position of the n th particle and V_m is called an m -body potential [9]. The first one-body term corresponds to an external potential, followed by the two-body term, the three-body term and so on. The simplest model that includes particle interaction is the pair potential truncating the expansion after the two-body term. A general feature of the pair potentials is that they favor close-packed structures which is unsuited to describe covalent bonds that take more open structures. In particular, pair potentials are completely inapplicable to strongly co-valent systems such as semiconductors [9]. In order to accommodate the description of covalent bonds the natural step is thus to include the next step of the expansion, the three-body terms, as we will see for the modeling of the graphene sheet C-C bonds and the Silicon sheet Si-Si bonds. For the interaction between the sheet and the substrate we can employ a Lennard Jones pair potential describing the non-bonded van der Waals interaction.

1.2.2 Lennard Jones

This sections is based on [[10], [11], [12]].

The Lennard-Jones (LJ) model is probably one of the most famous pair potentials used in MD simulations. LJ models the potential energy between two non-bonding atoms based solely on interatomic distance r . The model accounts for attractive forces arising from dipole-dipole, dipole-induced dipole and London interactions, and repulsive forces that capture the hard core (is this safe to say?) of overlapping wave functions at small distances. Thus it is assumes neutrally charged atoms and was originally proposed for noble gases. The classical 12-6 version of the model (referring to the power law of the repulsive and attractive forces respectively) reads

$$E = 4\epsilon \left[\left(\frac{\sigma}{r} \right)^{12} - \left(\frac{\sigma}{r} \right)^6 \right], \quad r < r_c, \quad (1.2)$$

where r is the interatomic distance with cut-off r_c , ϵ is the depth of the potential well and σ the distance where the potential is zero. By solving for the potential minimum ($dE/dr = 0$) we find the equilibrium distance to be $r_0 = \sigma 2^{1/6}$. This makes for an even clearer interpretation of σ which effectively sets the equilibrium distance between atoms, i.e. the dividing line for which the net force is repulsive or attractive. While the LJ model in many ways is an oversimplified model that is insufficient in its description of ... (get source and concrete examples) it is commonly used as a model for intermaterial interactions (between moving object and substrate) in friction studies [[8], [13], [4]].

1.2.3 Stillinger weber

This section is based on [[14], [15]]

The stillinger weber potential takes the form of a three body potential

$$E = \sum_i \sum_{j>i} \phi_2(r_{ij}) + \sum_i \sum_{j \neq i} \sum_{k>j} \phi_3(r_{ij}, r_{ik}, \theta_{ijk}),$$

where r_{ij} denotes the distance between atom i and j and θ_{ijk} the angle between bond ij and jk . The summations are over all neighbours j and k of atom i within a cut-off distance $r = a\sigma$.

The two-body term ϕ_2 builds from the LJ model with the addition of an exponential cutoff term

$$\phi_2(r_{ij}) = A_{ij}\epsilon_{ij} \left[B_{ij} \left(\frac{\sigma_{ij}}{r_{ij}} \right)^{p_{ij}} - \left(\frac{\sigma_{ij}}{r_{ij}} \right)^{q_{ij}} \right] \exp\left(\frac{\sigma_{ij}}{r_{ij} - a_{ij}\sigma_{ij}} \right). \quad (1.3)$$

The model parameters A , ϵ , B , σ , p , q and a comes with i, j indices to indicate that these parameters should be specified for each unique pair of atom types. However, in our case we will only provide a single value for each model parameter as we are exclusively dealing with Si-Si bonds. We see that the first term in eq. (1.3) is reminiscent of the LJ model in eq. (1.2) while the last term effectively drives the potential to zero at $r = a\sigma$, which is thus the chosen cut-off distance for the potential evaluation. With the model parameters for the Si-Si modelling (see table 1.1) the cut-off becomes ~ 3.8 Å.

The three body term includes an angle dependency as

$$\phi_3(r_{ij}, r_{ik}, \theta_{ijk}) = \lambda_{ijk} \epsilon_{ijk} \left[\cos \theta_{ijk} - \cos \theta_{0,ijk} \right]^2 \exp\left(\frac{\gamma_{ij}\sigma_{ij}}{r_{ij} - a_{ij}\sigma_{ij}} \right) \exp\left(\frac{\gamma_{ik}\sigma_{ik}}{r_{ik} - a_{ik}\sigma_{ik}} \right), \quad (1.4)$$

where $\theta_{0,ijk}$ is the equilibrium angle. The first term of eq. (1.4) includes an angle dependency analog to a harmonic oscillator based on a cosine angle distance from the equilibrium angle. The final two terms act again as a cut-off function by driving the potential to zero at $r_{ij} = a_{ij}\sigma_{ij}$ and $r_{ik} = a_{ik}\sigma_{ik}$ respectively.

The parameters used for the Si-Si bond modeling is displayed in table 1.1 along with an interpretation of each model parameter.

Table 1.1: Parameters for the stilliner weber potential used for intermolecular interactions in the silicon substrate.

Parameter	Value	Description
ϵ	2.1683	Individual depth of the potential well for each atom type pair/tiplets.
σ	2.0951	Distance for which the individual pair interactions has zero potential (analog to the LJ model).
a	1.80	The individual cut-off distance for each atom type pair.
λ	21.0	The overall depth of the three-body potential well.
γ	1.20	The shape of the three-body cut-off terms.
$\cos(\theta_0)$	-1/3	Cosine of equilibrium angle.
A	7.049556277	The overall depth of the two-body potential well.
B	0.6022245584	Scales the repulsion part of the two-body term.
p	4.0	The power dependency for the repulsion part of the two-body term.
q	0.0	The power dependency for the attraction part of the two-body term.
tol	0.0	LAMMPS: Option to define a different cut-off than the theoretical of $r = a\sigma$. $tol = 0$ refers to the theoretical being used.

1.2.4 Tersoff

This section is based on [[16], [9]].

The tersoff potential abandon the idea of a general n -body form and attempts instead to build the model on a more physics informed approach; The more neighbours an atom has the weaker the bonds will be. Thus it introduces the bond order (bond strength), that is environment specific and decrease with increasing bond coordination (number of neighbours for a given atom). The potential energy is taken to have the form

$$E = \sum_i E_i = \frac{1}{2} \sum_{i \neq j} V_{ij},$$

$$V_{ij} = f_C(r_{ij}) [f_R(r_{ij}) + b_{ij} f_A(r_{ij})],$$

where the total potential energy is decomposed into a bond energy V_{ij} . The indices i and j run over the atoms of the system with r_{ij} denoting the distance between atom i and j . Notice that the sum includes all combinations of i, j where $i \neq j$ meaning that the same bond is double counted which is the reason for the additional factor $1/2$. The reasoning behind comes from the asymmetry of the bond order $b_{ij} \neq b_{ji}$ leading to a $V_{ij} \neq V_{ji}$. The bond energy is composed of a repulsive term f_R , arising from overlapping wave functions, and an attractive term f_A associated with bonding. f_C is simply a smooth cut-off function to increase computational efficiency. b_{ij} represent the bond order, i.e. the strength of the bonds, which depends inversely on the number of bonds, the bond angles (θ_{ijk}) and optionally the relative bonds lengths (r_{ij}, r_{jk}). Notice that an additional cut-off term a_{ij} was originally multiplied to f_R as a way of including terms that limit the range of the interactions to the first neighbour shell. These kind of limitations is already included in b_{ij} for the attractive term f_A but is often omitted for the repulsive term f_R , and we do so to by setting $a_{ij} = 1$.

The cut-off function f_C goes from 1 to 0 over a small interval range $R \pm D$ as

$$f_C(r) = \begin{cases} 1 & r < R - D \\ \frac{1}{2} - \frac{1}{2} \sin\left(\frac{\pi}{2} \frac{r-R}{D}\right) & R - D < r < R + D, \\ 0 & r > R + D \end{cases}$$

which is continuous and differentiable for all r . R is usually chosen to include only the first neighbour shell. The repulsive and attractive terms f_R and f_A is modelled as an exponential function, similar to a morse potential,

$$\begin{aligned} f_R(r) &= A \exp(-\lambda_1 r), \\ f_A(r) &= -B \exp(-\lambda_2 r). \end{aligned}$$

The novel feature of the model lies in modeling of the bond order b_{ij} which includes three-body interactions by summing over a third atom $k \neq i, j$ within the cut-off $r_{ik} < R + D$ as shown in the following.

$$b_{ij} = (1 + \beta^n \zeta_{ij}^n)^{-\frac{1}{2n}} \quad (1.5)$$

$$\zeta_{ij} = \sum_{k \neq i, j} f_C(r_{ik}) g(\theta_{ijk}(r_{ij}, r_{ik})) \exp(\lambda_3^m (r_{ij} - r_{ik})^m) \quad (1.6)$$

$$g(\theta) = \gamma_{ijk} \left(1 + \frac{c^2}{d^2} - \frac{c^2}{[d^2 + (\cos \theta - \cos \theta_0)^2]} \right). \quad (1.7)$$

In eq. (1.7) $\zeta_{i,j}$ is an effective coordination and $g(\theta)$ captures angle dependency as it is minimized at the equilibrium angle $\theta = \theta_0$.

The parameters used to model the graphene C-C bonds is summarized in table 1.2

Table 1.2: Parameters for the tersoff potential used for intermolecular interactions in the graphene sheet

Parameter	Value	Description
m	3.0	Default (not used since $\lambda_3 = 0$)
γ	1.0	...
λ_3	0.0 \AA^{-1}	...
c	3.8049×10^4	Strength of the angular effect
d	4.3484	Determines the “sharpness” of the angular dependency
$\cos(\theta_0)$	-0.57058	Cosine of the equilibrium angle
n	0.72751	Power law exponent for the bond order dependency
β	1.5724×10^{-7}	...
λ_2	2.2119 \AA^{-1}	Decay of repulsion potential term
B	346.74 eV	Attractive potential term minimum at core ($r_{ij} = 0$).
R	1.95 \AA	Center distance for cut-off
D	0.15 \AA	Thickness of cut-off layers
λ_1	3.4879 \AA^{-1}	Decay of repulsion potential term
A	1393.6 eV	Repulsion potential term at core ($r_{ij} = 0$)

1.2.5 LAMMPS

1.2.6 Integration

Having defined a system of particles governed by interatomic potentials we need to move the system forward in time. By solving Newtons equations of motion we effectively do so by sampling the microcanonical ensemble characterized by a constant number of particles N , volume V and energy E , hence denoted NVE. Newtons equations of motion read

$$m_i \frac{d^2 \mathbf{r}_i}{dt^2} = \mathbf{F}_i = -\nabla U_i \quad (1.8)$$

where i is the particle index and m_i its mass, $\mathbf{r}_i = (x_i, y_i, z_i)$ the position, t is time, $\nabla_i = (\frac{\partial}{\partial x_i}, \frac{\partial}{\partial y_i}, \frac{\partial}{\partial z_i})$ and U_i the potential energy. In system the potential energy is a function of the particle positions of nearby particles depending on the specific potential in use. Since the forces defined by the potentials is conservative we expect

the energy of the solution to be conserved. We redefine eq.(1.8) in terms of two coupled first order differential equations

$$\dot{\mathbf{v}}_i(t) = \frac{\mathbf{F}}{m_i}, \quad \dot{\mathbf{r}}_i(t) = \mathbf{v}_i(t), \quad (1.9)$$

where $\dot{x} = dx/dt$ (Newton's notation) and $\mathbf{v} = (v_x, v_y, v_z)$ is velocity. Numerically we can solve the coupled equations (eq.(1.9)) by integrating over discrete timesteps. That is, we discretize the solution into temporal steps $t_k = t_0 + k \cdot \Delta t$ with time-step Δt .

1.2.6.1 Velocity Verlet

A common algorithm to integrate Newtons equation of motion (as formulated in eq.(1.9)) is the *velocity verlet*. We can derive the algorithm by the use of Taylor expansions. We begin by expanding the next-step position vector $\mathbf{r}_i(t + \Delta t)$ at time t

$$\mathbf{r}_i(t + \Delta t) = \mathbf{r}_i(t) + \dot{\mathbf{r}}_i(t)\Delta t + \frac{\ddot{\mathbf{r}}_i(t)}{2}\Delta t^2 + \mathcal{O}(\Delta t^3), \quad (1.10)$$

where $\ddot{\mathbf{r}} = d^2\mathbf{r}/dt^2$ and Δt^n is simply the relaxed notation for $(\Delta t)^n$. Similar we take the expansions of the next-step velocity vector $\mathbf{v}_i(t + \Delta t)$ at time t

$$\mathbf{v}_i(t + \Delta t) = \mathbf{v}_i(t) + \dot{\mathbf{v}}_i(t)\Delta t + \frac{\ddot{\mathbf{v}}_i(t)}{2}\Delta t^2 + \mathcal{O}(\Delta t^3). \quad (1.11)$$

Finally, by taking the expansion of $\dot{\mathbf{v}}_i(t + \Delta t)$ we can eliminate the $\ddot{\mathbf{v}}_i$ -term in eq. (1.11) and simplify it as shown in the following.

$$\begin{aligned} \dot{\mathbf{v}}_i(t + \Delta t) &= \dot{\mathbf{v}}_i(t) + \ddot{\mathbf{v}}_i(t)\Delta t + \mathcal{O}(\Delta t^2) \\ \frac{\ddot{\mathbf{v}}_i(t)}{2}\Delta t^2 &= \frac{\Delta t}{2} \left(\dot{\mathbf{v}}(t + \Delta t) - \dot{\mathbf{v}}_i(t) \right) + \mathcal{O}(\Delta t^3) \\ &\Downarrow \\ \mathbf{v}_i(t + \Delta t) &= \mathbf{v}_i(t) + \dot{\mathbf{v}}_i(t)\Delta t + \frac{\Delta t}{2} \left(\dot{\mathbf{v}}_i(t + \Delta t) - \dot{\mathbf{v}}_i(t) \right) + \mathcal{O}(\Delta t^3) \\ &= \mathbf{v}_i(t) + \frac{\Delta t}{2} \left(\dot{\mathbf{v}}_i(t) + \dot{\mathbf{v}}_i(t + \Delta t) \right) + \mathcal{O}(\Delta t^3). \end{aligned} \quad (1.12)$$

By combining eq. (1.10) and eq. (1.12) and using Newton's second equation $\dot{\mathbf{v}} = \mathbf{F}_i(t)/m_i$ and $\mathbf{v} = \dot{\mathbf{r}}$ we arrive at the final scheme

$$\begin{aligned} \mathbf{r}_i(t + \Delta t) &= \mathbf{r}_i(t) + \mathbf{v}_i(t)\Delta t + \frac{\mathbf{F}_i(t)}{2m_i}\Delta t^2 + \mathcal{O}(\Delta t^3), \\ \mathbf{v}_i(t + \Delta t) &= \mathbf{v}_i(t) + \frac{\mathbf{F}_i(t) + \mathbf{F}_i(t + \Delta t)}{2m_i}\Delta t + \mathcal{O}(\Delta t^3). \end{aligned}$$

The scheme will give a local error of order Δt^3 corresponding to a global error of Δt^2 . One of the most popular ways to implement this numerically is as stated in the following steps.

1. Calculate $v_{k+\frac{1}{2}} = v_k + \frac{F_k}{2m}\Delta t$.
2. Calculate $r_{k+1} = r_k + v_{k+\frac{1}{2}}\Delta t$.
3. Evaluate the force $F_{k+1} = F(r_{k+1})$.
4. Calculate $v_{k+1} = v_{k+\frac{1}{2}} + \frac{F_{k+1}}{2m}\Delta t$.

1.2.7 Thermostats

As we already mentioned above in Sec. 2, any kind of sliding friction involves mechanical work, some of which is then transformed into heat (the rest going into structural transformations, wear, etc.). The heat is then transported away by phonons (and electrons in the case of metallic sliders) and eventually dissipated to the environment [2].

Likewise all excitations generated in the simulations should be allowed to propagate in the system and disperse in the bulk of both sheet and substrate. Due to small simulation size these are likely to reflect back and ‘‘pile up’’ unphysically. Thus in order to avoid continuous heating and attain a steady state the (Joule) heat must be removed at a steady state. This is very the viscous damping of the Langevin equations enter the picture. It can be difficult to set the value γ for the magnitude of this damping. The unphysical introduction of heat sink can be mitigated by some modifications he mention, which is kind of next level I guess.

Langevin thermostat

In order to control the temperature of the system we introduce the so-called Langevin thermostat. This is a stochastic thermostat that modifies Newton's equation of motion such that solution lies in the canonical ensemble characterized by a constant number of particles N , constant volume V and constant temperature T , hence denoted NVT. The canonical ensemble system is represented by the finite system being in contact with an infinite heat bath of temperature T . The NVT ensemble is equivalent to sampling a system in thermodynamic equilibrium where the weight of each microscopic state is given by the Boltzmann factor $\exp[-E/(k_B T)]$.

The Langevin equation is the modified version of Newton's second law for a Brownian particle. A Brownian particle is a small particle suspended in liquid, e.g. pollen or dust, named after Robert Brown (1773–1858) who was the first to observe its jittery motion. The Langevin equation describes this motion as the combination of viscous drag force $-\gamma \mathbf{v}$, where γ is a positive friction coefficient and \mathbf{v} the velocity vector, and a random fluctuation force \mathbf{R} . The Langevin equation reads

$$m \frac{d\mathbf{v}}{dt} = -\gamma \mathbf{v} + \mathbf{R} \quad (1.13)$$

where m is the particle mass. This effectively describes the particle of interest, the Brownian particle, as being suspended in a sea of smaller particles. The collision with these smaller particles is modelled by the drag force and the fluctuation force. We notice that if the fluctuation force is excluded eq. (1.13) becomes

$$m \frac{d\mathbf{v}}{dt} = -\gamma \mathbf{v} \quad \Rightarrow \quad \mathbf{v}_i(t) = v(0) e^{-\frac{\gamma t}{m}},$$

where the solution shows that the Brownian particle will come to a complete stop after a long time $\mathbf{v}_i(t \rightarrow \infty) \rightarrow \mathbf{0}$. This is in violation with the equipartition theorem

$$\frac{1}{2} m \langle v^2 \rangle_{eq} = \frac{k_B T}{2},$$

and hence the fluctuation force is necessary to obtain the correct equilibrium.

The following calculations are done in one dimension in order to simplify the notation. We describe the statistical nature of the collisions as a sum of independent momentum transfers

$$\Delta P = \sum_i^N \delta p_i$$

where ΔP denotes the change of momentum after N momentum transfers δp_i from the environment to the Brownian particle. We assume the first and second moments $\langle \delta p \rangle = 0$ and $\langle \delta p^2 \rangle = \sigma^2$. When N is large the central limit theorem states that the random variable ΔP has a Gaussian distribution with $\langle P \rangle = 0$ and $\langle \Delta P^2 \rangle = N \sigma^2$. If we consider the momentum change ΔP over a discrete time Δt , where the number of collisions is proportional to time $N \propto \Delta t$, the corresponding fluctuation force $R = \Delta P / \Delta t$ will have a variance

$$\langle R^2 \rangle = \frac{\langle \Delta P^2 \rangle}{\Delta t^2} = \frac{N \sigma^2}{\Delta t^2} \propto \frac{1}{\Delta t}.$$

In a computer simulation we need to pick a random force $R(t)$ from a Gaussian distribution every time-step Δt . These forces will not be correlated as long as Δt is larger than the correlation time of the forces from the molecules which we will assume for this model (I think there exist corrections for this to refer to here). With this assumption we can write the correlation function as

$$\langle R(t)R(0) \rangle = \begin{cases} \frac{a}{\Delta t}, & |\Delta t| < \Delta t/2 \\ 0, & |\Delta t| > \Delta t/2, \end{cases} \quad (1.14)$$

where a is some strength of (...?). In the limit $\Delta t \rightarrow 0$ the correlation function becomes

$$\langle R(t)R(0) \rangle = a\delta(t), \quad (1.15)$$

where δ denotes the dirac delta function. This is valid for all spatial coordinates which will all be independent of each other. Since both the drag force and the fluctuation force originate from the molecular fluid, where the drag force $-\alpha\mathbf{v}$ is velocity dependent it is reasonable to assume that fluctuation force is independent of velocity, i.e. $\langle R_i v_j \rangle = 0$ for all cartesian indices i and j .

In the following we will attempt justify the Langevin equation (why it is like it is) and determine the relationship between the drag coefficient γ and the random force R .

From the Langevin equation eq. (1.13) we can compute the velocity autocorrelation function (Move to appendix?). We do this in one dimension for simplicity. We begin by multiplying by $(e^{\gamma t/m})/m$

$$\dot{v}(t)e^{\gamma t/m} + \frac{\gamma}{m}v(t)e^{\frac{\gamma t}{m}} = \frac{F}{m}e^{\frac{\gamma t}{m}},$$

and integrate from $t = -\infty$. By the use of integration by parts on the latter term on the left hand side we calculate the velocity

$$\begin{aligned} \int_{-\infty}^t dt' \dot{v}(t')e^{\frac{\gamma t'}{m}} + \frac{\gamma}{m}v(t)e^{\frac{\gamma t}{m}} &= \int_{-\infty}^t dt' e^{\frac{\gamma t'}{m}} \frac{F(t')}{m} \\ \int_{-\infty}^t dt' \dot{v}(t')e^{\frac{\gamma t'}{m}} + \left(\left[v(t')e^{\frac{\gamma t'}{m}} \right]_{-\infty}^t - \int_{-\infty}^t dt' \dot{v}(t')e^{\frac{\gamma t'}{m}} \right) &= \int_{-\infty}^t dt' e^{\frac{\gamma t'}{m}} \frac{F(t')}{m} \\ v(t) &= \int_{-\infty}^t dt' e^{\frac{-\gamma(t-t')}{m}} \frac{F(t')}{m}, \end{aligned}$$

where $e^{\frac{-\gamma t}{m}}$ plays the role of a response function. We can then calculate the autocorrelation

$$\begin{aligned} \langle v(t)v(0) \rangle &= \int_{-\infty}^t dt_1 \int_{-\infty}^0 dt_2 e^{\frac{t-t_1-t_2}{m}} \frac{\langle F(t_1)F(t_2) \rangle}{m^2} \\ &= \int_{-\infty}^t dt_1 \int_{-\infty}^0 dt_2 e^{\frac{t-t_1-t_2}{m}} \frac{a\delta(t_1-t_2)}{m^2} \\ &= \int_{-\infty}^0 dt_2 e^{\frac{t-2t_2}{m}} \frac{a}{m^2} = \frac{a}{2m\gamma} e^{-\frac{\gamma t}{m}}, \end{aligned}$$

where we used eq. (1.15) and the fact that the integration commutes with the average (we are allowed to flip the order). By comparing this with the equipartition theorem we get

$$\begin{aligned} \frac{1}{2}m\langle v^2 \rangle &= \frac{k_B T}{2} \\ \frac{1}{2}m\langle v(0)v(0) \rangle &= \frac{a}{4\gamma} = \frac{k_B T}{2} \\ a &= 2\gamma k_B T \end{aligned}$$

We notice the appearance of γ meaning that the magnitude of the fluctuations increase both with friction and temperature. Further we can integrate the velocity over time to get displacement $x(t)$ and show that the variance (show this? In appendix maybe?) is

$$\langle x^2(t) \rangle = \frac{2k_B T}{\gamma} \left(t - \frac{m}{\gamma} (1 - e^{-\gamma t/m}) \right),$$

where for $t \gg m/\gamma$ only the t -term survives yielding

$$\langle x^2(t) \rangle = 2k_B T t / \gamma.$$

In 1D, the diffusion constant D is related to the variance as $\langle x^2 \rangle = 2Dt$, meaning that this represents the Einstein relation $D = \mu k_B T$ with the mobility $\mu = 1/\gamma$.

when $t \ll m/\gamma$ we use the Taylor expansion $1 - e^{-x} \approx x - x^2/2$ for $x \ll 1$ to get

$$\langle x^2(t) \rangle = \frac{k_B T}{m} t^2$$

which exactly matches the thermal velocity

$$v_{\text{th}} \frac{\langle x^2(t) \rangle}{t^2} = \frac{k_B T}{m}$$

which follows from the equipartition theorem. The finite correlation time γ/m hence describe the crossover from the ballistic regime $\sqrt{\langle x^2(t) \rangle} \propto t$ to the diffusive regime $\sqrt{\langle x^2(t) \rangle} \propto \sqrt{t}$.

Introduce the fluctuation-dissipation theorem concept earlier since this is a motivation for the Langevin equation.

1.2.7.1 Implementing Langevin

The implementation of the Langevin equation into LAMMPS follows [17] and updates the force vector for each particle as

$$\begin{aligned} \mathbf{F} &= \mathbf{F}_c + \mathbf{F}_f + \mathbf{F}_r \\ &= -\nabla U - \gamma m \mathbf{v} + \sqrt{\frac{2k_B T m \gamma}{\Delta t}} \mathbf{h}(t) \end{aligned} \quad (1.16)$$

where \mathbf{F}_c is the conservative force computed via the usual inter-particle interactions described by the potential U , \mathbf{F}_f is the drag force and \mathbf{F}_r is the random fluctuation force where \mathbf{h} is a random vector drawn from a normal distribution with zero mean and unit variance. Notice that this generalized description of the Langevin equation deviates from the presentation in eq. (1.13) since we have added the conservative force \mathbf{F}_c , but also by the appearance of the mass in both the drag force and the fluctuation force due to the introduction of damping. It is beyond our scope to comprehend this. However, the fact that Δt now appears in the denominator for the random force variance $2k_B T m \gamma / \Delta t$ is due to the fact that we have discretized time. This is in agreement with the formulation in eq. (1.14). By applying eq. (1.16) we get the refined velocity Verlet scheme

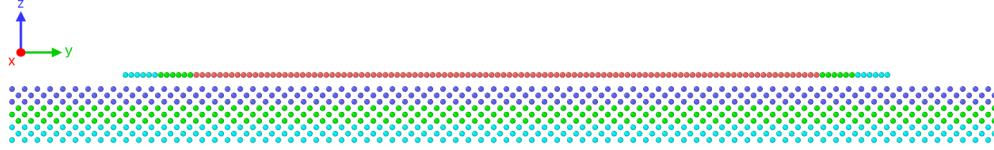
$$\begin{aligned} \mathbf{v}_i(t + \Delta t/2) &= \mathbf{v}_i(t) - \frac{\Delta t}{2} \left(\frac{\nabla_i U(t)}{m_i} + \gamma \mathbf{v}_i \right) + \sqrt{\frac{k_B T \gamma \Delta t}{2m_i}} \mathbf{h}_i \\ \mathbf{r}_i(t + \Delta t) &= \mathbf{r}_i(t) + \mathbf{v}_i(t + \Delta t/2) \Delta t \\ \mathbf{v}_i(t + \Delta t) &= \mathbf{v}_i(t + \Delta t/2) - \frac{\Delta t}{2} \left(\frac{\nabla_i U(t + \Delta t)}{m_i} + \gamma \mathbf{v}_i(t + \Delta t/2) \right) + \sqrt{\frac{k_B T \gamma \Delta t}{2m_i}} \mathbf{h}_i \end{aligned}$$

with new random vector \mathbf{h}_i for each particle and each update. Notice however, that LAMMPS only apply this scheme to the particle groups with the thermostat on.

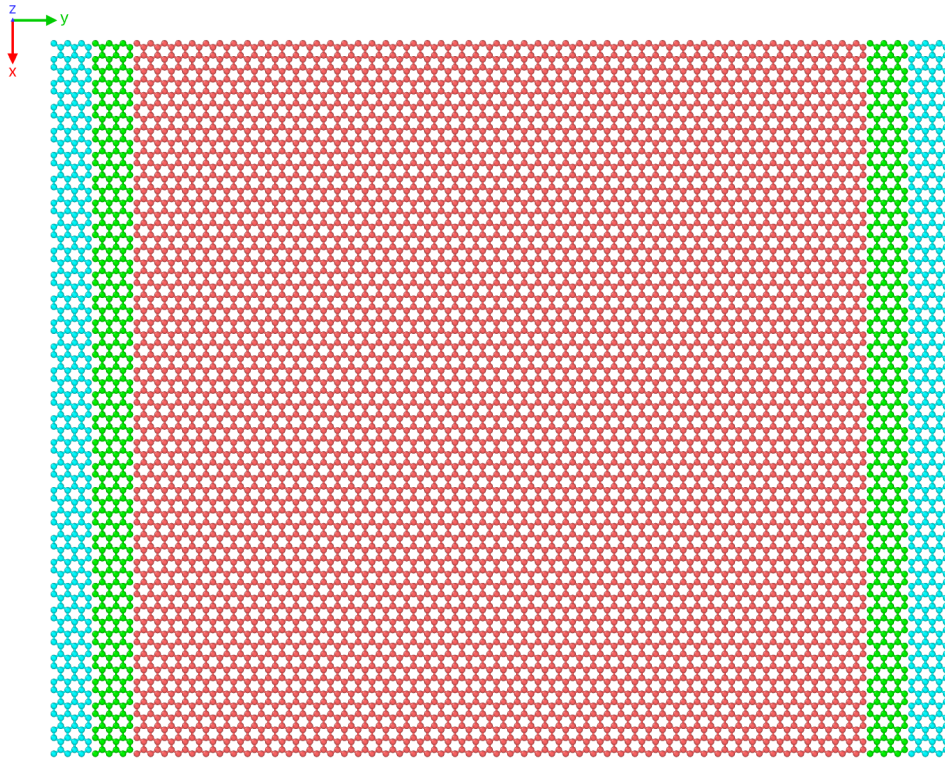
1.3 Defining the system

1.3.1 Groups: Sheet, pullblocks and substrate

Include figure of system to point out thermo layers and freeze layers.



(a) Side view



(b) Top view

Figure 1.5: System. 27456 atoms in total: 7272 with thermostat (orange), 7272 is locked (light blue) and the remaining 12912 just with NVE. (Get better colors)

Table 1.3: System atom count and region division.

Region	Total	Sub region	Sub total	NVE	NVT	Locked
Sheet	7800	Inner sheet	6360	6360	0	0
		Pull blocks	1440	0	720	720
Substrate	19656	Upper	6552	6552	0	0
		Middle	6552	0	6552	0
		Bottom	6552	0	0	6552
All	27456			12912	7272	7272

The sheet dimensions is

Table 1.4: Sheet dimensions

Group	x, y -dim	dim [\AA]	Area [\AA^2]
Full sheet	$x_S \times y_S$	$130.029 \times 163.219 \text{ \AA}$	21,223.203
Inner sheet	$x_S \times 81.40 \% y_S$	$130.029 \times 132.853 \text{ \AA}$	17,274.743
Pull blocks	$2 \times x_S \times 9.30 \% y_S$	$2 \times 130.029 \times 15.183 \text{ \AA}$	$2 \times 1,974.230$

1.3.1.1 Pressure reference

Find place to put this.

In order to relate the magntidue of the normal force in our friciton measurement we will use the pressure as a reference. We will use the pressure underneath a stiletto shoe as a worst case for human pressure execution underneath the shoes. From (source 1) it is reported that the diameter of a stiletto heeled shoe can be less than 1 cm. Hence a 80 kg man¹ standing on one stiletto heels (with all the weight on the heel) will result in a pressure corresponding diameter of

$$P = \frac{F}{A} = \frac{mg}{r^2\pi} = \frac{80 \text{ kg} \cdot 9.8 \frac{\text{m}}{\text{s}^2}}{\left(\frac{1 \times 10^{-2} \text{ m}}{2}\right)^2\pi} = 9.98 \text{ MPa}$$

While this is in itself a spectacular realization that is often used in introduction physics courses (source 2) to demonstrate the rather extreme pressure under a stiletto heel (greater than the foot of an elephant) (how many Atmos) this serves as a reasonable upperbound for human executed pressure. With a full sheet area of $\sim 21 \times 10^3 \text{ \AA}^2$ we can achieve a similar pressure of $\sim 10 \text{ MPa}$ with a normal force

$$F_N = 10 \text{ MPa} \cdot 21 \times 10^{-17} \text{ m}^2 = 2.10 \text{ nN}$$

Of course this pressure might be insufficient for various industrial purposes, but with no specific procedure in mind this serves as a decent reference point. Notice that if we consider a human foot with ares 113 cm^2 the pressure drops to a mere 70 kPa corresponding to $\sim 0.01 \text{ nN}$.

1.3.2 Creating sheets

We are going to create a 2D sheet graphene sheet.

1.3.2.1 Graphene

Graphene is a single layer of carbon atom, graphite is the bulk, arranged in a hexagonal lattice structure. We can describe the 2D crystal structure in terms of its primitive lattice vector and a basis. That is we populate each lattice site by the given basis and translate it to fill the whole plane by any linear combination of the lattice vectors

$$\mathbf{T}_{mn} = m\mathbf{a}_1 + n\mathbf{a}_2, \quad m, n \in \mathbb{N}.$$

For graphene we have the primitive lattice vectors

$$\mathbf{a}_1 = a \left(\frac{\sqrt{3}}{2}, -\frac{1}{2} \right), \quad \mathbf{a}_2 = a \left(\frac{\sqrt{3}}{2}, \frac{1}{2} \right), \quad |\mathbf{a}_1| = |\mathbf{a}_2| = 2.46 \text{ \AA}.$$

Notice that we deliberately excluded the third coordinate as we only consider a single graphene layer on not the bulk graphite consisting of multiple layers stacked on top of each other. The basis is

$$\left\{ (0, 0), \frac{a}{2} \left(\frac{1}{\sqrt{3}}, 1 \right) \right\}$$

¹Yes, a man can certainly wear stilleto heals.

It turns out that the spacing between atoms is equal for all pairs with an interatomic distance

$$\left| \frac{a}{2} \left(\frac{1}{\sqrt{3}}, 1 \right) \right| \approx 1.42 \text{ \AA}.$$

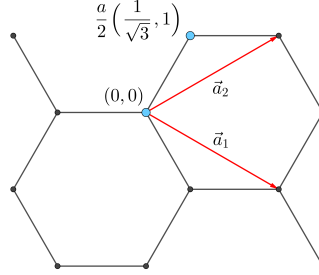


Figure 1.6: Graphene crystal structure with basis.

1.3.2.2 Indexing

In order to define the cut patterns applied to the graphene sheet we must define an indexing system. We must ensure that this gives a unique description of the atoms as we eventually want to pass a binary matrix, containing 0 for removed atom and 1 for present atom, that uniquely describes the sheet. We do this by letting the x-coordinate point to zigzag chains and the y-coordinate to the position along that chain. This is illustrated in figure 1.7. Other solutions might naturally involve the lattice vectors, but as these only can be used to translate to similar basis atoms a unfortunate duality is introduced as ones need to include the basis atom of choice into the indexing system. With the current system we notice that locality is somewhat preserved. That is, atom (i, j) is in the proximity of $\{(i+1, j), (i-1, j), (i, j+1), (i, j-1)\}$, but only three of them is categorized as nearest neighbours due to the hexagonal structure of the lattice. While $(i, j \pm 1)$ is always nearest neighbours the neighbour in the x-direction flip sides with incrementing y-coordinate. That is the nearest neighbours (NN) is decided as

$$\begin{aligned} j \text{ is even} &\rightarrow \text{NN} = \{(i+1, j), (i, j+1), (i, j-1)\}, \\ j \text{ is odd} &\rightarrow \text{NN} = \{(i-1, j), (i, j+1), (i, j-1)\}. \end{aligned}$$

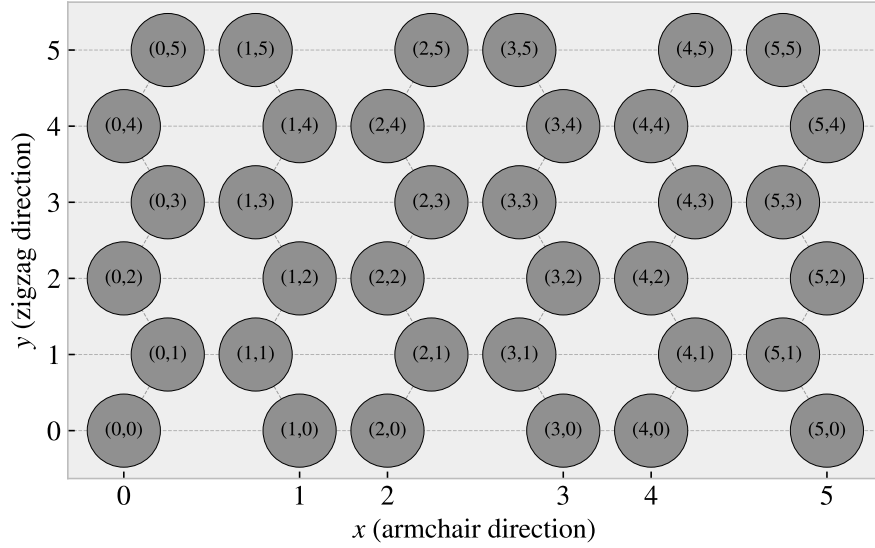


Figure 1.7: Graphene atom indexing

1.3.2.3 Removing atoms

As a mean to ease the formulation of cut patterns we introduce pseudo center element in each gap of the hexagonal honeycombs, see figure 1.8.

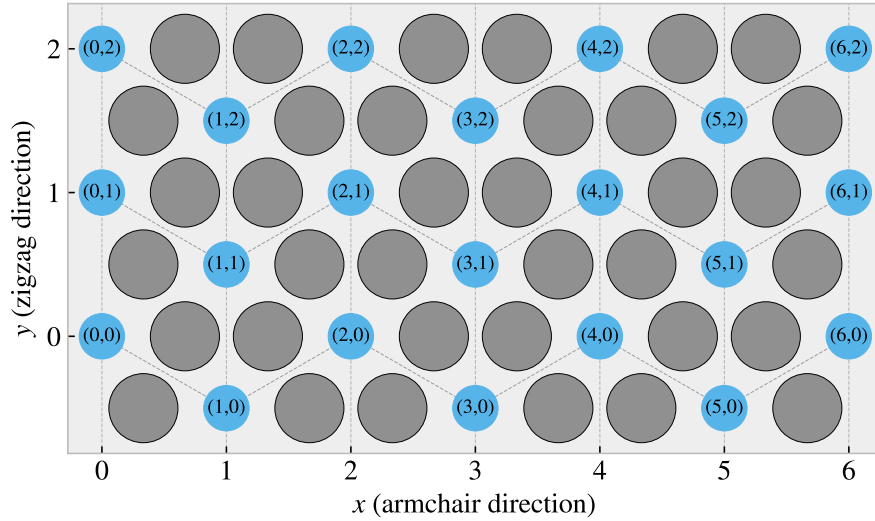


Figure 1.8: Graphene center indexing

Similar to the case of the indexing for the carbon atoms themselves the nearest neighbour center elements alternate with position, this time along the x-coordinate. Each center element has six nearest neighbours, in clock wise direction we can denote them: “up”, “upper right”, “lower right”, “down”, “lower left”, “upper left”. The “up” and “down” is always accessed as $(i, j \pm 1)$, but for even i the $(i + 1, j)$ index corresponds to the “lower right” neighbour while for odd i this corresponds to the “upper right” neighbour. This shifting applies for all left or right neighbours and the full neighbour list is illustrated in figure 1.9.

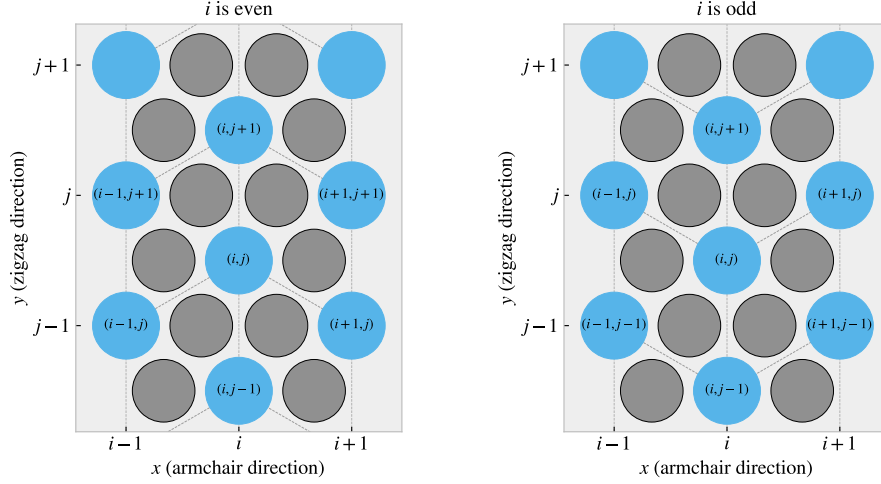


Figure 1.9: Graphene center elements directions

We define a cut pattern by connecting center elements into connected paths. As we walk element to element we remove atoms according to one of two rules

1. Remove intersection atoms: We remove the pair of atoms placed directly in the path we are walking. That is, when jumping to the “up” center element we remove the two upper atoms located in the local hexagon of atoms. This method is sensitive to the order of the center elements in the path.
2. Remove all surrounding atoms: We simply remove all atoms in the local hexagon surrounding each center element. This method is independent of the ordering of center elements in the path.

We notice that removing atoms using either of these rules will not guarantee a unique cut pattern. Rule 1 is the more sensitive to paths but we realize that, for an even i , we will remove the same five atoms following either of the following paths.

$$\begin{aligned}
 (i, j) &\rightarrow \underbrace{(i+1, j+1)}_{\text{upper right}} \rightarrow \underbrace{(i, j+1)}_{\text{up}} \rightarrow \underbrace{(i+1, j+2)}_{\text{upperright + up}} \rightarrow \underbrace{(i+1, j+1)}_{\text{upper right}} \\
 (i, j) &\rightarrow \underbrace{(i+1, j+1)}_{\text{upper right}} \rightarrow \underbrace{(i+1, j+2)}_{\text{upperright + up}} \rightarrow \underbrace{(i, j+1)}_{\text{up}}
 \end{aligned}$$

For rule 2 it is even more obvious that different paths can result in the same atoms being removed. This is the reason that we needed to define an indexing system for the atom position itself even though that all cuts generated manually will use the center element path as reference.

Illustrate some delete path?

1.3.3 Pull blocks

1.3.4 Kirigami inspired cut out patterns

1.3.4.1 Pop-up pattern

1.3.4.2 Honeycomb

1.3.4.3 Random walk

1.4 Fourier Transform (light)

Find out where to put this if necessary.

Fourier transform is a technique where we transform a function $f(t)$ of time to a function $F(k)$ of frequency. The Forward Fourier Transform is done as

$$F(k) = \int_{-\infty}^{\infty} f(t)e^{-2\pi i k x} dx$$

For any complex function $F(k)$ we can decompose it into magnitude $A(k)$ and phase $\phi(k)$

$$F(k) = A(k)e^{i\phi(k)}$$

Hence when performing a Forward Fourier transform on a time series we can determine the amplitude and phase as a function of frequency as

$$A(k) = |F(k)|^2, \quad \phi(k) = \Im \ln F(k)$$

1.4.1 Real life experimental procedures

From Introduction to Tribology, Second Edition, p. 526:

The surface force apparatus (SFA), the scanning tunneling microscopes (STM), and atomic force and friction force microscopes (AFM and FFM) are widely used in nanotribological and nanomechanics studies.

- Real life procedures to mimic in computation, for instance Atomic Force Microscopy (AFM) for friction measurements.
- Available technology for test of my findings if successful (possibilities for making the nano machine)

1.5 Machine Learning (ML)

- Feed forward fully connected
- CNN
- GAN (encoder + decoder)
- Genetic algorithm
- Using machine learning for inverse designs partly eliminate the black box problem. When a design is produced we can test it, and if it works we not rely on machine learning connections to verify it's relevance.
- However, using explainations techniques such as maybe t-SNE, Deep dream, LRP, Shapley values and linearizations, we can try to understand why the AI chose as it did. This can lead to an increased understanding of each design feature. Again this is not dependent on the complex network of the network as this can be tested and verified independently of the network.

1.5.1 Feed forward network / Neural networks

1.5.2 CNN for image recognition

1.5.3 GAN (encoder + deoder)

1.5.4 Inverse desing using machine learning

1.5.5 Prediction explanation

1.5.5.1 Shapley

1.5.5.2 Lineariations

1.5.5.3 LRP

1.5.5.4 t-SNE

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