

**FIRST PRINCIPLE CALCULATIONS OF DEFECT STRUCTURES
IN ZINC OXIDE**

By

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An undergraduate thesis submitted in partial fulfillment of
the requirements for the degree of

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FIRST PRINCIPLE CALCULATIONS OF DEFECT STRUCTURES IN ZINC OXIDE

Abstract

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March 2020

: Roland V. Sarmago

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Dedication

This dissertation/thesis is dedicated to my mother and father who
provided both emotional and financial support

Chapter One

INTRODUCTION

1.1 Purpose and Motivation

Describe the importance of defects in ZnO

1.2 Objectives

Study the mechanisms of different defects in ZnO

1.3 Outline

This is an example of how to cite [**Prades2009**]

Chapter Two

Review of Related Literature

2.1 Semiconductors

2.1.1 Properties

2.1.2 Applications of Semiconductors

2.1.3 Defects in Semiconductors

2.2 Zinc Oxide

describe ZnO in broad perspective

2.2.1 Crystal Structure

Consider different phases

2.2.2 Crystallographic Directions and Planes

2.2.3 Brillouin Zone Symmetry

2.2.4 Photoluminescence Properties

2.2.5 Defects

Chapter Three

THEORETICAL FRAMEWORK

3.1 Electronic Structure

The problem of electronic structure methods begins with the attempt to solve the general non-relativistic time-independent Schrödinger equation given as

$$\hat{\mathcal{H}}\Psi = E\Psi \tag{3.1}$$

where $\hat{\mathcal{H}}$ is the Hamiltonian operator for a system of electrons, Ψ is the electronic wavefunction and E is the energy of the system. Consider a single electron in three dimensional system, the Schrödinger equation can be expressed as

$$\hat{\mathcal{H}}\Psi_n = -\frac{\hbar^2}{2m} \left(\frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} + \frac{\partial^2}{\partial z^2} \right) \Psi_n + V\Psi_n = \epsilon_n\Psi_n \tag{3.2}$$

where m is the mass of electron, V is the effective potential energy and ϵ_n is the energy of electron in the orbital. The term orbital denotes the solution of the Schrödinger equation for a system of only one electron. This will be useful in later sections because this will allow to distinguish between the exact quantum state of a system of N interacting electrons from the approximate quantum state of N electrons in N orbitals, where each orbital is a solution to one-electron wavefunction in (3.2). If V is zero for the case of free electrons (i.e. non-interacting), then the orbital model is exact.

Since electrons are restricted by the potential inside the atom, the simplest way of solving (3.2) is by considering an infinite potential well. The electrons are confined inside a cube of length L where the potential V inside is zero and infinite at outside must satisfy the boundary condition

$$\Psi_n(L_x, L_y, L_z) = 0 \quad (3.3)$$

where L_x, L_y, L_z can be either 0 or L . The solution will have a sine dependence

$$\Psi_n(x, y, z) = \sqrt{\left(\frac{2}{L}\right)^3} \sin\left(\frac{n_x\pi}{L}x\right) \sin\left(\frac{n_y\pi}{L}y\right) \sin\left(\frac{n_z\pi}{L}z\right) \quad (3.4)$$

where n_x, n_y, n_z are integer quantum states. Provided that $k_i = n_i\pi/L$ where $i = x, y, \text{ or } z$; then the energy dispersion relation can be expressed as

$$\epsilon_k = \frac{\hbar^2}{2m}(k_x^2 + k_y^2 + k_z^2) = \frac{\hbar^2}{2m}k^2 \propto k^2 \quad (3.5)$$

Note that energy levels are discretized by the quantum states which arises from imposing the boundary conditions.

3.1.1 Electronic Band structure

Inside the crystal lattice, the periodic arrangement of atoms or ions causes the potential to be periodic which eventually gives rise to the formation of energy bands. The wavefunction Ψ will become periodic in space with a period L and must obey the Born-von Karman boundary condition [1]

$$\Psi_k(x, y, z) = \Psi_k(x + L, y, z) \quad (3.6)$$

and similarly for the y and z coordinates. It can be shown that wavefunctions satisfying (3.2) and (3.6) are the Bloch form of a travelling plane wave

$$\Psi_k(\vec{\mathbf{r}}) = u_k(\vec{\mathbf{r}}) \exp\left(i\vec{\mathbf{k}} \cdot \vec{\mathbf{r}}\right) \quad (3.7)$$

where $u_k(\vec{r})$ has the period of the crystal lattice with $u_k(\vec{r}) = u_k(\vec{r} + \vec{R})$. Here \vec{R} is the translation vector which can be simply thought as the periodicity expressed as vector. The Bloch expression can be written as

$$\begin{aligned}\Psi_k(\vec{r} + \vec{R}) &= u_k(\vec{r} + \vec{R}) \exp\left(i\vec{k} \cdot (\vec{r} + \vec{R})\right) \\ \Psi_k(\vec{r} + \vec{R}) &= u_k(\vec{r}) \exp\left(i\vec{k} \cdot \vec{r}\right) \exp\left(i\vec{k} \cdot \vec{R}\right) \\ \Psi_k(\vec{r} + \vec{R}) &= \Psi_k(\vec{r}) \exp\left(i\vec{k} \cdot \vec{R}\right)\end{aligned}\tag{3.8}$$

Notice that the wavefunction differs from the plane wave of free electrons only by a periodic modulation given by the new phase factor. This means that the electrons in the crystal lattice are treated as perturbed weakly by the periodic potential of the ion cores.

3.1.1.1 Band structure of free electron

A special case of periodicity is where the potential is set to zero, which is applicable for the free electrons. The wavefunction will be a plane wave

$$\Psi_k(\vec{r}) = \exp\left(i\vec{k} \cdot \vec{r}\right)\tag{3.9}$$

that represents travelling wave with a momentum $\vec{p} = \hbar\vec{k}$. The energy dispersion relation is still given by (3.5) but this time the allowed energy values are distributed essentially from zero to infinity. Figure 3.1 shows the parabolic dependence of energy with the wavevector k . Since the system is periodic in real space, it must be true for the reciprocal space, in this case by $2\pi/a$ where a is some lattice constant. Figure 3.1a shows the extended zone scheme where there are no restrictions on the values of wavevector \vec{k} . When wavevectors are outside the first Brillouin zone (BZ), they can be translated back to the first zone by subtracting a suitable reciprocal lattice vector. In mathematical sense [2]

$$\vec{k} + \vec{G} = \vec{k}'\tag{3.10}$$

where \vec{k}' is the unrestricted wavevector, \vec{k} is in the first Brillouin zone, and \vec{G} is the translational reciprocal lattice vector. The energy dispersion relation can always be written as

$$\begin{aligned}\epsilon(k_x, k_y, k_z) &= \frac{\hbar^2}{2m}(\vec{k} + \vec{G})^2 \\ &= \frac{\hbar^2}{2m}[(k_x + G_x)^2 + (k_y + G_y)^2 + (k_z + G_z)^2]\end{aligned}\quad (3.11)$$

Figure 3.1b shows the reduced zone scheme where the bands are folded into the first BZ by applying (3.10). Any energy state beyond the first BZ is the same to a state inside the first BZ with a different band index n .

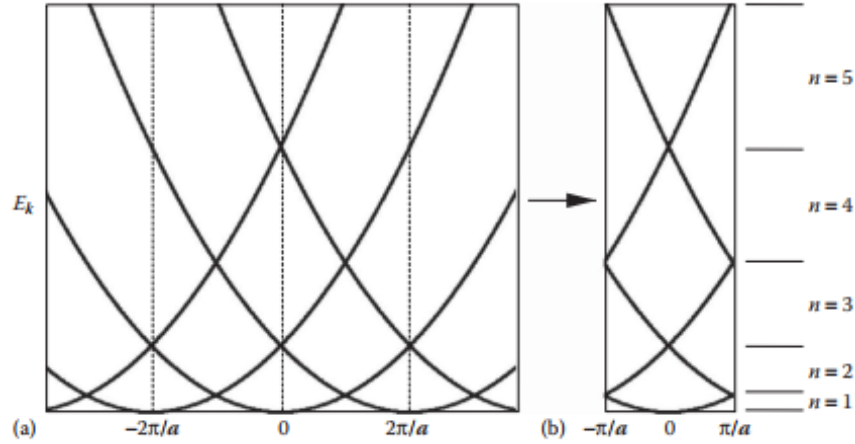


Figure 3.1 Free electron band structure

3.1.1.2 Band structure of electrons in solids

When atoms are very far from each other with no interaction, each electron occupies specific discrete orbitals such as 1s, 2p, 3d, etc. When they are brought closer enough, the outermost (valence) electrons interact with each other and will result in the energy level splitting. The innermost (core) electrons remain as they are, since they are closer to the nuclei and bounded by a deep potential well. For a solid containing a large N atoms, there will be N orbitals (i.e. N 3d-orbitals) trying to occupy the same energy level. Pauli's exclusion principle will prevent this from happening, hence what happens is there will be splitting of the energy

level that are closely spaced and this will eventually form a continuous band of energy levels. Figure 3.2 summarizes the evolution of energy levels as the atoms are brought together.

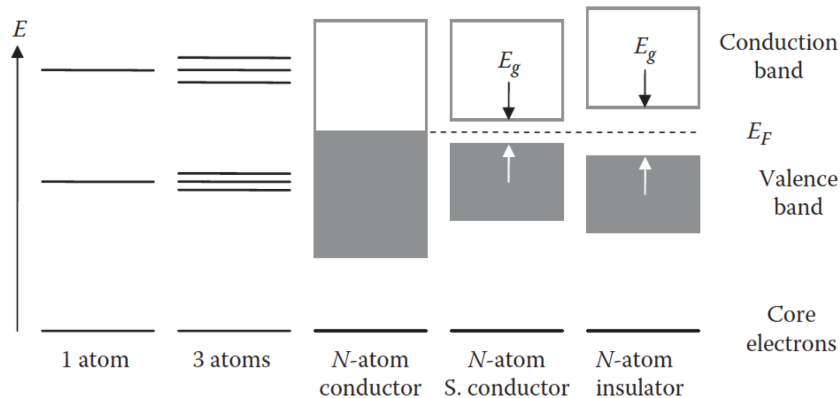


Figure 3.2 Formation of bands and band gaps when isolated atoms are brought closer together

Another interesting property of band structure is the formation of energy band gaps. This happens when the valence electrons interact with the periodic potential of the nuclei. Assuming a weak periodic potential, most of the band structure will not change very much, except possibly at the Brillouin zone boundaries with a wavevector of $\vec{k} = n\pi/a$. The orbitals with the wavevector at zone boundaries, chosen to be at high symmetry points, follow the Bragg diffraction condition and thus are diffracted. The valence electrons are scattered (or reflected) at the zone boundary in which the wavefunction is made up of equal plane waves travelling from the left and from the right. The wavefunction becomes a standing wave that resembles more of those bound states. Hence, there will be a forbidden region where travelling waves are not allowed. If sufficient energy is provided to the electron, they can overcome the binding potential.

The band gap is generally referred to the energy difference between the top of valence band, Valence band maximum (VBM), and the bottom of the conduction band, Conduction band minimum (CBM). If VBM and CBM coincide with each other, the material is said to be a conductor. Electrons can easily occupy the conduction band without any excitation, hence electrons are highly mobile that will lead to high current. For band gaps with a value

comparable to the quantity $k_B T$, where k_B is the Boltzmann constant and T is the absolute temperature near room temperature, then the material is semiconductor. If band gap is much larger than $k_B T$, then the material is insulator. However, this criterion is very loose because there are materials with large band gaps such as ZnO , $SrIn_2O_4$, that are categorized as semiconductors. These materials are generally called wide-band gap semiconductors. If the VBM and CBM are located in the same wavevector k , then the gap is direct. Otherwise, it is indirect.

3.1.2 Density of States

Another useful quantity in describing the electronic structure is the density of states (DOS). In general, the density of states can be defined as [3]

$$D(\epsilon) = 2 \sum_n \sum_k \delta(\epsilon - \epsilon_n(k)) \quad (3.12)$$

where for each band index n , the sum is over all allowed values of k lying inside the first Brillouin zone. The factor 2 comes from the allowed values of the spin quantum number for each allowed value of k . In the limit of large crystal, the k points are very close together, and the sum can be replaced by an integral. Since each allowed states will take up a volume of $(\Delta k)^3 = \pi^3/V$ where V is the volume of the solid in real space, it is convenient to write (3.12) as

$$D(\epsilon) = 2 \frac{V}{\pi^3} \sum_n \sum_k \delta(\epsilon - \epsilon_n(k)) (\Delta k)^3 \quad (3.13)$$

for in the limit of $V \rightarrow \infty$, $\Delta k \rightarrow 0$, it becomes

$$\lim_{V \rightarrow \infty} \frac{1}{V} D(\epsilon) = \frac{2}{\pi^3} \sum_n \int \delta(\epsilon - \epsilon_n(k)) d^3k \quad (3.14)$$

Usually, the total DOS is set to be the number of states per unit energy per unit volume.

The DOS can be projected in terms of the orbital contribution of each atoms. This can be expanded in a complete orthonormal basis as [4]

$$D(\epsilon) = \sum_i D_i(\epsilon) \quad (3.15)$$

$$= \sum_i \sum_n \int \langle \psi_n | \alpha | \psi_n \rangle \delta(\epsilon - \epsilon_n(k)) d^3k \quad (3.16)$$

where $D_i(\epsilon)$ is the projected density of states (PDOS) of orbital i with state α .

3.2 Many-body Quantum Mechanics

Despite the simplicity of Schrödinger equation in (3.1), solving it is a formidable task when dealing with many-electron systems. Analytical solutions to this equation only exist for the very simplest systems (i.e. hydrogenic atoms). Solving beyond '2 particle' system (electron and nucleus) is already intractable. In addition, solid state systems typically contains more than hundreds of particles, resulting in hundreds of simultaneous equations. Even the use of computational methods relies on a number of approximations just to make computations feasible enough. Hence, this section will discuss various levels of approximations without neglecting the parameter-free of first-principles calculations.

3.2.1 Time Independent Schrödinger Equation

The exact many-particle Hamiltonian consist of five operators which can be expressed as

$$\hat{\mathcal{H}} = \hat{\mathcal{T}}_n + \hat{\mathcal{T}}_e + \hat{\mathcal{V}}_{en} + \hat{\mathcal{V}}_{ee} + \hat{\mathcal{V}}_{nn} \quad (3.17)$$

where the $\hat{\mathcal{T}}$ and $\hat{\mathcal{V}}$ refer to kinetic energy and potential energy, respectively, and the labels e and n denotes the electronic and nuclear coordinates and their derivatives, respectively.

This equation can be expanded as

$$\begin{aligned} \hat{\mathcal{H}} = & -\frac{\hbar^2}{2} \sum_I \frac{\nabla_{\vec{R}_I}^2}{M_I} - \frac{\hbar^2}{2} \sum_i \frac{\nabla_{\vec{r}_i}^2}{m_e} \\ & - \frac{1}{4\pi\epsilon_0} \sum_{I,j} \frac{e^2 Z_I}{|\vec{R}_I - \vec{r}_j|} + \frac{1}{8\pi\epsilon_0} \sum_{i \neq j} \frac{e^2}{|\vec{r}_i - \vec{r}_j|} + \frac{1}{8\pi\epsilon_0} \sum_{I \neq J} \frac{e^2 Z_I Z_J}{|\vec{R}_I - \vec{R}_J|} \end{aligned} \quad (3.18)$$

where M_I is the mass of the I th nuclei (or usually ions) with charge Z_I located at site \vec{R}_I , and electrons have mass m_e located at site \vec{r} . The first and second terms are the kinetic energy of the atomic nuclei and electrons, respectively. The last three terms describe the Coulomb interaction between electrons and nuclei, between electrons and other electrons, and between nuclei and other nuclei.

3.2.2 Simplifying Assumptions

Solving (3.18) exactly is very impractical and not worth the effort. Hence, we resort to approximations in order to find acceptable eigenstates.

The first level of approximation is the Born-Oppenheimer approximation or the Adiabatic approximation. It begins with the observation that the mass of nuclei is much larger compared to the electron, as such one can assume that electrons moving in a potential much faster than the nuclei and that the nuclei can be treated as fixed or 'frozen' with respect to motion. As a consequence, the nuclear kinetic energy will be zero and the nuclear interaction with the electron cloud can be treated as an external parameter. Hence, the first term in (3.18) will vanish and the last term reduces to a constant which can be neglected. The third term will become the external potential. The Hamiltonian reduces to

$$\hat{\mathcal{H}} = \hat{\mathcal{T}} + \hat{\mathcal{V}} + \hat{\mathcal{V}}_{ext} \quad (3.19)$$

and using Hartree atomic units $\hbar = m_e = e = 4\pi/\epsilon_0 = 1$ for simplicity

$$\hat{\mathcal{H}} = -\frac{1}{2} \sum_i \nabla_{\vec{r}_i}^2 + \frac{1}{2} \sum_{i \neq j} \frac{1}{|\vec{r}_i - \vec{r}_j|} + \sum_{i,I} V_I(|\vec{R}_I - \vec{r}_i|) \quad (3.20)$$

3.2.3 Hartree Method

Since the second term in (3.20) includes electron-electron interaction which is difficult to evaluate, Hartree (1928) had proposed a simplified model where he treated each electrons to be independent and interacts with others in an averaged way [5]. This implies that each

electron does not recognize others as single entities but rather as a mean Coulomb field. The total energy will be sum of N numbers of one-electron energies

$$E = E_1 + E_2 + \cdots + E_N \quad (3.21)$$

then, the N -electron wavefunction can be approximated as a product of one-electron wavefunctions

$$\Psi = \Psi_1 \times \Psi_2 \times \cdots \times \Psi_N \quad (3.22)$$

Hartree model successfully predicts the ground-state energy of Hydrogen atom to be around -13.6 eV. However, for other systems, Hartree model produced crude estimations because it does not take into account the quantum mechanical effects such as antisymmetry principle and the Pauli's exclusion principle. Moreover, the model does not include the exchange and correlation energies of every interacting electrons in the actual systems.

3.2.4 Hartree-Fock Method

Due to the limitations of Hartree Model, Fock (1930) has taken into account the antisymmetric property of electron wavefunctions [6]. Pauli's exclusion principle posits that no two fermions can occupy the same quantum state because the wavefunction is antisymmetric upon particle exchange [7]. The many-electron wavefunction will be expressed in terms of Slater determinant [8]

$$\Psi = \frac{1}{\sqrt{N!}} \begin{vmatrix} \Psi_1(\vec{r}_1) & \Psi_2(\vec{r}_1) & \cdots & \Psi_N(\vec{r}_1) \\ \Psi_1(\vec{r}_2) & \Psi_2(\vec{r}_2) & \cdots & \Psi_N(\vec{r}_2) \\ \vdots & \vdots & \vdots & \vdots \\ \Psi_1(\vec{r}_N) & \Psi_2(\vec{r}_N) & \cdots & \Psi_N(\vec{r}_N) \end{vmatrix} \quad (3.23)$$

3.3 Density Functional Theory

3.3.1 Electron Density

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3.3.2.2 Second HK Theorem

3.3.3 Kohn Sham (KS) Formalism

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3.5.3 Hubbard U Correction

insert the symmetry points in IBZ.

Chapter Four

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This is sample text

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4.3.2.3 Irreducible Brillouin Zone

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4.3.3.1 Monkhorst-Pack method

4.3.3.2 Gamma Point Sampling

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5.4.7 bandstructure calculation

5.4.8 dos calculation

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APPENDIX

Appendix A

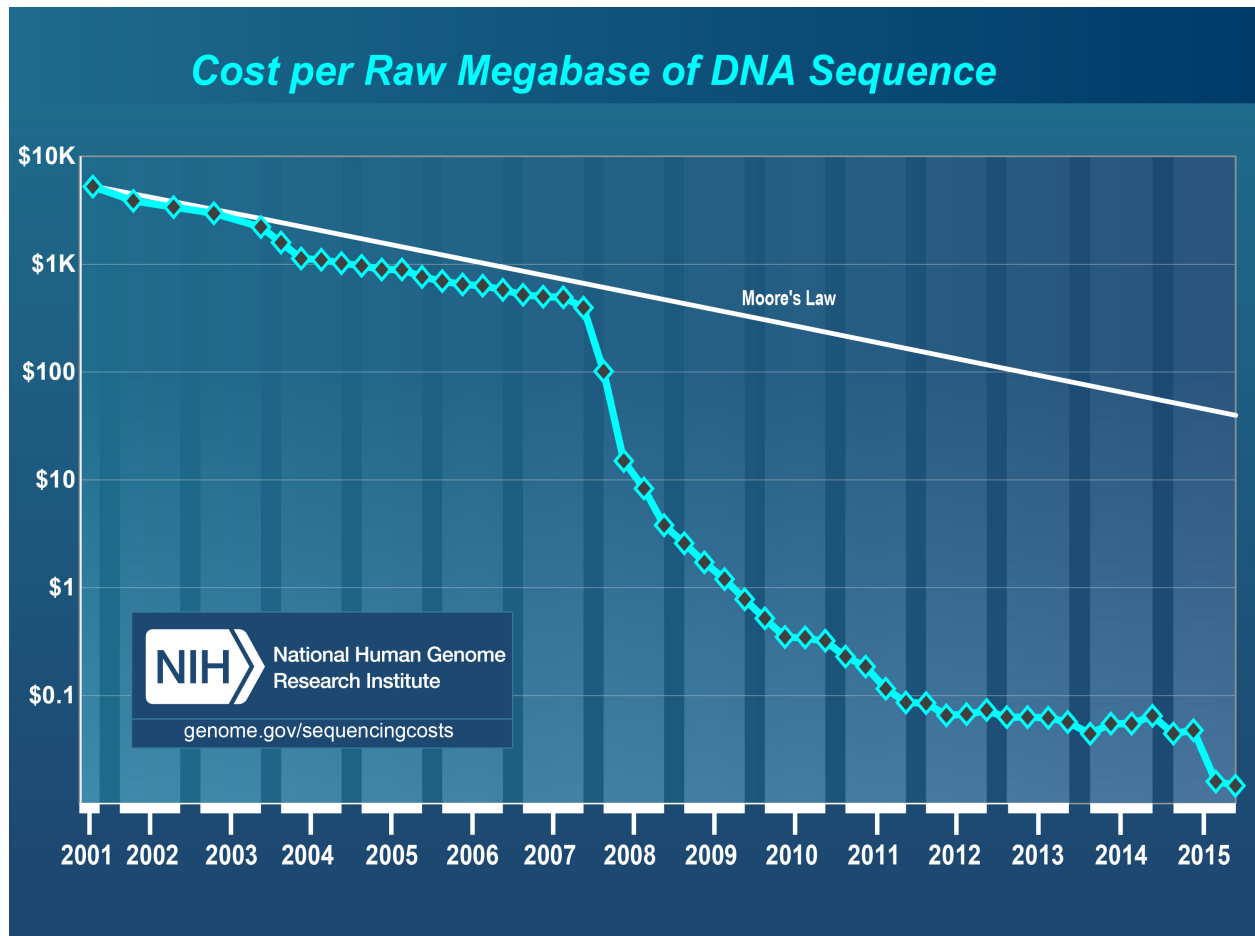


Figure A.1 Cost per raw megabase of DNA sequence from 2001 to 2015. Straight line - Moore's Law, blue curve - cost in US dollars, Y-axis scale is logarithmic. Graph reproduced from [wetterstrand2016]

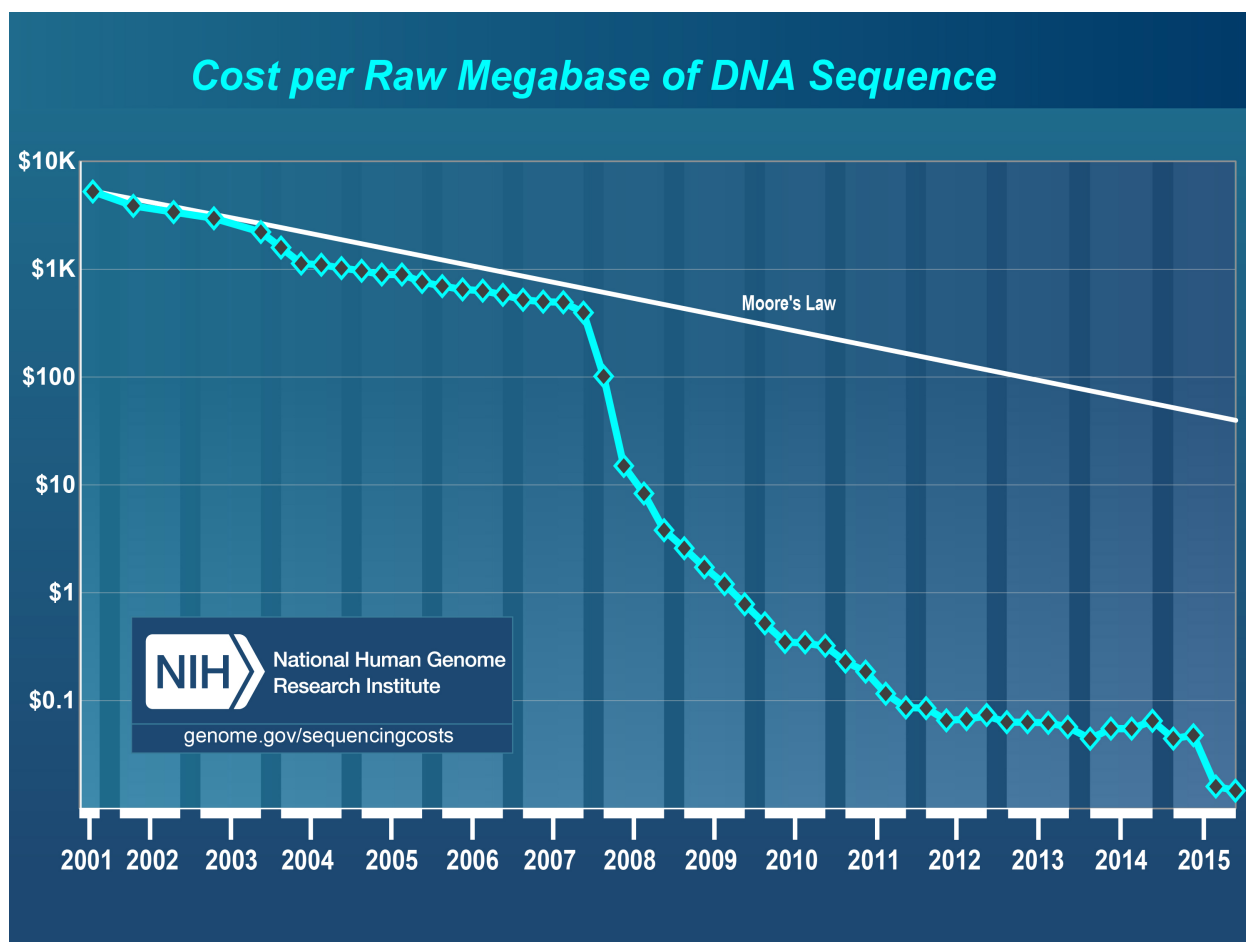


Figure A.2 Cost per raw megabase of DNA sequence from 2001 to 2015. Straight line - Moore's Law, blue curve - cost in US dollars, Y-axis scale is logarithmic. Graph reproduced from [wetterstrand2016]

Appendix B

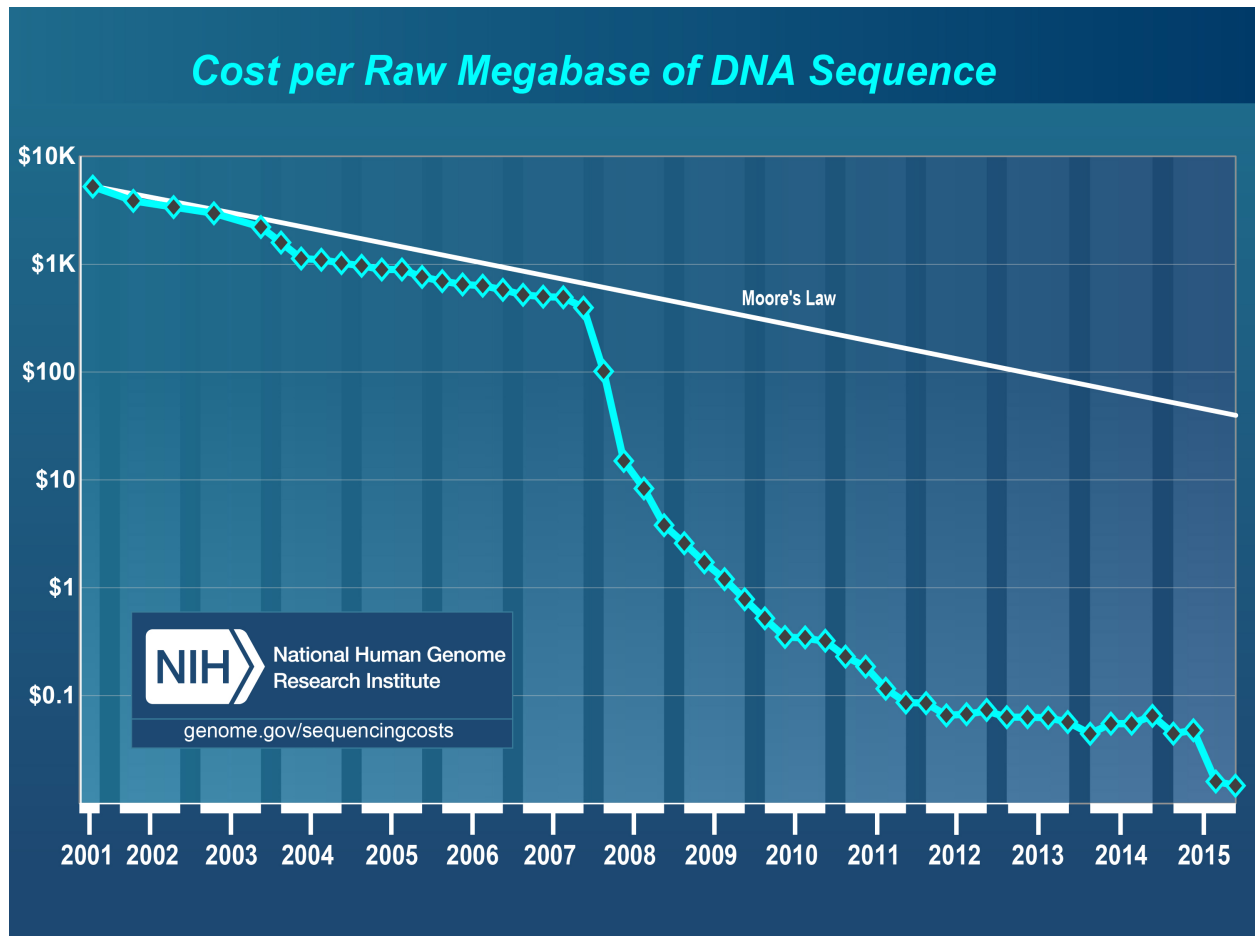


Figure B.1 Cost per raw megabase of DNA sequence from 2001 to 2015. Straight line - Moore's Law, blue curve - cost in US dollars, Y-axis scale is logarithmic. Graph reproduced from [wetterstrand2016]

Appendix C

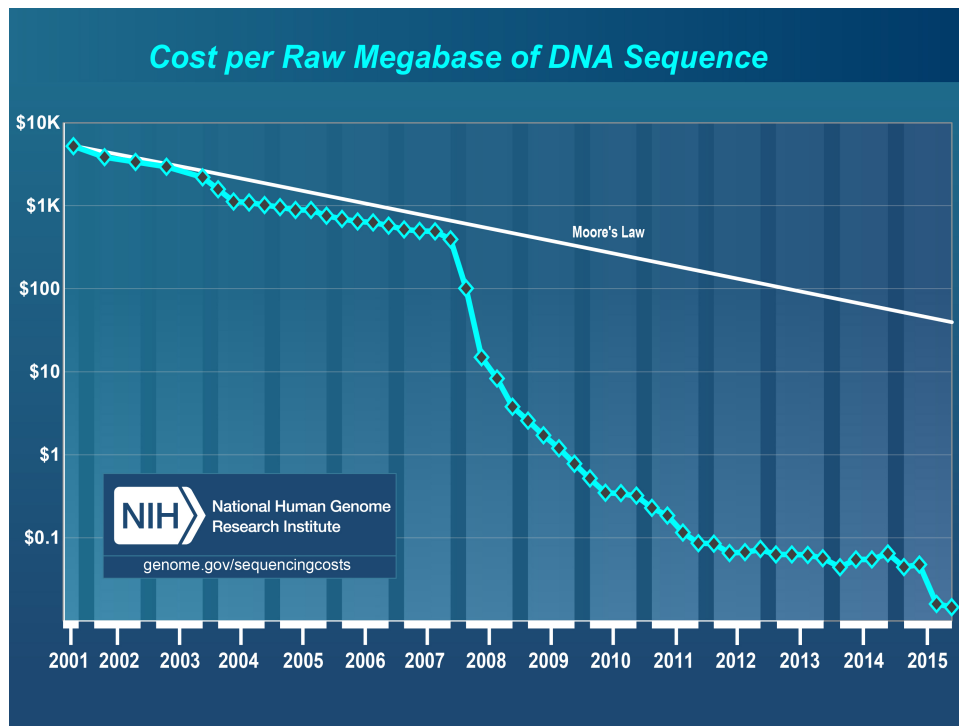


Figure C.1 Cost per raw megabase of DNA sequence from 2001 to 2015. Straight line - Moore's Law, blue curve - cost in US dollars, Y-axis scale is logarithmic. Graph reproduced from [wetterstrand2016]

Appendix D

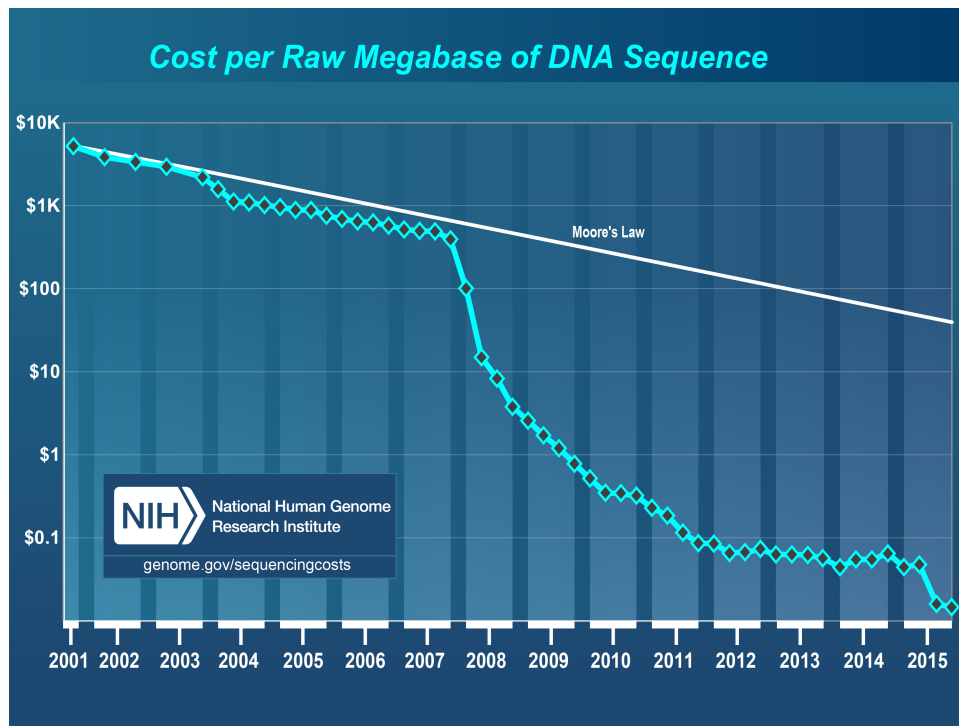


Figure D.1 Cost per raw megabase of DNA sequence from 2001 to 2015. Straight line - Moore's Law, blue curve - cost in US dollars, Y-axis scale is logarithmic. Graph reproduced from [wetterstrand2016]

Appendix E

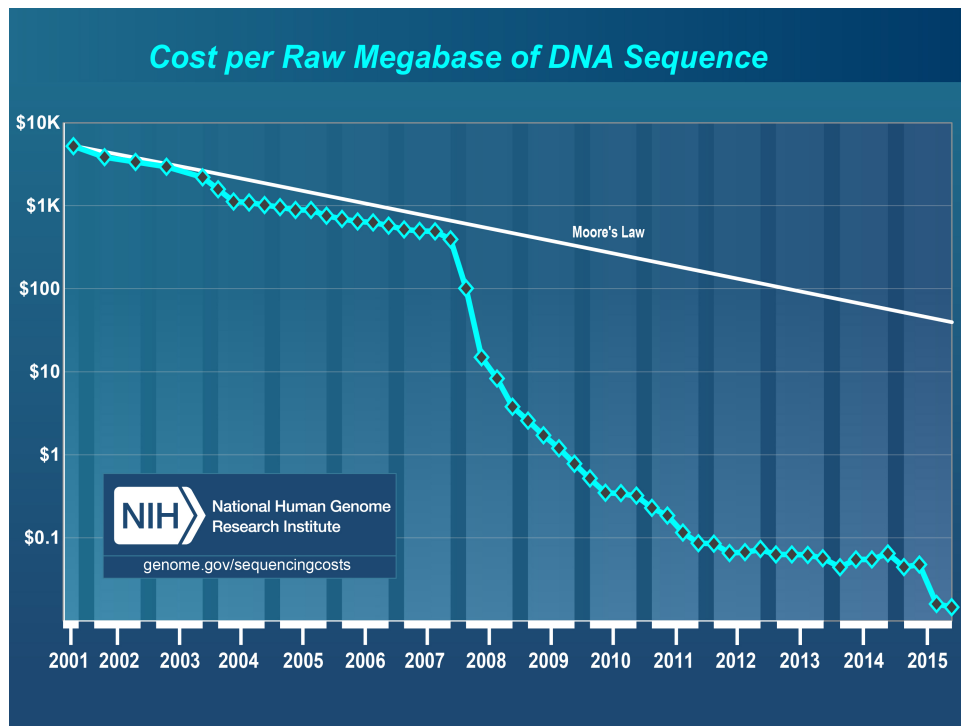


Figure E.1 Cost per raw megabase of DNA sequence from 2001 to 2015. Straight line - Moore's Law, blue curve - cost in US dollars, Y-axis scale is logarithmic. Graph reproduced from [wetterstrand2016]