

# Quantum Mechanics Made Simple: Lecture Notes

WENG CHO CHEW<sup>1</sup>

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<sup>1</sup>The author is with U of Illinois, Urbana-Champaign. He works part time at Hong Kong U this summer.



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# Preface

This set of supplementary lecture notes is the outgrowth of a course I taught, ECE 487, Quantum Electronics, at ECE Department, University of Illinois at Urbana-Champaign. It was intended to teach quantum mechanics to undergraduate students as well as graduate students. The primary text book for this course is Quantum Mechanics for Scientists and Engineers by D.A.B. Miller. I have learned a great deal by poring over Miller's book. But where I feel the book to be incomplete, I supplement them with my lecture notes. I try to reach into first principles as much as I could with these lecture notes. The only background needed for reading these notes is a background in undergraduate wave physics, and linear algebra.

I would still recommend using Miller's book as the primary text book for such a course, and use these notes as supplementary to teach this topic to undergraduates.

*Weng Cho CHEW*  
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# Chapter 1

## Introduction

### 1.1 Introduction

Quantum mechanics is an important intellectual achievement of the 20th century. It is one of the more sophisticated fields in physics that has affected our understanding of nano-meter length scale systems important for chemistry, materials, optics, and electronics. The existence of orbitals and energy levels in atoms can only be explained by quantum mechanics. Quantum mechanics can explain the behaviors of insulators, conductors, semi-conductors, and giant magneto-resistance. It can explain the quantization of light and its particle nature in addition to its wave nature. Quantum mechanics can also explain the radiation of hot body, and its change of color with respect to temperature. It explains the presence of holes and the transport of holes and electrons in electronic devices.

Quantum mechanics has played an important role in photonics, quantum electronics, and micro-electronics. But many more emerging technologies require the understanding of quantum mechanics; and hence, it is important that scientists and engineers understand quantum mechanics better. One area is nano-technologies due to the recent advent of nano-fabrication techniques. Consequently, nano-meter size systems are more common place. In electronics, as transistor devices become smaller, how the electrons move through the device is quite different from when the devices are bigger: nano-electronic transport is quite different from micro-electronic transport.

The quantization of electromagnetic field is important in the area of nano-optics and quantum optics. It explains how photons interact with atomic systems or materials. It also allows the use of electromagnetic or optical field to carry quantum information. Moreover, quantum mechanics is also needed to understand the interaction of photons with materials in solar cells, as well as many topics in material science.

When two objects are placed close together, they experience a force called the Casimir force that can only be explained by quantum mechanics. This is important for the understanding of micro/nano-electromechanical sensor systems (M/NEMS). Moreover, the understanding of spins is important in spintronics, another emerging technology where giant magneto-resistance, tunneling magneto-resistance, and spin transfer torque are being used.

Quantum mechanics is also giving rise to the areas of quantum information, quantum

communication, quantum cryptography, and quantum computing. It is seen that the richness of quantum physics will greatly affect the future generation technologies in many aspects.

## 1.2 Quantum Mechanics is Bizarre

The development of quantum mechanics is a great intellectual achievement, but at the same time, it is bizarre. The reason is that quantum mechanics is quite different from classical physics. The development of quantum mechanics is likened to watching two players having a game of chess, but the watchers have not a clue as to what the rules of the game are. By observations, and conjectures, finally the rules of the game are outlined. Often, equations are conjectured like conjurers pulling tricks out of a hat to match experimental observations. It is the interpretations of these equations that can be quite bizarre.

Quantum mechanics equations were postulated to explain experimental observations, but the deeper meanings of the equations often confused even the most gifted. Even though Einstein received the Nobel prize for his work on the photo-electric effect that confirmed that light energy is quantized, he himself was not totally at ease with the development of quantum mechanics as charted by the younger physicists. He was never comfortable with the probabilistic interpretation of quantum mechanics by Born and the Heisenberg uncertainty principle: “God doesn’t play dice,” was his statement assailing the probabilistic interpretation. He proposed “hidden variables” to explain the random nature of many experimental observations. He was thought of as the “old fool” by the younger physicists during his time.

Schrödinger came up with the bizarre “Schrödinger cat paradox” that showed the struggle that physicists had with quantum mechanics’s interpretation. But with today’s understanding of quantum mechanics, the paradox is a thing of yesteryear.

The latest twist to the interpretation in quantum mechanics is the parallel universe view that explains the multitude of outcomes of the prediction of quantum mechanics. All outcomes are possible, but with each outcome occurring in different universes that exist in parallel with respect to each other.

## 1.3 The Wave Nature of a Particle—Wave Particle Duality

The quantized nature of the energy of light was first proposed by Planck in 1900 to successfully explain the black body radiation. Einstein’s explanation of the photoelectric effect further asserts the quantized nature of light, or light as a photon.<sup>1</sup> However, it is well known that light is a wave since it can be shown to interfere as waves in the Newton ring experiment as far back as 1717.

The wave nature of an electron is revealed by the fact that when electrons pass through a crystal, they produce a diffraction pattern. That can only be explained by the wave nature

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<sup>1</sup>In the photoelectric effect, it was observed that electrons can be knocked off a piece of metal only if the light exceeded a certain frequency. Above that frequency, the electron gained some kinetic energy proportional to the excess frequency. Einstein then concluded that a packet of energy was associated with a photon that is proportional to its frequency.

of an electron. This experiment was done by Davisson and Germer in 1927. De Broglie hypothesized that the wavelength of an electron, when it behaves like a wave, is

$$\lambda = \frac{h}{p} \quad (1.3.1)$$

where  $h$  is the Planck's constant,  $p$  is the electron momentum,<sup>2</sup> and

$$h \approx 6.626 \times 10^{-34} \text{Joule} \cdot \text{second} \quad (1.3.2)$$

When an electron manifests as a wave, it is described by

$$\psi(z) \propto \exp(ikz) \quad (1.3.3)$$

where  $k = 2\pi/\lambda$ . Such a wave is a solution to<sup>3</sup>

$$\frac{\partial^2}{\partial z^2} \psi = -k^2 \psi \quad (1.3.4)$$

A generalization of this to three dimensions yields

$$\nabla^2 \psi(\mathbf{r}) = -k^2 \psi(\mathbf{r}) \quad (1.3.5)$$

We can define

$$p = \hbar k \quad (1.3.6)$$

where  $\hbar = h/(2\pi)$ .<sup>4</sup> Consequently, we arrive at an equation

$$-\frac{\hbar^2}{2m_0} \nabla^2 \psi(\mathbf{r}) = \frac{p^2}{2m_0} \psi(\mathbf{r}) \quad (1.3.7)$$

where

$$m_0 \approx 9.11 \times 10^{-31} \text{kg} \quad (1.3.8)$$

The expression  $p^2/(2m_0)$  is the kinetic energy of an electron. Hence, the above can be considered an energy conservation equation.

The Schrödinger equation is motivated by further energy balance that total energy is equal to the sum of potential energy and kinetic energy. Defining the potential energy to be  $V(\mathbf{r})$ , the energy balance equation becomes

$$\left[ -\frac{\hbar^2}{2m_0} \nabla^2 + V(\mathbf{r}) \right] \psi(\mathbf{r}) = E\psi(\mathbf{r}) \quad (1.3.9)$$

---

<sup>2</sup>Typical electron wavelengths are of the order of nanometers. Compared to 400 nm of wavelength of blue light, they are much smaller. Energetic electrons can have even smaller wavelengths. Hence, electron waves can be used to make electron microscope whose resolution is much higher than optical microscope.

<sup>3</sup>The wavefunction can be thought of as a "halo" that an electron carries that determine its underlying physical properties and how it interact with other systems.

<sup>4</sup>This is also called Dirac constant sometimes.

where  $E$  is the total energy of the system. The above is the time-independent Schrödinger equation. The *ad hoc* manner at which the above equation is arrived at usually bothers a beginner in the field. However, it predicts many experimental outcomes, as well as predicting the existence of electron orbitals inside an atom, and how electron would interact with other particles.

One can further modify the above equation in an *ad hoc* manner by noticing that other experimental finding shows that the energy of a photon is  $E = \hbar\omega$ . Hence, if we let

$$i\hbar\frac{\partial}{\partial t}\Psi(\mathbf{r}, t) = E\Psi(\mathbf{r}, t) \quad (1.3.10)$$

then

$$\Psi(\mathbf{r}, t) = e^{-i\omega t}\psi(\mathbf{r}, t) \quad (1.3.11)$$

Then we arrive at the time-dependent Schrödinger equation:

$$\left[-\frac{\hbar^2}{2m_0}\nabla^2 + V(\mathbf{r})\right]\psi(\mathbf{r}, t) = i\hbar\frac{\partial}{\partial t}\psi(\mathbf{r}, t) \quad (1.3.12)$$

Another disquieting fact about the above equation is that it is written in complex functions and numbers. In our prior experience with classical laws, they can all be written in real functions and numbers. We will later learn the reason for this.

Mind you, in the above, the frequency is not unique. We know that in classical physics, the potential  $V$  is not unique, and we can add a constant to it, and yet, the physics of the problem does not change. So, we can add a constant to both sides of the time-independent Schrödinger equation (1.3.10), and yet, the physics should not change. Then the total  $E$  on the right-hand side would change, and that would change the frequency we have arrived at in the time-dependent Schrödinger equation. We will explain how to resolve this dilemma later on. Just like potentials, in quantum mechanics, it is the difference of frequencies that matters in the final comparison with experiments, not the absolute frequencies.

The setting during which Schrödinger equation was postulated was replete with knowledge of classical mechanics. It will be prudent to review some classical mechanics knowledge next.







## Chapter 2

# Classical Mechanics

### 2.1 Introduction

Quantum mechanics cannot be derived from classical mechanics, but classical mechanics can inspire quantum mechanics. Quantum mechanics is richer and more sophisticated than classical mechanics. Quantum mechanics was developed during the period when physicists had rich knowledge of classical mechanics. In order to better understand how quantum mechanics was developed in this environment, it is better to understand some fundamental concepts in classical mechanics. Classical mechanics can be considered as a special case of quantum mechanics. We will review some classical mechanics concepts here.

In classical mechanics, a particle moving in the presence of potential<sup>1</sup>  $V(q)$  will experience a force given by

$$F(q) = -\frac{dV(q)}{dq} \quad (2.1.1)$$

where  $q$  represents the coordinate or the position of the particle. Hence, the particle can be described by the equations of motion

$$\frac{dp}{dt} = F(q) = -\frac{dV(q)}{dq}, \quad \frac{dq}{dt} = p/m \quad (2.1.2)$$

For example, when a particle is attached to a spring and moves along a frictionless surface, the force the particle experiences is  $F(q) = -kq$  where  $k$  is the spring constant. Then the equations of motion of this particle are

$$\frac{dp}{dt} = \dot{p} = -kq, \quad \frac{dq}{dt} = \dot{q} = p/m \quad (2.1.3)$$

Given  $p$  and  $q$  at some initial time  $t_0$ , one can integrate (2.1.2) or (2.1.3) to obtain  $p$  and  $q$  for all later time. A numerical analyst can think of that (2.1.2) or (2.1.3) can be solved by the

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<sup>1</sup>The potential here refers to potential energy.

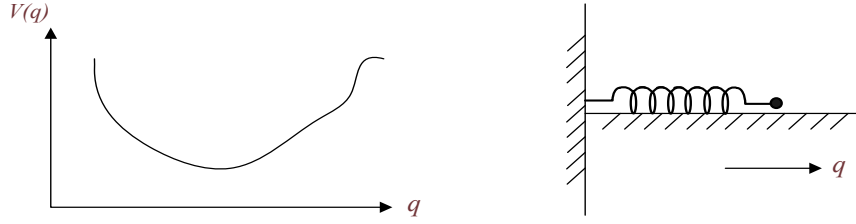


Figure 2.1: The left side shows a potential well in which a particle can be trapped. The right side shows a particle attached to a spring. The particle is subject to the force due to the spring, but it can also be described by the force due to a potential well.

finite difference method, where time-stepping can be used to find  $p$  and  $q$  for all later time. For instance, we can write the equations of motion more compactly as

$$\frac{d\mathbf{V}}{dt} = \mathbf{F}(\mathbf{V}) \quad (2.1.4)$$

where  $\mathbf{V} = [p, q]^t$ , and  $\mathbf{F}$  is a general vector function of  $\mathbf{V}$ . It can be nonlinear or linear; in the event if it is linear, then  $\mathbf{F}(\mathbf{V}) = \bar{\mathbf{A}} \cdot \mathbf{V}$ .

Using finite difference approximation, we can rewrite the above as

$$\begin{aligned} \mathbf{V}(t + \Delta t) - \mathbf{V}(t) &= \Delta t \mathbf{F}(\mathbf{V}(t)), \\ \mathbf{V}(t + \Delta t) &= \Delta t \mathbf{F}(\mathbf{V}(t)) + \mathbf{V}(t) \end{aligned} \quad (2.1.5)$$

The above can be used for time marching to derive the future values of  $\mathbf{V}$  from past values.

The above equations of motion are essentially derived using Newton's law. However, there exist other methods of deriving these equations of motion. Notice that only two variables  $p$  and  $q$  are sufficient to describe the state of a particle.

## 2.2 Lagrangian Formulation

Another way to derive the equations of motion for classical mechanics is via the use of the Lagrangian and the principle of least action. A Lagrangian is usually defined as the difference between the kinetic energy and the potential energy, i.e.,

$$L(\dot{q}, q) = T - V \quad (2.2.1)$$

where  $\dot{q}$  is the velocity. For a fixed  $t$ ,  $q$  and  $\dot{q}$  are independent variables, since  $\dot{q}$  cannot be derived from  $q$  if it is only known at one given  $t$ . The equations of motion is derived from the principle of least action which says that  $q(t)$  that satisfies the equations of motion between two times  $t_1$  and  $t_2$  should minimize the action integral

$$S = \int_{t_1}^{t_2} L(\dot{q}(t), q(t)) dt \quad (2.2.2)$$

Assuming that  $q(t_1)$  and  $q(t_2)$  are fixed, then the function  $q(t)$  between  $t_1$  and  $t_2$  should minimize  $S$ , the action. In other words, a first order perturbation in  $q$  from the optimal answer that minimizes  $S$  should give rise to second order error in  $S$ . Hence, taking the first variation of (2.2.2), we have

$$\delta S = \delta \int_{t_1}^{t_2} L(\dot{q}, q) dt = \int_{t_1}^{t_2} \delta L(\dot{q}, q) dt = \int_{t_1}^{t_2} \left( \delta \dot{q} \frac{\partial L}{\partial \dot{q}} + \delta q \frac{\partial L}{\partial q} \right) dt = 0 \quad (2.2.3)$$

In order to take the variation into the integrand, we have to assume that  $\delta L(\dot{q}, q)$  is taken with constant time. At constant time,  $\dot{q}$  and  $q$  are independent variables; hence, the partial derivatives in the next equality above follow. Using integration by parts on the first term, we have

$$\delta S = \delta q \frac{\partial L}{\partial \dot{q}} \Big|_{t_1}^{t_2} - \int_{t_1}^{t_2} \delta q \frac{d}{dt} \left( \frac{\partial L}{\partial \dot{q}} \right) dt + \int_{t_1}^{t_2} \delta q \frac{\partial L}{\partial q} dt = 0 \quad (2.2.4)$$

The first term vanishes because  $\delta q(t_1) = \delta q(t_2) = 0$  because  $q(t_1)$  and  $q(t_2)$  are fixed. Since  $\delta q(t)$  is arbitrary between  $t_1$  and  $t_2$ , we have

$$\frac{d}{dt} \left( \frac{\partial L}{\partial \dot{q}} \right) - \frac{\partial L}{\partial q} = 0 \quad (2.2.5)$$

The above is called the Lagrange equation, from which the equation of motion of a particle can be derived. The derivative of the Lagrangian with respect to the velocity  $\dot{q}$  is the momentum

$$p = \frac{\partial L}{\partial \dot{q}} \quad (2.2.6)$$

The derivative of the Lagrangian with respect to the coordinate  $q$  is the force. Hence

$$F = \frac{\partial L}{\partial q} \quad (2.2.7)$$

The above equation of motion is then

$$\dot{p} = F \quad (2.2.8)$$

Equation (2.2.6) can be inverted to express  $\dot{q}$  as a function of  $p$  and  $q$ , namely

$$\dot{q} = f(p, q) \quad (2.2.9)$$

The above two equations can be solved in tandem to find the time evolution of  $p$  and  $q$ .

For example, the kinetic energy  $T$  of a particle is given by

$$T = \frac{1}{2} m \dot{q}^2 \quad (2.2.10)$$

Then from (2.2.1), and the fact that  $V$  is independent of  $\dot{q}$ ,

$$p = \frac{\partial L}{\partial \dot{q}} = \frac{\partial T}{\partial \dot{q}} = m \dot{q} \quad (2.2.11)$$

or

$$\dot{q} = \frac{p}{m} \quad (2.2.12)$$

Also, from (2.2.1), (2.2.7), and (2.2.8), we have

$$\dot{p} = -\frac{\partial V}{\partial q} \quad (2.2.13)$$

The above pair, (2.2.12) and (2.2.13), form the equations of motion for this problem.

The above can be generalized to multidimensional problems. For example, for a one particle system in three dimensions,  $q_i$  has three degrees of freedom, and  $i = 1, 2, 3$ . (The  $q_i$  can represent  $x, y, z$  in Cartesian coordinates, but  $r, \theta, \phi$  in spherical coordinates.) But for  $N$  particles in three dimensions, there are  $3N$  degrees of freedom, and  $i = 1, \dots, 3N$ . The formulation can also be applied to particles constraint in motion. For instance, for  $N$  particles in three dimensions,  $q_i$  may run from  $i = 1, \dots, 3N - k$ , representing  $k$  constraints on the motion of the particles. This can happen, for example, if the particles are constraint to move in a manifold (surface), or a line (ring) embedded in a three dimensional space.

Going through similar derivation, we arrive at the equation of motion

$$\frac{d}{dt} \left( \frac{\partial L}{\partial \dot{q}_i} \right) - \frac{\partial L}{\partial q_i} = 0 \quad (2.2.14)$$

In general,  $q_i$  may not have a dimension of length, and it is called the generalized coordinate (also called conjugate coordinate). Also,  $\dot{q}_i$  may not have a dimension of velocity, and it is called the generalized velocity.

The derivative of the Lagrangian with respect to the generalized velocity is the generalized momentum (also called conjugate momentum), namely,

$$p_i = \frac{\partial L}{\partial \dot{q}_i} \quad (2.2.15)$$

The generalized momentum may not have a dimension of momentum. Hence, the equation of motion (2.2.14) can be written as

$$\dot{p}_i = \frac{\partial L}{\partial q_i} \quad (2.2.16)$$

Equation (2.2.15) can be inverted to yield an equation for  $\dot{q}_i$  as a function of the other variables. This equation can be used in tandem (2.2.16) as time-marching equations of motion.

## 2.3 Hamiltonian Formulation

For a multi-dimensional system, or a many particle system in multi-dimensions, the total time derivative of  $L$  is

$$\frac{dL}{dt} = \sum_i \left( \frac{\partial L}{\partial q_i} \dot{q}_i + \frac{\partial L}{\partial \dot{q}_i} \ddot{q}_i \right) \quad (2.3.1)$$

Since  $\partial L/\partial q_i = \frac{d}{dt}(\partial L/\partial \dot{q}_i)$  from the Lagrange equation, we have

$$\frac{dL}{dt} = \sum_i \left[ \frac{d}{dt} \left( \frac{\partial L}{\partial \dot{q}_i} \right) \dot{q}_i + \frac{\partial L}{\partial q_i} \ddot{q}_i \right] = \frac{d}{dt} \sum_i \left( \frac{\partial L}{\partial \dot{q}_i} \dot{q}_i \right) \quad (2.3.2)$$

or

$$\frac{d}{dt} \left( \sum_i \frac{\partial L}{\partial \dot{q}_i} \dot{q}_i - L \right) = 0 \quad (2.3.3)$$

The quantity

$$H = \sum_i \frac{\partial L}{\partial \dot{q}_i} \dot{q}_i - L \quad (2.3.4)$$

is known as the Hamiltonian of the system, and is a constant of motion, namely,  $dH/dt = 0$ . As shall be shown, the Hamiltonian represents the total energy of a system. It is a constant of motion because of the conservation of energy.

The Hamiltonian of the system, (2.3.4), can also be written, after using (2.2.15), as

$$H = \sum_i \dot{q}_i p_i - L \quad (2.3.5)$$

where  $p_i = \partial L/\partial \dot{q}_i$  is the generalized momentum. The first term has a dimension of energy, and in Cartesian coordinates, for a simple particle motion, it is easily seen that it is twice the kinetic energy. Hence, the above indicates that the Hamiltonian

$$H = T + V \quad (2.3.6)$$

The total variation of the Hamiltonian is

$$\begin{aligned} \delta H &= \delta \left( \sum_i p_i \dot{q}_i \right) - \delta L \\ &= \sum_i (\dot{q}_i \delta p_i + p_i \delta \dot{q}_i) - \sum_i \left( \frac{\partial L}{\partial q_i} \delta q_i + \frac{\partial L}{\partial \dot{q}_i} \delta \dot{q}_i \right) \end{aligned} \quad (2.3.7)$$

Using (2.2.15) and (2.2.16), we have

$$\begin{aligned} \delta H &= \sum_i (\dot{q}_i \delta p_i + p_i \delta \dot{q}_i) - \sum_i (\dot{p}_i \delta q_i + p_i \delta \dot{q}_i) \\ &= \sum_i (\dot{q}_i \delta p_i - \dot{p}_i \delta q_i) \end{aligned} \quad (2.3.8)$$

From the above, we gather that the Hamiltonian is a function of  $p_i$  and  $q_i$ . Taking the first variation of the Hamiltonian with respect to these variables, we have

$$\delta H = \sum_i \left( \frac{\partial H}{\partial p_i} \delta p_i + \frac{\partial H}{\partial q_i} \delta q_i \right) \quad (2.3.9)$$

Comparing the above with (2.3.8), we gather that

$$\dot{q}_i = \frac{\partial H}{\partial p_i} \quad (2.3.10)$$

$$\dot{p}_i = -\frac{\partial H}{\partial q_i} \quad (2.3.11)$$

These are the equations of motion known as the Hamiltonian equations.

The (2.3.4) is also known as the Legendre transformation. The original function  $L$  is a function of  $\dot{q}_i, q_i$ . Hence,  $\delta L$  depends on both  $\delta \dot{q}_i$  and  $\delta q_i$ . After the Legendre transformation,  $\delta H$  depends on the differential  $\delta p_i$  and  $\delta q_i$  as indicated by (2.3.8). This implies that  $H$  is a function of  $p_i$  and  $q_i$ . The equations of motion then can be written as in (2.3.10) and (2.3.11).

## 2.4 More on Hamiltonian

The Hamiltonian of a particle in classical mechanics is given by (2.3.6), and it is a function of  $p_i$  and  $q_i$ . For a non-relativistic particle in three dimensions, the kinetic energy

$$T = \frac{\mathbf{p} \cdot \mathbf{p}}{2m} \quad (2.4.1)$$

and the potential energy  $V$  is a function of  $\mathbf{q}$ . Hence, the Hamiltonian can be expressed as

$$H = \frac{\mathbf{p} \cdot \mathbf{p}}{2m} + V(\mathbf{q}) \quad (2.4.2)$$

in three dimensions. When an electromagnetic field is present, the Hamiltonian for an electron can be derived by letting the generalized momentum

$$p_i = m\dot{q}_i + eA_i \quad (2.4.3)$$

where  $e$  is the electron charge and  $A_i$  is component of the vector potential  $\mathbf{A}$ . Consequently, the Hamiltonian of an electron in the presence of an electromagnetic field is

$$H = \frac{(\mathbf{p} - e\mathbf{A}) \cdot (\mathbf{p} - e\mathbf{A})}{2m} + e\phi(\mathbf{q}) \quad (2.4.4)$$

The equation of motion of an electron in an electromagnetic field is governed by the Lorentz force law, which can be derived from the above Hamiltonian using the equations of motion provided by (2.3.10) and (2.3.11).

## 2.5 Poisson Bracket

Yet another way of expressing equations of motion in classical mechanics is via the use of Poisson bracket. This is interesting because Poisson bracket has a close quantum mechanics analogue. A Poisson bracket of two scalar variables  $u$  and  $v$  that are functions of  $q$  and  $p$  is defined as

$$\{u, v\} = \frac{\partial u}{\partial q} \frac{\partial v}{\partial p} - \frac{\partial v}{\partial q} \frac{\partial u}{\partial p} \quad (2.5.1)$$

In this notation, using (2.3.10) and (2.3.11),

$$\begin{aligned} \frac{du}{dt} &= \frac{\partial u}{\partial q} \frac{dq}{dt} + \frac{\partial u}{\partial p} \frac{dp}{dt} = \frac{\partial u}{\partial q} \frac{\partial H}{\partial p} - \frac{\partial u}{\partial p} \frac{\partial H}{\partial q} \\ &= \{u, H\} \end{aligned} \quad (2.5.2)$$

which is valid for any variable  $u$  that is a function of  $p$  and  $q$ . Hence, we have the equations of motion as

$$\dot{q} = \{q, H\}, \quad \dot{p} = \{p, H\} \quad (2.5.3)$$

in the Poisson bracket notation. As we shall see later, similar equations will appear in quantum mechanics.

The algebraic properties of Poisson bracket are

$$\{u, v\} = -\{v, u\} \quad (2.5.4)$$

$$\{u + v, w\} = \{u, w\} + \{v, w\} \quad (2.5.5)$$

$$\{uv, w\} = \{u, w\}v + u\{v, w\} \quad (2.5.6)$$

$$\{u, vw\} = \{u, v\}w + v\{u, w\} \quad (2.5.7)$$

$$\{\{u, v\}, w\} + \{\{v, w\}, u\} + \{\{w, u\}, v\} = 0 \quad (2.5.8)$$

These properties are antisymmetry, distributivity, associativity and Jacobi's identity.

If we define a commutator operation between two noncommuting operator  $\hat{u}$  and  $\hat{v}$  as

$$[\hat{u}, \hat{v}] = \hat{u}\hat{v} - \hat{v}\hat{u}, \quad (2.5.9)$$

it can be shown that the above commutator have the same algebraic properties as the Poisson bracket. An operator in quantum mechanics can be a matrix operator or a differential operator. In general, operators do not commute unless under very special circumstances.





## Chapter 3

# Quantum Mechanics—Some Preliminaries

### 3.1 Introduction

With some background in classical mechanics, we may motivate the Schrödinger equation in a more sanguine fashion. Experimental evidence indicated that small particles such as electrons behave quite strangely and cannot be described by classical mechanics alone. In classical mechanics, once we know  $p$  and  $q$  and their time derivatives (or  $\dot{p}$ ,  $\dot{q}$ ) of a particle at time  $t_0$ , one can integrate the equations of motion

$$\dot{p} = F, \quad \dot{q} = p/m \tag{3.1.1}$$

or use the finite difference method to find  $p$  and  $q$  at  $t_0 + \Delta t$ , and at all subsequent times.

In quantum mechanics, the use of two variables  $p$  and  $q$  and their derivatives is insufficient to describe the state of a particle and derive its future states. The state of a particle has to be more richly endowed and described by a wavefunction or state function  $\psi(q, t)$ . The state function (also known as a state vector) is a vector in the infinite dimensional space.

At this juncture, the state function or vector is analogous to when we study the control theory of a highly complex system. In the state variable approach, the state of a control system is described by the state vector, whose elements are variables that we think are important to capture the state of the system. For example, the state vector  $\mathbf{V}$ , describing the state of the factory, can contain variables that represent the number of people in a factory, the number of machines, the temperature of the rooms, the inventory in each room, etc. The state equation of this factory can then be written as

$$\frac{d}{dt}\mathbf{V}(t) = \overline{\mathbf{A}} \cdot \mathbf{V}(t) \tag{3.1.2}$$

It describes the time evolution of the factory. The matrix  $\overline{\mathbf{A}}$  causes the coupling between state variables as they evolve. It bears strong similarity to the time-dependent Schrödinger

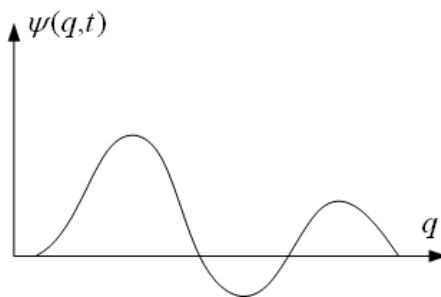


Figure 3.1: The state of a particle in quantum mechanics is described by a state function, which has infinitely many degrees of freedom.

equation, which is used to describe the time-evolution of the state function or the wavefunction of an electron. In the wavefunction, a complex number is assigned to each location in space.

In the Schrödinger equation, the wavefunction  $\psi(q, t)$  is a continuous function of the position variable  $q$  at any time instant  $t$ ; hence, it is described by infinitely many numbers, and has infinite degrees of freedom. The time evolution of the wavefunction  $\psi(q, t)$  is governed by the Schrödinger equation. It was motivated by experimental evidence and the works of many others such as Planck, Einstein, and de Broglie, who were aware of the wave nature of a particle and the dual wave-particle nature of light.

### 3.2 Probabilistic Interpretation of the wavefunction

The wavefunction of the Schrödinger equation has defied an acceptable interpretation for many years even though the Schrödinger equation was known to predict experimental outcomes. Some thought that it represented an electron cloud, and that perhaps, an electron, at the atomistic level, behaved like a charge cloud, and hence not a particle. The final, most accepted interpretation of this wavefunction (one that also agrees with experiments) is that its magnitude squared corresponds to the probabilistic density function. In other words, the probability of finding an electron in an interval  $[x, x + \Delta x]$  is equal to

$$|\psi(x, t)|^2 \Delta x \quad (3.2.1)$$

For the 3D case, the probability of finding an electron in a small volume  $\Delta V$  in the vicinity of the point  $\mathbf{r}$  is given by

$$|\psi(\mathbf{r}, t)|^2 \Delta V \quad (3.2.2)$$

Since the magnitude squared of the wavefunction represents a probability density function, it must satisfy the normalization condition of a probability density function, viz.,

$$\int dV |\psi(\mathbf{r}, t)|^2 = 1 \quad (3.2.3)$$

with its counterparts in 1D and 2D. The magnitude squared of this wavefunction is like some kind of “energy” that cannot be destroyed. Electrons cannot be destroyed and hence, charge conservation is upheld by the Schrödinger equation.

Motivated by the conservation of the “energy” of the wavefunction, we shall consider an “energy” conserving system where the classical Hamiltonian will be a constant of motion. In this case, there is no “energy” loss from the system. The Schrödinger equation that governs the time evolution of the wavefunction  $\psi$  is

$$\hat{H}\psi = i\hbar \frac{d\psi}{dt} \quad (3.2.4)$$

where  $\hat{H}$  is the Hamiltonian operator.<sup>1</sup> One can solve (3.2.4) formally to obtain

$$\psi(t) = e^{-i\frac{\hat{H}}{\hbar}t}\psi(t=0) \quad (3.2.5)$$

Since the above is a function of an operator, it has meaning only if this function acts on the eigenvectors of the operator  $\hat{H}$ . It can be shown easily that if  $\bar{\mathbf{A}} \cdot \mathbf{V}_i = \lambda_i \mathbf{V}_i$ ,

$$\exp(\bar{\mathbf{A}}) \cdot \mathbf{V}_i = \exp(\lambda_i) \mathbf{V}_i \quad (3.2.6)$$

If  $\hat{H}$  is a Hermitian operator, then there exists eigenfunctions, or special wavefunctions,  $\psi_n$ , such that

$$\hat{H}\psi_n = E_n\psi_n \quad (3.2.7)$$

where  $E_n$  is purely real. In this case, the time evolution of  $\psi_n$  from (3.2.5) is

$$\psi_n(t) = e^{-i\frac{E_n}{\hbar}t}\psi_n(t=0) = e^{-i\omega_n t}\psi_n(t=0) \quad (3.2.8)$$

In the above,  $E_n = \hbar\omega_n$ , or the energy  $E_n$  is related to frequency  $\omega_n$  via the reduced Planck constant  $\hbar$ . The reduced Planck constant is related to the Planck constant by  $\hbar = h/(2\pi)$  and  $h = 6.626068 \times 10^{-34}$  J s.

Scalar variables that are measurable in classical mechanics, such as  $p$  and  $q$ , are known as observables in quantum mechanics. They are elevated from scalar variables to operators in quantum mechanics, denoted by a “^” symbol here. In classical mechanics, for a one particle system, the Hamiltonian is given by

$$H = T + V = \frac{p^2}{2m} + V \quad (3.2.9)$$

The Hamiltonian contains the information from which the equations of motion for the particle can be derived. But in quantum mechanics, this is not sufficient, and  $H$  becomes an operator

$$\hat{H} = \frac{\hat{p}^2}{2m} + \hat{V} \quad (3.2.10)$$

---

<sup>1</sup>Rightfully, one should use the bra and ket notation to write this equation as  $\hat{H}|\psi\rangle = i\hbar \frac{d}{dt}|\psi\rangle$ . In the less rigorous notation in (3.2.4), we will assume that  $\hat{H}$  is in the representation in which the state vector  $\psi$  is in. That is if  $\psi$  is in coordinate space representation,  $\hat{H}$  is also in coordinates space representation.

This operator works in tandem with a wavefunction  $\psi$  to describe the state of the particle. The operator acts on a wavefunction  $\psi(t)$ , where in the coordinate  $q$  representation, is  $\psi(q, t)$ . When  $\psi(q, t)$  is an eigenfunction with energy  $E_n$ , it can be expressed as

$$\psi_n(q, t) = \psi_n(q)e^{-i\omega_n t} \quad (3.2.11)$$

where  $E_n = \hbar\omega_n$ . The Schrödinger equation for  $\psi_n(q)$  then becomes<sup>2</sup>

$$\hat{H}\psi_n(q) = \left( \frac{\hat{p}^2}{2m} + \hat{V} \right) \psi_n(q) = E_n\psi_n(q) \quad (3.2.12)$$

For simplicity, we consider an electron moving in free space where it has only a constant kinetic energy but not influenced by any potential energy. In other words, there is no force acting on the electron. In this case,  $\hat{V} = 0$ , and this equation becomes

$$\frac{\hat{p}^2}{2m}\psi_n(q) = E_n\psi_n(q) \quad (3.2.13)$$

It has been observed by de Broglie that the momentum of a particle, such as an electron which behaves like a wave, has a momentum

$$p = \hbar k \quad (3.2.14)$$

where  $k = 2\pi/\lambda$  is the wavenumber of the wavefunction. This motivates that the operator  $\hat{p}$  can be expressed by

$$\hat{p} = -i\hbar \frac{d}{dq} \quad (3.2.15)$$

in the coordinate space representation. This is chosen so that if an electron is described by a state function  $\psi(q) = c_1 e^{ikq}$ , then  $\hat{p}\psi(q) = \hbar k\psi(q)$ . The above motivation for the form of the operator  $\hat{p}$  is highly heuristic. We will see other reasons for the form of  $\hat{p}$  when we study the correspondence principle and the Heisenberg picture.

Equation (3.2.13) for a free particle is then

$$-\frac{\hbar^2}{2m} \frac{d^2}{dq^2} \psi_n(q) = E_n \psi_n(q) \quad (3.2.16)$$

Since this is a constant coefficient ordinary differential equation, the solution is of the form

$$\psi_n(q) = e^{\pm ikq} \quad (3.2.17)$$

which when used in (3.2.16), yields

$$\frac{\hbar^2 k^2}{2m} = E_n \quad (3.2.18)$$

Namely, the kinetic energy  $T$  of the particle is given by

$$T = \frac{\hbar^2 k^2}{2m} \quad (3.2.19)$$

---

<sup>2</sup>From the Schrödinger equation in coordinate space,  $\hat{V}$  turns out to be a scalar operator.

where  $p = \hbar k$  is in agreement with de Broglie's finding.

In many problems, the operator  $\hat{V}$  is a scalar operator in coordinate space representation which is a scalar function of position  $V(q)$ . This potential traps the particle within it acting as a potential well. In general, the Schrödinger equation for a particle becomes

$$\left[ -\frac{\hbar^2}{2m} \frac{\partial^2}{\partial q^2} + V(q) \right] \psi(q, t) = i\hbar \frac{\partial}{\partial t} \psi(q, t) \quad (3.2.20)$$

For a particular eigenstate with energy  $E_n$  as indicated by (3.2.11), it becomes

$$\left[ -\frac{\hbar^2}{2m} \frac{d^2}{dq^2} + V(q) \right] \psi_n(q) = E_n \psi_n(q) \quad (3.2.21)$$

The above is an eigenvalue problem with eigenvalue  $E_n$  and eigenfunction  $\psi_n(q)$ . These eigenstates are also known as stationary states, because they have a time dependence indicated by (3.2.11). Hence, their probability density functions  $|\psi_n(q, t)|^2$  are time independent.

These eigenfunctions correspond to trapped modes in the potential well defined by  $V(q)$  very much like trapped guided modes in a dielectric waveguide. These modes are usually countable and they can be indexed by the index  $n$ .

In the special case of a particle in free space, or the absence of the potential well, the particle or electron is not trapped and it is free to assume any energy or momentum indexed by the continuous variable  $k$ . In (3.2.18), the index for the energy should rightfully be  $k$  and the eigenfunctions are uncountably infinite. Moreover, the above can be generalized to two and three dimensional cases.

### 3.3 Simple Examples of Time Independent Schrödinger Equation

At this juncture, we have enough knowledge to study some simple solutions of time-independent Schrödinger equation such as a particle in a box, a particle impinging on a potential barrier, and a particle in a finite potential well.

#### 3.3.1 Particle in a 1D Box

Consider the Schrödinger equation for the 1D case where the potential  $V(x)$  is defined to be a function with zero value for  $0 < x < a$  (inside the box) and infinite value outside this range. The Schrödinger equation is given by

$$\left[ -\frac{\hbar^2}{2m} \frac{d^2}{dx^2} + V(x) \right] \psi(x) = E\psi(x) \quad (3.3.1)$$

where we have replaced  $q$  with  $x$ . Since  $V(x)$  is infinite outside the box,  $\psi(x)$  has to be zero. Inside the well,  $V(x) = 0$  and the above equation has a general solution of the form

$$\psi(x) = A \sin(kx) + B \cos(kx) \quad (3.3.2)$$

The boundary conditions are that  $\psi(x = 0) = 0$  and  $\psi(x = a) = 0$ . For this reason, a viable solution for  $\psi(x)$  is

$$\psi(x) = A \sin(kx) \quad (3.3.3)$$

where  $k = n\pi/a$ ,  $n = 1, \dots, \infty$ . There are infinitely many eigensolutions for this problem. For each chosen  $n$ , the corresponding energy of the solution is

$$E_n = \frac{(\hbar n\pi/a)^2}{2m} \quad (3.3.4)$$

These energy values are the eigenvalues of the problem, with the corresponding eigenfunctions given by (3.3.3) with the appropriate  $k$ . It is seen that the more energetic the electron is (high  $E_n$  values), the larger the number of oscillations the wavefunction has inside the box. The solutions that are highly oscillatory have higher  $k$  values, and hence, higher momentum or higher kinetic energy. The solutions which are even about the center of the box are said to have even parity, while those that are odd have odd parity.

One other thing to be noted is that the magnitude squared of the wavefunction above represents the probability density function. Hence, it has to be normalized. The normalized version of the wavefunction is

$$\psi(x) = \sqrt{2/a} \sin(n\pi x/a) \quad (3.3.5)$$

Moreover, these eigenfunctions are orthonormal to each other, viz.,

$$\int_0^a dx \psi_n^*(x) \psi_m(x) = \delta_{nm} \quad (3.3.6)$$

The orthogonality is the generalization of the fact that for a Hermitian matrix system, where the eigenvectors are given by

$$\bar{\mathbf{H}} \cdot \mathbf{V}_i = \lambda_i \mathbf{V}_i \quad (3.3.7)$$

then it can be proven easily that

$$\mathbf{V}_j^\dagger \cdot \mathbf{V}_i = C_j \delta_{ij} \quad (3.3.8)$$

Moreover, the eigenvalues are real.

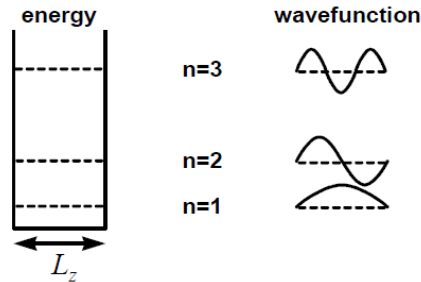


Figure 3.2: The wavefunctions of an electron trapped in a 1D box (from DAB Miller).

### 3.3.2 Particle Scattering by a Barrier

In the previous example, it is manifestly an eigenvalue problem since the solution can be found only at discrete values of  $E_n$ . The electron is trapped inside the box. However, in an open region problem where the electron is free to roam, the energy of the electron  $E$  can be arbitrary. We can assume that the potential profile is such that  $V(x) = 0$  for  $x < 0$  while  $V(x) = V_0$  for  $x > 0$ . The energy of the electron is such that  $0 < E < V_0$ . On the left side, we assume an electron coming in from  $-\infty$  with the wavefunction described by  $A \exp(ik_1x)$ . When this wavefunction hits the potential barrier, a reflected wave will be present, and the general solution on the left side of the barrier is given by

$$\psi_1(x) = A_1 e^{ik_1x} + B_1 e^{-ik_1x} \quad (3.3.9)$$

where  $(\hbar k_1)^2/(2m) = E$  is the kinetic energy of the incident electron. On the right side, however, the Schrödinger equation to be satisfied is

$$\left[ -\frac{\hbar^2}{2m} \frac{d^2}{dx^2} \right] \psi_2(x) = (E - V_0) \psi_2(x) \quad (3.3.10)$$

The solution of the transmitted wave on the right is

$$\psi_2(x) = A_2 e^{ik_2x} \quad (3.3.11)$$

where

$$k_2 = \sqrt{2m(E - V_0)}/\hbar \quad (3.3.12)$$

Given the known incident wave amplitude  $A_1$ , we can match the boundary conditions at  $x = 0$  to find the reflected wave amplitude  $B_1$  and the transmitted wave amplitude  $A_2$ . By eyeballing the Schrödinger equation (3.3.1), we can arrive at the requisite boundary conditions are that  $\psi$  and  $\partial\psi/\partial x$  are continuous at  $x = 0$ .

Since  $E < V_0$ ,  $k_2$  is pure imaginary, and the wave is evanescent and decays when  $x \rightarrow \infty$ . This effect is known as tunneling. The electron has a nonzero probability of being found inside the barrier, albeit with decreasing probability into the barrier. The larger  $V_0$  is compared to  $E$ , the more rapidly decaying is the wavefunction into the barrier.

However, if the electron is energetic enough so that  $E > V_0$ ,  $k_2$  becomes real, and then the wavefunction is no more evanescent. It penetrates into the barrier; it can be found even a long way from the boundary.

It is to be noted that the wavefunction in this case cannot be normalized as the above represents a fictitious situation of an electron roaming over infinite space. The above example illustrates the wave physics at the barrier.

### 3.3.3 Particle in a Potential Well

If the potential profile is such that

$$V(x) = \begin{cases} V_1 & \text{if } x < -a/2, \text{ region 1} \\ V_2 = 0 & \text{if } |x| < a/2, \text{ region 2} \\ V_3 & \text{if } x > a/2, \text{ region 3} \end{cases} \quad (3.3.13)$$

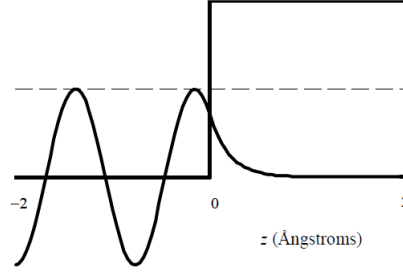


Figure 3.3: Scattering of the electron wavefunction by a 1D barrier (from DAB Miller).

then there can be trapped modes (or states) inside the well represented by standing waves, whereas outside the well, the waves are evanescent for eigenmodes for which  $E < V_1$  and  $E < V_3$ .

The wavefunction for  $|x| < a/2$  can be expressed as

$$\psi_2(x) = A_2 \sin(k_2 x) + B_2 \cos(k_2 x) \quad (3.3.14)$$

where  $k_2 = \sqrt{2mE}/\hbar$ . In region 1 to the left, the wavefunction is

$$\psi_1(x) = A_1 e^{\alpha_1 x} \quad (3.3.15)$$

where  $\alpha_1 = \sqrt{2m(V_1 - E)}/\hbar$ . The wave has to decay in the left direction. Similar, in region 3 to the right, the wavefunction is

$$\psi_3(x) = B_3 e^{-\alpha_3 x} \quad (3.3.16)$$

where  $\alpha_3 = \sqrt{2m(V_3 - E)}/\hbar$ . It has to decay in the right direction. Four boundary conditions can be imposed at  $x = \pm a/2$  to eliminate the four unknowns  $A_1$ ,  $A_2$ ,  $B_2$ , and  $B_3$ . However, non-trivial eigensolutions can only exist at selected values of  $E$  which are the eigenvalues of the Schrödinger equation. The eigenequation from which the eigenvalues can be derived is a transcendental equation.

To illustrate this point, we impose that  $\psi$  is continuous at  $x = \pm a/2$  to arrive at the following two equations:

$$A_1 e^{-\alpha_1 a/2} = -A_2 \sin(k_2 a/2) + B_2 \cos(k_2 a/2) \quad (3.3.17)$$

$$A_3 e^{-\alpha_3 a/2} = A_2 \sin(k_2 a/2) + B_2 \cos(k_2 a/2) \quad (3.3.18)$$

We further impose that  $\partial\psi/\partial x$  is continuous at  $x = \pm a$  to arrive at the following two equations:

$$\alpha_1 A_1 e^{-\alpha_1 a/2} = k_2 A_2 \cos(k_2 a/2) + k_2 B_2 \sin(k_2 a/2) \quad (3.3.19)$$

$$-\alpha_3 A_3 e^{-\alpha_3 a/2} = k_2 A_2 \cos(k_2 a/2) - k_2 B_2 \sin(k_2 a/2) \quad (3.3.20)$$



The above four equations form a matrix equation

$$\overline{\mathbf{M}} \cdot \mathbf{V} = 0 \tag{3.3.21}$$

where  $\mathbf{V} = [A_1, A_2, B_2, B_3]^t$ , and the elements of  $\overline{\mathbf{M}}$ , which depend on  $E$ , can be gleaned off the above equations. Non-trivial solutions exist for  $\mathbf{V}$  only if

$$\det(\overline{\mathbf{M}}(E)) = |\overline{\mathbf{M}}(E)| = 0 \tag{3.3.22}$$

The above is the transcendental eigenequation from which the eigenvalues  $E$  can be derived. The nontrivial solutions for  $\mathbf{V}$  are in the null space of  $\overline{\mathbf{M}}$ . Having known  $\mathbf{V} = [A_1, A_2, B_2, B_3]^t$ , the eigenfunctions of the Schrödinger equation can be constructed. Notice that  $\mathbf{V}$  is known to an arbitrary multiplicative constant. The normalization of the eigenfunction will pin down the value of this constant.

When the potential well is symmetric such that  $V_1 = V_3 = V_0$ , then the solutions can be decomposed into odd and even solutions about  $x = 0$ . In this case, either  $A_2 = 0$  for even modes, or  $B_2 = 0$  for odd modes. Furthermore,  $A_1 = \pm B_3$  for these modes. The problem then has two unknowns, and two boundary conditions at one of the interfaces suffice to deduce the eigenequation.

The above problem is analogous to the 1D dielectric waveguide problem in classical electromagnetics. The plot of the eigenmodes and their energy levels are shown in Figure 3.4. It is an interesting example showing that a trapped electron exists with different energy levels. The more energetic the electron is, the higher the energy level. For the case when  $E > V_0$ , the wavefunction is not evanescent outside the well, and the electron is free to roam outside the well.

Modern technology has allowed the engineering of nano-structures so small that a quantum well can be fabricated. Quantum well technology is one of the emerging nano-technologies. The energy level of a quantum well can also be engineered so that laser technology of different wavelengths can be fabricated.

In the case of a hydrogen atom, the Coulomb potential around the proton at the nucleus yields a potential well described by  $-q^2/(\epsilon 4\pi r)$ . This well can trap an electron into various eigenstates, yielding different electronic orbitals. The Schrödinger equation has predicted the energy levels of a hydrogen atom with great success.

### 3.4 The Quantum Harmonic Oscillator

The pendulum, a particle attached to a spring, or many vibrations in atoms and molecules can be described as a harmonic oscillator. Hence, the harmonic oscillator is one of the most important examples in quantum mechanics. Its quantum mechanical version can be described by the 1D Schrödinger equation.

The classical equation for a harmonic oscillator is given by

$$\frac{d^2 z}{dt^2} = -\omega^2 z \tag{3.4.1}$$

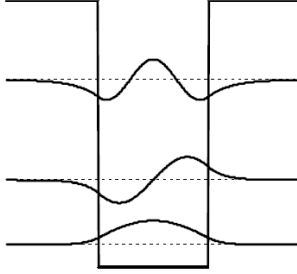


Figure 3.4: Trapped modes representing the electron wavefunctions inside a 1D potential well (from DAB Miller).

where  $\omega = \sqrt{K/m}$  where  $K$  is the spring constant, and  $m$  is the mass of the particle. The potential energy of a particle attached to a spring is given by

$$V(z) = \frac{1}{2}m\omega^2 z^2 \quad (3.4.2)$$

Consequently, the above potential energy can be used in the Schrödinger equation to describe the trapping of wave modes. The kinetic energy of the particle is described by a term proportional to the square of the momentum operator. Hence, the corresponding 1D Schrödinger equation is

$$\left[ -\frac{\hbar^2}{2m} \frac{d^2}{dz^2} + \frac{1}{2}m\omega^2 z^2 \right] \psi_n(z) = E_n \psi_n(z) \quad (3.4.3)$$

with a parabolic potential well. It turns out that this equation has closed-form solutions, yielding the wavefunction for an eigenstate given by

$$\psi_n(z) = \sqrt{\frac{1}{2^n n!}} \sqrt{\frac{m\omega}{\pi \hbar}} e^{-\frac{m\omega}{2\hbar} z^2} H_n \left( \sqrt{\frac{m\omega}{\hbar}} z \right) \quad (3.4.4)$$

where  $H_n(x)$  is a Hermite polynomial, and the wavefunction is Gaussian tapered. The energy of the eigenstate is given by

$$E_n = \left( n + \frac{1}{2} \right) \hbar\omega \quad (3.4.5)$$

The energy levels are equally spaced  $\hbar\omega$  apart. Even the lowest energy state, the ground state, has a nonzero energy of  $\hbar\omega/2$ . The higher energy states correspond to larger amplitudes of oscillation, and vice versa for the lower energy states. In order to kick the quantum harmonic oscillator from the low energy state to a level above, it needs a packet of energy of  $\hbar\omega$ , the quantization energy of a photon. The physics of quantized electromagnetic oscillations (photons) and quantized mechanical oscillations (phonons) is intimately related to the quantum harmonic oscillator.

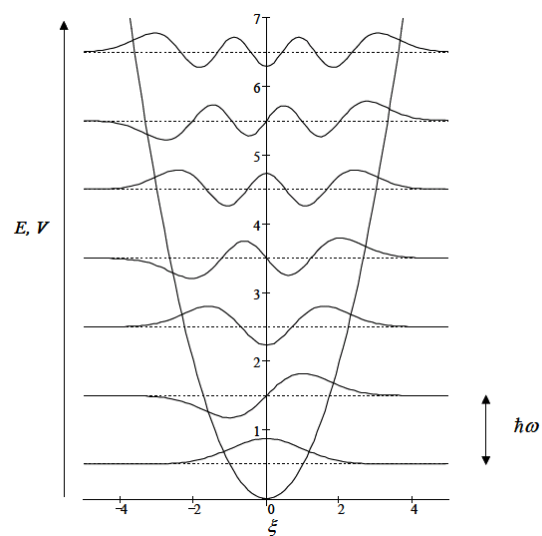


Figure 3.5: Sketch of the eigenstates, energy levels, and the potential well of a quantum harmonic oscillator (picture from DAB Miller).



## Chapter 4

# Time-Dependent Schrödinger Equation

### 4.1 Introduction

Each eigenstate of Schrödinger equation has its own time dependence of  $\exp(-i\omega_n t)$ . When we consider one eigenstate alone, its time dependence is unimportant, as the time dependence disappears when we convert the wavefunction into a probability density function by taking its amplitude square. Moreover, the absolute frequency of an eigenstate is arbitrary. Hence, the probability density function is quite uninteresting. This is an antithesis to the classical harmonic oscillator where the position of the particle moves with respect to time.

However, when a quantum state is described by a wavefunction which is a linear superposition of two eigenstates, it is important that we take into account their individual frequency value and time dependence. The two eigenstates will “beat” with each other to produce a difference in frequencies when we take the amplitude square of the wavefunction.

### 4.2 Quantum States in the Time Domain

Consider a quantum state which is a linear superposition of two eigenstates

$$\psi(\mathbf{r}, t) = c_a e^{-i\omega_a t} \psi_a(\mathbf{r}) + c_b e^{-i\omega_b t} \psi_b(\mathbf{r}) \quad (4.2.1)$$

where  $c_a$  and  $c_b$  are properly chosen to normalize the corresponding probability density function. Then the probability function is

$$|\psi(\mathbf{r}, t)|^2 = |c_a|^2 |\psi_a(\mathbf{r})|^2 + |c_b|^2 |\psi_b(\mathbf{r})|^2 + 2\Re \left[ c_a c_b^* \psi_a(\mathbf{r}) \psi_b^*(\mathbf{r}) e^{-i(\omega_a - \omega_b)t} \right] \quad (4.2.2)$$

It is clearly time varying. We are seeing the motion of the particle through the potential well.

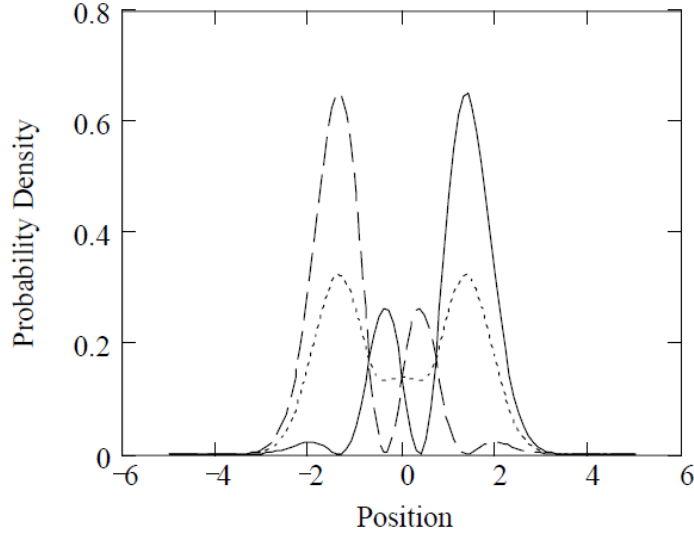


Figure 4.1: The time evolution of the linear superposition of the first two eigenstates of the simple quantum harmonic oscillator. (i) Beginning of a period (solid line). (ii) One quarter and three quarters way through the period (dotted line). (iii) Half way through the period (picture from DAB Miller).

### 4.3 Coherent State

The eigenstates of a quantum harmonic oscillator do not resemble its classical state. In order to see the classical state emerging from the quantum harmonic oscillator eigenstates, we need to take a judicious linear superposition of them. Such a state is called the coherent state.<sup>1</sup> The coherent state is given by

$$\psi_N(\xi, t) = \sum_{n=0}^{\infty} c_{Nn} e^{-i(n+1/2)\omega t} \psi_n(\xi) \quad (4.3.1)$$

where  $\xi = \sqrt{m\omega/\hbar}z$ , and  $c_{Nn} = \sqrt{N^n e^{-N}/n!}$ .<sup>2</sup> It forms a wave packet that emulates the motion of the classical harmonic oscillators such as a pendulum. The total energy of this wave packet is given by  $(N + 1/2)\hbar\omega$ . The larger  $N$  is, the more energy the particle has, and the closer is the coherent state to a classical state.

When a particle is traveling in vacuum, it should be described by a localized wavefunction, such as a Gaussian wave packet. The study of the Gaussian wave packet is given in Appendix C. It can be seen that the classical limit emerges when the momentum of the particle becomes very large.

<sup>1</sup>The state was derived by R. Glauber who won the Nobel Prize in 2005 for his contribution.

<sup>2</sup>The coefficient is also written as  $e^{-|\alpha|^2/2} \frac{\alpha^n}{\sqrt{n!}}$  where  $N = |\alpha|^2$ .

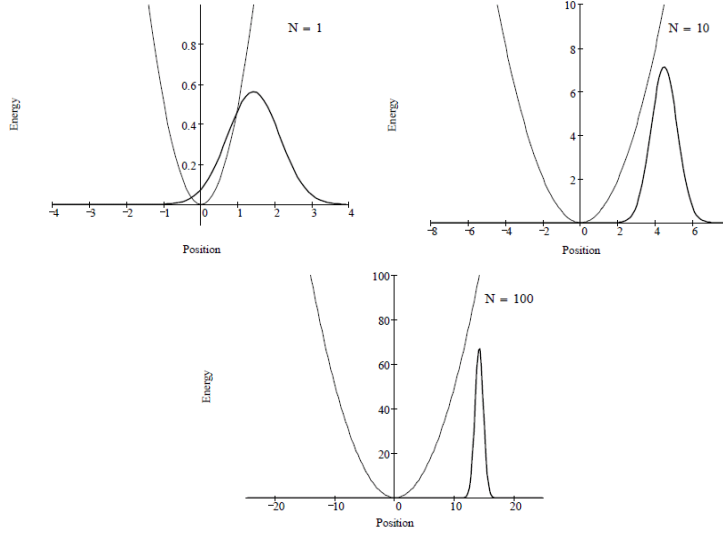


Figure 4.2: Coherent state for different values of  $N$  (picture from DAB Miller).

## 4.4 Measurement Hypothesis and Expectation Value

The wavefunction of a quantum system can be written as a linear superposition of the stationary states

$$\psi(\mathbf{r}, t) = \sum_n c_n(t) \psi_n(r) \quad (4.4.1)$$

The amplitude square of this wavefunction should integrate to one due to its probabilistic interpretation. In other words,

$$\int_V |\psi(\mathbf{r}, t)|^2 d\mathbf{r} = \int_V \sum_n c_n(t) \psi_n(\mathbf{r}) \sum_{n'} c_{n'}^*(t) \psi_{n'}^*(\mathbf{r}) = 1 \quad (4.4.2)$$

Exchanging the order of integration and summations, we arrive at

$$\int_V |\psi(\mathbf{r}, t)|^2 d\mathbf{r} = \sum_n \sum_{n'} c_n(t) c_{n'}^*(t) \int_V \psi_n(r) \psi_{n'}^*(r) \quad (4.4.3)$$

Using the orthonormality of the eigenstates, viz.,  $\int_V \psi_n(r) \psi_{n'}^*(r) = \delta_{nn'}$ , we arrive at the fact that

$$\sum_n |c_n(t)|^2 = 1 \quad (4.4.4)$$

Since the amplitude squares add up to one, they can be assigned probabilistic interpretation as well. Hence,  $P_n = |c_n(t)|^2$  represents the probability of finding the electron in eigenstate  $n$ .

The quantum measurement hypothesis states that before the measurement, the electron lives as a linear superposition of different eigenstates. After the measurement, the electron collapses into one of the eigenstates, say the  $n$  eigenstate, with probability proportional to  $|c_n(t)|^2$ .

The above can go further by saying that an electron is described by a wavefunction where its position is indeterminate: it can be found any place where the wavefunction is nonzero. Its probability of being found is proportional to the amplitude square of the wavefunction. However, once the measurement is made to determine the location of the electron, the electron's position collapses to one location as is discovered by the experiment. Its position is determinate after the measurement.

Due to this probabilistic interpretation, the expected energy of the quantum system is given by

$$\langle E \rangle = \sum_n E_n P_n = \sum_n E_n |c_n(t)|^2 \quad (4.4.5)$$

The above is the expectation value of  $E$ . Notice that for the eigenstate  $c_n(t) \propto e^{-i\omega_n t}$  and hence,  $|c_n(t)|^2$  is time independent. The above expectation of the energy is a constant. This means that the quantum system is energy conserving, namely, it cannot lose energy, for instance, due to radiation. This is because this is an isolated quantum system. In order for this system to lose energy, it has to be coupled with other quantum systems, which is a subject of later study.

For the coherent state, the expectation value of  $E$  is

$$\langle E \rangle = \sum_{n=0}^{\infty} E_n \frac{N^n e^{-N}}{n!} \quad (4.4.6)$$

$$= \hbar\omega \left[ \sum_{n=0}^{\infty} n \frac{N^n e^{-N}}{n!} \right] + \frac{\hbar\omega}{2} \quad (4.4.7)$$

$$= \left( N + \frac{1}{2} \right) \hbar\omega \quad (4.4.8)$$

Hence, the energy of the coherent state is that of  $N$  photons. The larger  $N$  is, the closer it is to a classical state.

The above expected energy value can also be found by taking the expectation of the Hamiltonian operator, viz.,

$$\langle E \rangle = \int d\mathbf{r} \psi^*(\mathbf{r}) \hat{H} \psi(\mathbf{r}) \quad (4.4.9)$$

The above equality can be easily shown by substituting (4.4.1) into the right-hand side and making use of the fact that  $\psi_n(\mathbf{r})$  is the eigenfunction of the Hamiltonian operator

$$\hat{H} = -\frac{\hbar^2}{2m} \nabla^2 + V(\mathbf{r}) \quad (4.4.10)$$



In general, a quantity that is measurable, like the energy  $E$ , is replaced by an operator  $\hat{H}$  in quantum mechanics. Such operators are called observables, such as momentum, position, etc. The relation between the classical value of an observable and its quantum mechanical counterpart, the operator, is via the expectation relation, viz.,

$$\langle \hat{A} \rangle = \int d\mathbf{r} \psi^*(\mathbf{r}) \hat{A} \psi(\mathbf{r}) \quad (4.4.11)$$

The above reduces to a scalar number. It is also the bridge between classical quantities and quantum mechanical entities.

#### 4.4.1 Time Evolution of the Hamiltonian Operator

The Schrödinger equation can be written as

$$\frac{d\psi(\mathbf{r}, t)}{dt} = -\frac{i\hat{H}}{\hbar} \psi(\mathbf{r}, t) \quad (4.4.12)$$

Assuming that  $\hat{H}$  is a constant (which is not), we can integrate the above to yield

$$\psi(\mathbf{r}, t) = e^{-\frac{i\hat{H}}{\hbar}t} \psi(\mathbf{r}, t=0) \quad (4.4.13)$$

The above is a function of an operator which is also an operator. A function of an operator has meaning only when it acts on the eigenvectors of the operator.

For example, we have a polynomial function

$$f(x) = a_0 + a_1x + a_2x^2 + a_3x^3 + \dots \quad (4.4.14)$$

If  $\bar{\mathbf{A}}$  is an operator, say a matrix operator, then the function of an operator is defined to be

$$f(\bar{\mathbf{A}}) = a_0\bar{\mathbf{I}} + a_1\bar{\mathbf{A}} + a_2\bar{\mathbf{A}}^2 + a_3\bar{\mathbf{A}}^3 + \dots \quad (4.4.15)$$

which is itself an operator as is obvious from the above definition. When the above acts on an eigenvector of  $\bar{\mathbf{A}}$  such that  $\bar{\mathbf{A}} \cdot \mathbf{v}_i = \lambda_i \mathbf{v}_i$ , then

$$f(\bar{\mathbf{A}}) \cdot \mathbf{v}_i = a_0\mathbf{v}_i + a_1\bar{\mathbf{A}} \cdot \mathbf{v}_i + a_2\bar{\mathbf{A}}^2 \cdot \mathbf{v}_i + a_3\bar{\mathbf{A}}^3 \cdot \mathbf{v}_i + \dots \quad (4.4.16)$$

$$= (a_0 + a_1\lambda_i + a_2\lambda_i^2 + a_3\lambda_i^3 + \dots) \cdot \mathbf{v}_i \quad (4.4.17)$$

$$= f(\lambda_i) \mathbf{v}_i \quad (4.4.18)$$

The idea can be extended to operators of infinite dimensions if we know the eigenfunctions of the operator. If the eigenfunction  $\psi_n(\mathbf{r}, t)$  is such that  $\hat{H}\psi_n(\mathbf{r}, t) = E_n\psi_n(\mathbf{r}, t)$  then (4.4.13) becomes

$$\psi_n(\mathbf{r}, t) = e^{-\frac{i\hat{H}}{\hbar}t} \psi_n(\mathbf{r}, t=0) = e^{-\frac{iE_n}{\hbar}t} \psi_n(\mathbf{r}, t=0) \quad (4.4.19)$$

The above is a solution to (4.4.12). When we write

$$\psi(\mathbf{r}, t) = \sum_n c_n \psi_n(\mathbf{r}, t) \quad (4.4.20)$$

$$\psi(\mathbf{r}, t) = e^{-\frac{i\hat{H}}{\hbar}t}\psi(\mathbf{r}, t=0) = \sum_n c_n e^{-\frac{iE_n}{\hbar}t}\psi_n(\mathbf{r}, t=0) \quad (4.4.21)$$

It can be easily shown that the above, which is derived from (4.4.13) is a solution to (4.4.12). Hence, (4.4.13) is the formal solution to (4.4.12).

#### 4.4.2 Uncertainty Principle

The uncertainty principle is asserted by the fact that Fourier transform of a Gaussian function is another Gaussian function. For instance, if we have a wave packet that is formed by superposing waves with different  $k$  or momentum, we can express it as

$$\psi(z) = \int_{-\infty}^{\infty} \tilde{\psi}(k) e^{ikz} dk \quad (4.4.22)$$

where

$$\tilde{\psi}(k) = C e^{-\left(\frac{k-k_0}{2\Delta k}\right)^2} \quad (4.4.23)$$

The above represents a Gaussian-tapered function with a spread of wavenumbers centered around  $k = k_0$ . It can be Fourier transformed in closed form yielding that<sup>3</sup>

$$\psi(z) = D e^{-(\Delta k)^2 z^2} \quad (4.4.24)$$

The probability density functions are proportional to the amplitude square of the above. The first pulse has a standard deviation of  $\Delta k$  while the standard deviation of the second pulse is  $\Delta z = 1/(2\Delta k)$ . The product of these two standard deviations yields

$$\Delta k \Delta z = \frac{1}{2} \quad (4.4.25)$$

or

$$\Delta p \Delta z = \frac{\hbar}{2} \quad (4.4.26)$$

#### 4.4.3 Particle Current

It turns out that even though the state of an electron is defined by a wavefunction, other equations do not see this wavefunction. For instance, Maxwell's equations, will produce electric field from sources, but they will only produce the electric field from the charge cloud and the current produced by the charge cloud. We see that for stationary states of a trapped electron in a potential well, the charge cloud is static. Hence, it will not radiate according to electromagnetic theory. This resolves the conflict in early days as to why the electron, supposedly "orbiting" around the nucleus of an atom, does not radiate.

<sup>3</sup>The Fourier transform of a Gaussian is a Gaussian: it can be proved by using the identity that  $I = \int_{-\infty}^{\infty} dx e^{-x^2} = \sqrt{\pi}$ . This can be proved by noticing that  $I^2 = \int_{-\infty}^{\infty} dx e^{-x^2} \int_{-\infty}^{\infty} dy e^{-y^2} = \iint_{-\infty}^{\infty} dx dy e^{-(x^2+y^2)} = 2\pi \int_0^{\infty} \rho d\rho e^{-\rho^2} = \pi$ . The last integral can be integrated in closed form.

However, when an electron is in more than one stationary state, the charge cloud is time varying, and can potentially couple to an external electric field and radiate.<sup>4</sup> For conservation of charge, we have the relation that

$$\frac{\partial \rho_p}{\partial t} = -\nabla \cdot \mathbf{J}_p \quad (4.4.27)$$

where  $\rho_p$  is the particle density, and  $\mathbf{J}_p$  is the particle current density. The particle density  $\rho_p(\mathbf{r}, t) = |\psi(\mathbf{r}, t)|^2$ . We can take the time derivative of  $\rho_p$  to yield

$$\frac{\partial}{\partial t} [\psi^*(\mathbf{r}, t)\psi(\mathbf{r}, t)] = \frac{\partial \psi^*(\mathbf{r}, t)}{\partial t} \psi(\mathbf{r}, t) + \psi^*(\mathbf{r}, t) \frac{\partial \psi(\mathbf{r}, t)}{\partial t} \quad (4.4.28)$$

We can use Schrödinger equation to replace the time derivatives on the right-hand side to get

$$\frac{\partial}{\partial t} [\psi^*(\mathbf{r}, t)\psi(\mathbf{r}, t)] = -\frac{i}{\hbar} (\psi^* \hat{H} \psi - \psi \hat{H}^* \psi^*) \quad (4.4.29)$$

Substituting in for the definition of the Hamiltonian operator, we have further that

$$\frac{\partial}{\partial t} [\psi^*(\mathbf{r}, t)\psi(\mathbf{r}, t)] = \frac{i\hbar}{2m} (\psi^* \nabla^2 \psi - \psi \nabla^2 \psi^*) \quad (4.4.30)$$

$$\frac{\partial}{\partial t} [\psi^*(\mathbf{r}, t)\psi(\mathbf{r}, t)] = -\frac{i\hbar}{2m} \nabla \cdot (\psi \nabla \psi^* - \psi^* \nabla \psi) \quad (4.4.31)$$

Hence, we can define the particle current as

$$\mathbf{J}_p = \frac{i\hbar}{2m} (\psi \nabla \psi^* - \psi^* \nabla \psi) \quad (4.4.32)$$

When a stationary eigenstate is substituted into the above, due to the product of the function and its complex conjugate above, the particle current becomes time independent. Hence, stationary states can only give rise to non-time-varying current (static or DC current), and such current does not radiate according to electromagnetic theory.

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<sup>4</sup>We will learn how to treat such coupling later on.



## Chapter 5

# Mathematical Preliminaries

### 5.1 A Function is a Vector

A function is actually an equivalence of a vector. In linear algebra, we denote a vector as  $\mathbf{v} = [v_1, v_2, v_3, \dots, v_N]^t$ . A short-hand notation for this vector is  $v_j$  where  $j = 1, \dots, N$ . But for a function  $f(x)$ , the countably finite index  $j$  in  $v_j$  is now replaced by an uncountably infinite (nondenumerable) set of indices denoted by  $x$ . Hence, we can think of  $f(x)$  as a vector in an infinite dimensional vector space. An inner product in linear algebra is written as

$$\mathbf{V}^\dagger \cdot \mathbf{U} = \sum_{j=1}^N V_j^* U_j \quad (5.1.1)$$

The analogue of the above for functions is

$$\langle f|g \rangle = \int_{-\infty}^{\infty} dx f(x)^* g(x) \quad (5.1.2)$$

The above is the inner product between two functions which are actually vectors.<sup>1</sup> The left-hand side is the compact notation for the inner product between two state vectors known as Dirac's bra-ket notation. The  $\langle f|$  is the "bra" while  $|g\rangle$  is the "ket". The "bra" can be thought of as the conjugate transpose of a vector, while the "ket" is analogous to an ordinary vector.

The set of functions, just as a set of vectors, can then span a vector space.<sup>2</sup> A vector space for which an inner product is defined as above is known as an inner product space. Moreover, an inner product space which is complete is a Hilbert space. Hilbert spaces can be infinite dimensional. The above inner product facilitates the definition of a norm since  $\langle f|f \rangle$

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<sup>1</sup>In the mathematics literature, the inner product is often denoted as  $\langle f, g \rangle$  or  $(f, g)$ .

<sup>2</sup>When any vector in the space can be represented as a linear superposition of the set of independent vectors, the set of vectors are said to span the vector space. They also form the basis set for spanning the space.

is a positive definite number. Hence, the norm of a vector can be defined to be

$$\|f\| = (\langle f|f \rangle)^{1/2} \quad (5.1.3)$$

It is the measure of the length of the vector. We can use this norm to define the distance between two vectors to be

$$d(f, g) = \|f - g\| \quad (5.1.4)$$

In general, if we have a set of orthonormal eigenfunctions,  $\{\psi_n(x), n = 1, \dots, \infty\}$ , that spans a linear vector space, we can expand an arbitrary function in the same space as

$$g(x) = \sum_n d_n \psi_n(x) \quad (5.1.5)$$

The set  $\{\psi_n(x), n = 1, \dots, \infty\}$  also forms the orthonormal basis or the orthonormal basis set for spanning the vector space. A member of the set is known as a basis function or a basis vector.<sup>3</sup> Eigenfunctions of an operator can be used as basis functions.

The above can be written using Dirac's notation as

$$|g\rangle = \sum_n d_n |\psi_n\rangle \quad (5.1.6)$$

Inner product the above with  $\langle\psi_m|$  from the left, we arrive at that

$$\langle\psi_m|g\rangle = \sum_n d_n \langle\psi_m|\psi_n\rangle \quad (5.1.7)$$

Using the orthonormality of the eigenfunction such that  $\langle\psi_m|\psi_n\rangle = \delta_{nm}$ , the above yields that

$$d_m = \langle\psi_m|g\rangle \quad (5.1.8)$$

Consequently, we have

$$|g\rangle = \sum_n |\psi_n\rangle \langle\psi_n|g\rangle \quad (5.1.9)$$

We can identify the operator

$$\hat{I} = \sum_n |\psi_n\rangle \langle\psi_n| \quad (5.1.10)$$

as an identity operator since when it operates on a vector, it returns the same vector. We can often construct an identity operator of a vector space once we have identify the set of orthonormal vectors that span the space. The product  $|\psi_n\rangle \langle\psi_n|$  in Dirac's notation, is analogous to outer product between two vectors  $\mathbf{U} \cdot \mathbf{V}^\dagger$  in linear algebraic notation. This

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<sup>3</sup>We will use "function" and "vector" interchangeably since they are the same.

outer product produces a matrix, and in the case of infinite dimensional linear vector space, produces an operator.

For instance, if we are in a 3-space (the short for 3D vector space), the unit vectors that span the 3-space are  $\{\hat{a}_1, \hat{a}_2, \hat{a}_3\}$ . They are complete and orthonormal. Hence, the identity operator in a 3-space is<sup>4</sup>

$$\bar{\mathbf{I}} = \sum_{i=1}^3 \hat{a}_i \hat{a}_i \tag{5.1.11}$$

To be strictly correct, the above is an outer product between two vectors, and a transpose sign should be attached to one of the unit vectors. But it is customary in the physics literature to ignore this transpose sign. Written in terms of  $x, y, z$  notation, the identity operator becomes

$$\bar{\mathbf{I}} = \hat{x}\hat{x} + \hat{y}\hat{y} + \hat{z}\hat{z} \tag{5.1.12}$$

One can easily be convinced that  $\bar{\mathbf{I}} \cdot \mathbf{A} = \mathbf{A}$ , confirming the above.<sup>5</sup>

In an  $N$  dimensional vector space spanned by a set of orthonormal vectors  $\{\mathbf{U}_n, n = 1, \dots, N\}$ , the identity operator is formed by their outer products; namely,

$$\bar{\mathbf{I}} = \sum_{n=1}^N \mathbf{U}_n \cdot \mathbf{U}_n^\dagger \tag{5.1.13}$$

One can easily show that when  $\bar{\mathbf{I}}$  operates on a member of the vector space, it returns the same member. Namely,

$$\bar{\mathbf{I}} \cdot \mathbf{V} = \sum_{n=1}^N \mathbf{U}_n \cdot \mathbf{U}_n^\dagger \cdot \mathbf{V} = \sum_{n=1}^N \mathbf{U}_n (\mathbf{U}_n^\dagger \cdot \mathbf{V}) = \mathbf{V} \tag{5.1.14}$$

The last equality follows because the second last term is just an orthonormal eigenvector expansion of the vector  $\mathbf{V}$ .

A vector space is also defined by a set of axioms. For  $\mathbf{u}, \mathbf{v}, \mathbf{w}$  that belong to a vector space, the following axioms hold:

Associativity of addition  $\mathbf{u} + (\mathbf{v} + \mathbf{w}) = (\mathbf{u} + \mathbf{v}) + \mathbf{w}$  (5.1.15)

Commutativity of addition  $\mathbf{u} + \mathbf{v} = \mathbf{v} + \mathbf{u}$  (5.1.16)

Identity element of addition  $\mathbf{v} + \mathbf{0} = \mathbf{v}$  (5.1.17)

Inverse elements of addition  $\mathbf{v} + (-\mathbf{v}) = \mathbf{0}$  (5.1.18)

Distributivity of scalar multiplication over vectors  $a(\mathbf{u} + \mathbf{v}) = a\mathbf{u} + a\mathbf{v}$  (5.1.19)

Distributivity of scalar multiplication by a vector  $(a + b)\mathbf{v} = a\mathbf{v} + b\mathbf{v}$  (5.1.20)

Compatibility of scalar multiplication  $a(b\mathbf{v}) = (ab)\mathbf{v}$  (5.1.21)

Identity element of scalar multiplication  $1\mathbf{v} = \mathbf{v}$  (5.1.22)

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<sup>4</sup>Throughout these lecture notes, the hat (“^”) symbol is used to denote an operator, but here, it denotes a unit vector.

<sup>5</sup>It is to be noted that an outer product in matrix notation is  $\mathbf{U} \cdot \mathbf{V}^\dagger$ , while in physics notation for 3 space, it is often just written as  $\mathbf{EB}$ . In Dirac notation, an outer product is  $|\psi\rangle\langle\phi|$ .

## 5.2 Operators

An operator maps vectors from one space to vectors in another space. It is denoted mathematically as  $\hat{A} : V \rightarrow W$  where  $\hat{A}$  is the operator, while  $V$  and  $W$  are two different vector spaces:  $V$  is the domain space, while  $W$  is the range space. In linear algebra, the operator is a matrix operator. In Hilbert spaces, it can be a differential operator such as

$$g(x) = \frac{d}{dx}f(x) \quad (5.2.1)$$

It can be an integral operator such as a Fourier transform operator

$$g(k) = \int_{-\infty}^{\infty} dx e^{ikx} f(x) \quad (5.2.2)$$

In Dirac notation,

$$|g\rangle = \hat{A}|f\rangle \quad (5.2.3)$$

Linear operators are defined such that

$$\hat{L}(c_1|g_1\rangle + c_2|g_2\rangle) = c_1\hat{L}|g_1\rangle + c_2\hat{L}|g_2\rangle \quad (5.2.4)$$

It is quite clear that matrix operators satisfy the above, and hence, they are linear. They are also known as linear maps. In general, like matrix operators, linear operators are not commutative; namely

$$\hat{A}\hat{B} \neq \hat{B}\hat{A} \quad (5.2.5)$$

### 5.2.1 Matrix Representation of an Operator

An operator equation can be written as<sup>6</sup>

$$|g\rangle = \hat{A}|f\rangle \quad (5.2.6)$$

We can convert the above into a matrix equation by inserting an identity operator on the right-hand side to give

$$|g\rangle = \sum_n \hat{A}|\psi_n\rangle\langle\psi_n|f\rangle \quad (5.2.7)$$

Furthermore, we can multiply the above from the left by the basis vector  $\langle\psi_m|$ ,  $m = 1, \dots, \infty$  to yield<sup>7</sup>

$$\langle\psi_m|g\rangle = \sum_n \langle\psi_m|\hat{A}|\psi_n\rangle\langle\psi_n|f\rangle, \quad m = 1, \dots, \infty \quad (5.2.8)$$

<sup>6</sup>In the mathematics literature, this is often just denoted as  $\hat{A}f = g$ .

<sup>7</sup>This process is called testing or weighting, and  $\psi_m$  is called the testing or weighting function.



The above is an infinite dimensional matrix equation which can be written as

$$\mathbf{g} = \overline{\mathbf{A}} \cdot \mathbf{f} \quad (5.2.9)$$

where

$$[\mathbf{A}]_{mn} = \langle \psi_m | \hat{A} | \psi_n \rangle \quad (5.2.10)$$

$$[\mathbf{g}]_m = \langle \psi_m | g \rangle \quad (5.2.11)$$

$$[\mathbf{f}]_n = \langle \psi_n | f \rangle \quad (5.2.12)$$

The matrix equation can be solved approximately by truncating its size to  $N \times N$ , or without truncation, it can be solved iteratively.

The matrix denoted by

$$A_{mn} = \langle \psi_m | \hat{A} | \psi_n \rangle \quad (5.2.13)$$

is the matrix representation of the operator  $\hat{A}$ . By the same token,  $\langle \psi_m | g \rangle$  and  $\langle \psi_n | f \rangle$  are the vector representations of the functions  $g$  and  $f$  respectively.

In the above, we have assumed that the range space and the domain space of the operator are the same, and hence, they can be spanned by the same basis set. For a Hermitian operator, this is usually the case. However, for some operators where the range space and the domain space are different, we may choose to test (5.2.7) with a different set of basis functions.

## 5.2.2 Bilinear Expansion of an Operator

We have seen how the use of the identity operator allows us to expand a function in terms of a set of basis functions as in (5.1.9). The same can be done with an operator. Pre- and post-multiply an operator with the identity operator as given by (5.1.10), we have

$$\hat{A} = \sum_n \sum_m |\psi_n\rangle \langle \psi_n | \hat{A} | \psi_m\rangle \langle \psi_m | \quad (5.2.14)$$

The above can be rewritten as

$$\hat{A} = \sum_n \sum_m |\psi_n\rangle A_{nm} \langle \psi_m | = \sum_n \sum_m A_{nm} |\psi_n\rangle \langle \psi_m | \quad (5.2.15)$$

where  $A_{nm}$  is the matrix representation of the operator  $\hat{A}$ . The above is the bilinear expansion of the operator in terms of orthonormal functions. Notice that the expansion of an identity operator given by (5.1.10) is a bilinear expansion.

## 5.2.3 Trace of an Operator

The trace of a matrix operator is defined to be the sum of its diagonal elements; namely,

$$\text{tr}(\overline{\mathbf{M}}) = \sum_{i=1}^N M_{ii} \quad (5.2.16)$$

If an operator  $\hat{A}$  has matrix representation given by  $\langle\psi_i|\hat{A}|\psi_j\rangle$ , the trace of the operator  $\hat{A}$  is defined to be

$$\text{tr}(\hat{A}) = \sum_i \langle\psi_i|\hat{A}|\psi_i\rangle \quad (5.2.17)$$

It can be shown that the trace of an operator is independent of the basis used for its representation. To this end, we insert the identity operator

$$\hat{I} = \sum_m |\phi_m\rangle\langle\phi_m| \quad (5.2.18)$$

into (5.2.17) to get

$$\text{tr}(\hat{A}) = \sum_i \sum_m \langle\psi_i|\phi_m\rangle\langle\phi_m|\hat{A}|\psi_i\rangle \quad (5.2.19)$$

Exchanging the order of summation above, and the order of the two scalar numbers in the summand, we have

$$\text{tr}(\hat{A}) = \sum_m \sum_i \langle\phi_m|\hat{A}|\psi_i\rangle\langle\psi_i|\phi_m\rangle \quad (5.2.20)$$

The inner summation reduces to an identity operator which can be removed, and the above becomes

$$\text{tr}(\hat{A}) = \sum_m \langle\phi_m|\hat{A}|\phi_m\rangle \quad (5.2.21)$$

This is the trace of  $\hat{A}$  using another basis set that is complete and orthonormal. Hence, the trace of an operator is invariant to the choice of basis.

If we choose a basis function that is the eigenfunction of  $\hat{A}$  such that  $\hat{A}|\psi_n\rangle = \lambda_n|\psi_n\rangle$ , then (5.2.19) becomes

$$\text{tr}(\hat{A}) = \sum_i \sum_m \lambda_i \langle\psi_i|\phi_m\rangle\langle\phi_m|\psi_i\rangle = \sum_i \lambda_i \quad (5.2.22)$$

Hence, the trace of an operator is also the sum of its eigenvalues.

It can also be shown that

$$\text{tr}(\hat{A}\hat{B}) = \text{tr}(\hat{B}\hat{A}) \quad (5.2.23)$$

This is quite easy to show for matrix operators since

$$\text{tr}(\bar{\mathbf{A}} \cdot \bar{\mathbf{B}}) = \sum_i \sum_j A_{ij} B_{ji} = \sum_j \sum_i B_{ji} A_{ij} \quad (5.2.24)$$

Trace is usually used in quantum mechanics as an alternative way to write the expectation value of an operator. As mentioned before, for a quantum system in a state defined by the state function  $|\psi\rangle$ , the expectation value of a quantum operator in such a state is

$$\langle\psi|\hat{A}|\psi\rangle \quad (5.2.25)$$

For denumerable indices, the above is analogous to

$$\mathbf{U}^\dagger \cdot \bar{\mathbf{A}} \cdot \mathbf{U} = \sum_n \sum_m U_n^* A_{nm} U_m = \sum_n \sum_m A_{nm} U_m U_n^* = \text{tr}(\bar{\mathbf{A}} \cdot \mathbf{U} \cdot \mathbf{U}^\dagger) \quad (5.2.26)$$

where  $U_m U_n^*$  is the outer product of two vectors  $U_m$  and  $U_n$ . Converting the above to Dirac notation, we have

$$\langle \psi | \hat{A} | \psi \rangle = \text{tr}\{\hat{A} | \psi \rangle \langle \psi | \} \quad (5.2.27)$$

The operator

$$\hat{\rho} = | \psi \rangle \langle \psi | \quad (5.2.28)$$

is known as the density operator. It is used in quantum mechanics to denote the state of a quantum system as an alternative to the state vector.

## 5.2.4 Unitary Operators

Unitary operators satisfy the property that

$$\bar{\mathbf{U}}^\dagger \cdot \bar{\mathbf{U}} = \bar{\mathbf{I}} \quad (5.2.29)$$

Unitary operators are operators that when acted on a vector does not change its length. In matrix notation, a unitary transform is

$$\bar{\mathbf{U}} \cdot \mathbf{V} = \mathbf{V}' \quad (5.2.30)$$

The length square of the new vector  $\mathbf{V}'$  is defined to be

$$\mathbf{V}'^\dagger \cdot \bar{\mathbf{U}}^\dagger \cdot \bar{\mathbf{U}} \cdot \mathbf{V} = \mathbf{V}'^\dagger \cdot \mathbf{V} \quad (5.2.31)$$

Making use of (5.2.29), the above implies that

$$\mathbf{V}'^\dagger \cdot \mathbf{V} = \mathbf{V}'^\dagger \cdot \mathbf{V}' \quad (5.2.32)$$

or that the length of the vector has not changed. For unitary matrix, it is clear that  $\bar{\mathbf{U}}^\dagger = \bar{\mathbf{U}}^{-1}$ . The above can be rewritten using Dirac notation. Since in quantum mechanics,  $\langle \psi | \psi \rangle = 1$ , the length of the vector is 1. The time evolution operator by integrating Schrödinger equation given below

$$\hat{\tau} = e^{-\frac{i}{\hbar} \hat{H} t} \quad (5.2.33)$$

where  $\hat{H}$  is the Hamiltonian operator, is a unitary operator, since the state vector it acts on cannot change its length.

Since many basis set comprises orthonormal basis vectors, a unitary matrix is needed for the change of basis from one set to another.

### 5.2.5 Hermitian Operators

Hermitian operators appear frequently in quantum mechanics since they have real eigenvalues. As such, the expectation value of such operators are real valued, so that it can be connected to measurable or observable quantities in the real world.

A Hermitian matrix is one where its conjugate transpose (also called Hermitian transpose or Hermitian conjugate) is itself. For operators, the term “adjoint” is often used. A Hermitian operator is also called a self-adjoint operator. The notation for this is

$$\hat{M}^\dagger = \hat{M} \quad (5.2.34)$$

For matrix operator with denumerable indices, the above is the same as

$$M_{ij}^* = M_{ji} \quad (5.2.35)$$

For operators with nondenumerable indices, it is better to define the adjoint of the operator via inner products. Using the rule for conjugate transpose, we have

$$\langle f | \hat{M} | g \rangle^* = \langle f | \hat{M} | g \rangle^\dagger = \langle g | \hat{M}^\dagger | f \rangle \quad (5.2.36)$$

The first equality follows because the above are scalar quantities: hence, the conjugate transpose of a scalar is the same as its conjugation. To obtain the second equality, we have used the rule that

$$(\overline{\mathbf{A} \cdot \mathbf{B} \cdot \mathbf{C}})^t = \overline{\mathbf{C}^t \cdot \mathbf{B}^t \cdot \mathbf{A}^t} \quad (5.2.37)$$

in linear algebra, where the matrices need not be square. Hence,  $\overline{\mathbf{A}}$  and  $\overline{\mathbf{C}}$  can be  $1 \times N$  and  $N \times 1$  respectively representing vectors. This identity immediately implies that

$$(\overline{\mathbf{A} \cdot \mathbf{B} \cdot \mathbf{C}})^\dagger = \overline{\mathbf{C}^\dagger \cdot \mathbf{B}^\dagger \cdot \mathbf{A}^\dagger} \quad (5.2.38)$$

These rules can easily be shown to apply to inner product of operators and vectors under the Dirac notation. Hence, (5.2.36) defines the adjoint of the operator  $\hat{M}^\dagger$  when  $f$  and  $g$  are arbitrary. It does so by using just inner products.

A Hermitian operator has real eigenvalues and their eigenvectors are orthogonal. The proof is analogous to the proof for matrix operators, for which we will provide next. Given

$$\overline{\mathbf{M}} \cdot \mathbf{v}_i = \lambda_i \mathbf{v}_i \quad (5.2.39)$$

$$\overline{\mathbf{M}} \cdot \mathbf{v}_j = \lambda_j \mathbf{v}_j \quad (5.2.40)$$

Dot-multiply the first equation by  $\mathbf{v}_j^\dagger$  from the left, and likewise for the second equation with  $\mathbf{v}_i^\dagger$ , we have

$$\mathbf{v}_j^\dagger \cdot \overline{\mathbf{M}} \cdot \mathbf{v}_i = \lambda_i \mathbf{v}_j^\dagger \cdot \mathbf{v}_i \quad (5.2.41)$$

$$\mathbf{v}_i^\dagger \cdot \overline{\mathbf{M}} \cdot \mathbf{v}_j = \lambda_j \mathbf{v}_i^\dagger \cdot \mathbf{v}_j \quad (5.2.42)$$

We can take the conjugate transpose of the first equation, using the rule in (5.2.38), the left-hand side becomes the same as that for the second equation after making use of the Hermitian property of the matrix  $\bar{\mathbf{M}}$ . On subtracting the two equations, the following equation ensues:

$$0 = (\lambda_i^* - \lambda_j) \mathbf{v}_i^\dagger \cdot \mathbf{v}_j \quad (5.2.43)$$

If  $i = j$ , on the right-hand side, we have  $\mathbf{v}_i^\dagger \cdot \mathbf{v}_i = |\mathbf{v}_i|^2$  which is a positive definite number. The above can be zero only if  $\lambda_i^* = \lambda_i$  or that  $\lambda_i$  is real. On the other hand, if  $\lambda_i \neq \lambda_j$ , it is necessary that  $\mathbf{v}_i^\dagger \cdot \mathbf{v}_j = 0$ . In other words,

$$\mathbf{v}_i^\dagger \cdot \mathbf{v}_j = C_n \delta_{ij} \quad (5.2.44)$$

The above proof can be repeated using Dirac notation, and the conclusion will be the same. The eigenvectors of a Hermitian operator are also complete in the space that the operator acts on. It is obvious in the finite dimensional case, but not so obvious in the infinite dimensional space.

We can use the rule expressed in (5.2.36) to see if an operator is Hermitian. For instance, the momentum operator is  $\hat{p}$ . Using (5.2.36), we have

$$\langle f | \hat{p} | g \rangle^* = \langle g | \hat{p}^\dagger | f \rangle \quad (5.2.45)$$

The above defines the adjoint of the operator  $\hat{p}$ . Writing the above explicitly in 1D space using coordinate space representation,<sup>8</sup> we have on the left-hand side

$$\int_{-\infty}^{\infty} dx f(x) i\hbar \frac{d}{dx} g^*(x) = \int_{-\infty}^{\infty} dx g^*(x) \left( -i\hbar \frac{d}{dx} \right) f(x) \quad (5.2.46)$$

We have arrived at the form of the right-hand side by using integration by parts, and assuming that the functions are vanishing at infinity. By comparing the above, we identify that

$$\hat{p}^\dagger = -i\hbar \frac{d}{dx} \quad (5.2.47)$$

We see that  $(\hat{p})^\dagger = \hat{p}$  implying that it is Hermitian. The eigenfunction of the momentum operator is  $f_k(x) \propto e^{ikx}$ . It is quite clear that  $\hat{p}f_k(x) = \hbar k f_k(x)$ , and hence, its eigenvalues are also real, as is required of a Hermitian operator. The above can be generalized to 3D. It can also be shown that the kinetic energy operator

$$\hat{T} = \hat{p}^2 / 2m = -\frac{\hbar \nabla^2}{2m} \quad (5.2.48)$$

is Hermitian. In 3D,  $\hat{p}^2$  is proportional to the Laplacian operator  $\nabla^2$  in coordinate space representation. We can use integration by parts in 3D to show that the above operator is also Hermitian. Using the fact that

$$\nabla \cdot [f(\mathbf{r}) \nabla g^*(\mathbf{r})] = f(\mathbf{r}) \nabla^2 g^*(\mathbf{r}) + \nabla f(\mathbf{r}) \nabla g^*(\mathbf{r}) \quad (5.2.49)$$

---

<sup>8</sup>A note is in order here regarding the term “coordinate space”, since there is only one Hilbert space. It is understood that when the term “coordinate space representation” is used, it means that the Hilbert space is represented in the space where the basis is the coordinate space basis.

we have in coordinate space representation

$$\int_V d\mathbf{r} f(\mathbf{r}) \nabla^2 g^*(\mathbf{r}) = \int_V d\mathbf{r} [\nabla \cdot (f(\mathbf{r}) \nabla g^*(\mathbf{r})) - \nabla f(\mathbf{r}) \nabla g^*(\mathbf{r})] = - \int_V d\mathbf{r} \nabla f(\mathbf{r}) \nabla g^*(\mathbf{r}) \quad (5.2.50)$$

The last form is symmetric between  $f$  and  $g$ , and hence, we can easily show that

$$\int_V d\mathbf{r} g^*(\mathbf{r}) \nabla^2 f(\mathbf{r}) = - \int_V d\mathbf{r} \nabla g^*(\mathbf{r}) \nabla f(\mathbf{r}) \quad (5.2.51)$$

Consequently, we can show that

$$\begin{aligned} \langle f | \hat{\nabla}^2 | g \rangle^* &= \left[ \int_V d\mathbf{r} f^*(\mathbf{r}) \nabla^2 g(\mathbf{r}) \right]^* \\ &= \int_V d\mathbf{r} g^*(\mathbf{r}) \nabla^2 f(\mathbf{r}) \\ &= \langle g | \hat{\nabla}^2 | f \rangle \end{aligned} \quad (5.2.52)$$

indicating the Hermitian property of the Laplacian operator. Consequently, the  $\hat{T}$  operator is also Hermitian or self-adjoint.

If we have chosen  $f = g$  in the above, (5.2.51) is always negative, implying that the Laplacian operator is a negative definite operator.

### 5.3 Identity Operator in a Continuum Space

We have seen the derivation of the identity operator when the basis functions are indexed by integers, such as the eigenfunctions of an infinite potential well, or those of a harmonic oscillator in a parabolic potential well. In many situations, we are required to work with indices that are a continuum (nondenumerable) like the coordinate space indices  $x, y, z$ . When an electron is freely roaming, its energy values are also a continuum.

We have seen how we have used denumerable eigenfunctions, such as  $\psi_n$  to project a function into its vector representation. Namely,

$$f_n = \langle \psi_n | f \rangle \quad (5.3.1)$$

Very often, a denumerable basis function is just written as  $\langle n |$ , and the above becomes

$$f_n = \langle n | f \rangle \quad (5.3.2)$$

Similar, we can think of a coordinate (or position) basis function  $\langle p_x |$  whose property is that its inner product with  $|f\rangle$  yields the value of the function  $f$  at position  $x$ ; namely,

$$f(x) = \langle p_x | f \rangle \quad (5.3.3)$$

The above is often abbreviated as

$$f(x) = \langle x | f \rangle \quad (5.3.4)$$

Assuming that this basis  $|x\rangle$  is complete and orthogonal in  $0 < x < a$ , then we can define an identity operator such that

$$\bar{\mathbf{I}} = \int_0^a dx' |x'\rangle \langle x'| \quad (5.3.5)$$

so that

$$|f\rangle = \int_0^a dx' |x'\rangle \langle x'| f\rangle = \int_0^a dx' |x'\rangle f(x') \quad (5.3.6)$$

Notice that the above is quite different from the identity operator when the basis functions are denumerable, which is

$$\bar{\mathbf{I}} = \sum_{n=1}^{\infty} |f_n\rangle \langle f_n| \quad (5.3.7)$$

Taking the product of (5.3.6) with  $\langle x|$ , we have

$$\langle x|f\rangle = f(x) = \int_0^a dx' \langle x|x'\rangle f(x') \quad (5.3.8)$$

Hence, in order for (5.3.5) to be an identity operator, the basis function must satisfy

$$\langle x|x'\rangle = \delta(x - x') \quad (5.3.9)$$

Notice that in the denumerable case, the orthonormal relation of the basis function is expressed in terms of Kronecker delta function  $\delta_{ij}$ , but in the continuum case, the equivalent relationship is expressed in terms of the Dirac delta function  $\delta(x - x')$ .

We note that  $|x\rangle$  is not orthonormal in the strict sense that (5.3.9) is infinite when  $x = x'$ . The identity operator (5.3.5) is different from (5.3.7) because of the extra weight  $dx$  in the integral summation. One may think that  $|x\rangle\sqrt{dx'}$  as analogous to the orthonormal vector of the countable case.

The inner product between two vectors is written as

$$\langle f|g\rangle = \int_0^a dx \langle f|x\rangle \langle x|g\rangle = \int_0^a dx f^*(x)g(x) \quad (5.3.10)$$

where we have inserted the identity operator (5.3.5) in the first expression above to get the second expression, and making use of (5.3.4) to get the third expression. Furthermore, we note that  $\langle f|x\rangle$  is the complex conjugate of  $\langle x|f\rangle$ , since  $\langle f|$  and  $|x\rangle$  are conjugate transpose of  $|f\rangle$  and  $\langle x|$ , respectively.

A vector  $|f\rangle$  can have other representations. For example, we can define a set of “orthonormal” vectors  $\langle k|$  such that

$$\langle k|f\rangle = f(k), \quad (5.3.11)$$

where  $f(k)$  is the Fourier transform of  $f(x)$  via the relationship

$$f(x) = \frac{1}{\sqrt{2\pi}} \int_{-\infty}^{\infty} dk e^{ikx} f(k) \quad (5.3.12)$$

The above can be written as

$$f(x) = \langle x|f \rangle = \int_{-\infty}^{\infty} dk \langle x|k \rangle \langle k|f \rangle \quad (5.3.13)$$

where we have defined the identity operator

$$I = \int_{-\infty}^{\infty} dk |k \rangle \langle k| \quad (5.3.14)$$

Comparing (5.3.11), (5.3.12) and (5.3.13), we deduce that

$$\langle x|k \rangle = \frac{1}{\sqrt{2\pi}} e^{ikx} \quad (5.3.15)$$

In other words, the coordinate space representation of the vector  $|k \rangle$  is given by (5.3.15). Since  $\hbar k$  in quantum mechanics is related to the momentum of a particle,  $f(k) = \langle k|f \rangle$  is also called the momentum representation of the vector  $|f \rangle$ , while  $f(x) = \langle x|f \rangle$  is the coordinate space representation of the vector  $|f \rangle$ .

From the above, we can derive Parseval's theorem

$$\begin{aligned} \langle f|g \rangle &= \int_0^a dx \langle f|x \rangle \langle x|g \rangle = \int_0^a dx f^*(x) g(x) \\ &= \int_{-\infty}^{\infty} dk \langle f|k \rangle \langle k|g \rangle = \int_{-\infty}^{\infty} dk f^*(k) g(k) \end{aligned} \quad (5.3.16)$$

Parseval's theorem is the statement that the inner product between two vectors is invariant with respect to the basis that represent it.

The above can be generalized to 3-space where a position basis function is  $\langle \mathbf{r}|$ , with the property that

$$f(\mathbf{r}) = \langle \mathbf{r}|f \rangle \quad (5.3.17)$$

The identity operator can then be expressed as

$$\bar{\mathbf{I}} = \int_V d\mathbf{r}' |\mathbf{r}' \rangle \langle \mathbf{r}'| \quad (5.3.18)$$

And the “orthonormality” relationship is

$$\langle \mathbf{r}|\mathbf{r}' \rangle = \delta(\mathbf{r} - \mathbf{r}') \quad (5.3.19)$$

Similar basis can be defined for the 3D  $\mathbf{k}$  space.



## 5.4 Changing Between Representations

If we have an operator equation such as

$$\hat{p}|\psi\rangle = |g\rangle \quad (5.4.1)$$

The above equation is analogous to the equation  $\bar{\mathbf{P}} \cdot \mathbf{f} = \mathbf{g}$  in matrix algebra. The explicit forms and the representations of the matrix operator  $\bar{\mathbf{P}}$  and the vector  $\mathbf{f}$  are not given. In order to express the above explicitly, we first insert the identity operator to transform (5.4.1) to

$$\int dx \hat{p}|x\rangle\langle x|\psi\rangle = |g\rangle \quad (5.4.2)$$

Then multiply the above from the left by  $\langle x'|$ , we have

$$\int dx \langle x'|\hat{p}|x\rangle\langle x|\psi\rangle = \langle x'|g\rangle \quad (5.4.3)$$

$$\int dx p(x', x)\psi(x) = g(x') \quad (5.4.4)$$

The above is the coordinate space representation of (5.4.1). It is valid for any operator in a 1D space. But if  $\hat{p}$  is the momentum operator, we know that the above equation in coordinate space representation should be

$$-i\hbar \frac{d}{dx}\psi(x) = g(x') \quad (5.4.5)$$

Hence we conclude that

$$p(x', x) = -i\hbar\delta(x' - x)\frac{d}{dx}\psi(x) \quad (5.4.6)$$

Therefore, a differential operator is a highly localized operator, or a quasi-diagonal operator.

Next, we will study the action of the momentum operator  $\hat{p}$  on a momentum eigenstate  $|k\rangle$ ; namely,

$$\hat{p}|k\rangle \quad (5.4.7)$$

Transforming the above to coordinate space representation similar to what have been done above, we have

$$\int dx \delta(x' - x) \left(-i\hbar \frac{d}{dx}\right) \langle x|k\rangle = \left(-i\hbar \frac{d}{dx}\right) \langle x|k\rangle \quad (5.4.8)$$

The expression for  $\langle x|k\rangle$  is given in equation (5.3.15). Hence, the above becomes

$$\hbar k \langle x|k\rangle \quad (5.4.9)$$

Therefore, we see that  $|k\rangle$  is an eigenvector of the momentum operator  $\hat{p}$ . In other words,

$$\hat{p}|k\rangle = \hbar k|k\rangle \quad (5.4.10)$$

In fact, any function of  $\hat{p}$ , when operating on a momentum eigenstate gives

$$f(\hat{p})|k\rangle = f(\hbar k)|k\rangle \quad (5.4.11)$$

We can repeat (5.4.2) and (5.4.3) using the basis  $|k\rangle$  to arrive at

$$\int dk \langle k'|\hat{p}|k\rangle \langle k|\psi\rangle = \langle k'|g\rangle \quad (5.4.12)$$

Since  $|k\rangle$  is an eigenvector of  $\hat{p}$ , the above becomes

$$\int dk \hbar k \langle k'|k\rangle \langle k|\psi\rangle = \int dk \hbar k \delta(k' - k) \langle k|\psi\rangle = \hbar k' \langle k'|\psi\rangle = \langle k'|g\rangle \quad (5.4.13)$$

which is just a scalar equation.

### 5.4.1 The Coordinate Basis Function

Using the coordinate basis function  $|\mathbf{r}\rangle$ , we can write many of our previously derived identities in coordinate space representation. Say if the functions defined in (5.1.9) and (5.1.10) are functions of 3-space, indexed by  $\mathbf{r}$ , then multiplying (5.1.9) from the left by  $\langle \mathbf{r}|$ ,

$$\langle \mathbf{r}|g\rangle = g(\mathbf{r}) = \sum_n \langle \mathbf{r}|\psi_n\rangle \langle \psi_n|g\rangle = \sum_n \psi_n(\mathbf{r}) \langle \psi_n|g\rangle \quad (5.4.14)$$

We can pre- and post-multiply the identity operator defined in (5.1.10) by  $\langle \mathbf{r}|$  and  $|\mathbf{r}'\rangle$  respectively, we can identify the operator

$$\langle \mathbf{r}|\hat{I}|\mathbf{r}'\rangle = \langle \mathbf{r}|\mathbf{r}'\rangle = \delta(\mathbf{r} - \mathbf{r}') = \sum_n \langle \mathbf{r}|\psi_n\rangle \langle \psi_n|\mathbf{r}'\rangle = \sum_n \psi_n(\mathbf{r}) \psi_n^*(\mathbf{r}') \quad (5.4.15)$$

where we have noted that  $\langle \psi_n|\mathbf{r}'\rangle = \langle \mathbf{r}'|\psi_n\rangle^*$  since they are conjugate transpose of each other. The above is the bilinear eigenfunction expansion of the Dirac delta function. We can also apply the above to the bilinear expansion of an operator given by (5.2.15). Going through similar operations, we have

$$\langle \mathbf{r}|\hat{A}|\mathbf{r}'\rangle = A(\mathbf{r}, \mathbf{r}') = \sum_n \sum_m A_{nm} \psi_n(\mathbf{r}) \psi_m^*(\mathbf{r}') \quad (5.4.16)$$

We have let  $A(\mathbf{r}, \mathbf{r}')$  be the coordinate space representation of the operator  $\hat{A}$ . The above is just the bilinear expansion in coordinate space representation.

We can also choose the momentum basis set  $|\mathbf{k}\rangle$  (or any other basis sets) and project the previously obtained identities in momentum representation (or any other representations).

## 5.5 Commutation of Operators

Operators, like matrices, are non-commuting for most parts. In other words,

$$\hat{A}\hat{B} \neq \hat{B}\hat{A} \quad (5.5.1)$$

The commutator is defined to be

$$[\hat{A}, \hat{B}] = \hat{A}\hat{B} - \hat{B}\hat{A} \quad (5.5.2)$$

If  $\hat{A}$  and  $\hat{B}$  are commuting, then

$$[\hat{A}, \hat{B}] = 0 \quad (5.5.3)$$

But for most parts,

$$[\hat{A}, \hat{B}] = i\hat{C} \quad (5.5.4)$$

It can be shown easily that  $\hat{C}$  is a Hermitian operator if  $\hat{A}$  and  $\hat{B}$  are Hermitian.

If  $\hat{A}$  and  $\hat{B}$  are commuting, they share the same set of eigenfunctions, but not the same eigenvalues. Say if  $|\psi_n\rangle$  is an eigenfunction of  $\hat{B}$ , with eigenvalue  $B_n$ , then

$$\hat{B}\hat{A}|\psi_n\rangle = \hat{A}\hat{B}|\psi_n\rangle = \hat{A}B_n|\psi_n\rangle = B_n\hat{A}|\psi_n\rangle \quad (5.5.5)$$

From the above, it is clear that  $\hat{A}|\psi_n\rangle$  is also eigenfunction of  $\hat{B}$  with the same eigenvalue  $B_n$ . Therefore,  $\hat{A}|\psi_n\rangle$  is proportional to  $|\psi_n\rangle$ , or that  $\hat{A}|\psi_n\rangle = F_n|\psi_n\rangle$ . In other words,  $|\psi_n\rangle$  is also an eigenfunction of  $\hat{A}$ , but with a different eigenvalue.

## 5.6 Expectation Value and Eigenvalue of Operators

The expectation value of a random variable is a concept from probability. In quantum mechanics, a measurable quantity or an observable in the real world is a random variable. For each observable in the real world, there is a corresponding operator in the quantum world.

The expectation value of an operator for a quantum system in state  $|f\rangle$  is defined to be

$$\langle \hat{A} \rangle = \langle f | \hat{A} | f \rangle \quad (5.6.1)$$

It is prudent to understand this expectation value in terms of the eigenstates and eigenvalues of  $\hat{A}$ . Letting

$$|f\rangle = \sum_n |\psi_n\rangle f_n \quad (5.6.2)$$

then (5.6.1) can be rewritten as

$$\langle \hat{A} \rangle = \sum_n \sum_m \langle \psi_n | \hat{A} | \psi_m \rangle f_n^* f_m \quad (5.6.3)$$

When  $\psi_m$  is the eigenstate of  $\hat{A}$  with eigenvalue  $A_m$ , the above can be written as

$$\langle \hat{A} \rangle = \sum_n \sum_m A_m \langle \psi_n | \psi_m \rangle f_n^* f_m = \sum_n A_n |f_n|^2 \quad (5.6.4)$$

where  $|f_n|^2$  is the probability of finding the state in eigenstate  $n$ . Hence, the above is the statistical average or the expectation value of the eigenvalue  $A_n$ . The eigenvalue  $A_n$  can be thought of as a random variable. Hence, the expectation value of an operator is the average value of its eigenvalue when the quantum state is in state  $f$ . Therefore, it is customary to denote

$$\langle \hat{A} \rangle = \langle A \rangle = \langle f | \hat{A} | f \rangle = \bar{A} \quad (5.6.5)$$

where the scalar variable  $A$  denotes the eigenvalue of  $\hat{A}$ , which is a random variable, while the angular brackets on the left-hand side  $\langle A \rangle$ , indicate statistical average. An overbar in the last equality above is often used as a short-hand for angular brackets to denote an average of a random variable.

The above indicates that if we prepare a quantum state that is exactly the eigenstate of the quantum operator  $\hat{A}$ , then its expectation value in this quantum eigenstate is just the eigenvalue of the eigenstate. In a sense, we can “measure” the eigenvalue of the quantum operator through such an experiment. Namely,

$$\langle \hat{A} \rangle = \langle \psi_n | \hat{A} | \psi_n \rangle = A_n \quad (5.6.6)$$

where  $A_n$  is the eigenvalue of the eigenstate.

We see in the previous section that when two operators commute, they share the same eigenfunctions but with different eigenvalues. Therefore, if we prepare an eigenstate shared by these two operators, their respective eigenvalues can be “measured” exactly. In other words,

$$\langle \hat{A} \rangle = \langle \psi_n | \hat{A} | \psi_n \rangle = A_n \quad (5.6.7)$$

$$\langle \hat{B} \rangle = \langle \psi_n | \hat{B} | \psi_n \rangle = B_n \quad (5.6.8)$$

On the other hand, if the two operators do not commute, they do not share the same set of eigenstates. If we prepare an eigenstate  $\psi_n$  that is the eigenstate of  $\hat{A}$ , but it is not the eigenstate of  $\hat{B}$ . However, we can expand the eigenstate  $\psi_n$  in terms of the eigenstates of the operator  $\hat{B}$ ; namely,

$$|\psi_n\rangle = \sum_i a_i |\phi_i\rangle \quad (5.6.9)$$

where  $|\phi_i\rangle$  are the eigenstates of  $\hat{B}$ . We now see that the expectation value of  $\hat{B}$  due to the eigenstate  $|\psi_n\rangle$  is

$$\langle \hat{B} \rangle = \langle \psi_n | \hat{B} | \psi_n \rangle = \sum_{ij} a_i^* a_j \langle \phi_i | \hat{B} | \phi_j \rangle = \sum_i |a_i|^2 B_i \quad (5.6.10)$$

The expectation value of  $\hat{B}$  is due to a range of eigenvalues of  $\hat{B}$ , and not just due to one single eigenvalue like before. This is the gist of the uncertainty principle. When two operators do not commute, the precise measurement of the eigenvalue of one operator implies that the measurement of the second operator will involve a range of eigenvalues.

## 5.7 Generalized Uncertainty Principle

We have seen the Heisenberg uncertainty principle expressed for uncertainty in momentum  $\Delta p$  and position  $\Delta x$  as  $\Delta p \Delta x \geq \hbar/2$ . In this case,  $p$  and  $x$  are both observables in the real world. In quantum mechanics, all observables are replaced by operators. To connect the operators to real world observables, we take the expectation values of the operators. That is

$$\bar{A} = \langle A \rangle = \langle f | \hat{A} | f \rangle \quad (5.7.1)$$

where  $A$  is a scalar variable representing the eigenvalue of  $\hat{A}$  and  $f$  is a state vector that defines that state the quantum system is in. The expectation value of an operator also gives the statistical mean of the observable expressed as the mean of the eigenvalue of the corresponding operator.

We can define the deviation from its mean by the operator

$$\Delta \hat{A} = \hat{A} - \bar{A} \quad (5.7.2)$$

On the right-hand side, the first term is an operator while the second term is the mean.<sup>9</sup> The above operator has zero expectation value or zero mean.

We can study the variance of the operator  $(\Delta \hat{A})^2$  which has nonzero expectation value. Namely,

$$\begin{aligned} \langle f | (\Delta \hat{A})^2 | f \rangle &= \langle f | (\hat{A} - \bar{A})^2 | f \rangle \\ &= \langle (A - \bar{A})^2 \rangle \\ &= \langle A^2 - 2\bar{A}\hat{A} + \bar{A}^2 \rangle \\ &= \langle A^2 \rangle - 2\bar{A}\langle A \rangle + \bar{A}^2 \\ &= \langle A^2 \rangle - \bar{A}^2 = \langle (\Delta A)^2 \rangle = \overline{(\Delta A)^2} = \sigma_A^2 \end{aligned} \quad (5.7.3)$$

The above is just the definition of variance of random variable  $A$  as in statistics. The standard deviation is obtained by taking the square root of the variance to get<sup>10</sup>

$$\sigma_A = \sqrt{\overline{(\Delta A)^2}} \quad (5.7.4)$$

We can derive a similar expression for  $\sigma_B$ , the standard deviation for  $B$ .

<sup>9</sup>To be strictly correct, we should multiply the second term by the identity operator, but this is usually understood.

<sup>10</sup>It is tempting to denote the standard deviation as  $\Delta A$  but this could be confusing in view of (5.7.2).

The generalized uncertainty principle is obtained by using the Schwartz inequality:

$$\left| \int f^* g dx \right|^2 \leq \left( \int |f|^2 dx \right) \left( \int |g|^2 dx \right) \quad (5.7.5)$$

We can rewrite the above using Dirac notation to get

$$|\langle f|g \rangle|^2 \leq \langle f|f \rangle \langle g|g \rangle \quad (5.7.6)$$

The above is the generalization of the cosine inequality we have for 3-vectors<sup>11</sup>

$$|\mathbf{A} \cdot \mathbf{B}|^2 = |\mathbf{A}|^2 |\mathbf{B}|^2 \cos^2 \theta \leq |\mathbf{A}|^2 |\mathbf{B}|^2 \quad (5.7.7)$$

It can be generalized to  $N$ -vectors or vectors in  $N$  dimensional space, and then to vectors in infinite dimensional space if the integrals converge. If we define, for a quantum system in state  $\psi$ ,

$$|f\rangle = (\hat{A} - \bar{A}) |\psi\rangle = \hat{a} |\psi\rangle \quad (5.7.8)$$

$$|g\rangle = (\hat{B} - \bar{B}) |\psi\rangle = \hat{b} |\psi\rangle \quad (5.7.9)$$

where  $\hat{a} = \hat{A} - \bar{A}$  and  $\hat{b} = \hat{B} - \bar{B}$ . Then

$$\langle f|f \rangle = \langle \psi | \hat{a}^2 | \psi \rangle = \left\langle \psi \left| (\hat{A} - \bar{A})^2 \right| \psi \right\rangle = \overline{(\Delta A)^2} \quad (5.7.10)$$

$$\langle g|g \rangle = \langle \psi | \hat{b}^2 | \psi \rangle = \left\langle \psi \left| (\hat{B} - \bar{B})^2 \right| \psi \right\rangle = \overline{(\Delta B)^2} \quad (5.7.11)$$

Using the inequality in (5.7.6), we have

$$\left( \overline{(\Delta A)^2} \right) \left( \overline{(\Delta B)^2} \right) \geq |\langle \psi | \hat{a} \hat{b} | \psi \rangle|^2 \quad (5.7.12)$$

It can be shown easily that

$$[\hat{a}, \hat{b}] = i\hat{C} \quad (5.7.13)$$

Furthermore,

$$\hat{a} \hat{b} = \frac{\hat{a} \hat{b} + \hat{b} \hat{a}}{2} + \frac{\hat{a} \hat{b} - \hat{b} \hat{a}}{2} = \frac{\hat{a} \hat{b} + \hat{b} \hat{a}}{2} + i \frac{\hat{C}}{2} \quad (5.7.14)$$

Taking the expectation value of the above, we have

$$\langle \hat{a} \hat{b} \rangle = \left\langle \frac{\hat{a} \hat{b} + \hat{b} \hat{a}}{2} \right\rangle + \frac{i}{2} \langle \hat{C} \rangle \quad (5.7.15)$$

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<sup>11</sup>The integral for the inner products above can be generalized to 3D space.

Since  $\hat{a}$  and  $\hat{b}$  are Hermitian, the first term on the right-hand side is also Hermitian. Hence, its expectation value is purely real, while the second term is purely imaginary. Therefore, the amplitude square of the above is

$$|\langle \hat{a}\hat{b} \rangle|^2 = \left| \left\langle \frac{\hat{a}\hat{b} + \hat{b}\hat{a}}{2} \right\rangle \right|^2 + \frac{1}{4} |\langle \hat{C} \rangle|^2 \geq \frac{1}{4} |\langle \hat{C} \rangle|^2 \quad (5.7.16)$$

Using the above in (5.7.12), we have

$$\left( \overline{(\Delta A)^2} \right) \left( \overline{(\Delta B)^2} \right) \geq \frac{1}{4} |\langle \hat{C} \rangle|^2 \quad (5.7.17)$$

The above is the generalized uncertainty principle for two observables  $A$  and  $B$ . We can take the square root of the above to get

$$\sigma_A \sigma_B \geq \frac{1}{2} |\langle \hat{C} \rangle| \quad (5.7.18)$$

## 5.8 Time Evolution of the Expectation Value of an Operator

The expectation value of an operator is given by

$$\langle \hat{A} \rangle = \langle \psi(t) | \hat{A} | \psi(t) \rangle \quad (5.8.1)$$

Taking the time derivative of the above yields

$$\partial_t \langle \hat{A} \rangle = \langle \partial_t \psi(t) | \hat{A} | \psi(t) \rangle + \langle \psi(t) | \hat{A} | \partial_t \psi(t) \rangle \quad (5.8.2)$$

Assuming that  $\hat{A}$  is time independent, and using the fact that

$$i\hbar \partial_t | \psi(t) \rangle = i\hbar | \partial_t \psi(t) \rangle = \hat{H} | \psi(t) \rangle \quad (5.8.3)$$

we have

$$\begin{aligned} \partial_t \langle \hat{A} \rangle &= \frac{i}{\hbar} \left\langle \left[ \hat{H} \psi(t) \right] | \hat{A} | \psi(t) \right\rangle - \frac{i}{\hbar} \langle \psi(t) | \hat{A} | \hat{H} \psi(t) \rangle \\ &= \frac{i}{\hbar} \langle \psi(t) | \hat{H} \hat{A} | \psi(t) \rangle - \frac{i}{\hbar} \langle \psi(t) | \hat{A} \hat{H} | \psi(t) \rangle \\ &= \frac{i}{\hbar} \langle \psi(t) | \hat{H} \hat{A} - \hat{A} \hat{H} | \psi(t) \rangle \end{aligned} \quad (5.8.4)$$

or

$$i\hbar \partial_t \langle \hat{A} \rangle = \langle \hat{A} \hat{H} - \hat{H} \hat{A} \rangle = \left\langle \left[ \hat{A}, \hat{H} \right] \right\rangle \quad (5.8.5)$$

In other words, if  $\hat{A}$  commutes with  $\hat{H}$ , its expectation value is time independent, or it is a constant of motion. The operator  $\hat{A}$  in this case represents an observable that is conserved, such as linear momentum or angular momentum. If such an operator commutes with

the Hamiltonian of the quantum system, the observable that corresponds to the operator is conserved.

Furthermore, it can be shown that if  $\hat{A}$  is time dependent, we need only to augment the above equation as

$$i\hbar\partial_t\langle\hat{A}\rangle = \left\langle\left[\hat{A}, \hat{H}\right]\right\rangle + i\hbar\langle\partial_t\hat{A}\rangle \quad (5.8.6)$$

## 5.9 Periodic Boundary Condition

One technique often used to create countably infinite modes is to use periodic boundary condition. In this method, we require that the wavefunction satisfies the boundary condition that  $\psi(x=0) = \psi(x=L)$ . This can be satisfied easily by a traveling wave

$$\psi(x) = Ce^{ikx} \quad (5.9.1)$$

The periodic boundary condition implies that

$$e^{ikL} = 1 \quad (5.9.2)$$

or that

$$k = k_n = \frac{2n\pi}{L} \quad (5.9.3)$$

where we have indexed  $k$  with subscript  $n$ , and  $n$  is all integer values on the real line from  $-\infty$  to  $\infty$ . Furthermore, we can pick  $C = 1/\sqrt{L}$  to normalize the above function. It is quite easy to show that with

$$\psi_n(x) = \sqrt{\frac{1}{L}}e^{i2n\pi x/L} \quad (5.9.4)$$

then

$$\langle\psi_n|\psi_m\rangle = \delta_{nm} \quad (5.9.5)$$

Hence, given any function  $f(x)$  defined between  $0 < x < L$ , we can expand

$$f(x) = \sum_{n=-\infty}^{\infty} f_n\psi_n(x) = \sum_{n=-\infty}^{\infty} f_n\sqrt{\frac{1}{L}}e^{i2n\pi x/L} \quad (5.9.6)$$

The above is just a Fourier series expansion of the function  $f(x)$ . The Fourier coefficient  $f_n$  can be found by using the orthonormality of  $\psi_n$ , to get

$$f_n = \langle\psi_n|f\rangle = \sqrt{\frac{1}{L}}\int_0^L dx e^{-i2n\pi x/L} f(x) \quad (5.9.7)$$

In the above, since  $\psi_n(x)$  is a periodic function with period  $L$ ,  $f(x)$  is also a periodic function with period  $L$ .



Fourier transform can be derived from Fourier series expansion. To this end, we look at (5.9.3), and note that the  $k_n$  values are equally spaced by  $2\pi/L$  on the real  $k$  line. The spacing becomes increasingly small as  $L \rightarrow \infty$ . We can define the spacing to be  $\Delta k = 2\pi/L$ . Consequently, in view of changing a summation to an integral, we can rewrite (5.9.6) more suggestively as

$$f(x) = \sum_{n=-\infty}^{\infty} \Delta k f_n \frac{L}{2\pi} \sqrt{\frac{1}{L}} e^{i2n\pi x/L} = \frac{1}{2\pi} \sum_{n=-\infty}^{\infty} \Delta k \tilde{f}(k_n) e^{ik_n x} \quad (5.9.8)$$

where we have defined  $\tilde{f}(k_n) = f_n \sqrt{L}$ . From (5.9.7), we have

$$\tilde{f}(k_n) = \sqrt{L} f_n = \int_0^L dx e^{-ik_n x} f(x) \quad (5.9.9)$$

In the limit when  $L$  tends to infinity,  $\Delta k \rightarrow 0$ , the sum in (5.9.8) becomes an integral; namely

$$f(x) = \frac{1}{2\pi} \int_{-\infty}^{\infty} dk \tilde{f}(k) e^{ikx} \quad (5.9.10)$$

The above is just a Fourier inverse transform integral with  $\tilde{f}(k)$  from (5.9.9)

$$\tilde{f}(k) = \int_0^{\infty} dx e^{-ikx} f(x) \quad (5.9.11)$$

which is a Fourier transform. Here,  $f(x)$  is originally a periodic function, but now the period  $L$  is infinite. Also, if  $f(x) \rightarrow 0$  when  $x \rightarrow \infty$ , then by its infinite periodicity,  $f(x) \rightarrow 0$  when  $x \rightarrow -\infty$ . We can replace the semi-infinite integral above with an infinite integral:

$$\tilde{f}(k) = \int_{-\infty}^{\infty} dx e^{-ikx} f(x) \quad (5.9.12)$$

The above, (5.9.10) and (5.9.12), form a Fourier transform pair.



## Chapter 6

# Approximate Methods in Quantum Mechanics

### 6.1 Introduction

There are many problems in quantum mechanics where closed form or simple solutions cannot be found. What we have shown as examples are usually problems that have simple solutions, such as a infinite potential well, harmonic oscillator, or a finite potential well. Even the case of finite potential well requires the solution of a transcendental equation. Nevertheless, these are called textbook problems, because they are wonderful examples in textbooks to teach beginning students on the subject matter. But in general, most quantum mechanics problems do not have closed form solution. We have to resort to approximate or numerical methods to solve such problems.

### 6.2 Use of an Approximate Subspace

We have seen that in general, a Schrödinger equation problem can be cast into a matrix equation by projecting it into a space spanned by countably infinite orthonormal basis functions. With such a basis, we define an identity operator

$$\hat{I} = \sum_{n=1}^{\infty} |\psi_n\rangle\langle\psi_n| \quad (6.2.1)$$

Given a general time independent Schrödinger equation, which is

$$\hat{H}|\psi\rangle = E|\psi\rangle \quad (6.2.2)$$

we can first insert an identity operator between  $\hat{H}$  and  $|\psi\rangle$ . Then the above becomes

$$\sum_{n=1}^{\infty} \hat{H}|\psi_n\rangle\langle\psi_n|\psi\rangle = E|\psi\rangle \quad (6.2.3)$$

Testing the above by multiplying it from the left with  $\langle \psi_m |$ ,  $m = 1, \dots, \infty$ , we have<sup>1</sup>

$$\sum_{n=1}^{\infty} \langle \psi_m | \hat{H} | \psi_n \rangle \langle \psi_n | \psi \rangle = E \langle \psi_m | \psi \rangle, \quad m = 1, \dots, \infty \quad (6.2.4)$$

The above is a matrix equation of the form

$$\sum_{n=1}^{\infty} H_{mn} f_n = E f_m, \quad m = 1, \dots, \infty \quad (6.2.5)$$

It corresponds to an infinite dimension matrix, but it can be truncated to form a finite matrix system

$$\bar{\mathbf{H}} \cdot \mathbf{f} = E \mathbf{f} \quad (6.2.6)$$

The above is a matrix eigenvalue problem that can be solved numerically. If  $\bar{\mathbf{H}}$  is  $N \times N$ , it will have  $N$  eigenvalues and  $N$  eigenvectors. If we choose our basis functions appropriately, and are only interested in certain eigenvalues, which represent the energy levels of the stationary states, we need only to pick a small set of basis functions. Therefore,  $N$  can be a small number. This is important because the number of computer operations needed to find the eigenvectors and eigenvalues of the above problem is proportional to  $N^3$ .

In solving the above, we will have to form the matrix elements  $H_{mn}$ . For an infinite potential well, whose well bottom has been distorted, we can use the eigenfunctions of the flat-bottom well as our basis functions. Then, the matrix elements  $H_{mn}$  are explicitly given as

$$H_{mn} = \int_0^L dx \psi_m^*(x) \left( -\frac{\hbar}{2m_o} \frac{d^2}{dx^2} + V(x) \right) \psi_n(x) \quad (6.2.7)$$

We can use integration by parts to convert the term that involves second derivative to a form that involves only first derivatives. That is

$$H_{mn} = \int_0^L dx \left( \frac{\hbar}{2m_o} \frac{d}{dx} \psi_m^*(x) \frac{d}{dx} \psi_n(x) + \psi_m^*(x) V(x) \psi_n(x) \right) \quad (6.2.8)$$

which is a more symmetric form. In the above, we have used the boundary condition of the eigenfunctions to remove the extraneous term resulting from integration by parts.

Modern numerical methods have allowed us to use flexible basis set to expand our unknown functions. For example, we can let

$$|\psi\rangle = \sum_{n=1}^N a_n |f_n\rangle \quad (6.2.9)$$

where  $f_n(x)$  need not be orthogonal or complete in Hilbert space, but they can approximate the unknown function  $\psi(x)$  well. Also,  $f_n(x), n = 1, \dots, N$  spans an  $N$  dimensional vector

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<sup>1</sup>This procedure is called testing or weighting in the mathematics literature,  $\langle \psi_m |$  is called the testing or weighting function.

space that is the subspace of the Hilbert space. We assume that the solution in the subspace is a good approximation of the solution in the original infinite dimensional Hilbert space. Examples of such finite basis set are triangle functions shown in Figure 6.1.

Using (6.2.9) in (6.2.2), we have

$$\sum_{n=1}^N \hat{H}|f_n\rangle a_n = E \sum_{n=1}^N a_n |f_n\rangle \quad (6.2.10)$$

Multiplying the above from  $\langle f_m|$  from the left, we have

$$\sum_{n=1}^N \langle f_m|\hat{H}|f_n\rangle a_n = E \sum_{n=1}^N \langle f_m|f_n\rangle a_n \quad (6.2.11)$$

The above is a matrix system of the form

$$\tilde{\mathbf{H}} \cdot \mathbf{a} = E \bar{\mathbf{B}} \cdot \mathbf{a} \quad (6.2.12)$$

In the above,

$$\begin{aligned} [\tilde{\mathbf{H}}]_{mn} &= \langle f_m|\hat{H}|f_n\rangle \\ [\bar{\mathbf{B}}]_{mn} &= \langle f_m|f_n\rangle \\ [\mathbf{a}]_n &= a_n \end{aligned} \quad (6.2.13)$$

The equation (6.2.12) above is a generalized eigenvalue problem with eigenvalue  $E$  and eigenvector  $\mathbf{a}$ . The difference of it from (6.2.6) is that a matrix  $\bar{\mathbf{B}}$  is on the right-hand side. But numerical software to seek the eigenvalues and eigenvectors for such generalized eigenvalue problems are widely available.

The approximate subspace method is varyingly known as the subspace projection method, Galerkin's method, Petrov-Galerkin method, weighted-residual method, and the method of moments. The finite element method falls under the same category.

### 6.3 Time Independent Perturbation Theory

Let us assume that we can solve a problem when it is unperturbed, such as the infinite potential well problem. We like to solve the problem when, say there is a small electric field in the well indicated by a potential with a gradient, or that the well is imperfect due to some impurities indicated by a small bump (see Figure 6.2). The Schrödinger equation for the perturbed system can be written as

$$(\hat{H}_0 + \gamma \hat{H}_p)|\phi\rangle = E|\phi\rangle \quad (6.3.1)$$

where  $\hat{H}_0$  is the Hamiltonian of the unperturbed system whose solution is known, and  $\gamma \hat{H}_p$  is due to the small perturbation where  $\gamma$  is a small parameter. Here,  $\hat{H}_0$  can be the Hamiltonian of the infinite potential well, for instance. In the above equation,  $|\phi\rangle$  and  $E$  are both

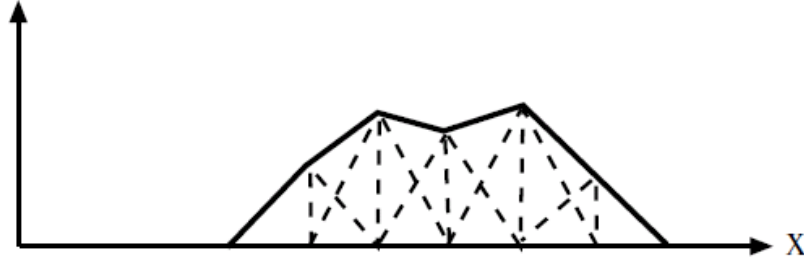


Figure 6.1: The triangle functions for a piecewise linear approximation of a function. This is a basis that is not orthogonal but yet can be used to seek approximate solutions to (6.2.12).

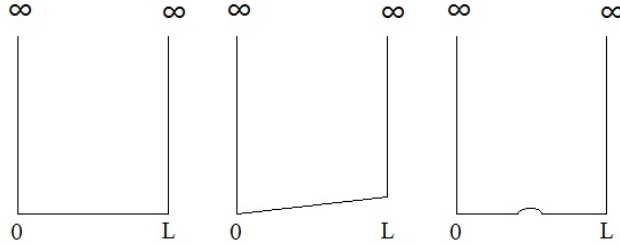


Figure 6.2: The infinite potential well on the left represents the unperturbed problem. The middle figure represents a perturbation due to a tiny electric field. The right figure represents a perturbation due to imperfection in fabrication or impurities.

unknowns, but we can write them in a perturbation series or expansion, namely

$$|\phi\rangle = |\phi^{(0)}\rangle + \gamma|\phi^{(1)}\rangle + \gamma^2|\phi^{(2)}\rangle + \dots \quad (6.3.2)$$

$$E = E^{(0)} + \gamma E^{(1)} + \gamma^2 E^{(2)} + \dots \quad (6.3.3)$$

Upon substituting the above series into (6.3.1), we obtain

$$\begin{aligned} & (\hat{H}_0 + \gamma\hat{H}_p) \left( |\phi^{(0)}\rangle + \gamma|\phi^{(1)}\rangle + \gamma^2|\phi^{(2)}\rangle + \dots \right) \\ &= \left( E^{(0)} + \gamma E^{(1)} + \gamma^2 E^{(2)} + \dots \right) \left( |\phi^{(0)}\rangle + \gamma|\phi^{(1)}\rangle + \gamma^2|\phi^{(2)}\rangle + \dots \right) \end{aligned} \quad (6.3.4)$$

The left-hand side of (6.3.4) can be expanded and rewritten on a power series in  $\gamma$

$$a_0 + a_1\gamma + a_2\gamma^2 + \dots \quad (6.3.5)$$

while the right-hand side is similarly written as

$$b_0 + b_1\gamma + b_2\gamma^2 + \dots \quad (6.3.6)$$

These two power series in  $\gamma$  are equal only if  $a_i = b_i$ ,  $i = 0, 1, \dots, \infty$ .<sup>2</sup>

Equating the coefficients of the power series on both sides of (6.3.4) we have the following equations:

Zeroth Order:

$$\hat{H}_0|\phi^{(0)}\rangle = E^{(0)}|\phi^{(0)}\rangle \quad (6.3.7)$$

First Order:

$$\hat{H}_0|\phi^{(1)}\rangle + \hat{H}_p|\phi^{(0)}\rangle = E^{(0)}|\phi^{(1)}\rangle + E^{(1)}|\phi^{(0)}\rangle \quad (6.3.8)$$

Second Order:

$$\hat{H}_0|\phi^{(2)}\rangle + \hat{H}_p|\phi^{(1)}\rangle = E^{(0)}|\phi^{(2)}\rangle + E^{(1)}|\phi^{(1)}\rangle + E^{(2)}|\phi^{(0)}\rangle \quad (6.3.9)$$

We assume that the zeroth order equation is known in terms of an eigenstate  $|\psi_m\rangle$  with energy  $E_m$ . In other words

$$|\phi^{(0)}\rangle = |\psi_m\rangle, \quad E^{(0)} = E_m \quad (6.3.10)$$

We will use this knowledge to solve the first order equation (6.3.8) above.

Before we proceed further, a note is in order regarding the uniqueness of the eigenvalue problem (6.3.1). An eigenvector is known only within a multiplicative factor. Hence, its length is indeterminate. This non-uniqueness in its length manifests in the non-uniqueness of the value of the perturbation series (6.3.2) as we shall see later. To achieve uniqueness, it is best to pin down the length of the total eigenvector given by (6.3.2). We fix the length of the eigenvector  $|\phi\rangle$  by requiring that

$$\langle\psi_m|\phi\rangle = 1 \quad (6.3.11)$$

With this requirement, we substitute (6.3.2) into the above. Since  $\langle\psi_m|\phi^{(0)}\rangle = 1$ , because  $|\phi^{(0)}\rangle = |\psi_m\rangle$ , it is easy to show that  $\langle\psi_m|\phi^{(i)}\rangle = 0$ ,  $i > 0$ . As a consequence,  $|\phi^{(i)}\rangle$  is orthogonal to  $|\psi_m\rangle$ . The perturbation series is not necessarily normalized, but it can be normalized later after the series has been calculated.

Next, to find the first order corrections to the eigenvalue and the eigenvector, we move the unknowns  $|\phi^{(1)}\rangle$  to the left of (6.3.8). We then have

$$\left(\hat{H}_0 - E_m\right)|\phi^{(1)}\rangle = E^{(1)}|\psi_m\rangle - \hat{H}_p|\psi_m\rangle \quad (6.3.12)$$

where we have made use of (6.3.10). Notice that the above equation is non-unique since the operator  $\left(\hat{H}_0 - E_m\right)$  has a null space with a null space vector  $|\psi_m\rangle$ .

Testing the above equation with  $\langle\psi_m|$ , we have

$$\langle\psi_m|\hat{H}_0 - E_m|\phi^{(1)}\rangle = E^{(1)} - \langle\psi_m|\hat{H}_p|\psi_m\rangle \quad (6.3.13)$$

---

<sup>2</sup>This can be easily proved by setting  $\gamma = 0$ , which immediately implies that  $a_0 = b_0$ . By differentiating the power series again and let  $\gamma = 0$ , we prove that  $a_i = b_i$  etc.

But  $\langle \psi_m | (\hat{H}_0 - E_m) = 0$  because  $\hat{H}_0 |\psi_m\rangle = E_m |\psi_m\rangle$ . Hence, the above gives

$$E^{(1)} = \langle \psi_m | \hat{H}_p | \psi_m \rangle \quad (6.3.14)$$

the first order correction to the energy of the perturbed system.

Therefore, one of the two unknowns in (6.3.12) is found. The remaining unknown  $|\phi^{(1)}\rangle$  can be expanded in terms of the eigenfunctions of the unperturbed system, that is

$$|\phi^{(1)}\rangle = \sum_n a_n^{(1)} |\psi_n\rangle \quad (6.3.15)$$

First, testing the equation (6.3.12) with  $\langle \psi_i |$ , we have

$$\langle \psi_i | \hat{H}_0 - E_m | \phi^{(1)} \rangle = E^{(1)} \langle \psi_i | \psi_m \rangle - \langle \psi_i | \hat{H}_p | \psi_m \rangle \quad (6.3.16)$$

Upon substituting (6.3.15) into the above, the left-hand side evaluates to

$$\langle \psi_i | E_i - E_m | \phi^{(1)} \rangle = (E_i - E_m) \sum_n \langle \psi_i | a_n^{(1)} | \psi_n \rangle = (E_i - E_m) a_i^{(1)} \quad (6.3.17)$$

The right-hand side of (6.3.16), for  $i \neq m$ , is just  $-\langle \psi_i | \hat{H}_p | \psi_m \rangle$ . Hence

$$a_i^{(1)} = \frac{\langle \psi_i | \hat{H}_p | \psi_m \rangle}{E_m - E_i}, i \neq m \quad (6.3.18)$$

When  $i = m$ , (6.3.17) evaluates to zero, while the right hand side of (6.3.16) also evaluates to zero because of (6.3.14). Hence  $a_m^{(1)}$  is undefined. We choose  $a_m^{(1)} = 0$  for a number of reasons: It makes the correction term unique since  $|\psi^{(1)}\rangle$  is orthogonal to  $|\psi^{(0)}\rangle$ . It makes the normalization of the eigenvector  $|\phi\rangle$  accurate to second order even though the correction is first order. It will also make the second order corrections much simpler to find.

To find the second order corrections, we rewrite (6.3.9) with the unknown  $|\phi^{(2)}\rangle$  on the left hand side. Then (6.3.9) becomes

$$\left( \hat{H}_0 - E_m \right) |\phi^{(2)}\rangle = E^{(1)} |\phi^{(1)}\rangle + E^{(2)} |\psi_m\rangle - \hat{H}_p |\phi^{(1)}\rangle \quad (6.3.19)$$

Testing the above with  $\langle \psi_m |$ , the left hand side becomes zero as before.<sup>3</sup> Since we have made  $|\phi^{(1)}\rangle$  orthogonal to  $|\psi_m\rangle$ , on the right-hand side, only the last two terms remain. Consequently,

$$0 = E^{(2)} - \langle \psi_m | \hat{H}_p | \phi^{(1)} \rangle \quad (6.3.20)$$

or

$$E^{(2)} = \langle \psi_m | \hat{H}_p | \phi^{(1)} \rangle \quad (6.3.21)$$

---

<sup>3</sup>Again, the above is non-unique for the same reason cited for (6.3.12).



Letting

$$|\phi^{(2)}\rangle = \sum_n a_n^{(2)} |\psi_n\rangle \quad (6.3.22)$$

and substituting it into the left side of (6.3.19), testing with  $\langle\psi_i|$ , we have

$$(E_i - E_m)a_i^{(2)} = E^{(1)}\langle\psi_i|\phi^{(1)}\rangle + E^{(2)}\delta_{im} - \langle\psi_i|\hat{H}_p|\phi^{(1)}\rangle \quad (6.3.23)$$

Therefore, for  $i \neq m$ ,

$$a_i^{(2)} = \frac{\langle\psi_i|\hat{H}_p|\phi^{(1)}\rangle}{E_m - E_i} - \frac{E^{(1)}a_i^{(1)}}{E_m - E_i} \quad (6.3.24)$$

When  $i = m$ , both sides of (6.3.23) vanish, and  $a_m^{(2)}$  is undefined. Again, we pick  $a_m^{(2)} = 0$  to obtain a unique solution.

The above procedure can be generalized to arbitrary order. By induction, we notice that the equivalence of (6.3.9) to  $p$ -th order is

$$\hat{H}_0|\phi^{(p)}\rangle + \hat{H}_p|\phi^{(p-1)}\rangle = E^{(0)}|\phi^{(p)}\rangle + E^{(1)}|\phi^{(p-1)}\rangle + \dots + E^{(p)}|\phi^{(0)}\rangle \quad (6.3.25)$$

The above can be rewritten as

$$(\hat{H}_0 - E^{(0)})|\phi^{(p)}\rangle = E^{(1)}|\phi^{(p-1)}\rangle + \dots + E^{(p)}|\phi^{(0)}\rangle - \hat{H}_p|\phi^{(p-1)}\rangle \quad (6.3.26)$$

Testing the above with  $\langle\psi_m|$  gives

$$E^{(p)} = \langle\psi_m|\hat{H}_p|\phi^{(p-1)}\rangle \quad (6.3.27)$$

Letting

$$|\phi^{(p)}\rangle = \sum_n a_n^{(p)} |\psi_n\rangle \quad (6.3.28)$$

in (6.3.25), and testing with  $\langle\psi_i|$ , one can easily show, for  $i \neq m$ , and after noting that  $E^{(0)} = E_m$ ,  $|\phi^{(0)}\rangle = |\psi_m\rangle$ , that

$$a_i^{(p)} = \frac{1}{E_m - E_i} \left( \langle\psi_i|\hat{H}_p|\phi^{(p-1)}\rangle - E^{(1)}a_i^{(p-1)} - E^{(2)}a_i^{(p-2)} + \dots - E^{(p-1)}a_i^{(1)} \right) \quad (6.3.29)$$

It is to be noted that with modern advent of computer technology, and given the availability of numerical methods, the calculation of perturbation theory to very high order is laborious and not necessary. However, a perturbation correction can give us insight on how a small change in the Hamiltonian can change the solution.

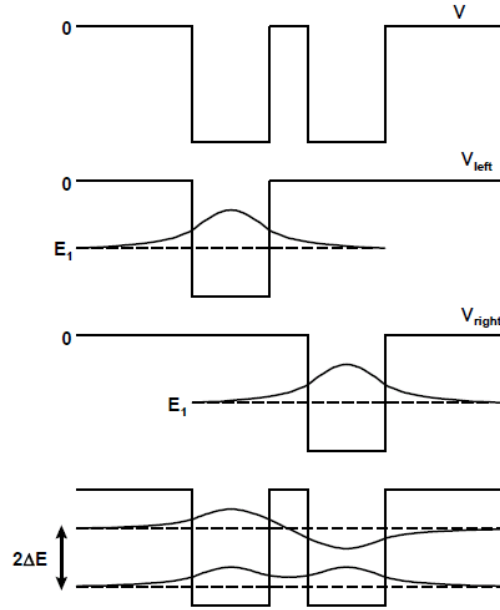


Figure 6.3: The tight binding model can be used to find the approximate eigenstates of two quantum wells that are weakly coupled to each other. The eigenstates of the isolated quantum wells are known (Figure is from DAB Miller).

## 6.4 Tight Binding Model

The tight binding model is a way to find the solution to two weakly coupled quantum systems when the solution to each isolated quantum system is known (see Figure 6.3).

To solve this system, we can use the subspace projection method of nonorthogonal basis to derive the relevant matrix equation. Given the Schrödinger equation

$$\hat{H} |\psi\rangle = E |\psi\rangle \quad (6.4.1)$$

we let

$$|\psi\rangle = \sum_{i=1}^2 a_i |\psi_i\rangle \quad (6.4.2)$$

where  $|\psi_1\rangle$  is the eigenfunction of the ground state of well 1 in isolation and  $|\psi_2\rangle$  is the eigenfunction of the ground state of well 2. Going through the same procedure we have had before, we arrive at

$$\sum_{i=1}^2 a_i \langle \psi_j | \hat{H} | \psi_i \rangle = E \sum_{i=1}^2 a_i \langle \psi_j | \psi_i \rangle \quad (6.4.3)$$

The above is the same as the matrix equation

$$\overline{\mathbf{H}} \cdot \mathbf{a} = E \overline{\mathbf{B}} \cdot \mathbf{a} \quad (6.4.4)$$

where

$$\left[ \overline{\mathbf{H}} \right]_{ij} = \langle \psi_i | \hat{H} | \psi_j \rangle \quad (6.4.5)$$

$$\left[ \overline{\mathbf{B}} \right]_{ij} = \langle \psi_i | \psi_j \rangle \quad (6.4.6)$$

It is easy to show that  $\overline{\mathbf{H}}$  and  $\overline{\mathbf{B}}$  are Hermitian matrices.<sup>4</sup> In the above, we can assume that  $|\psi_i\rangle$  is quasi-orthonormal, so that

$$\langle \psi_i | \psi_j \rangle \simeq \delta_{ij} \quad (6.4.7)$$

Then we can write the system (6.4.4) as

$$\begin{bmatrix} E_1 & \Delta E \\ \Delta E^* & E_1 \end{bmatrix} \begin{bmatrix} a_1 \\ a_2 \end{bmatrix} = E \begin{bmatrix} 1 & \Delta B \\ \Delta B^* & 1 \end{bmatrix} \begin{bmatrix} a_1 \\ a_2 \end{bmatrix} \quad (6.4.8)$$

The off diagonal elements are conjugate of each other because the matrix is Hermitian. Moreover,  $\Delta E \ll E_1$  and  $\Delta B \ll 1$ . These terms are small because the overlap between the eigenfunctions  $\psi_1$  and  $\psi_2$  is small. The further apart the wells are, the smaller these terms would be, as the overlap is smaller. When the wells are infinite far apart, the off diagonal terms vanish and the two quantum systems are uncoupled from each other. The two eigenvalues are then degenerate.

The eigenvalue of (6.4.8) above is obtained by solving

$$\det \begin{bmatrix} E_1 - E & \Delta E - E\Delta B \\ \Delta E^* - E\Delta B^* & E_1 - E \end{bmatrix} = 0 \quad (6.4.9)$$

or

$$(E_1 - E)^2 - |\Delta E - E\Delta B|^2 = 0 \quad (6.4.10)$$

Since the  $E\Delta B$  on the left-hand side of the above is small, we do not incur much error by replacing it with  $E_1\Delta B$ . Upon solving the above, we have

$$E \cong E_1 \pm |\Delta E_1 - E\Delta B|^2 \quad (6.4.11)$$

Note that the degenerate eigenvalues are now split into two non-degenerate eigenvalues. The corresponding eigenvectors can be found, and the eigenfunctions are of the form

$$\psi_- = \frac{1}{\sqrt{2}} (\psi_1 + \psi_2), \quad \psi_+ = \frac{1}{\sqrt{2}} (\psi_1 - \psi_2) \quad (6.4.12)$$

The eigenfunctions are due to odd and even coupling between the isolated quantum states of the two wells as they are brought closer together.

<sup>4</sup>It can be easily shown that the matrix representation of a Hermitian operator remains Hermitian.

### 6.4.1 Variational Method

Given an eigenvalue problem denoted by

$$\hat{H} |\phi\rangle = E |\phi\rangle \quad (6.4.13)$$

we can convert the above into a matrix generalized eigenvalue problem

$$\bar{\mathbf{H}} \cdot \mathbf{a} = E \bar{\mathbf{B}} \cdot \mathbf{a} \quad (6.4.14)$$

We next convert the above into a functional by multiplying it with  $\bar{\mathbf{a}}^\dagger$  to obtain

$$\bar{\mathbf{a}}^\dagger \cdot \bar{\mathbf{H}} \cdot \mathbf{a} = E \bar{\mathbf{a}}^\dagger \cdot \bar{\mathbf{B}} \cdot \mathbf{a} \quad (6.4.15)$$

By setting  $\mathbf{a} = \mathbf{a}_0 + \delta\mathbf{a}$ ,  $E = E_0 + \delta E$ , where  $\mathbf{a}_0$  and  $E_0$  are the exact solution to (6.4.14), and then taking the first variation of (6.4.15) we have

$$\begin{aligned} \delta\bar{\mathbf{a}}^\dagger \cdot \bar{\mathbf{H}} \cdot \mathbf{a}_0 + \bar{\mathbf{a}}_0^\dagger \cdot \bar{\mathbf{H}} \cdot \delta\mathbf{a} &= \delta E \bar{\mathbf{a}}_0^\dagger \cdot \bar{\mathbf{B}} \cdot \mathbf{a}_0 + E_0 \cdot \delta\bar{\mathbf{a}}^\dagger \cdot \bar{\mathbf{B}} \cdot \mathbf{a}_0 \\ &+ E_0 \bar{\mathbf{a}}_0^\dagger \cdot \bar{\mathbf{B}} \cdot \delta\mathbf{a} \end{aligned} \quad (6.4.16)$$

If  $\mathbf{a}_0$  and  $E_0$  are exact solution to (6.4.16), they satisfy

$$\bar{\mathbf{H}} \cdot \mathbf{a}_0 = E_0 \bar{\mathbf{B}} \cdot \mathbf{a}_0 \quad (6.4.17)$$

$$\bar{\mathbf{a}}_0^\dagger \cdot \bar{\mathbf{H}} = E_0 \bar{\mathbf{a}}_0^\dagger \cdot \bar{\mathbf{B}} \quad (6.4.18)$$

Then the term multiplied by  $\delta\bar{\mathbf{a}}^\dagger$  and  $\delta\mathbf{a}$  cancel each other on both sides and we have, to first order,

$$0 = \delta E \bar{\mathbf{a}}_0^\dagger \cdot \bar{\mathbf{B}} \cdot \mathbf{a}_0 \quad (6.4.19)$$

The above implies that  $\delta E$  is zero. It means that if we were to use (6.4.15) to find  $E$  by writing it as

$$E = \frac{\bar{\mathbf{a}}^\dagger \cdot \bar{\mathbf{H}} \cdot \mathbf{a}}{\bar{\mathbf{a}}^\dagger \cdot \bar{\mathbf{B}} \cdot \mathbf{a}} \quad (6.4.20)$$

and if we substitute inexact or approximate value for  $\mathbf{a} = \mathbf{a}_0 + \delta\mathbf{a}$ , the value we obtain for  $E$  is second order accurate since  $\delta E = 0$ . The above is known as the Rayleigh quotient. To begin with, we need an estimate of the solution that is quite close to the exact solution so that  $\delta\mathbf{a}$  is small. In other words, it allows one to substitute in approximate value of  $\mathbf{a}$  that is just first-order accurate, and yet get an estimate of the eigenvalue or energy  $E$  that is second-order accurate.

A more general form of the Rayleigh quotient is to test (6.4.13) with  $\langle\phi|$  and convert it into a functional. The corresponding Rayleigh quotient is

$$E = \frac{\langle\phi|\hat{H}|\phi\rangle}{\langle\phi|\phi\rangle} \quad (6.4.21)$$

It assumes that  $|\phi\rangle$  is not normalized. In the above, it can be proved by the same token that an approximate, first-order accurate value of  $|\phi\rangle$  can yield a second-order accurate estimate of  $E$ . Rayleigh quotient is great for estimating the eigenvalues when only approximate eigenfunctions are known. These eigenfunctions are also called trial functions. A systematic way to estimate these eigenfunctions is known as the Rayleigh-Ritz method.

### 6.4.2 Time Dependent Perturbation Theory

Now, we assume that the perturbing Hamiltonian is a function of time. In other words, we have a Schrödinger equation where

$$i\hbar \frac{\partial}{\partial t} |\Psi\rangle = \hat{H} |\Psi\rangle \quad (6.4.22)$$

and

$$\hat{H} = \hat{H}_0 + \gamma \hat{H}_p(t) \quad (6.4.23)$$

where  $\hat{H}_p(t)$  is time varying.<sup>5</sup> It can come from an external electric field that is time-harmonic, for instance. We also assume that we can solve the unperturbed system in terms of stationary states, namely

$$\hat{H}_0 |\Psi_n\rangle = E_n |\Psi_n\rangle \quad (6.4.24)$$

The total time dependent wavefunction can be expanded in terms of the stationary states, namely

$$|\Psi\rangle = \sum_n a_n(t) e^{-i\omega_n t} |\Psi_n\rangle \quad (6.4.25)$$

where  $\hbar\omega_n = E_n$ , and  $a_n(t)$  is a function of time.

Substituting the above into (6.4.22), we have

$$\frac{\partial}{\partial t} |\Psi\rangle = \sum_n (\dot{a}_n(t) - i\omega_n a_n(t)) e^{-i\omega_n t} |\Psi_n\rangle = \frac{1}{i\hbar} \hat{H} \sum_n a_n(t) e^{-i\omega_n t} |\Psi_n\rangle \quad (6.4.26)$$

$$= \frac{1}{i\hbar} \sum_n a_n(t) e^{-i\omega_n t} \left( E_n |\Psi_n\rangle + \gamma \hat{H}_p(t) |\Psi_n\rangle \right) \quad (6.4.27)$$

where  $\dot{a}_n(t) = \partial a_n(t) / \partial t$ .

Since  $E_n = \hbar\omega_n$ , the  $E_n$  term on the right-hand side cancels the  $\omega_n$  term on the left-hand side. Finally we have

$$\sum_n \dot{a}_n(t) e^{-i\omega_n t} |\Psi_n\rangle = \frac{\gamma}{i\hbar} \sum_n a_n(t) e^{-i\omega_n t} \hat{H}_p(t) |\Psi_n\rangle \quad (6.4.28)$$

To simplify the equation further, we test it with  $\langle \Psi_q |$  to get

$$\dot{a}_q(t) e^{-i\omega_q t} = \frac{\gamma}{i\hbar} \sum_n a_n(t) e^{-i\omega_n t} \langle \Psi_q | \hat{H}_p(t) | \Psi_n \rangle \quad (6.4.29)$$

The above equation is exact as of this point. To solve it further, we expand

$$a_n = a_n^{(0)} + \gamma a_n^{(1)} + \gamma^2 a_n^{(2)} \dots \quad (6.4.30)$$

---

<sup>5</sup>To be precise, the Hamiltonian of a system has to be a constant of motion, since it represents the conservation of energy. But when a quantum system is coupled to another quantum system to which energy is transferred, the first Hamiltonian may not be a constant anymore. But we can assume that the perturbing Hamiltonian is small.

Using the above in (6.4.29), and matching terms of like orders, we have

$$\dot{a}_q^{(0)}(t) = 0 \quad (6.4.31)$$

or that  $a_q^{(0)}$  is a constant. By matching terms of the first order, we have

$$\dot{a}_q^{(1)}(t) = \frac{1}{i\hbar} \sum_n a_n^{(0)} e^{i\omega_{qn}t} \langle \psi_q | \hat{H}_p(t) | \psi_n \rangle \quad (6.4.32)$$

where  $\omega_{qn} = \omega_q - \omega_n$ . To make things simpler, we can assume that the electron is in the  $m$ -th stationary state before we turn on the perturbation  $\hat{H}_p(t)$ . Namely,

$$|\psi\rangle = a_m^{(0)} e^{-i\omega_m t} |\psi_m\rangle = e^{-i\omega_m t} |\psi_m\rangle \quad (6.4.33)$$

In the above, since if only one state exists,  $a_m^{(0)} = 1$  by normalization. Then, (6.4.32) becomes

$$\dot{a}_q^{(1)}(t) = \frac{1}{i\hbar} a_m^{(0)} e^{i\omega_{qm}t} \langle \psi_q | \hat{H}_p(t) | \psi_m \rangle \quad (6.4.34)$$

We assume the perturbing Hamiltonian to be of the form:

$$\hat{H}_p(t) = \begin{cases} 0, & t < 0 \\ \hat{H}_{p0}(e^{-i\omega t} + e^{i\omega t}), & 0 < t < t_0 \\ 0, & t > t_0 \end{cases} \quad (6.4.35)$$

so that it is turned on only for a time window  $0 < t < t_0$ . Then for  $t > t_0$ , by integrating (6.4.34), we have

$$a_q^{(1)} = \frac{1}{i\hbar} \int_0^{t_0} \langle \psi_q | \hat{H}_p(t') | \psi_m \rangle e^{i\omega_{qm}t'} dt' \quad (6.4.36)$$

$$= \frac{1}{i\hbar} \langle \psi_q | \hat{H}_{p0} | \psi_m \rangle \int_0^{t_0} \left( e^{i(\omega_{qm}-\omega)t'} + e^{i(\omega_{qm}+\omega)t'} \right) dt' \quad (6.4.37)$$

$$= -\frac{1}{i\hbar} \langle \psi_q | \hat{H}_{p0} | \psi_m \rangle \left( \frac{e^{i(\omega_{qm}-\omega)t_0} - 1}{\omega_{qm} - \omega} + \frac{e^{i(\omega_{qm}+\omega)t_0} - 1}{\omega_{qm} + \omega} \right) \quad (6.4.38)$$

$$= -\frac{t_0}{i\hbar} \langle \psi_q | \hat{H}_{p0} | \psi_m \rangle \left( e^{\frac{i(\omega_{qm}-\omega)t_0}{2}} \operatorname{sinc} \left( \frac{(\omega_{qm} - \omega)t_0}{2} \right) + e^{\frac{i(\omega_{qm}+\omega)t_0}{2}} \operatorname{sinc} \left( \frac{(\omega_{qm} + \omega)t_0}{2} \right) \right) \quad (6.4.39)$$

The probability of finding the state in  $q$ -state is

$$P(q) = |a_q^{(1)}|^2 \quad (6.4.40)$$

or

$$P(q) \approx \frac{t_0^2}{\hbar^2} |\langle \psi_q | \hat{H}_{p0} | \psi_m \rangle|^2 \left( \operatorname{sinc}^2 \left( \frac{(\omega_{qm} - \omega)t_0}{2} \right) + \operatorname{sinc}^2 \left( \frac{(\omega_{qm} + \omega)t_0}{2} \right) \right) \quad (6.4.41)$$

where we have ignored the cross terms since the sinc functions in (6.4.39) are highly peaked at  $\omega = \pm\omega_{qm}$  when  $t_0 \rightarrow \infty$ ; and hence their cross terms are small.

The above equation means that the probability of transition from the  $m$  to the  $q$  eigenstate is high only if

$$\omega = \pm\omega_{qm} = \pm(\omega_q - \omega_m) \quad (6.4.42)$$

or that

$$\hbar\omega = \pm(\hbar\omega_q - \hbar\omega_m) = \pm(E_q - E_m) \quad (6.4.43)$$

The “+” sign corresponds to when the electron jumps from a low energy  $E_m$  to high energy  $E_q$  requiring the absorption of a photon with energy  $\hbar\omega = (E_q - E_m)$ . The “-” sign in (6.4.43) corresponds to the electron dropping from a high energy state  $E_m$  to a low energy state  $E_q$  emitting a photon with energy  $\hbar\omega = (E_m - E_q)$ . This is known as the Fermi’s golden rule.





# Chapter 7

## Quantum Mechanics in Crystals

### 7.1 Introduction

In a crystal, the atoms are located in a periodic lattice. Hence, when an electron wave propagates on a lattice, it is propagating in a periodic structure. However, the real world is not that simple. Usually, there are more than one electron traveling in a lattice. The electrons see each other's electric field or potentials. There will be electron-electron interaction that has to be accounted for. Moreover, electrons are fermions meaning that they obey Pauli's exclusion principle. Two electrons cannot be in the same state simultaneously. There are many-body effects, but we will ignore them here. The many-body effect can be lumped approximately into an effective potential. We will assume that the effective potential is still periodic, even though this may not be true.

Another effect is that as an electron moves through a lattice, the attractive force between the electron and a nucleus of the atoms distort the lattice locations. This gives rise to lattice vibrations called phonons. We will ignore electron-phonon coupling here.

In a periodic lattice, the atoms are located at the lattice points given by

$$\mathbf{R}_L = n_1 \mathbf{a}_1 + n_2 \mathbf{a}_2 + n_3 \mathbf{a}_3 \quad (7.1.1)$$

where  $\mathbf{a}_1$ ,  $\mathbf{a}_2$ ,  $\mathbf{a}_3$  are lattice vectors, and  $n_1$ ,  $n_2$ ,  $n_3$  are integers. Every lattice point can be related to every other lattice point by a proper choice of  $n_1$ ,  $n_2$ , and  $n_3$ .

Assuming that an electron sees a periodic potential in a lattice, then

$$V_P(\mathbf{r} + \mathbf{R}_L) = V_P(\mathbf{r}) \quad (7.1.2)$$

The single electron sees this periodic potential and its wavefunction satisfies

$$\left[ -\frac{\hbar^2}{2m_e} \nabla^2 + V_P(\mathbf{r}) \right] \psi(\mathbf{r}) = E\psi(\mathbf{r}) \quad (7.1.3)$$

We will study next the kind of wave that can propagate in this periodic potential, known as Bloch-Floquet waves.<sup>1</sup>

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<sup>1</sup>In the subsequent discussions, we will replace  $m_e$  with  $m$ , with implicit understanding that this is the effective mass.

## 7.2 Bloch-Floquet Waves

Let us study the case of a wave propagating in a 1-D periodic structure, where  $V(x)$  may look like that of Figure 7.1.

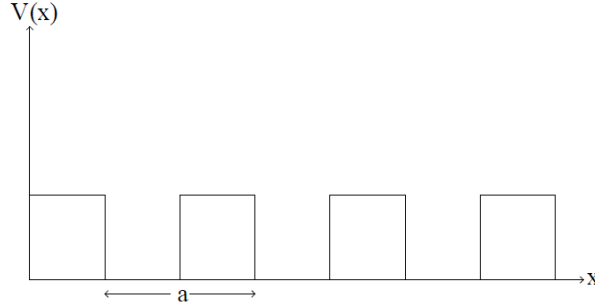


Figure 7.1: A 1D periodic structure of a potential profile for  $V(x)$  where a Bloch-Floquet wave can travel on it.

The periodicity of the potential is  $a$ . Let us assume that the wave that can exist on such a structure is of the form

$$\psi(k, x) = e^{ikx} u_p(x) \quad (7.2.1)$$

This wavefunction must satisfy the 1-D Schrödinger equation that

$$\left[ -\frac{\hbar^2}{2m} \frac{d^2}{dx^2} + V_P(x) \right] \psi(k, x) = E\psi(k, x) \quad (7.2.2)$$

In order for  $\psi(k, x)$  to satisfy the above,  $u_p(x)$  has to satisfy a certain condition. Since the above is a periodic structure, if one translate the coordinate such that  $x \Rightarrow x + na$ , then  $V_P(x) \Rightarrow V_P(x + na) = V_P(x)$  which remains unchanged. Letting  $x \Rightarrow x + na$  in (7.2.1), we have

$$\psi(k, x) \Rightarrow \psi(k, x + na) = e^{ikx} e^{ikna} u_p(x + na) \quad (7.2.3)$$

This new wavefunction must be a solution of Schrödinger equation as well. In order for it to satisfy (7.2.2), we require that

$$u_p(x + na) = u_p(x), \quad \text{for all } n. \quad (7.2.4)$$

In other words,  $u_p(x)$  is a periodic function of  $x$ . Since  $e^{ikna}$  is just a constant, it is clear that  $\psi(x + na)$  satisfies the same equation as (7.2.2). Hence, the form of (7.2.1) that satisfies (7.2.2) has to be such that  $u_p(x)$  is a periodic function. Equation (7.2.1) represents a traveling wave modulated by a periodic function  $u_p(x)$ . However,  $u_p(x)$  is more rapidly varying than  $e^{ikx}$ . Such a wave is a Bloch wave or a Bloch-Floquet wave.

Noting that the second derivative of (7.2.1) is given by

$$\frac{d^2}{dx^2} [e^{ikx} u_p(x)] = -k^2 e^{ikx} u_p(x) + 2ike^{ikx} u_p'(x) + e^{ikx} u_p''(x), \quad (7.2.5)$$

substituting (7.2.1) into (7.2.2), and using the above, we have

$$-\frac{\hbar^2}{2m} [-k^2 u_p(x) + 2ik u_p'(x) + u_p''(x)] + V_p(x) u_p(x) = E u_p(x) \quad (7.2.6)$$

In the above, if  $V_p(x) = 0$ , a simple solution to the above is  $u_p(x)$  a constant with  $E = (\hbar k)^2/(2m)$ . Then the Bloch-Floquet wave becomes a simple traveling wave. Otherwise, we can rewrite (7.2.6) as

$$\left\{ -\frac{\hbar^2}{2m} \left[ \frac{d^2}{dx^2} + 2ik \frac{d}{dx} - k^2 \right] + V_p(x) \right\} u_p(x) = E u_p(x) \quad (7.2.7)$$

with  $u_p(x)$  satisfying the periodic boundary condition within period  $a$ . The above is a simple eigenvalue problem with eigenvector  $u_p(x)$  and eigenvalue  $E$ . We can use matrix method to convert the above into a matrix equation

$$\sum_n H_{mn} a_n = E \sum_n B_{mn} a_n \quad (7.2.8)$$

by letting  $u_p(x) = \sum_{n=1}^N a_n \phi_n(x)$ , and testing with  $\phi_m(x)$ . One also ensures that the choice of basis is such that the periodic boundary condition is satisfied. Here,  $\phi_n(x)$  are not necessarily orthogonal. In the above,

$$H_{mn} = \langle \phi_m | \hat{H}_p | \phi_n \rangle, \quad B_{mn} = \langle \phi_m | \phi_n \rangle \quad (7.2.9)$$

where

$$\hat{H}_p = -\frac{\hbar^2}{2m} \left[ \frac{d^2}{dx^2} + 2ik \frac{d}{dx} - k^2 \right] + V_p(x) \quad (7.2.10)$$

Notice that  $\hat{H}_p$  is a function of  $k$ , and so  $H_{mn}$  is a function of  $k$ . Therefore, (7.2.8) can be written as a generalized eigenvalue problem

$$\bar{\mathbf{H}}(k) \cdot \mathbf{a} = E \bar{\mathbf{B}} \cdot \mathbf{a} \quad (7.2.11)$$

In the above, we have explicitly denoted that  $\bar{\mathbf{H}}(k)$  is a function of  $k$ . Hence, we have to fix  $k$  in order to solve the eigenvalue problem (7.2.7), or the generalized matrix eigenvalue problem (7.2.11). Therefore, the eigenvalue  $E$  obtained is a function of  $k$ , the wavenumber in the Bloch-Floquet wave (7.2.1). In general, a  $E(k)$  plot will look like that in Figure 7.2.

For each  $k$  value, there are many possible values of  $E$ . Also, in (7.2.1), if we let  $k \Rightarrow k + \frac{2p\pi}{a}$ , where  $p$  is an integer value, then, (7.2.1) becomes

$$\psi(k, x) \Rightarrow \psi(k + 2p\pi/a, x) = e^{ikx} u_p(x) e^{\frac{2ip\pi}{a}x} = e^{ikx} \tilde{u}_p(x) \quad (7.2.12)$$

Notice that  $e^{\frac{2ip\pi}{a}x}$  is a periodic function; hence, it can be lumped with  $u_p(x)$  to form a new  $\tilde{u}_p(x)$  periodic function. They satisfy the same equation (7.2.7) with the same  $k$  value. Therefore,  $u_p(x)$  and  $\tilde{u}_p(x)$  share the same set of eigenvalues. They are the same eigenfunctions

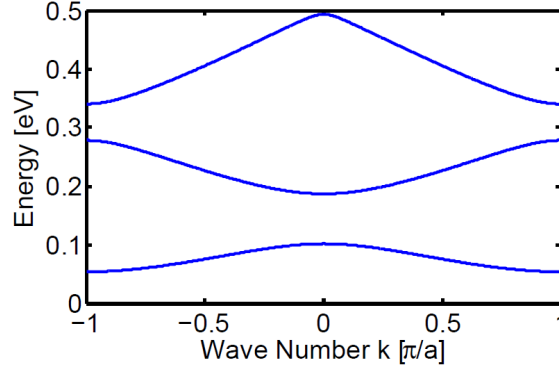


Figure 7.2: The E-k diagram or band structure diagram of a Bloch-Floquet wave propagating in a 1D periodic structure.

within a multiplicative constant. Consequently,  $E(k)$  is a periodic function in  $k$  with period  $\frac{2\pi}{a}$ . Each of this period is called the Brillouin zone. The zone that is centered about the origin is called the first Brillouin zone.

Furthermore, since everything in the operator of (7.2.2) is real,  $\psi^*(k, x)$  is also a solution of it whenever  $\psi(k, x)$  is a solution. But

$$\psi^*(k, x) = e^{-ikx} u_p^*(x) \quad (7.2.13)$$

is also a Bloch-Floquet wave with  $-k$  wavenumber. In fact,  $u_p^*(x)$  satisfies an equation that is the conjugate of (7.2.7). But since the operator in (7.2.7) is Hermitian, its matrix representation remains Hermitian. Furthermore, the eigenvalues are real, and remain unchanged even when the equation is conjugated.

Since  $u_p(x)$  is periodic function with period  $a$ , it can be expanded as a Fourier series, namely

$$u_p(x) = \sum_{n=-\infty}^{\infty} a_n e^{in\pi x/(2a)} \quad (7.2.14)$$

### 7.3 Bloch-Floquet Theorem for 3D

In three dimensions, the Bloch-Floquet wave looks like

$$\psi(\mathbf{k}, \mathbf{r}) = e^{i\mathbf{k}\cdot\mathbf{r}} u_p(\mathbf{r}) \quad (7.3.1)$$

with the property that

$$u_p(\mathbf{r} + \mathbf{R}_L) = u_p(\mathbf{r}) \quad (7.3.2)$$

We can expand  $u_p(\mathbf{r})$  as a generalized Fourier series

$$u_p(\mathbf{r}) = \sum_{\mathbf{G}} a_{\mathbf{G}} e^{i\mathbf{G}\cdot\mathbf{r}} \quad (7.3.3)$$

where  $\mathbf{G}$  represents points on the reciprocal lattice denoted by

$$\mathbf{G} = l_1 \mathbf{b}_1 + l_2 \mathbf{b}_2 + l_3 \mathbf{b}_3 \quad (7.3.4)$$

with the property that

$$e^{i\mathbf{G}\cdot\mathbf{R}_L} = 1 \quad (7.3.5)$$

where

$$\mathbf{R}_L = n_1 \mathbf{a}_1 + n_2 \mathbf{a}_2 + n_3 \mathbf{a}_3 \quad (7.3.6)$$

In the above,  $l_i$  and  $n_i$ ,  $i = 1, 2, 3$  are integers. The above property indicates that  $u_p(\mathbf{r})$  in (7.3.3) is periodic as indicated by (7.3.2). The summation in (7.3.2) is over all possible values of  $\mathbf{G}$  or  $l_1$ ,  $l_2$ , and  $l_3$ .

One can solve for  $\mathbf{b}_1, \mathbf{b}_2$  and  $\mathbf{b}_3$  to yield

$$\mathbf{b}_1 = \frac{\mathbf{a}_2 \times \mathbf{a}_3}{\mathbf{a}_1 \cdot (\mathbf{a}_2 \times \mathbf{a}_3)} 2\pi, \quad \mathbf{b}_2 = \frac{\mathbf{a}_3 \times \mathbf{a}_1}{\mathbf{a}_1 \cdot (\mathbf{a}_2 \times \mathbf{a}_3)} 2\pi, \quad \mathbf{b}_3 = \frac{\mathbf{a}_1 \times \mathbf{a}_2}{\mathbf{a}_1 \cdot (\mathbf{a}_2 \times \mathbf{a}_3)} 2\pi \quad (7.3.7)$$

for this  $\mathbf{G}$ . Then by back substitution, (7.3.5) is satisfied. Furthermore,

$$\psi(\mathbf{k} + \mathbf{G}, \mathbf{r}) = e^{i\mathbf{k}\cdot\mathbf{r}} e^{i\mathbf{G}\cdot\mathbf{r}} u_p(\mathbf{r}) \quad (7.3.8)$$

where

$$e^{i\mathbf{G}\cdot(\mathbf{r}+\mathbf{R}_L)} = e^{i\mathbf{G}\cdot\mathbf{r}} \quad (7.3.9)$$

Since it is a periodic function on the lattice, the  $e^{i\mathbf{G}\cdot\mathbf{r}}$  term can be lumped with the unit cell function  $u_p(\mathbf{r})$ ; therefore, the new Bloch-Floquet wave is

$$\psi(\mathbf{k} + \mathbf{G}, \mathbf{r}) = e^{i\mathbf{k}\cdot\mathbf{r}} \tilde{u}_p(\mathbf{r}) \quad (7.3.10)$$

It is of the same form as (7.3.1); and hence will yield the same eigenvalues as  $\psi(\mathbf{k}, \mathbf{r})$ . Consequently, the lattice points are defined by the lattice vector  $\mathbf{R}_L$ , and the reciprocal lattice points are defined by the reciprocal lattice vector  $\mathbf{G}$ . The Brillouin zone is defined by the reciprocal lattice points. The first Brillouin zone is the first unique zone centered around the origin. These zones in 3D can be quite complicated as shown in Figure 7.3.

The band structure diagram is rather complicated and is usually plot with  $E$  versus  $k$  vector along prescribed lines in the first Brillouin zone as shown in Figure 7.4.

We can take yet another viewpoint of the existence of band structure in a crystal. We have studied the tight-binding model of two identical quantum wells. This can be extended to two identical atoms as shown in Figure 7.5. The figure shows how the energy of the trapped modes of the atoms are split versus the spacing of the atoms (as we have found on our tight binding model in the previous chapter). The closer the spacing, the stronger is the coupling between the modes, and the larger is the split in the level of the degenerate modes.

If we have a system with billions of atoms, each atom has its energy levels due to trapped modes in the atom. If the inter-atomic distance is large, there is no interaction between the atoms, and energy levels will be many billion times degenerate. However, when the interatomic spacing becomes smaller, the wavefunctions from one atom will overlap with the

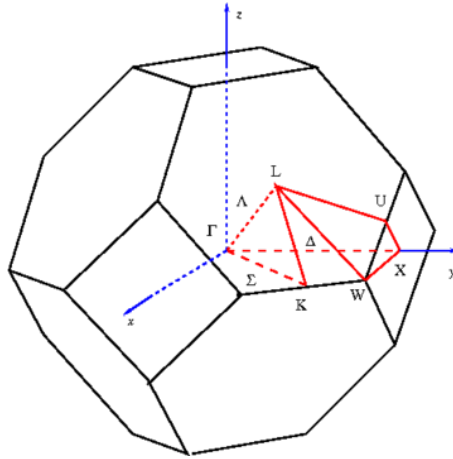


Figure 7.3: The first Brillouin zone of the FCC (face centered cubic) lattice (from Wikipedia).

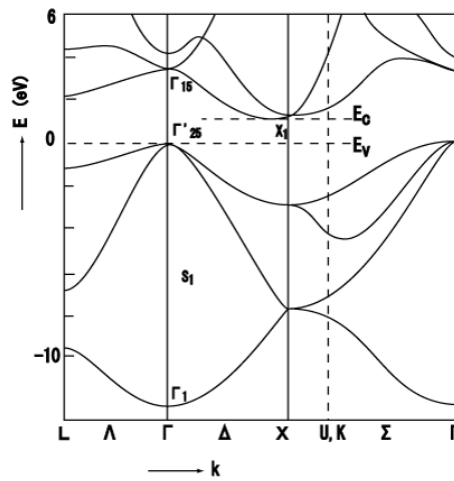


Figure 7.4: The band structure diagram of Silicon along prescribed lines in the first Brillouin zone (from Warwick Physics Dept.).

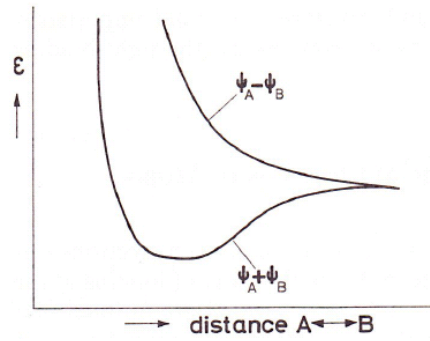


Figure 7.5: The energy levels of two atoms versus the atomic spacing. (from Seeger, Semiconductor Physics).

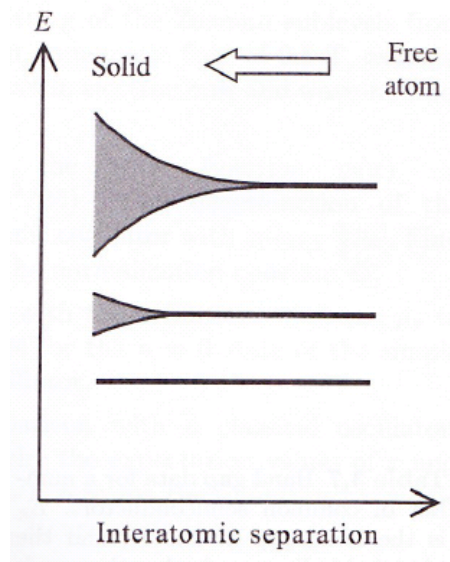


Figure 7.6: The genesis of band structure as the interatomic spacing of the atoms of the lattice becomes smaller and smaller. (from Fox, Quantum Optics).

neighboring atoms through tunneling. The interatomic coupling, like the example of the tight-binding model, will split the energy levels and make them into a continuous band of energy. This give rise to energy band where energy states exist and energy band where no states exist, yielding the band structures we have seen.

As we have seen in the simple case of the tight-binding model study of two quantum wells, the degenerate modes split into even and odd mode coupling between them. In one case, the modes are in phase, and the other case, the modes are out of phase with respect to each other. When multitude of mode coupling prevails, there could be infinitely many possible phase relations between the atomic modes, giving rise to Bloch-Floquet wave with different  $\exp(i\mathbf{k} \cdot \mathbf{r})$  across the lattice. The Bloch-Floquet wave establishes the phase relationship between different trapped modes of the atoms as they couple to each other.

## 7.4 Effective Mass Schrödinger Equation

The bottom of many conduction band is parabolic. In the vicinity of  $k = 0$ , we can write the  $E$ - $k$  relation as

$$E_k = \frac{\hbar^2 k^2}{2m_e} + V \quad (7.4.1)$$

where  $m_e$  is the effective mass chosen to fit the curvature of the parabola, and  $V$  is the energy at the bottom of the conduction band. We can write the Bloch-Floquet wave as a superposition of waves with different energies or  $k$  values, namely

$$\psi(\mathbf{r}, t) = \sum_k c_k u_k(\mathbf{r}) e^{i\mathbf{k} \cdot \mathbf{r}} e^{-i\omega_k t}, \quad k \simeq 0, \quad \hbar\omega_k = E_k \quad (7.4.2)$$

When  $k \simeq 0$ , we can approximate

$$u_k(\mathbf{r}) = u_0(\mathbf{r}) \quad (7.4.3)$$

so that

$$\begin{aligned} \psi(\mathbf{r}, t) &= u_0(\mathbf{r}) \sum_k c_k e^{i\mathbf{k} \cdot \mathbf{r}} e^{-i\omega_k t} \\ &= u_0(\mathbf{r}) \psi_e(\mathbf{r}, t) \end{aligned} \quad (7.4.4)$$

where  $\psi_e(\mathbf{r}, t)$  is the envelope function, namely

$$\psi_e(\mathbf{r}, t) = \sum_k c_k e^{i\mathbf{k} \cdot \mathbf{r}} e^{-i\omega_k t} \quad (7.4.5)$$



We can show that

$$\begin{aligned}
 i\hbar \frac{\partial \psi_e(\mathbf{r}, t)}{\partial t} &= \sum_k c_k E_k e^{i\mathbf{k}\cdot\mathbf{r} - i\omega_k t} \\
 &= \sum_k c_k \left( \frac{\hbar^2 k^2}{2m_e} + V \right) e^{i\mathbf{k}\cdot\mathbf{r} - i\omega_k t} \\
 &= \left( -\frac{\hbar^2 \nabla^2}{2m_e} + V \right) \sum_k c_k e^{i\mathbf{k}\cdot\mathbf{r} - i\omega_k t}
 \end{aligned} \tag{7.4.6}$$

or

$$i\hbar \frac{\partial \psi_e(\mathbf{r}, t)}{\partial t} = \left( -\frac{\hbar^2 \nabla^2}{2m_e} + V \right) \psi_e(\mathbf{r}, t) \tag{7.4.7}$$

The above is the effective mass Schrödinger equation. It is valid when  $k$  is small, or  $e^{i\mathbf{k}\cdot\mathbf{r}}$  is slowly varying compared to the lattice spacing  $a$ . In the above, we can assume that  $V(\mathbf{r})$  is slowly varying, and that the effective mass  $m_e(\mathbf{r})$  is also a slowly varying function of  $\mathbf{r}$ . Then the above can be correctly written as

$$-\frac{\hbar^2}{2} \nabla \cdot \frac{1}{m_e(\mathbf{r})} \nabla \psi_e(\mathbf{r}, t) + V(\mathbf{r}) \psi_e(\mathbf{r}, t) = i\hbar \frac{\partial}{\partial t} \psi_e(\mathbf{r}, t) \tag{7.4.8}$$

The essence of the above approximation is that if the material properties are slowly varying with respect to the lattice spacing, an electron “sees” locally an environment of an infinite lattice. Consequently, if there are two domains with different  $V$  and  $m_e$ , the above equation induces the boundary conditions at the interface of the domains as

$$\psi_{e1}(\mathbf{r}) = \psi_{e2}(\mathbf{r}), \quad \mathbf{r} \in S \tag{7.4.9}$$

$$\frac{1}{m_{e1}} \frac{\partial}{\partial n} \psi_{e1}(\mathbf{r}) = \frac{1}{m_{e2}} \frac{\partial}{\partial n} \psi_{e2}(\mathbf{r}), \quad \mathbf{r} \in S \tag{7.4.10}$$

where  $S$  is the interface separating the two domains and  $\partial/\partial n$  refers to normal derivative. Even though  $V(\mathbf{r})$  and  $m_e(\mathbf{r})$  have jump discontinuities at such an interface, we assume that the discontinuities are still slowly varying compared to the lattice spacings.

## 7.5 Density of States (DOS)

We can use periodic boundary condition to discretize the  $k$  vector in the Bloch-Floquet wave so that it is countable, or that

$$\mathbf{k}_L = l_1 \frac{2\pi}{L_1} \hat{a}_1 + l_2 \frac{2\pi}{L_2} \hat{a}_2 + l_3 \frac{2\pi}{L_3} \hat{a}_3 \tag{7.5.1}$$

In one dimension, it is

$$k_l = l \frac{2\pi}{L} = l\Delta k \tag{7.5.2}$$

where  $\Delta k = \frac{2\pi}{L}$ . When we count the number of discrete modes, we have 1 mode per  $\Delta k$  on the  $k_l$  real line. Or the density of states is

$$g_{1D}(k) = 1/\Delta k = L/2\pi \quad (7.5.3)$$

The number of states is proportional to the length of  $L$ . In 2D, the density of states can be easily derived to be

$$g_{2D}(k) = L_1 L_2 / (2\pi)^2 = A / (2\pi)^2 \quad (7.5.4)$$

In 3D, it becomes

$$g_{3D}(k) = L_1 L_2 L_3 / (2\pi)^3 = V / (2\pi)^3 \quad (7.5.5)$$

We can normalize the density of states with respect to length  $L$ , area  $A$ , and volume  $V$  for 1D, 2D, and 3D, respectively. Then

$$g_{1D}(k) = 1/2\pi, \quad g_{2D}(k) = 1/(2\pi)^2, \quad g_{3D}(k) = 1/(2\pi)^3 \quad (7.5.6)$$

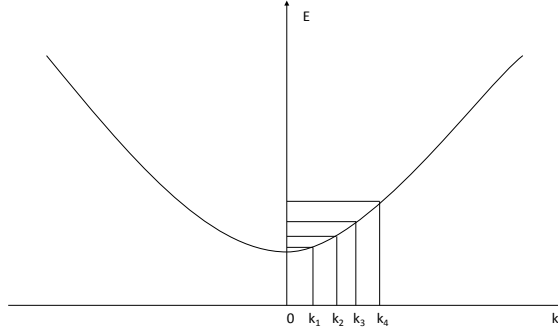


Figure 7.7: If every state on the  $k$  line corresponds to a state on the  $E$  line, then the density of states per unit length on the  $E$  line is different from that on the  $k$  line.

We may want to count the density of states per unit energy instead as shown by the above  $E$ - $k$  diagram. Even though the states are evenly spaced on the  $k$  line, they are not evenly spread on the  $E$  line, thereby, altering the density of states per  $\Delta E$ . In general, we can write, in 3D,

$$\begin{aligned} g(E) dE &= 2g(k) d^3k = 2g(k) 4\pi k^2 dk \\ &= 8\pi g(k) k^2 \frac{dk}{dE} dE \end{aligned} \quad (7.5.7)$$

where we put a factor of 2 for two electron states per energy state.

Therefore, the relation between DOS on the  $E$  line and the  $k$  line is

$$g(E) = 8\pi g(k) k^2 \frac{dk}{dE} \quad (7.5.8)$$

Since

$$E = \frac{\hbar^2 k^2}{2m_e} + V \quad (7.5.9)$$

for a free electron, we have

$$k = \sqrt{E - V} \frac{\sqrt{2m_e}}{\hbar} \quad (7.5.10)$$

$$\frac{dk}{dE} = \sqrt{\frac{m_e}{2\hbar^2}} \frac{1}{\sqrt{E - V}} \quad (7.5.11)$$

and

$$g(E) = \frac{1}{2\pi^2} \left( \frac{2m_e}{\hbar^2} \right)^{\frac{3}{2}} (E - V)^{\frac{1}{2}} \quad (7.5.12)$$

The above is the DOS in the conduction band of a bulk material.

## 7.6 DOS in a Quantum Well

In a 1D quantum well, the potential varies as a function of  $z$ , but is independent of  $x$  and  $y$ . Then the effective mass Schrödinger equation becomes

$$-\frac{\hbar^2}{2} \nabla \frac{1}{m_e(z)} \nabla \psi(\mathbf{r}) + V(z)\psi(\mathbf{r}) = E\psi(\mathbf{r}) \quad (7.6.1)$$

We can further separate the  $z$  variation from the  $x$  and  $y$  variations to get

$$-\frac{\hbar^2}{2m_e} \nabla_s^2 \psi(\mathbf{r}) - \frac{\hbar^2}{2} \frac{\partial}{\partial z} \frac{1}{m_e} \frac{\partial}{\partial z} \psi(\mathbf{r}) + V(z)\psi(\mathbf{r}) = E\psi(\mathbf{r}) \quad (7.6.2)$$

where  $\nabla_s^2 = \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2}$ .

By separation of variables, we let

$$\psi(\mathbf{r}) = \psi_n(z)\psi_s(\mathbf{r}_s) \quad (7.6.3)$$

where  $\mathbf{r}_s = \hat{x}x + \hat{y}y$ . Then, we can show that

$$-\frac{\hbar^2}{2m_e} \nabla_s^2 \psi_s(\mathbf{r}_s) = E_s \psi_s(\mathbf{r}_s) \quad (7.6.4)$$

and

$$-\frac{\hbar^2}{2} \frac{d}{dz} \frac{1}{m_e(z)} \frac{d}{dz} \psi_n(z) + V(z)\psi_n(z) = E_n \psi_n(z) \quad (7.6.5)$$

In the above

$$E = E_s + E_n \quad (7.6.6)$$

Equation (7.6.5) can be solved for the eigenvalues  $E_n$ . For example, it can be the 1-D finite potential well problem, or the infinite potential well problem. Alternatively, it can have a complicated  $V(z)$  and an  $m_e(z)$  that are amenable to numerical methods only.

In the  $x$ - $y$  direction, we let

$$\psi(\mathbf{r}_s) \propto e^{i\mathbf{k}_s \cdot \mathbf{r}_s} \quad (7.6.7)$$

where  $\mathbf{k}_s = \hat{x}k_x + \hat{y}k_y$ . Then,

$$E_s = \frac{\hbar^2 k_s^2}{2m_e} \quad (7.6.8)$$

where  $k_s^2 = k_x^2 + k_y^2$ . Therefore, the total  $E$ - $k_s$  diagram looks like that shown in Figure 7.8

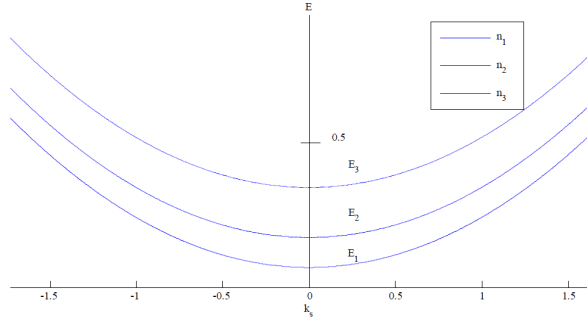


Figure 7.8: The subband  $E$ - $k_s$  relation of a quantum well.

In general,  $g_{2D}(k_s) = \frac{1}{(2\pi)^2}$ , and we have

$$g_{2D}(E_s)dE_s = 2g_{2D}(k_s)2\pi k_s \frac{dk_s}{dE_s} dE_s \quad (7.6.9)$$

Finally, we get

$$g_{2D}(E_s) = \frac{m_e}{\pi\hbar^2} \quad (7.6.10)$$

which is a constant independent of  $E_s$ .

Hence, if we march along the  $E$ -line from 0 upward, the DOS as a function of total  $E$  is shown in Figure 7.9.

If the 1D quantum well is an infinite potential well, the energy is given by

$$E_n = \frac{\hbar^2}{2m_e} \left( \frac{n\pi}{L} \right)^2 \quad (7.6.11)$$

But density of states for a 3D bulk material is

$$g_{3D}(E) = \frac{1}{2\pi^2} \left( \frac{2m_e}{\hbar^2} \right)^{\frac{3}{2}} E^{\frac{1}{2}} \quad (7.6.12)$$

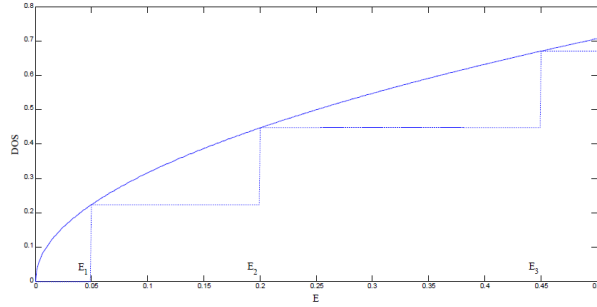


Figure 7.9: The density of states of a 2D quantum well compared to that of a 3D bulk material.

If we evaluate this at  $E = E_n$ , we have

$$g_{3D}(E_n) = \frac{m_e}{\pi \hbar^2} \left( \frac{n}{L} \right) = g_{2D} \frac{n}{L} \tag{7.6.13}$$

Therefore,  $\frac{g_{2D}}{L}$  touches the  $g_{3D}$  at  $E_n$  values as shown in Figure 7.9.

When  $L$  becomes large, so that the quantization level in the quantum well becomes finer, the DOS plot versus total  $E$  is shown in Figure 7.10. The DOS for a quantum well resembles that of a bulk material.

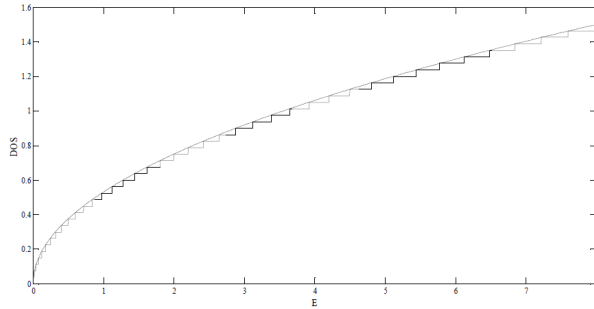


Figure 7.10: The density of states of a 2D quantum well compared to that of a 3D bulk material when the width of the quantum well  $L$  is large.



# Chapter 8

## Angular Momentum

### 8.1 Introduction

In quantum mechanics, we learn that travelling wave  $e^{ikx}$  carries a linear momentum proportional to  $\hbar k$ . Hence, the more oscillatory the wave is, the higher the momentum it carries. However, when an electron is trapped inside the potential of an atom, the wave is not a linearly travelling wave, but a wave that swirls around the nucleus of the atom. Hence, the wave carries mostly angular momentum. The magnitude of this angular momentum, intuitively, is proportional to the angular variation of the wavefunction.

If we take Schrödinger's Equation in free space, it is given by

$$-\frac{\hbar^2}{2m_0}\nabla^2\psi(\mathbf{r}) = E\psi(\mathbf{r}) \quad (8.1.1)$$

We can rewrite this as

$$(\nabla^2 + k_0^2)\psi(\mathbf{r}) = 0 \quad (8.1.2)$$

where  $k_0^2 = 2m_0E/\hbar^2$ . The above is the Helmholtz wave equation whose solution is well known. In cylindrical coordinates, it is

$$\left(\frac{1}{\rho}\frac{\partial}{\partial\rho}\rho\frac{\partial}{\partial\rho} + \frac{1}{\rho^2}\frac{\partial^2}{\partial\phi^2} + \frac{\partial^2}{\partial z^2} + k_0^2\right)\psi(\rho, \phi, z) = 0 \quad (8.1.3)$$

where  $\rho = \sqrt{x^2 + y^2}$ ,  $\phi = \tan^{-1}(y/x)$ . In spherical coordinates, it is

$$\left(\frac{1}{r^2}\frac{\partial}{\partial r}r^2\frac{\partial}{\partial r} + \frac{1}{r^2\sin\theta}\frac{\partial}{\partial\theta}\sin\theta\frac{\partial}{\partial\theta} + \frac{1}{r^2\sin^2\theta}\frac{\partial^2}{\partial\phi^2} + k_0^2\right)\psi(r, \theta, \phi) = 0 \quad (8.1.4)$$

where  $r = \sqrt{x^2 + y^2 + z^2}$ ,  $\theta = \cos^{-1}(z/r)$ ,  $\phi = \tan^{-1}(y/x)$ .

### 8.1.1 Electron Trapped in a Pill Box

We will first solve the above equation in cylindrical coordinates to gain physical insight into the wave-angular momentum relationship. Hence, we study the case of an electron trapped inside a pill box. The solutions to Helmholtz equation can be found by the separation of variables. For example, to solve (8.1.3), we let

$$\psi(\rho, \phi, z) = R(\rho)\Phi(\phi)Z(z) \quad (8.1.5)$$

Using (8.1.5) in (8.1.3) yields

$$\begin{aligned} \Phi(\phi)Z(z)\frac{1}{\rho}\frac{\partial}{\partial\rho}\rho\frac{\partial}{\partial\rho}R(\rho) + R(\rho)Z(z)\frac{1}{\rho^2}\frac{\partial^2}{\partial\phi^2}\Phi(\phi) \\ + R(\rho)\Phi(\phi)\frac{\partial^2}{\partial z^2}Z(z) + k_0^2R(\rho)\Phi(\phi)Z(z) = 0 \end{aligned} \quad (8.1.6)$$

To solve the above, we propose two eigenvalue problems,

$$\frac{\partial^2}{\partial z^2}Z(z) = -k_z^2Z(z) \quad (8.1.7)$$

$$\frac{\partial^2}{\partial\phi^2}\Phi(\phi) = -n^2\Phi(\phi) \quad (8.1.8)$$

The solutions to the above are easily found:

$$Z(z) = e^{\pm ik_z z} \quad (8.1.9)$$

$$\Phi(\phi) = e^{\pm in\phi} \quad (8.1.10)$$

By so doing, Equation (8.1.6) becomes

$$\frac{1}{\rho}\frac{d}{d\rho}\rho\frac{d}{d\rho}R(\rho) - \frac{n^2}{\rho^2}R(\rho) + k_\rho^2R(\rho) = 0 \quad (8.1.11)$$

where

$$k_\rho^2 = k_0^2 - k_z^2 \quad (8.1.12)$$

Equation (8.1.11) is the Bessel equation where solutions are Bessel function  $J_n(k_\rho\rho)$  and Neumann function  $Y_n(k_\rho\rho)$ . Here,  $n$  is the order of these functions. So the general solution (8.1.5) becomes

$$\psi(\rho, \phi, z) = [AJ_n(k_\rho\rho) + BY_n(k_\rho\rho)]e^{\pm in\phi}e^{\pm ik_z z} \quad (8.1.13)$$

We can now put an electron in a pill box whose wall is an infinite potential as shown in Figure 8.1. Since  $Y_n(k_\rho\rho) \rightarrow \infty$ , when  $\rho \rightarrow 0$ , we set  $B = 0$  to eliminate the Neumann function. Also,  $\psi(\rho, \phi, z)$  has to be regular and finite inside the pill box. Matching boundary conditions at the two ends of the pill box, we have

$$\psi(\rho, \phi, z) = AJ_n(k_\rho\rho)e^{\pm in\phi}\sin\left(\frac{p\pi}{L}z\right) \quad (8.1.14)$$



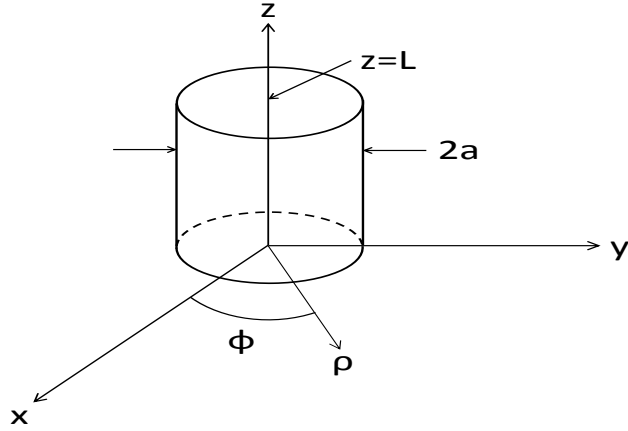


Figure 8.1: The geometry of a cylindrical pill box where an electron is trapped.

Table 8.1: The first few zeros of Bessel functions:

$m$	$\zeta_{0m}$	$\zeta_{1m}$	$\zeta_{2m}$	$\zeta_{3m}$	$\zeta_{4m}$	$\zeta_{5m}$
1	2.4048	3.8317	5.1356	6.3802	7.5883	8.7715
2	5.5201	7.0156	8.4172	9.7610	11.0647	12.3386
3	8.6537	10.1735	11.6198	13.0152	14.3725	15.7002
4	11.7915	13.3237	14.7960	16.2235	17.6160	18.9801
5	14.9309	16.4706	17.9598	19.4094	20.8269	22.2178

where  $k_z = p\pi/L$ , for all integer  $p$ .

The boundary condition at  $\rho = a$  is that  $\psi(a, \phi, z) = 0$ , or that

$$J_n(k_\rho a) = 0 \tag{8.1.15}$$

The zeros of Bessel functions are tabulated and given by

$$J_n(\zeta_{nm}) = 0 \tag{8.1.16}$$

The first few zeros are given in Table 8.1. Hence, we obtain that

$$k_\rho = \zeta_{nm}/a \tag{8.1.17}$$

So (8.1.14) becomes

$$\psi(\rho, \phi, z) = AJ_n\left(\zeta_{nm}\frac{\rho}{a}\right)e^{\pm in\phi}\sin\left(\frac{p\pi}{L}z\right) \tag{8.1.18}$$

In the above, we see a standing wave  $\sin(p\pi z/L)$  in the  $z$  direction, a standing wave represented by  $J_n(\zeta_{nm}\frac{\rho}{a})$  in the  $\rho$  direction, and a travelling wave  $e^{\pm in\phi}$  in the  $\phi$  direction. The

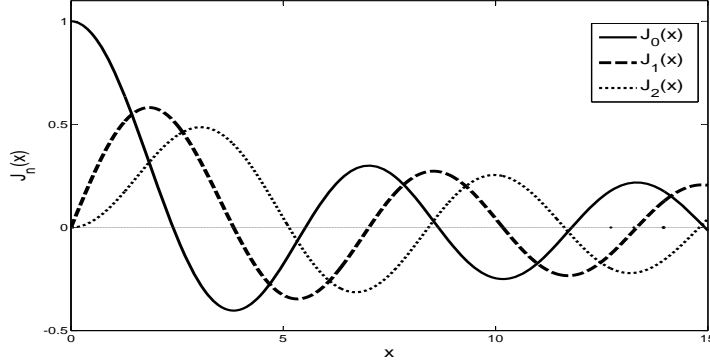


Figure 8.2: The Bessel function of different order as a function of its argument. Notice that the higher the order, the smaller the Bessel function is at the origin.

larger  $n$  is, the more rapidly varying is the travelling wave as a function of  $\phi$ , and the more angular momentum it carries.

In fact, plots of  $J_n(x)$  versus  $x$  for different  $n$  is shown in Figure 8.2.<sup>1</sup> We notice that the larger  $n$  is, the more rapidly varying  $e^{\pm im\phi}$  is, and the more the Bessel function  $J_n(x)$  is pulled away from the origin. This is characteristic of a particle swirling around in a circle. The centrifugal force is keeping the particle away from the origin.

From the fact that

$$k_0^2 = k_\rho^2 + k_z^2 = \left(\frac{\zeta_{nm}}{a}\right)^2 + \left(\frac{p\pi}{L}\right)^2 = \frac{2m_0E}{\hbar^2} \quad (8.1.19)$$

we can derive that the energy levels of the trapped electron are given by

$$E_{nmp} = \frac{\hbar^2}{2m_0} \left[ \left(\frac{\zeta_{nm}}{a}\right)^2 + \left(\frac{p\pi}{L}\right)^2 \right] \quad (8.1.20)$$

Hence, the energy levels of the trapped electron assume quantized values. This is the characteristic of the wave nature of the electron.

### 8.1.2 Electron Trapped in a Spherical Box

To obtain different physical insight with different wavefunctions, next, we place the electron in a spherical box with infinite potential at the wall. We also solve this problem by separation of variables by letting

$$\psi(r, \theta, \phi) = R(r) \Theta(\theta) \Phi(\phi) \quad (8.1.21)$$

<sup>1</sup>Ways to compute  $J_n(x)$  is well documented and is available in programming toolbox like Matlab.

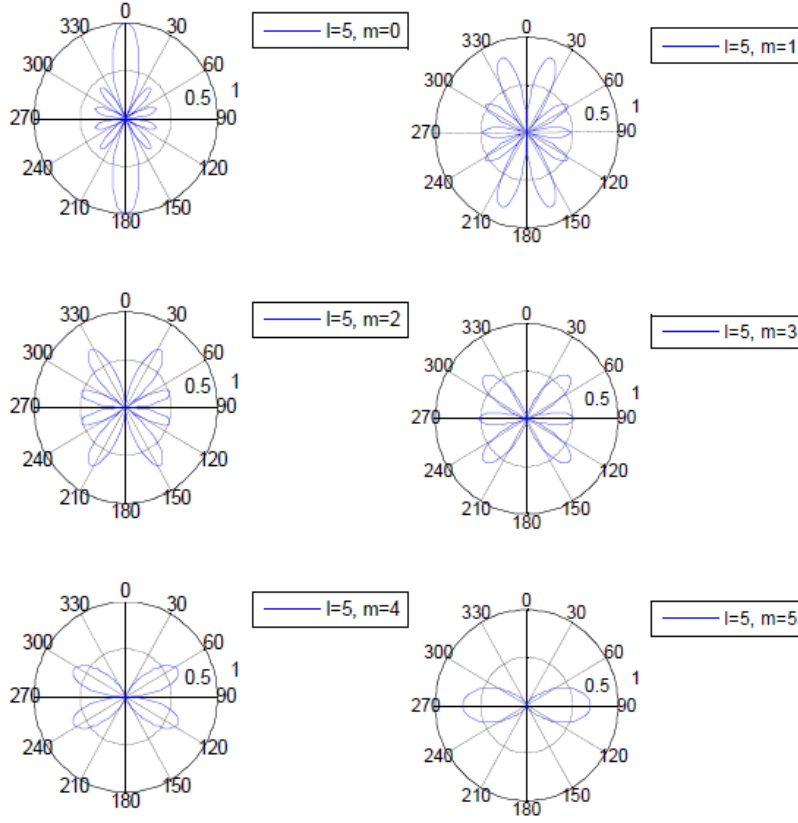


Figure 8.3: Legendre Polynomial with argument  $\cos \theta$  plotted for  $l = 5$  and different  $m$  values.

Using (8.1.21) in (8.1.4), we have

$$\Theta(\theta) \Phi(\phi) \frac{1}{r^2} \frac{\partial}{\partial r} r^2 \frac{\partial}{\partial r} R(r) + R(r) \Phi(\phi) \frac{1}{r^2 \sin \theta} \frac{\partial}{\partial \theta} \left( \sin \theta \frac{\partial}{\partial \theta} \Theta(\theta) \right) + R(r) \Theta(\theta) \frac{1}{r^2 \sin^2 \theta} \frac{\partial^2}{\partial \phi^2} \Phi(\phi) + k_0^2 R(r) \Theta(\theta) \Phi(\phi) = 0 \quad (8.1.22)$$

We propose another eigenvalue problem such that

$$\frac{\partial^2}{\partial \phi^2} \Phi(\phi) = -m^2 \Phi(\phi) \quad (8.1.23)$$

Using the above in (8.1.22),  $\Phi(\phi)$  can be canceled in the equation. Then we propose another eigenvalue problem such that

$$\frac{1}{\sin \theta} \frac{\partial}{\partial \theta} \sin \theta \frac{\partial}{\partial \theta} \Theta(\theta) - \frac{m^2}{\sin^2 \theta} \Theta(\theta) = -l(l+1) \Theta(\theta) \quad (8.1.24)$$

The solution to (8.1.23) is  $e^{\pm im\phi}$  while the solution to (8.1.24) is in terms of associated Legendre polynomials,  $P_l^m(x)$ , or in terms of  $\theta$ ,<sup>2</sup>

$$\Theta(\theta) = P_l^m(\cos\theta), \quad -l \leq m \leq l \quad (8.1.25)$$

Eventually, the equation for  $R(r)$  is

$$\frac{1}{r^2} \frac{d}{dr} r^2 \frac{d}{dr} R(r) + \left[ k_0^2 - \frac{l(l+1)}{r^2} \right] R(r) = 0 \quad (8.1.26)$$

The last equation is the spherical Bessel equation whose solution is

$$R(r) = A j_l(k_0 r) + B h_l^{(1)}(k_0 r) \quad (8.1.27)$$

In the above,  $j_l(x)$  represents a spherical Bessel function of order  $l$ , which is regular at  $x = 0$ , while  $h_l^{(1)}(x)$  represents a spherical Hankel function of the first kind of order  $l$ , which is singular at  $x = 0$ .

When we put this solution inside a spherical box, we set  $B = 0$  since  $h_l^{(1)}(k_0 r)$  is singular at  $r = 0$ . Then the general solution to (8.1.4) inside a box is

$$\psi(r, \theta, \phi) = A j_l(k_0 r) P_l^m(\cos\theta) e^{\pm im\phi}, \quad -l \leq m \leq l \quad (8.1.28)$$

Assuming the radius of the sphere is  $a$ , to satisfy the boundary condition, we require that

$$j_l(k_0 a) = 0, \quad (8.1.29)$$

Since  $j_l(\zeta_{lp}) = 0$ , we deduce that  $k_0 = \zeta_{lp}/a$ . The first few zeros of the spherical Bessel functions are given in Table 8.2.

Table 8.2: The first few zeros of the spherical Bessel functions:

$p$	$\zeta_{0p}$	$\zeta_{1p}$	$\zeta_{2p}$	$\zeta_{3p}$	$\zeta_{4p}$
1	3.142	4.493	5.763	6.988	8.183
2	6.283	7.725	9.095	10.417	11.705
3	9.425	10.904	12.323	13.698	15.040
4	12.566	14.066	15.515	16.924	18.301

Therefore, the general solution becomes

$$\psi(r, \theta, \phi) = A j_l\left(\zeta_{lp} \frac{r}{a}\right) P_l^m(\cos\theta) e^{\pm im\phi}, \quad -l \leq m \leq l \quad (8.1.30)$$

In the above,  $e^{\pm im\phi}$  represents a traveling wave in the  $\phi$  direction,  $P_l^m(\cos\theta)$  represents a standing wave in the  $\theta$  direction. When  $l$  becomes large,  $P_l^m(x)$  is a higher order polynomial

<sup>2</sup>The other solution to (8.1.24) is the associate Legendre function of the second kind,  $Q_l^n(\cos\theta)$  but this function is singular for  $0 \leq \theta \leq \pi$ , and hence, is not admissible as a solution.

and  $P_l^m(\cos \theta)$  is a rapidly varying function of  $\theta$ . In this case,  $e^{\pm im\phi}$ ,  $-l \leq m \leq l$  could also be a rapidly varying function of  $\phi$  implying the presence of high angular momentum. From Figure 8.3, it is seen that the larger  $m$  is, the more slowly varying is the associated Legendre polynomial. The variations from  $P_l^m(\cos \theta)$  and  $e^{\pm im\phi}$  compensate each other for a fixed  $l$  so as to keep the “sum” of the angular momentum a constant when  $l$  is fixed.

The function  $j_l(x)$  is well known and is plotted in Figure 8.4, showing that when  $n$  is large, the centrifugal force of the particle pulls it away from the origin as in the pill-box case in cylindrical coordinates. From the fact that

$$k_0^2 = \left( \frac{\zeta_{lp}}{a} \right)^2 = \frac{2m_0 E}{\hbar^2} \quad (8.1.31)$$

we deduce that the energy levels of the trapped electron are

$$E_{lp} = \frac{\hbar^2}{2m_0} \left( \frac{\zeta_{lp}}{a} \right)^2 \quad (8.1.32)$$

This is again quantized due to the wave nature of the electron.

The above cases illustrate the trapping of an electron by a pill box and a spherical box. In nature, the positive charge of a nucleus forms a potential that can trap an electron. In the case of the hydrogen atom, closed form solutions can be obtained for the trapping of an electron by the Coulomb potential of the positive nucleus. This has been documented in many quantum mechanics books.

## 8.2 Mathematics of Angular Momentum

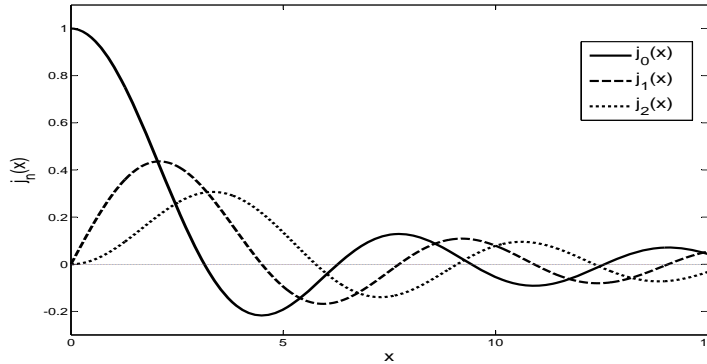


Figure 8.4: Spherical Bessel function of different order versus its argument. The higher order functions are smaller around the origin.

We have shown previously that an electron trapped inside a pill box or a spherical box does display angular momentum. Next, we derive the operator that is related to angular momentum. In classical mechanics, the angular momentum is given by

$$\mathbf{L} = \mathbf{r} \times \mathbf{p} \quad (8.2.1)$$

where  $\mathbf{p}$  is the linear momentum. Notice that by taking the cross product of  $\mathbf{p}$  with  $\mathbf{r}$ , only the circumferential component of the linear momentum is extracted. Written explicitly, it is

$$L_x = yp_z - zp_y, \quad L_y = zp_x - xp_z, \quad L_z = xp_y - yp_x \quad (8.2.2)$$

We raise these to operators, and write them as, in coordinate space representation,

$$\hat{L}_x = -i\hbar \left( y \frac{\partial}{\partial z} - z \frac{\partial}{\partial y} \right), \quad \hat{L}_y = -i\hbar \left( z \frac{\partial}{\partial x} - x \frac{\partial}{\partial z} \right), \quad \hat{L}_z = -i\hbar \left( x \frac{\partial}{\partial y} - y \frac{\partial}{\partial x} \right), \quad (8.2.3)$$

$x, y, z$  can be raised to be operators as well, but in coordinate space representation, they are just scalars. Actually, we can write (8.2.1) as the operator

$$\hat{\mathbf{L}} = \hat{\mathbf{r}} \times \hat{\mathbf{p}} \quad (8.2.4)$$

The hat signs above indicate that they are operators and not unit vectors. The coordinate space representation of the  $\hat{\mathbf{r}}$  operator is just  $\mathbf{r}$ . The coordinate space representation representation of the momentum operator  $\hat{\mathbf{p}}$  has been derived before. Consequently, the above operators can be shown to satisfy the following commutation relations

$$[\hat{L}_x, \hat{L}_y] = i\hbar \hat{L}_z, \quad [\hat{L}_y, \hat{L}_z] = i\hbar \hat{L}_x, \quad [\hat{L}_z, \hat{L}_x] = i\hbar \hat{L}_y \quad (8.2.5)$$

By transforming to spherical coordinates, the operators become

$$\hat{L}_x = i\hbar \left( \sin \phi \frac{\partial}{\partial \theta} + \cot \theta \cos \phi \frac{\partial}{\partial \phi} \right) \quad (8.2.6)$$

$$\hat{L}_y = i\hbar \left( -\cos \phi \frac{\partial}{\partial \theta} + \cot \theta \sin \phi \frac{\partial}{\partial \phi} \right) \quad (8.2.7)$$

$$\hat{L}_z = -i\hbar \frac{\partial}{\partial \phi} \quad (8.2.8)$$

The eigenfunction of the  $\hat{L}_z$  operator is  $\Phi = e^{\pm im\phi}$  so that

$$\hat{L}_z \Phi(\phi) = m\hbar \Phi(\phi) \quad (8.2.9)$$

### 8.3 $\hat{L}^2$ Operator

The  $\hat{L}^2$  operator is defined as

$$\hat{L}^2 = \hat{L}_x^2 + \hat{L}_y^2 + \hat{L}_z^2 \quad (8.3.1)$$

It can be shown that

$$\hat{L}^2 = -\hbar^2 \nabla_{\theta, \phi}^2 \quad (8.3.2)$$

where  $\theta$  and  $\phi$  mean only variation in these variables is involved, and

$$\nabla_{\theta, \phi}^2 = \left[ \frac{1}{\sin \theta} \frac{\partial}{\partial \theta} \sin \theta \frac{\partial}{\partial \theta} + \frac{1}{\sin^2 \theta} \frac{\partial^2}{\partial \phi^2} \right] \quad (8.3.3)$$

The above is just the circumferential variation of the Laplacian operator in (8.1.4). It is quite easy to show that the above commutes with  $\hat{L}_z$ . Hence,  $\hat{L}^2$  commutes with  $\hat{L}_z$ . By symmetry,  $\hat{L}^2$  also commutes with  $\hat{L}_x$  and  $\hat{L}_y$ . We can find the eigenfunction of the  $\hat{L}^2$  operator by first finding that of  $\nabla_{\theta, \phi}^2$ . Traditionally, this eigenfunction and eigenvalue problem is expressed as

$$\nabla_{\theta, \phi}^2 Y_{lm}(\theta, \phi) = -l(l+1)Y_{lm}(\theta, \phi), \quad -l \leq m \leq l \quad (8.3.4)$$

Its solution is similar to the method of solving for the solution of Helmholtz wave equation in spherical coordinates. By the separation of variables, the solution is

$$Y_{lm}(\theta, \phi) = C P_l^m(\cos \theta) e^{im\phi}, \quad -l \leq m \leq l \quad (8.3.5)$$

which can be normalized to yield

$$Y_{lm}(\theta, \phi) = (-1)^m \sqrt{\frac{(2l+1)(l-m)!}{4\pi(l+m)!}} P_l^m(\cos \theta) e^{im\phi}, \quad -l \leq m \leq l \quad (8.3.6)$$

In Hilbert space, this eigenfunction is denoted as  $|l, m\rangle$ . We can define another operator

$$\hat{L}_s^2 = \hat{L}^2 - \hat{L}_z^2 = \hat{L}_x^2 + \hat{L}_y^2 \quad (8.3.7)$$

It is quite clear that  $\hat{L}_z$  commutes with  $\hat{L}_s^2$  and  $\hat{L}_s^2$  commutes with  $\hat{L}^2$ . Hence,  $\hat{L}_z$ ,  $\hat{L}_s^2$  and  $\hat{L}^2$  share the same eigenfunctions  $|l, m\rangle$ . Therefore,

$$\hat{L}_s^2 |l, m\rangle = \hbar^2 [l(l+1) - m^2] |l, m\rangle \quad (8.3.8)$$

A pictorial illustration of this concept is given in Figure 8.5. It is to be noted that this picture can be drawn because  $\hat{L}_z$  commutes with  $\hat{L}_s^2$ . Thus their eigenvalues can be “measured” uniquely and simultaneously. However,  $\hat{L}_z$  does not commute with  $\hat{L}_x$  or  $\hat{L}_y$ . In fact, from the picture of  $|Y_{lm}|^2$ , the function is completely indeterminate in the  $\phi$  direction. Hence, the eigenvalues of  $\hat{L}_x$  or  $\hat{L}_y$  are indeterminate or smeared out for a given  $Y_{lm}$  with a determinate eigenvalue of  $\hat{L}_z$ .

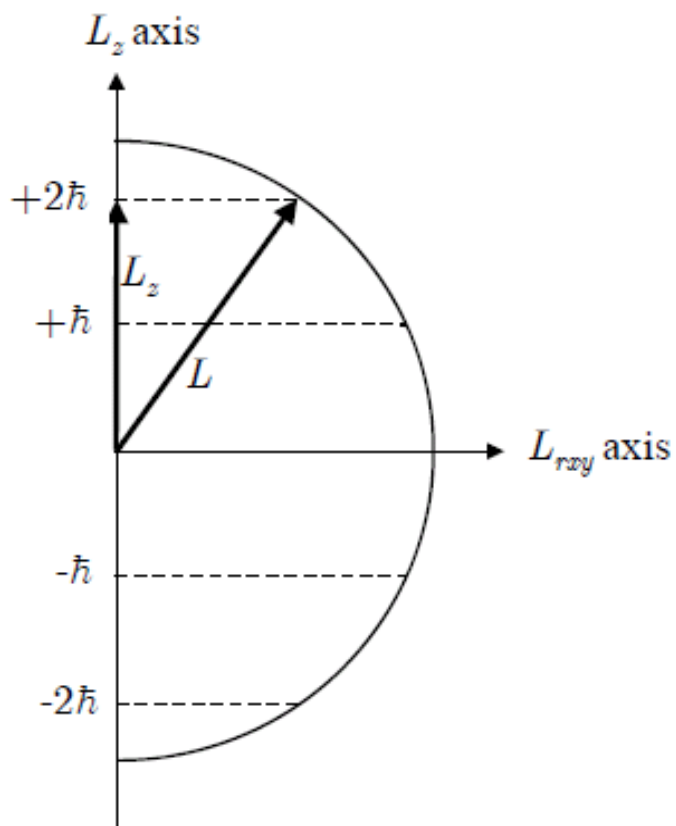


Figure 8.5: A picture showing the angular momentum components for different tilt angles of the total angular momentum  $L$  (from DAB Miller).



# Chapter 9

## Spin

### 9.1 Introduction

Spin is a special property of atomic or subatomic particles that has no classical analogue. Electron has spin. We can think of it as being due to the self spinning of the electron, but we should not let our imagination run further than that. Spin of an electron gives it a spin angular momentum in addition to the orbital angular momentum it possesses. The spin also endows an electron with a magnetic dipole moment that causes it to interact with a magnetic field.

The spin of some particles is found to have binary values of “spin up” and “spin down” experimentally by the famous Stern-Gerlach experiment. This binary nature, as we shall see, fits nicely in the mathematical structure of angular momentum in quantum mechanics, but it cannot be described by a wavefunction or wave mechanics. Instead, it can be represented by matrix mechanics.

### 9.2 Spin Operators

We have seen that the  $z$  component of the orbital angular momentum, represented by the operator  $\hat{L}_z$ , is quantized to be  $m\hbar$  where  $-l \leq m \leq l$ ,  $l$  being an integer related to the total angular momentum square operator  $\hat{L}^2$  with eigenvalue  $l(l+1)\hbar^2$ .

It can be shown that the relationship between the total angular momentum number  $l$  and the  $z$ -component of the angular number  $m$  is not restricted to orbital angular momenta. It can be established for all quantum mechanical angular momenta, as is shown in Appendix A. A more general framework for angular momentum is that for  $\hat{J}^2 = \hat{J}_x^2 + \hat{J}_y^2 + \hat{J}_z^2$ , an operator that represents the square of the total angular momentum, and  $\hat{J}_x$ ,  $\hat{J}_y$ ,  $\hat{J}_z$ , operators that represent the  $x$ ,  $y$ , and  $z$  components of angular momenta, then

$$\hat{J}^2 |L, M\rangle = L(L+1)\hbar^2 |L, M\rangle \quad (9.2.1)$$

$$\hat{J}_z |L, M\rangle = M\hbar |L, M\rangle, \quad -L \leq M \leq L \quad (9.2.2)$$

We have proved the above results for orbital angular momentum by using wave mechanics and wavefunctions, but they can be proven for general angular momentum by using rotational symmetry of 3D coordinate space, and mathematics of raising and lowering operators.

Spin angular momentum operators also fit under the framework of general angular momentum operator, and can be thought of as a special case of the above framework. Hence, the above can also be applied to spin  $\frac{1}{2}$  angular momentum where  $L = \frac{1}{2}$ ,  $-\frac{1}{2} \leq M \leq \frac{1}{2}$ . The values of  $M$  has to be spaced by value one apart, hence,  $M = \pm\frac{1}{2}$ .

For spins, we let  $\hat{S}$  represent the total angular momentum operator, while  $\hat{S}_z$  represents the  $z$  component of the spin angular momentum. Applying (9.2.1) and (9.2.2) to spin  $\frac{1}{2}$  angular momentum, we have

$$\hat{S}^2 \left| \frac{1}{2}, \pm\frac{1}{2} \right\rangle = \frac{3}{4}\hbar^2 \left| \frac{1}{2}, \pm\frac{1}{2} \right\rangle, \quad L = \frac{1}{2}, \quad M = \pm\frac{1}{2} \quad (9.2.3)$$

$$\hat{S}_z \left| \frac{1}{2}, \pm\frac{1}{2} \right\rangle = \pm\frac{1}{2}\hbar \left| \frac{1}{2}, \pm\frac{1}{2} \right\rangle \quad (9.2.4)$$

As a result, the corresponding  $z$  component of the spin angular momentum, represented by the operator  $\hat{S}_z$ , has only two eigenvalues and two eigenstates: an up state with angular momentum of  $\frac{1}{2}\hbar$  and a down state with angular momentum of  $-\frac{1}{2}\hbar$ .<sup>1</sup>

The corresponding  $x$  and  $y$  components of the spin angular momentum can be represented by operators  $\hat{S}_x$  and  $\hat{S}_y$ . Together with  $\hat{S}_z$ , they satisfy the following commutation relations,

$$\left[ \hat{S}_x, \hat{S}_y \right] = i\hbar\hat{S}_z, \quad \left[ \hat{S}_y, \hat{S}_z \right] = i\hbar\hat{S}_x, \quad \left[ \hat{S}_z, \hat{S}_x \right] = i\hbar\hat{S}_y \quad (9.2.5)$$

The above is similar to the commutation relations satisfied by  $\hat{L}_x$ ,  $\hat{L}_y$ , and  $\hat{L}_z$ , where they have been motivated by wave mechanics. However, as has been shown in Appendix A, that if an operator is to represent an angular momentum, then their  $x$ ,  $y$ , and  $z$  components have to satisfy the above commutation relations by rotational symmetry of the 3D coordinate space.

It is expedient at this point to define Pauli spin matrices given by

$$\hat{\sigma}_x = \frac{2}{\hbar}\hat{S}_x, \quad \hat{\sigma}_y = \frac{2}{\hbar}\hat{S}_y, \quad \hat{\sigma}_z = \frac{2}{\hbar}\hat{S}_z \quad (9.2.6)$$

with the commutation relation

$$\left[ \hat{\sigma}_x, \hat{\sigma}_y \right] = 2i\hat{\sigma}_z, \quad \left[ \hat{\sigma}_y, \hat{\sigma}_z \right] = 2i\hat{\sigma}_x, \quad \left[ \hat{\sigma}_z, \hat{\sigma}_x \right] = 2i\hat{\sigma}_y \quad (9.2.7)$$

Since there are only two spin states, they can be represented by a two dimensional column vector. The spin operators can in turn be represented by  $2 \times 2$  matrices. Hence, we have

$$|\uparrow\rangle = \begin{bmatrix} 1 \\ 0 \end{bmatrix}, \quad |\downarrow\rangle = \begin{bmatrix} 0 \\ 1 \end{bmatrix} \quad (9.2.8)$$

Then, the matrix representation of  $\hat{\sigma}_z$  is

$$\hat{\sigma}_z = \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix} \quad (9.2.9)$$

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<sup>1</sup>This was discovered in silver atoms by the famous Stern-Gerlach experiment.

such that

$$\hat{\sigma}_z |\uparrow\rangle = |\uparrow\rangle, \quad \hat{\sigma}_z |\downarrow\rangle = -|\downarrow\rangle \quad (9.2.10)$$

The other Pauli matrices can be found to be

$$\hat{\sigma}_x = \begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix}, \quad \hat{\sigma}_y = \begin{bmatrix} 0 & -i \\ i & 0 \end{bmatrix} \quad (9.2.11)$$

It can be shown that

$$\hat{\sigma}_x^2 = \hat{\sigma}_y^2 = \hat{\sigma}_z^2 = \bar{\mathbf{I}} = \begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} \quad (9.2.12)$$

Hence, we can define a  $\hat{\sigma}^2$  operator such that

$$\hat{\sigma}^2 = \hat{\sigma}_x^2 + \hat{\sigma}_y^2 + \hat{\sigma}_z^2 = 3\bar{\mathbf{I}} \quad (9.2.13)$$

By the same token, the  $\hat{S}^2$  operator, from (9.2.6) and the above, can be evaluated to be

$$\hat{S}^2 = \hat{S}_x^2 + \hat{S}_y^2 + \hat{S}_z^2 = \frac{\hbar^2}{4} \hat{\sigma}^2 = \frac{3}{4} \hbar \bar{\mathbf{I}} \quad (9.2.14)$$

In the above, the quantum mechanics of spin half particle such as an electron is described by matrix mechanics.

### 9.3 The Bloch Sphere

We have seen that the eigenstates of  $\hat{\sigma}_z$  are

$$|\uparrow\rangle = \begin{bmatrix} 1 \\ 0 \end{bmatrix}, \quad |\downarrow\rangle = \begin{bmatrix} 0 \\ 1 \end{bmatrix} \quad (9.3.1)$$

In the two dimensional space in which spin states are represented, these two states are orthonormal and complete. Hence, the eigenstates of the  $\hat{\sigma}_x$  and  $\hat{\sigma}_y$  operators can be expanded as well in terms of  $|\uparrow\rangle$  and  $|\downarrow\rangle$  states. Consequently, a general spin state can be written as

$$|s\rangle = a_\uparrow |\uparrow\rangle + a_\downarrow |\downarrow\rangle = \cos\left(\frac{\theta}{2}\right) |\uparrow\rangle + e^{i\phi} \sin\left(\frac{\theta}{2}\right) |\downarrow\rangle \quad (9.3.2)$$

We have chosen  $a_\uparrow$  and  $a_\downarrow$  judiciously so that the wavefunction is clearly normalized. The relative phase between  $a_\uparrow$  and  $a_\downarrow$  is in  $e^{i\phi}$  as absolute phase is unimportant. Let

$$\boldsymbol{\sigma} = \boldsymbol{i}\hat{\sigma}_x + \boldsymbol{j}\hat{\sigma}_y + \boldsymbol{k}\hat{\sigma}_z \quad (9.3.3)$$

be an operator that represents a vector of the spin angular momentum, where  $\boldsymbol{i}$ ,  $\boldsymbol{j}$ , and  $\boldsymbol{k}$  are unit vectors in the  $x$ ,  $y$ , and  $z$  directions. Then the expectation value of this operator for a given quantum state  $|s\rangle$  in (9.3.2) should point in a vectorial direction. This is a basic

tenet of quantum mechanics since  $\boldsymbol{\sigma}$  represents angular momentum, which is an observable. Consequently, we find the vector

$$\mathbf{P}_s = \langle s | \boldsymbol{\sigma} | s \rangle = \mathbf{i} \langle s | \hat{\sigma}_x | s \rangle + \mathbf{j} \langle s | \hat{\sigma}_y | s \rangle + \mathbf{k} \langle s | \hat{\sigma}_z | s \rangle \quad (9.3.4)$$

After substituting  $|s\rangle$  from (9.3.2) into the above, we can show that

$$\mathbf{P}_s = \mathbf{i} \sin \theta \cos \phi + \mathbf{j} \sin \theta \sin \phi + \mathbf{k} \cos \theta \quad (9.3.5)$$

The vector  $\mathbf{P}_s$  maps out a sphere called the Bloch sphere with corresponding  $\theta$  and  $\phi$  values as shown in Figure 9.1.

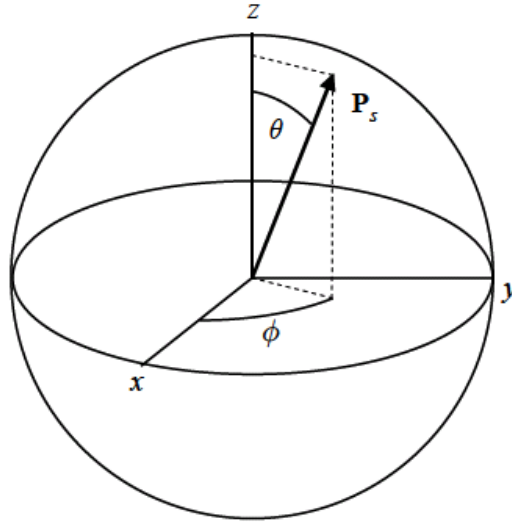


Figure 9.1: The Bloch sphere showing the vector  $\mathbf{P}_s$ .

## 9.4 Spinor

In general, we need a two component vector to describe both the wavefunction and the spin state of an electron, namely,

$$|\psi\rangle = \begin{bmatrix} |\psi_{up}\rangle \\ |\psi_{down}\rangle \end{bmatrix} = |\psi_{up}\rangle \begin{bmatrix} 1 \\ 0 \end{bmatrix} + |\psi_{down}\rangle \begin{bmatrix} 0 \\ 1 \end{bmatrix} \quad (9.4.1)$$

These wave functions with two components are known as spinors. The vector space for spinors consists of the vector space for the electron wavefunction multiplying with the two-dimensional vector space for the spin. This is known as a direct product space.

## 9.5 Pauli Equation

The classical angular momentum of a particle moving in a circle is

$$L = m_0 v r \quad (9.5.1)$$

which can also be written as a vector

$$\mathbf{L} = \mathbf{r} \times \mathbf{p} = \mathbf{r} \times m_0 \mathbf{v} \quad (9.5.2)$$

If this particle is an electron, it will complete an orbit in  $2\pi r/v$  time, and we can denote the current due to this orbiting electron as

$$I = -\frac{ev}{2\pi r} \quad (9.5.3)$$

A circulating current loop produces a magnetic dipole moment

$$\boldsymbol{\mu}_d = I \mathbf{a} \quad (9.5.4)$$

where  $|\mathbf{a}|$  is the area of the current loop, the vector  $\mathbf{a}$  points in the direction normal to the surface  $|\mathbf{a}|$ . For circulating electron, this moment is

$$|\boldsymbol{\mu}_d| = -\frac{ev}{2\pi r} \pi r^2 = -\frac{evr}{2} \quad (9.5.5)$$

We can then write

$$\boldsymbol{\mu}_d = -\frac{e\mathbf{r} \times \mathbf{v}}{2} = -\frac{e\mathbf{L}}{2m_0} \quad (9.5.6)$$

The interaction energy of a magnetic dipole and a background magnetic field  $\mathbf{B}$  is given by

$$E_u = -\boldsymbol{\mu}_d \cdot \mathbf{B} \quad (9.5.7)$$

Assume that  $\mathbf{B} = \hat{z}B_z$ , then

$$E_u = \frac{e}{2m_0} B_z L_z = \frac{e}{2m_0} B_z m \hbar \quad (9.5.8)$$

since the eigenvalue of  $\hat{L}_z$  is  $m\hbar$  and the expectation value of  $\hat{L}_z$  can be precisely  $m\hbar$ . Equation (9.5.8) can also be written as

$$E_u = m\mu_B B_z \quad (9.5.9)$$

where  $\mu_B$  is a Bohr magneton defined to be

$$\mu_B = \frac{e\hbar}{2m_0} \quad (9.5.10)$$

The Hamiltonian of quantum mechanics is inspired by the classical Hamiltonian mechanics. The Hamiltonian represents the total energy of the system. If an applied external magnetic

field contribution to a small change in the total energy, we can add a perturbation Hamiltonian in accordance to (9.5.8) and let

$$\hat{H}_p = \frac{eB_z}{2m_0} \hat{L}_z \quad (9.5.11)$$

Originally, for angular momentum with number  $l$ , the  $z$  component of the angular momentum is  $m\hbar$  with  $(2l + 1)$  values for  $m$  that are degenerate. The applied B field will split this degeneracy into  $2l + 1$  distinct energy values. This is known as the Zeeman effect.

Equation (9.5.11) can be written as

$$\hat{H}_p = \frac{e}{2m_0} \hat{\mathbf{L}} \cdot \mathbf{B} \quad (9.5.12)$$

When the spin angular momentum is added, the perturbing Hamiltonian becomes

$$\hat{H}_p = \frac{e}{2m_0} (\hat{\mathbf{L}} + g\hat{\mathbf{S}}) \cdot \mathbf{B} \quad (9.5.13)$$

where  $g$  is the gyromagnetic factor found to be approximately 2 experimentally.

In general, the Hamiltonian for an electron in the presence of an electromagnetic field is

$$\hat{H}_p = \frac{1}{2m_0} (\hat{\mathbf{p}} - e\mathbf{A})^2 + V + \frac{e}{2m_0} (\hat{\mathbf{L}} + g\hat{\mathbf{S}}) \cdot \mathbf{B} \quad (9.5.14)$$

The above Hamiltonian is again motivated by the Hamiltonian for classical mechanics for an electron in the presence of an electromagnetic field, which is

$$\hat{H}_p = \frac{1}{2m_0} [\mathbf{p} - e\mathbf{A}(\mathbf{r})]^2 + V(\mathbf{r}) \quad (9.5.15)$$

One can show, using the classical equation of motion for Hamiltonian mechanics, that Lorentz force law emerges from the above. Equation (9.5.14) is motivated by (9.5.15) where the momentum  $\mathbf{p}$  in (9.5.15) is elevated to be an operator  $\hat{\mathbf{p}}$  in (9.5.14). The vector potential  $\mathbf{A}(\mathbf{r})$ , however remains as a classical variable. As a result, (9.5.14) is a semi-classical Hamiltonian.

# Chapter 10

## Identical Particles

### 10.1 Introduction

The physics of identical particles is another triumph in quantum mechanics where the phenomenon is not observed in the classical world. When particles are identical or indistinguishable, their wavefunctions have to assume a certain form to satisfy a certain symmetry when the positions of the two particles are swapped. In quantum mechanics, it is the magnitude square of a wavefunction that has physical meaning. If the particles are indistinguishable, the magnitude square of a wavefunction does not change when the two particles are swapped.

In general, we can write the general wavefunction for two particles as

$$\psi_{tp}(\mathbf{r}_1, \mathbf{r}_2) = c_{12}\psi_a(\mathbf{r}_1)\psi_b(\mathbf{r}_2) + c_{21}\psi_a(\mathbf{r}_2)\psi_b(\mathbf{r}_1) \quad (10.1.1)$$

where  $\psi_\eta(\mathbf{r}_i)$  is a one-particle wavefunction for a particle in state  $\eta$  described by coordinate  $\mathbf{r}_i$ .

But when  $\mathbf{r}_1 \rightleftharpoons \mathbf{r}_2$ , the assumption is that the magnitude square of the wavefunction remains unchanged. Namely,

$$|\psi_{tp}(\mathbf{r}_1, \mathbf{r}_2)|^2 = |\psi_{tp}(\mathbf{r}_2, \mathbf{r}_1)|^2 \quad (10.1.2)$$

which means that

$$\psi_{tp}(\mathbf{r}_1, \mathbf{r}_2) = \gamma\psi_{tp}(\mathbf{r}_2, \mathbf{r}_1) \quad (10.1.3)$$

where  $\gamma = e^{i\theta}$ . Using (10.1.1) in (10.1.3), we have

$$c_{12}\psi_a(\mathbf{r}_1)\psi_b(\mathbf{r}_2) + c_{21}\psi_a(\mathbf{r}_2)\psi_b(\mathbf{r}_1) = \gamma [c_{12}\psi_a(\mathbf{r}_2)\psi_b(\mathbf{r}_1) + c_{21}\psi_a(\mathbf{r}_1)\psi_b(\mathbf{r}_2)] \quad (10.1.4)$$

From the above, we get

$$c_{21} = \gamma c_{12}, \quad c_{12} = \gamma c_{21} \quad (10.1.5)$$

Consequently,

$$c_{12} = \gamma c_{21} = \gamma^2 c_{12} \quad (10.1.6)$$

or  $\gamma^2 = 1$ ,  $\gamma = \pm 1$ ,  $c_{21} = \pm c_{12}$ , and

$$\psi_{tp}(\mathbf{r}_1, \mathbf{r}_2) = c_{12} [\psi_a(\mathbf{r}_1)\psi_b(\mathbf{r}_2) \pm \psi_a(\mathbf{r}_2)\psi_b(\mathbf{r}_1)] \quad (10.1.7)$$

When the “+” sign is chosen in (10.1.7), it corresponds to the wavefunction of two identical bosons. When the “−” sign is chosen, it corresponds to the wavefunction of two identical fermions.

## 10.2 Pauli Exclusion Principle

For fermions, the two-particle wavefunction is given by

$$\psi_{tp}(\mathbf{r}_1, \mathbf{r}_2) = c [\psi_a(\mathbf{r}_1)\psi_b(\mathbf{r}_2) - \psi_a(\mathbf{r}_2)\psi_b(\mathbf{r}_1)] \quad (10.2.1)$$

when  $a = b$ , the above vanishes, meaning that two fermions cannot be in the same state. The above vanishes when  $\mathbf{r}_1 = \mathbf{r}_2$ , meaning that two fermions cannot be in the same position simultaneously. This known as the Pauli exclusion principle. No two identical electrons can be in the same state nor the same position simultaneously. However, two electrons with different spins, up spin and down spin, are considered different, and they can be in the same state, like in the same orbital in an atom. Also, the sign of the total wavefunction,  $\psi_{tp}(\mathbf{r}_1, \mathbf{r}_2)$ , changes sign when the positions of the two particles are swapped.

## 10.3 Exchange Energy

When we have two electrons, each will see the Coulomb potential of the other electron. Hence, the Hamiltonian is given by

$$\hat{H} = -\frac{\hbar^2}{2m_e} (\nabla_1^2 + \nabla_2^2) + \frac{e^2}{4\pi\epsilon_0 |\mathbf{r}_1 - \mathbf{r}_2|} \quad (10.3.1)$$

The two-particle fermion wavefunction can be written as

$$|\psi_{tp}\rangle = \frac{1}{\sqrt{2}} [|1, a\rangle |2, b\rangle - |2, a\rangle |1, b\rangle] \quad (10.3.2)$$

where  $\frac{1}{\sqrt{2}}$  ensures that  $|\psi_{tp}\rangle$  is also normalized. In the above,  $|i, \eta\rangle$  is the Hilbert space representation of the wavefunction  $\psi_\eta(\mathbf{r}_i)$ . The above does not solve the Schrödinger equation corresponding to the above Hamiltonian, but it can be used to estimate the eigenvalue or the energy of the Hamiltonian in the spirit of the variational method described previously, using the Rayleigh quotient.

To find the corresponding expectation value of the energy, we evaluate

$$\langle E \rangle = \langle \psi_{tp} | \hat{H} | \psi_{tp} \rangle \quad (10.3.3)$$



Since the function is normalized, the above is also the Rayleigh quotient for the energy. When expanded to four terms, gives

$$\begin{aligned} \langle E \rangle = \frac{1}{2} & \left[ \langle 1, a | \langle 2, b | \hat{H} | 1, a \rangle | 2, b \rangle + \langle 2, a | \langle 1, b | \hat{H} | 2, a \rangle | 1, b \rangle \right. \\ & \left. - \langle 1, a | \langle 2, b | \hat{H} | 2, a \rangle | 1, b \rangle - \langle 2, a | \langle 1, b | \hat{H} | 1, a \rangle | 2, b \rangle \right] \end{aligned} \quad (10.3.4)$$

The first two terms above are equal to each other. Specifically,

$$\begin{aligned} & \langle 1, a | \langle 2, b | \hat{H} | 1, a \rangle | 2, b \rangle \\ &= \langle 1, a | \langle 2, b | \frac{-\hbar^2}{2m_0} \nabla_1^2 | 1, a \rangle | 2, b \rangle + \langle 1, a | \langle 2, b | \frac{-\hbar^2}{2m_0} \nabla_2^2 | 1, a \rangle | 2, b \rangle \\ & \quad + \langle 1, a | \langle 2, b | \frac{e^2}{4\pi\epsilon_0 |\mathbf{r}_1 - \mathbf{r}_2|} | 1, a \rangle | 2, b \rangle \\ &= \langle 1, a | \frac{-\hbar^2}{2m_0} \nabla_1^2 | 1, a \rangle + \langle 2, b | \frac{-\hbar^2}{2m_0} \nabla_2^2 | 2, b \rangle \\ & \quad + \langle 1, a | \langle 2, b | \frac{e^2}{4\pi\epsilon_0 |\mathbf{r}_1 - \mathbf{r}_2|} | 1, a \rangle | 2, b \rangle \\ &= E_{KE,a} + E_{KE,b} + E_{PE,ab} \end{aligned} \quad (10.3.5)$$

$E_{KE,a}$  is the kinetic energy of particle 1 in state  $a$ , while  $E_{KE,b}$  is the kinetic energy of particle 2 in state  $b$ . The  $E_{PE,ab}$  term can be derived explicitly to yield

$$E_{PE,ab} = e^2 \int d\mathbf{r} d\mathbf{r}' \frac{|\psi_a(\mathbf{r})|^2 |\psi_b(\mathbf{r}')|^2}{4\pi\epsilon_0 |\mathbf{r} - \mathbf{r}'|} \quad (10.3.6)$$

This is exactly the same as the potential energy from the second term on the right-hand side of (10.3.4) if we were to write it explicitly. Consequently, we have

$$E_{KE,a} + E_{KE,b} + E_{PE,ab} = \frac{1}{2} \left[ \langle 1, a | \langle 2, b | \hat{H} | 1, a \rangle | 2, b \rangle + \langle 2, a | \langle 1, b | \hat{H} | 2, a \rangle | 1, b \rangle \right] \quad (10.3.7)$$

However, the third and fourth terms on the right-hand side of (10.3.4) have no classical analog. They are there because of the presence of the second term in (10.1.7) be it for bosons or fermions. The second term is needed because the particles are indistinguishable or identical if we exchange their positions.

Due to the Hermiticity of the Hamiltonian operator, the fourth term on the right-hand side of (10.3.4) is the complex conjugate of the third term. Consequently, these two terms can be written as

$$E_{EX,ab} = -\frac{1}{2} \left[ \langle 1, a | \langle 2, b | \hat{H} | 2, a \rangle | 1, b \rangle + \left( \langle 1, a | \langle 2, b | \hat{H} | 2, a \rangle | 1, b \rangle \right)^* \right] \quad (10.3.8)$$

$$= -\Re \left[ \int d\mathbf{r} d\mathbf{r}' \psi_a^*(\mathbf{r}) \psi_b^*(\mathbf{r}') \hat{H} \psi_a(\mathbf{r}') \psi_b(\mathbf{r}) \right] \quad (10.3.9)$$

The exchange energy is proportional to the overlap between the wavefunctions  $\psi_a(\mathbf{r})$  and  $\psi_b(\mathbf{r})$ .

## 10.4 Extension to More Than Two Particles

In the following, we will discuss methods to construct basis functions for different kinds of particles. These basis functions does not solve the Schrödinger equation yet, but they have to satisfy certain symmetry conditions depending on the kind of particles they represent.

### 1. Non-identical Particle Case:

Let us assume that we have  $N$  particles, and  $M$  modes to fit this  $N$  particles. We can construct a state for non-identical particles that looks like

$$|\psi_{\text{diff}}\rangle = |1, a\rangle |2, b\rangle |3, c\rangle \dots |N, n\rangle \quad (10.4.1)$$

In terms of basis function, we may express the above as

$$|\psi_{ab\dots n}\rangle = |1, a\rangle |2, b\rangle |3, c\rangle \dots |N, n\rangle \quad (10.4.2)$$

or

$$\psi_{ab\dots n}(\mathbf{r}_1, \mathbf{r}_2, \dots, \mathbf{r}_N) = \psi_a(\mathbf{r}_1)\psi_b(\mathbf{r}_2)\dots\psi_n(\mathbf{r}_N) \quad (10.4.3)$$

We can fit the  $N$  particles in  $n$  modes, and these  $n$  modes can be repeating or non-repeating. For non-repeating case, it is necessary for  $M \geq N$ .

However, the above wavefunction cannot be used for bosons and fermions, as we will get a new wavefunction when we swap the positions of two particles. But bosons and fermions are indistinguishable particles. We will consider them separately.

### 2. Boson Case:

For the  $N$  boson particle case, we can write the legitimate wavefunction, which can be used as a basis function, as

$$|\psi_{\text{identical-bosons}}\rangle \propto \sum_{\hat{P}} \hat{P} |1, a\rangle |2, b\rangle |3, c\rangle \dots |N, n\rangle \quad (10.4.4)$$

where  $\hat{P}$  is a permutation operator, and the above summation is over all possible permutations of the coordinate  $\mathbf{r}_i$  over the one-particle eigenstates  $a, b, c, \dots, n$ . The above wavefunction remains unchange when we permute the positions of two particles, because for every  $|1, a\rangle \dots |i, l\rangle \dots |j, p\rangle \dots |N, n\rangle$ , there is a  $|1, a\rangle \dots |j, l\rangle \dots |i, p\rangle \dots |N, n\rangle$  in the above summation. Hence, swapping of  $i$  and  $j$  will not change the sign of the above wavefunction. The above can also be written as a basis function as

$$|\psi_{ab\dots n}\rangle \propto \sum_{\hat{P}} \hat{P} |1, a\rangle |2, b\rangle |3, c\rangle \dots |N, n\rangle \quad (10.4.5)$$

### 3. Fermion Case:

For the  $N$  fermion case, we can write the wavefunction, which can be used as a basis function, as

$$|\psi_{\text{identical-fermion}}\rangle = \frac{1}{\sqrt{N!}} \sum_{\hat{P}=1}^{N!} \pm \hat{P} |1, a\rangle |2, b\rangle |3, c\rangle \dots |N, n\rangle \quad (10.4.6)$$

where the “+” sign is chosen for even permutation while the “−” sign is chosen for odd permutation. A permutation involves a unique pairwise exchange of two particles. The permutation is even or odd depending on the number of pairwise exchanges that have taken place.

Therefore, given a term  $|1, a\rangle \cdots |i, l\rangle \cdots |j, p\rangle \cdots |N, n\rangle$ , there always exists another term:  $-|1, a\rangle \cdots |j, l\rangle \cdots |i, p\rangle \cdots |N, n\rangle$  in the above summation since they differ by one permutation. If  $i = j$ , the two terms cancel each other implying that they cannot be in the same position. Likewise all the terms in the sum cancel each other since every term that contains  $i$  and  $j$  can be paired up with every other terms in the sum. Moreover, If  $l = p$ , all terms in the summation above cancel as well implying that they cannot be in the same mode or state. Therefore, the above is a legitimate basis function that represents the fermions as it obeys Pauli’s exclusion principle. Also, there is a sign change when the position of two particles are swapped.

Another way of writing the fermion case is

$$|\psi_{\text{identical-fermion}}\rangle = \frac{1}{\sqrt{N!}} \begin{vmatrix} |1, a\rangle & |2, a\rangle & \cdots & |N, a\rangle \\ |1, b\rangle & |2, b\rangle & \cdots & |N, b\rangle \\ \vdots & \vdots & \ddots & \vdots \\ |1, n\rangle & |2, n\rangle & \cdots & |N, n\rangle \end{vmatrix} \quad (10.4.7)$$

The above is known as the Slater determinant. When two particles are in the same position, two columns are the same, and the above determinant is zero. When two states are the same, two rows are the same, and the determinant of the above is again zero, implying that these two cases are not allowed. Also when two columns are swapped, the sign of the determinant changes, because it corresponds to two particles exchanging positions. In terms of basis function, we can express the above as

$$\begin{aligned} |\psi_{ab\dots n}\rangle &= \frac{1}{\sqrt{N!}} \sum_{\hat{P}=1}^{N!} \pm \hat{P} |1, a\rangle |2, b\rangle |3, c\rangle \cdots |N, n\rangle \\ &= \frac{1}{\sqrt{N!}} \begin{vmatrix} |1, a\rangle & |2, a\rangle & \cdots & |N, a\rangle \\ |1, b\rangle & |2, b\rangle & \cdots & |N, b\rangle \\ \vdots & \vdots & \ddots & \vdots \\ |1, n\rangle & |2, n\rangle & \cdots & |N, n\rangle \end{vmatrix} \end{aligned} \quad (10.4.8)$$

## 10.5 Counting the Number of Basis states

Usually, the one particle eigenfunction has infinitely many possible states. In the finite basis method, we may want to choose a subset of this eigenstates, say  $M$  states. At this point, if we have  $N$  particles and  $M$  states to put this  $N$  particles in, it may be prudent to count how many possible basis states there are. Such counting scheme will also be used in the appendix to derive the thermal distribution functions.

### 1. Non-identical Particles:

The first particle has  $M$  state to choose from in (10.4.2), the second particle also has  $M$  states to choose from, and eventually, there are  $M^N$  states possible for (10.4.2), since repetition of the states is allowed.

$$N_{basis} = M^N \quad (10.5.1)$$

## 2. Identical Bosons:

In this case, the number of particles that can be put into one state is not limited. Since we need to permute the state (10.4.2) to arrive at new states to be linear superposed in the manner of (10.4.5), the ordering of the particles within the states is not important.

We will first derive the number of permutations when ordering is important. One can imagine the  $M$  states to be equivalent to  $M$  bins. The first particle has  $M$  possible bins to fit into. However, once the particle is put into one of the bins, since ordering is important, repetition is allowed, there are two ways to put in the second particle in the same bin, before or after the first particle. In other words, the first particle splits the bin into two compartments, allowing two ways to put two particles in the same bin. Together with the other  $M - 1$  bins, there are  $M + 1$  ways to insert the second particle into the bins. In other words, adding a new particle always adds a new compartment for the next particle to fit in. By the same token, the third particle has  $M + 2$  ways to fit into the bins etc. Finally the number of possible ways is

$$M(M + 1)(M + 2)\dots(M + N - 1) \quad (10.5.2)$$

One can verify that the above gives the correct answer of  $N!$  if only one bin is available for  $N$  particles. Since ordering is unimportant, we have

$$N_{basis} = \frac{(M + N - 1)!}{N!(M - 1)!} \quad (10.5.3)$$

## 3. Identical Fermions:

For fermions, the first particle has  $M$  states to choose from. The second particle has  $M - 1$  states to choose from, since repetition is not allowed. Also, ordering is not important since all permutations are used in the summation (10.4.8). Consequently,

$$N_{basis} = \frac{M!}{N!(M - N)!} \quad (10.5.4)$$

# 10.6 Examples

## 1. Non-identical Particles:

Say if we have two electrons with different spins, the distinct states are

$$|1, a\rangle|2, a\rangle, |1, b\rangle|2, b\rangle, |1, a\rangle|2, b\rangle, |1, b\rangle|2, a\rangle \quad (10.6.1)$$

The above is in agreement with  $M = 2$ ,  $N = 2$ ,  $M^N = 2^2 = 4$ .

### 2. Identical Bosons:

Consider the  ${}^4\text{He}$  (helium four) atoms which are bosons. Then the possible boson states are

$$|1, a\rangle|2, a\rangle, |1, b\rangle|2, b\rangle, \frac{1}{\sqrt{2}}(|1, a\rangle|2, b\rangle + |1, b\rangle|2, a\rangle) \quad (10.6.2)$$

Again this is in agreement with  $M = 2$ ,  $N = 2$ ,  $\frac{(M+N-1)!}{N!(M-1)!} = \frac{3!}{2!1!} = 3$ .

### 3. Identical Fermions:

If we have two identical fermions, the only state is

$$\frac{1}{\sqrt{2}}(|1, a\rangle|2, b\rangle - |1, b\rangle|2, a\rangle) \quad (10.6.3)$$

Again this is in agreement with  $M = 2$ ,  $N = 2$ ,  $\frac{M!}{N!(M-N)!} = \frac{2!}{2!0!} = 1$ .

In the above simple examples, we see that for the boson case, there are more states with particles in the same state compared to the other cases. In the boson case, there are two out of three states where the particles are in the same state (see (10.6.2)). In the non-identical particle case, there are two out of four such states. In the fermion case, there is none. This is indicative of the fact that bosons like to cluster together.

## 10.7 Thermal Distribution Functions

There are some important thermal distribution functions in understanding semiconductor physics. These are the Maxwell-Boltzmann distribution, the Fermi-Dirac distribution, and the Bose-Einstein distribution. These distribution functions are derived via statistical mechanics, and the results will just quoted here. The first-principles derivations are given in Appendix B.

### 1. Maxwell-Boltzmann Distribution:

For this distribution, the number of particles  $N(E)$  in a given energy state is given by

$$N(E) = e^{-\frac{(E-\mu)}{k_B T}} \quad (10.7.1)$$

where  $\mu$  is the chemical potential,  $k_B$  is the Boltzmann constant, and  $T$  is temperature in Kelvin.

### 2. Fermi-Dirac Distribution:

This is given by

$$N(E) = \frac{1}{1 + e^{\frac{E-\mu}{k_B T}}} \quad (10.7.2)$$

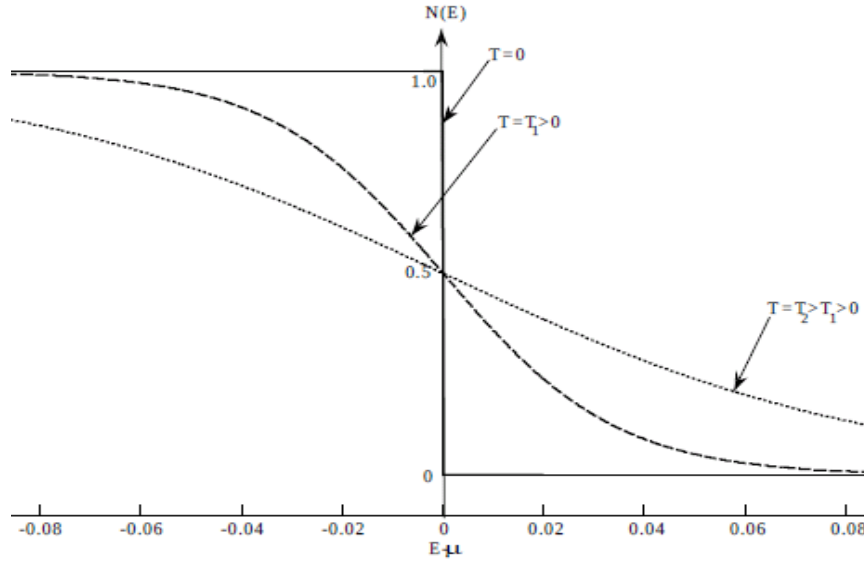


Figure 10.1: Fermi-Dirac Distribution for changing temperature  $T$ .

This is the most important distribution function for semiconductors. The chemical potential  $\mu$  is loosely called the Fermi level. To be precise,  $\mu$  is the same as the Fermi level at absolute zero temperature. When  $T$  is zero, the distribution function looks like as shown in Figure 10.1. The electrons are frozen into the ground state for  $E < \mu$ . When  $T > 0$ , some of the electrons from the states where  $E < \mu$  are dislodged by thermal agitation into states where  $E > \mu$ . This phenomenon is more pronounced as  $T$  increases. This distribution also explains the physical character of semiconductors.

In semiconductors, the Fermi level is midway in between the valence band and the conduction band. When  $T = 0$ , all the electrons are frozen in the valence band, and the semiconductor cannot conduct electricity, as there are no free electrons in the conduction band. When  $T > 0$ , some electrons in the valence band are dislodged into the conduction band. This gives rise to electrons in the conduction band, and holes in the valence band. They contribute to the flow of electric current and the semiconductor starts to conduct. The conductivity of the semi-conductor material increases with increasing temperature.

### 3. Bose-Einstein Distribution:

The Bose-Einstein distribution is for bosons, and it is given by

$$N(E) = \frac{1}{e^{(E-\mu)/(k_B T)} - 1} \quad (10.7.3)$$

This distribution has a divergence when  $E = \mu$ . It also reflects the fact that bosons like to cluster together. When the temperature is low, they condense to around the chemical potential  $\mu$ .

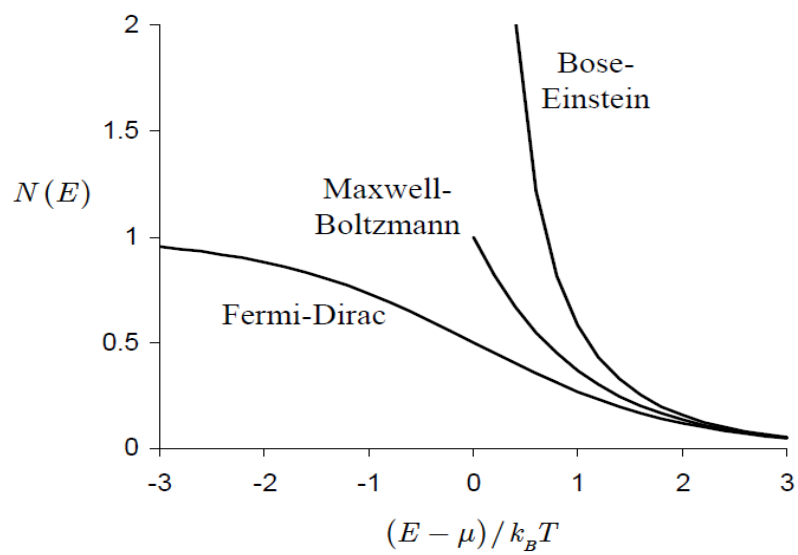


Figure 10.2: Comparison of different thermal distribution functions. For high energy states, they are similar to each other (from DAB Miller).





# Chapter 11

## Density Matrix

### 11.1 Pure and Mixed States

Given a quantum state described by

$$|\psi\rangle = a_1|\psi_1\rangle + a_2|\psi_2\rangle \quad (11.1.1)$$

which is a linear superposition of two states, the interpretation of quantum mechanics is that the particle is in a linear superposition of the two states, and one does not know what state the particle is in until the measurement is done. The measurement collapses the particle into one of the two quantum states. For instance, the two states could be the spin states of an atom. The famous Stern-Gerlach experiment separates the spin states into the up state and the down state. Such a state indicated in (11.1.1) is known as a pure quantum state. The phase relationship between  $a_1$  and  $a_2$  is maintained precisely. When the quantum state is in such a state, we say that the two states are coherent.

However, there are states where the particle has already collapsed into state 1 or state 2 even before the measurement. Or these two states are entirely uncorrelated or incoherent. The particle is either in state 1 or state 2. The probability of finding the particle in state 1 is  $|a_1|^2$  while that for state 2 is  $|a_2|^2$ . The phase relationship between  $a_1$  and  $a_2$  is completely random. States for which there is no coherence between  $a_1$  and  $a_2$  are known as mixed states. If we take an ensemble average of the measurement outcomes of the mixed states, they satisfy the aforementioned probability property.

Nevertheless, it is difficult for a particle to be in a pure quantum state if it involves a linear superposition of many quantum states unless the quantum system is completely isolated. Coupling to the thermal bath can destroy the coherence between these states. Hence, there are states for which partial coherence still exists, and the linear superposition of quantum states de-coheres with respect to each other, before the linear superposition of states becomes entirely a mixed state. Coherence lifetime is used to measure when a linear superposition of quantum states, which starts out being a coherent sum, becomes an incoherent superposition.

## 11.2 Density Operator

An elegant way to represent a quantum state that is either in the pure state, mixed state, or partial coherent state, is via the density operator

$$\hat{\rho} = |\psi\rangle\langle\psi| \quad (11.2.1)$$

We can easily show that the expectation of the operator  $\hat{A}$  in this state is

$$\langle\hat{A}\rangle = \langle\psi|\hat{A}|\psi\rangle = \text{tr}\left(\hat{A}\hat{\rho}\right) \quad (11.2.2)$$

This has been shown in Section 5.2.3. Therefore, knowing the density operator is equivalent to knowing the quantum state of a system which is denoted by state  $|\psi\rangle$ .

Say if we start with (11.1.1) for a quantum state, using it in (11.2.1) will yield

$$\hat{\rho} = |a_1|^2|\psi_1\rangle\langle\psi_1| + |a_2|^2|\psi_2\rangle\langle\psi_2| + a_1a_2^*|\psi_1\rangle\langle\psi_2| + a_2a_1^*|\psi_2\rangle\langle\psi_1| \quad (11.2.3)$$

Notice that only relative phases between  $a_1$  and  $a_2$  are needed to form the above. Absolute phase has no meaning in quantum mechanics.

For a mixed state, we assume that the off diagonal terms involving  $a_1^*a_2$  and  $a_2^*a_1$  will either time average or ensemble average to zero since they are incoherent. Hence, we suppose that

$$\begin{aligned} \hat{\rho} &= |a_1|^2|\psi_1\rangle\langle\psi_1| + |a_2|^2|\psi_2\rangle\langle\psi_2| \\ &= p_1|\psi_1\rangle\langle\psi_1| + p_2|\psi_2\rangle\langle\psi_2| \end{aligned} \quad (11.2.4)$$

The form (11.2.4) can be thought of as having been ensemble averaged or time averaged.<sup>1</sup> For a general mixed state we can further write the density operator as

$$\hat{\rho} = \sum_j p_j |\psi_j\rangle\langle\psi_j| \quad (11.2.5)$$

where  $|\psi_j\rangle$  are pure quantum states but not necessarily stationary states, and  $p_j$  is the probability of finding the quantum system in state  $|\psi_j\rangle$ .

When we find the expectation value of an operator  $\hat{A}$  with a mixed state, it can be thought of as having been ensemble averaged. For example,

$$\begin{aligned} \langle\hat{A}\rangle &= \text{tr}(\hat{A}\hat{\rho}) = \sum_j p_j \text{tr}(\hat{A}|\psi_j\rangle\langle\psi_j|) \\ &= \sum_j p_j \langle\psi_j|\hat{A}|\psi_j\rangle \end{aligned} \quad (11.2.6)$$

The trace of the matrix representation of an operator is independent of the basis states chosen to represent the operator as explained in Chapter 5. It will be interesting to see how the density operator changes under the change of basis state. To this end, we let

$$|\psi_j\rangle = \sum_n c_n^{(j)}(t)|\phi_n\rangle \quad (11.2.7)$$

---

<sup>1</sup>Statistical processes whose time average and ensemble average are equivalent to each other are known as ergodic processes.

Using this in (11.2.5), we have

$$\begin{aligned}\hat{\rho} &= \sum_j p_j \sum_{n,m} c_n^j |\phi_n\rangle \langle \phi_m| (c_m^j)^* \\ &= \sum_{n,m} |\phi_n\rangle \langle \phi_m| \sum_j p_j c_n^j (c_m^j)^*\end{aligned}\quad (11.2.8)$$

We can find the  $uv$  element of the matrix representation of  $\hat{\rho}$  as

$$\rho_{uv} = \langle \phi_u | \hat{\rho} | \phi_v \rangle = \sum_j p_j c_u^{(j)} [c_v^{(j)}]^* = \overline{c_u c_v^*} \quad (11.2.9)$$

where the overbar stands for ensemble average. From the above, it is quite clear that

$$\rho_{uv} = \rho_{vu}^* \quad (11.2.10)$$

or that the density matrix is Hermitian. The Hermiticity property of the density operator is obvious from (11.1.1) and (11.2.5).

Since a trace of an operator is independent of the basis state in which it is represented, the trace of the density operator is defined as

$$\text{tr}(\hat{\rho}) = \sum_u \langle \phi_u | \hat{\rho} | \phi_u \rangle \quad (11.2.11)$$

where  $\phi_u$  constitutes an orthonormal basis.

For a pure quantum state,

$$\hat{\rho}^2 = |\psi\rangle \langle \psi | \psi\rangle \langle \psi| = |\psi\rangle \langle \psi| = \hat{\rho} \quad (11.2.12)$$

Therefore,

$$\text{tr}(\hat{\rho}^2) = \text{tr}(\hat{\rho}) = 1 \quad (11.2.13)$$

The trace of the above is 1 is obviated by setting  $\hat{A}$  to 1 in Equation (11.2.6).

For a mixed state

$$\hat{\rho}^2 = \sum_i \sum_j p_i p_j |\psi_i\rangle \langle \psi_i | \psi_j\rangle \langle \psi_j| \quad (11.2.14)$$

Taking the trace of the above, we have

$$\begin{aligned}\text{tr}(\hat{\rho}^2) &= \sum_n \langle \phi_n | \hat{\rho}^2 | \phi_n \rangle = \sum_n \sum_i \sum_j p_i p_j \langle \phi_n | \psi_i\rangle \langle \psi_i | \psi_j\rangle \langle \psi_j | \phi_n \rangle \\ &= \sum_i \sum_j p_i p_j |\langle \psi_i | \psi_j \rangle|^2 \leq \sum_i \sum_j p_i p_j = \left( \sum_i p_i \right)^2 = 1\end{aligned}\quad (11.2.15)$$

In the above, we did not assume that  $|\psi_j\rangle$  are orthogonal, but they are normalized. The inequality follows from that  $|\langle \psi_i | \psi_j \rangle|^2 \leq 1$ . Hence, the trace of the density operator has the property that

$$\text{tr}(\hat{\rho}^2) < 1, \quad \text{mixed state} \quad (11.2.16)$$

$$\text{tr}(\hat{\rho}^2) = 1, \quad \text{pure state} \quad (11.2.17)$$

### 11.3 Time Evolution of the Matrix Element of an Operator

The matrix representation of an operator is given by

$$A_{mn}(t) = \langle \psi_m(t) | \hat{A} | \psi_n(t) \rangle \quad (11.3.1)$$

Taking the time derivative of the above, we have

$$\partial_t A_{mn}(t) = \langle \partial_t \psi_m(t) | \hat{A} | \psi_n(t) \rangle + \langle \psi_m(t) | \hat{A} | \partial_t \psi_n(t) \rangle \quad (11.3.2)$$

assuming that  $\hat{A}$  is time independent, using the fact that

$$i\hbar \partial_t | \psi_m(t) \rangle = i\hbar | \partial_t \psi_m(t) \rangle = \hat{H} | \psi_m(t) \rangle \quad (11.3.3)$$

we have

$$\begin{aligned} \partial_t A_{mn}(t) &= \frac{i}{\hbar} \left\langle \left[ \hat{H} \psi_m(t) \right] | \hat{A} | \psi_n(t) \right\rangle - \frac{i}{\hbar} \langle \psi_m(t) | \hat{A} | \hat{H} \psi_n(t) \rangle \\ &= \frac{i}{\hbar} \langle \psi_m(t) | \hat{H} \hat{A} | \psi_n(t) \rangle - \frac{i}{\hbar} \langle \psi_m(t) | \hat{A} \hat{H} | \psi_n(t) \rangle \\ &= \frac{i}{\hbar} \langle \psi_m(t) | \hat{H} \hat{A} - \hat{A} \hat{H} | \psi_n(t) \rangle \end{aligned} \quad (11.3.4)$$

or

$$i\hbar \partial_t A_{mn}(t) = (\hat{A} \hat{H} - \hat{H} \hat{A})_{mn} = \left[ \hat{A}, \hat{H} \right]_{mn} \quad (11.3.5)$$

The above applies to expectation value of  $\hat{A}$  as well, as shown in Chapter 5, where

$$\langle \hat{A} \rangle = \langle \psi(t) | \hat{A} | \psi(t) \rangle \quad (11.3.6)$$

and

$$i\hbar \partial_t \langle \hat{A} \rangle = \langle \hat{A} \hat{H} - \hat{H} \hat{A} \rangle = \left\langle \left[ \hat{A}, \hat{H} \right] \right\rangle \quad (11.3.7)$$

In other words, if  $\hat{A}$  commutes with  $\hat{H}$ , its expectation value is time independent. In the above derivation, we have assumed that  $\hat{A}$  is independent of time.

The time evolution of the density operator can be derived, similar to the time evolution of the expectation value of an operator. However, the density operator is in general time dependent, whereas operators considered previously are time independent; and hence, the derivation is slightly different. Given

$$\hat{\rho} = |\psi\rangle\langle\psi| \quad (11.3.8)$$

then

$$\partial_t \hat{\rho} = |\partial_t \psi\rangle\langle\psi| + |\psi\rangle\langle\partial_t \psi| \quad (11.3.9)$$

Using (11.3.3), we have

$$\partial_t \hat{\rho} = -\frac{i}{\hbar} \hat{H} |\psi\rangle \langle \psi| + \frac{i}{\hbar} |\psi\rangle \langle \psi| \hat{H} = \frac{i}{\hbar} [\hat{\rho} \hat{H} - \hat{H} \hat{\rho}] = \frac{i}{\hbar} [\hat{\rho}, \hat{H}] \quad (11.3.10)$$

There is a sign difference between (11.3.7) and (11.3.10). The above holds if  $\hat{\rho}$  is in a mixed state as well

$$\hat{\rho} = \sum_j p_j |\psi_j\rangle \langle \psi_j| \quad (11.3.11)$$

This follows from the linearity of the above equations and that  $p_j$  is time-independent.

Equation (11.3.10) can be easily generalized to matrix elements when the basis functions used for seeking the matrix representation are from time-independent basis. Then,

$$\partial_t \rho_{mn} = \frac{i}{\hbar} [\hat{\rho} \hat{H} - \hat{H} \hat{\rho}]_{mn} = \frac{i}{\hbar} [\hat{\rho}, \hat{H}]_{mn} \quad (11.3.12)$$

where  $A_{mn} = \langle \psi_m | \hat{A} | \psi_n \rangle$ , and  $\psi_i$  is from a time-independent basis. By inserting identity operator defined using the same basis set, we can convert the above into wholly a matrix form:

$$\partial_t \bar{\rho} = \frac{i}{\hbar} [\bar{\rho}, \bar{\mathbf{H}}] \quad (11.3.13)$$

where  $\bar{\rho}$  and  $\bar{\mathbf{H}}$  are the matrix representations of  $\hat{\rho}$  and  $\hat{H}$ , respectively. We often do not distinguish the matrix representations  $\bar{\rho}$  and  $\bar{\mathbf{H}}$ , and operator representations  $\hat{\rho}$  and  $\hat{H}$ , as mathematically, they are the same. The above equation replaces the Schrödinger equation as the equation of motion for the quantum system.

## 11.4 Interaction of Light with Two-Level Atomic System

Two-level systems are encountered in nuclear magnetic resonance when the spin states of a particle interact with a magnetic field. Many more complex systems, for simplicity, can be approximated by a two-level system. When an electric field (optical field) interacts with an atom and causes an atomic transition between two energy levels, while the other energy levels are far away, the system can be approximated by a simpler two-level system. Also, approximate two-level systems are often used to represent the value “0” and “1” in quantum computing as we shall see later.

One assumes that before the system is perturbed, the atom is described by a simple two-level system with two stationary states bearing energies  $E_1$  and  $E_2$ . In this case, the perturbing Hamiltonian is assumed to be

$$\hat{H}_p = ezE = -E\hat{\mu} \quad (11.4.1)$$

The matrix representation of the dipole moment in terms of two stationary states before perturbation is given by<sup>2</sup>

$$\mu_{mn} = -e\langle\psi_m|\hat{z}|\psi_n\rangle \quad (11.4.2)$$

where  $\hat{z}$  is the position operator, and

$$\left(\hat{H}_p\right)_{mn} = H_{p,mn} = -E\mu_{mn} \quad (11.4.3)$$

In the above  $\mu_{11} = \mu_{22} = 0$ , so are  $H_{p11} = H_{p22} = 0$ , because  $\psi_m$  and  $\psi_n$  have different parities. In other words, if one of them has odd parity, the other one has even parity. Then in the coordinate representation, the function  $|\psi_i|^2 z$  will have odd parity, and its integral evaluates to zero. But for the off diagonal elements, the integrands will have even parity, and their integrals evaluate to nonzero values. The matrix elements above are formed with stationary states, and  $\mu_{12}$  can be made pure real so that  $\mu_{12} = \mu_{21} = \mu_d$ . Hence, the perturbing Hamiltonian becomes

$$\hat{H}_p = \begin{bmatrix} 0 & -E\mu_d \\ -E\mu_d & 0 \end{bmatrix} \quad (11.4.4)$$

The above can be simplified to, and approximated by a two-level system with the use of simple matrix mechanics. This is also very much in the spirit of a finite-basis method, where a small number of basis can form a good approximation of the quantum system. The unperturbed Hamiltonian,  $\hat{H}_0$  is assumed to have eigenvalues  $E_1$  and  $E_2$  with respect to these two stationary states. The total Hamiltonian is then

$$\hat{H} = \hat{H}_0 + \hat{H}_p = \begin{bmatrix} E_1 & -E\mu_d \\ -E\mu_d & E_2 \end{bmatrix} \quad (11.4.5)$$

The equation of motion according to Schrödinger equation is

$$i\hbar\partial_t|\psi\rangle = \hat{H}|\psi\rangle \quad (11.4.6)$$

In the above, the perturbing Hamiltonian introduces the off diagonal terms. They are responsible for causing the transition of eigenstates between state 1 and state 2. If the off diagonal terms are absent, the stationary states will remain in their respective states without transition. In the above finite basis approximation to the Hamiltonian, the basis used consists of the two eigenstates of the original unperturbed problem. This is a subspace approximation method, whereby an infinite dimensional Hilbert space has been replaced by a two-dimensional subspace.

The density matrix which is the matrix representation of the density operator is

$$\hat{\rho} = \begin{bmatrix} \rho_{11} & \rho_{12} \\ \rho_{21} & \rho_{22} \end{bmatrix} \quad (11.4.7)$$

---

<sup>2</sup>This is an approximation since after perturbation, the quantum state is not necessarily describable by just two stationary states. We saw that in the time-dependent perturbation theory, all eigenstates of the system was necessary to approximate the perturbed system.

With the notation of (11.2.3), the density matrix is

$$\hat{\rho} = \begin{bmatrix} |a_1(t)|^2 & a_1(t)a_2^*(t) \\ a_2(t)a_1^*(t) & |a_2(t)|^2 \end{bmatrix} \quad (11.4.8)$$

In (11.2.3), we assume here that the two states are stationary states (they need not be), and the time variations of the states have been put into  $a_1(t)$  and  $a_2(t)$ . Hence, the off diagonal terms can be rapidly varying functions of time due to the interference between the two stationary states.

The expectation value of the dipole moment is

$$\langle \hat{\mu} \rangle = \text{tr}(\hat{\rho}\hat{\mu}) \quad (11.4.9)$$

where

$$\hat{\rho}\hat{\mu} = \begin{bmatrix} \rho_{11} & \rho_{12} \\ \rho_{21} & \rho_{22} \end{bmatrix} \begin{bmatrix} 0 & \mu_d \\ \mu_d & 0 \end{bmatrix} = \begin{bmatrix} \rho_{12}\mu_d & \rho_{11}\mu_d \\ \rho_{22}\mu_d & \rho_{21}\mu_d \end{bmatrix} \quad (11.4.10)$$

Hence

$$\langle \hat{\mu} \rangle = \mu_d(\rho_{12} + \rho_{21}) = 2\mu_d\Re(\rho_{12}) \quad (11.4.11)$$

because  $\rho_{12} = \rho_{21}^*$ .

Next, we derive the equation of motion for the density matrix. Since

$$\begin{aligned} \frac{d\hat{\rho}}{dt} &= \frac{i}{\hbar} [\hat{\rho}\hat{H} - \hat{H}\hat{\rho}] \\ &= \frac{i}{\hbar} \left( \begin{bmatrix} \rho_{11} & \rho_{12} \\ \rho_{21} & \rho_{22} \end{bmatrix} \begin{bmatrix} E_1 & -E\mu_d \\ -E\mu_d & E_2 \end{bmatrix} - \begin{bmatrix} E_1 & -E\mu_d \\ -E\mu_d & E_2 \end{bmatrix} \begin{bmatrix} \rho_{11} & \rho_{12} \\ \rho_{21} & \rho_{22} \end{bmatrix} \right) \\ &= \frac{i}{\hbar} \begin{bmatrix} -E\mu_d(\rho_{12} - \rho_{21}) & -E\mu_d(\rho_{11} - \rho_{22}) + (E_2 - E_1)\rho_{12} \\ -E\mu_d(\rho_{22} - \rho_{11}) + (E_1 - E_2)\rho_{21} & -E\mu_d(\rho_{21} - \rho_{12}) \end{bmatrix} \end{aligned} \quad (11.4.12)$$

From the above

$$\frac{d\rho_{21}}{dt} = \frac{i}{\hbar} [(\rho_{11} - \rho_{22})E\mu_d - (E_2 - E_1)\rho_{21}] = -i\omega_{21}\rho_{21} + i\frac{\mu_d E}{\hbar}(\rho_{11} - \rho_{22}) \quad (11.4.13)$$

where  $\hbar\omega_{21} = E_2 - E_1$ . We need not write down the equation for  $\rho_{12}$  since  $\rho_{12} = \rho_{21}^*$ .

From (11.4.12)

$$\frac{d\rho_{11}}{dt} = -\frac{i}{\hbar}E\mu_d(\rho_{12} - \rho_{21}), \quad \frac{d\rho_{22}}{dt} = -\frac{i}{\hbar}E\mu_d(\rho_{21} - \rho_{12}) \quad (11.4.14)$$

From them, we get

$$\frac{d}{dt}(\rho_{11} - \rho_{22}) = -\frac{i}{\hbar}E\mu_d(\rho_{12} - \rho_{21} - \rho_{21} + \rho_{12}) = \frac{2i}{\hbar}E\mu_d(\rho_{21} - \rho_{21}^*) = \frac{-4}{\hbar}E\mu_d\Im(\rho_{21}) \quad (11.4.15)$$

since  $\rho_{12} = \rho_{21}^*$ . The above equations, (11.4.13) and (11.4.15), form two coupled equations from which the unknowns,  $\rho_{11} - \rho_{22}$  and  $\rho_{12}$  can be solved for.

At his point, we are going to add phenomenological terms to the above equations to make them agree with experimental observations. The above two level system, so far, has been isolated from a thermal bath. It more aptly describes the two level system of a single atom in isolation. In actuality, there will be a collection of atoms that will be similarly excited by the electric field. These atoms will be coupled to other thermal sources such as by collision with other atoms or with the wall of the container. The density matrix can be used to represent the ensemble averaged state of the collection of these atoms. Here, for a collection of atoms,  $\rho_{11} - \rho_{22}$  is the fractional population difference in electrons of the atoms between state 1 and state 2. But when coupled to a thermal bath, there will be some electrons excited to level 2 even in the absence of an external electric field. Hence, we modify (11.4.15) to account for the coupling to a thermal bath to become

$$\frac{d}{dt}\Delta\rho = \frac{2i}{\hbar}E\mu_d(\rho_{21} - \rho_{21}^*) - \frac{\Delta\rho - \Delta\rho_o}{T_1} \quad (11.4.16)$$

where  $\Delta\rho = \rho_{11} - \rho_{22}$ ,  $\Delta\rho_o = (\rho_{11} - \rho_{22})_o$  is the quiescent steady state value of  $\Delta\rho$  in the absence of external driving field. The transient value  $\Delta\rho$  is assumed to relax to the quiescent value  $\Delta\rho_o$  in time  $T_1$ .

Also,  $\rho_{21} = C e^{-i\omega_{21}t}$  in the absence of external driving field. But if the two level system is coupled to an external heat bath the phase coherence between the two energy states will be lost and average to zero. We can describe this by adding a term to (11.4.13) to account for dephasing, or

$$\frac{d\rho_{21}}{dt} = -i\omega_{21}\rho_{21} + i\frac{\mu_d}{\hbar}E\Delta\rho - \frac{\rho_{21}}{T_2} \quad (11.4.17)$$

$T_2$  is the dephasing time which is usually shorter than  $T_1$ .

At this point, it is prudent to discuss the role of the exciting field, which is of the form

$$E(t) = E_o \cos \omega t = \frac{E_o}{2} [e^{i\omega t} + e^{-i\omega t}] \quad (11.4.18)$$

When  $\Delta\rho$  is slowly varying in (11.4.17), the driving term, which is the second term on the right-hand side in (11.4.17) has two rotating signals  $e^{i\omega t}$  and  $e^{-i\omega t}$ . The response,  $\rho_{21}$  will have two rotating signals as well. Hence,  $\rho_{21} = C e^{-i\omega t} + c.c.$ <sup>3</sup> when driven by a sinusoidal source (11.4.18). Therefore, we define

$$\beta_{21}(t) = \rho_{21}(t) e^{i\omega t} = C + C^* e^{2i\omega t} \quad (11.4.19)$$

where  $C$  is slowly varying so that  $\beta_{21}(t)$  is slowly varying when we ignore the rapid term. Using this in (11.4.16), we have

$$\frac{d}{dt}\Delta\rho = \frac{i}{\hbar}E_o\mu_d(\beta_{21} - \beta_{12}^*) - \frac{\Delta\rho - \Delta\rho_o}{T_1} \quad (11.4.20)$$

---

<sup>3</sup>where c.c. stands for complex conjugate.



In the above, we have kept only the slowest varying term of the first term on the right-hand side. The rapid term is assumed to only result in a small response in  $\Delta\rho$ . This is known as the rotating wave approximation.

From (11.4.18), we deduce that

$$\frac{d\rho_{21}(t)}{dt} = \left[ \frac{d}{dt} \beta_{21}(t) \right] e^{-i\omega t} - i\omega \beta_{21}(t) e^{-i\omega t} \quad (11.4.21)$$

Consequently, (11.4.13) becomes

$$\frac{d}{dt} \beta_{21} = i \Delta\omega \beta_{21} + i \frac{\mu_d E_o}{2\hbar} \Delta\rho - \frac{\beta_{21}}{T_2} \quad (11.4.22)$$

where  $\Delta\omega = \omega - \omega_{21}$ , and we have kept only the slowly varying terms.

Previously, we have found that

$$\langle \hat{\mu} \rangle = \mu_d (\rho_{12} + \rho_{21}) = \mu_d (\beta_{12} e^{i\omega t} + \beta_{21} e^{-i\omega t}) \quad (11.4.23)$$

after using (11.4.19). Therefore

$$\langle \hat{\mu} \rangle = 2\mu_d [\Re(\beta_{21}) \cos \omega t + \Im(\beta_{21}) \sin \omega t] \quad (11.4.24)$$

since  $\beta_{21} = \beta_{12}^*$ .

In steady state,  $d\Delta\rho/dt = 0$  in (11.4.20). Furthermore,  $\Delta\rho$  in (11.4.22) is a constant in steady state, and hence,  $\beta_{21}$ , must tend to a constant in steady state with  $d\beta_{21}/dt = 0$ . Defining  $\Omega = \mu_d E_o / (2\hbar)$ , we have from (11.4.20) and (11.4.22) that

$$0 = i\Delta\omega \beta_{21} + i\Omega \Delta\rho - \beta_{21}/T_2 \quad (11.4.25)$$

$$0 = -\Delta\Omega \Im(\beta_{21}) - (\Delta\rho - \Delta\rho_o)/T_1 \quad (11.4.26)$$

Taking the real and imaginary parts of (11.4.25), we arrive at

$$0 = -\Delta\omega \Im(\beta_{21}) - \Re(\beta_{21})/T_2 \quad (11.4.27)$$

$$0 = \Delta\omega \Re(\beta_{21}) + \Omega \Delta\rho - \Im(\beta_{21})/T_2 \quad (11.4.28)$$

Equations (11.4.26)-(11.4.28) constitute three equations with three unknowns  $\Delta\rho$ ,  $\Im(\beta_{21})$ ,  $\Re(\beta_{21})$ , from which they can be solved for. Therefore

$$\Delta\rho = \Delta\rho_o \frac{1 + \Delta\omega^2 T_2^2}{1 + \Delta\omega^2 T_2^2 + 4\Omega^2 T_1 T_2} \quad (11.4.29)$$

$$\Re(\beta_{21}) = -\frac{\Delta\omega \Omega T_2^2 \Delta\rho_o}{1 + \Delta\omega^2 T_2^2 + 4\Omega^2 T_1 T_2} \quad (11.4.30)$$

$$\Im(\beta_{21}) = -\frac{\Omega T_2 \Delta\rho_o}{1 + \Delta\omega^2 T_2^2 + 4\Omega^2 T_1 T_2} \quad (11.4.31)$$

If there are  $N$  atoms, then we define

$$\Delta N = N\Delta\rho, \quad \Delta N_o = N\Delta\rho_o \quad (11.4.32)$$

and

$$\Delta N = \Delta N_o \frac{1 + \Delta\omega^2 T_2^2}{1 + \Delta\omega^2 T_2^2 + 4\Omega^2 T_1 T_2} \quad (11.4.33)$$

In electromagnetics, we have

$$P = \epsilon_o \chi E \quad (11.4.34)$$

in phasor representation. In time domain

$$\begin{aligned} P &= \Re e [\epsilon_o (\chi' + i\chi'') E_o e^{-i\omega t}] \\ &= \epsilon_o (\chi' \cos \omega t + \chi'' \sin \omega t) E_o \end{aligned} \quad (11.4.35)$$

But

$$P = N \langle \hat{\mu} \rangle \quad (11.4.36)$$

From (11.4.24), we have

$$P = 2N\mu_d [\Re e (\beta_{21}) \cos \omega t + \Im m (\beta_{21}) \sin \omega t] \quad (11.4.37)$$

Comparing (11.4.35) and (11.4.37), we have

$$\begin{aligned} \chi' &= \frac{2N\mu_d}{\epsilon_o E_o} \Re e (\beta_{21}) \\ &= -\frac{\mu_d^2 T_2 \Delta N_o}{\epsilon_o \hbar} \frac{\Delta\omega T_2}{1 + (\Delta\omega)^2 T_2^2 + 4\Omega^2 T_1 T_2} \end{aligned} \quad (11.4.38)$$

$$\begin{aligned} \chi'' &= \frac{2N\mu_d}{\epsilon_o E} \Im m (\beta_{21}) \\ &= \frac{\mu_d^2 T_2 \Delta N_o}{\epsilon_o \hbar} \frac{1}{1 + (\Delta\omega)^2 T_2^2 + 4\Omega^2 T_1 T_2} \end{aligned} \quad (11.4.39)$$

When the applied electric field is small, we can ignore the  $4\Omega^2 T_1 T_2$  term to get

$$\chi'(\omega) = \frac{\mu_d^2 T_2 \Delta N_o}{\epsilon_o \hbar} \frac{(\omega_{21} - \omega) T_2}{1 + (\omega - \omega_{21})^2 T_2^2} \quad (11.4.40)$$

$$\chi''(\omega) = \frac{\mu_d^2 T_2 \Delta N_o}{\epsilon_o \hbar} \frac{1}{1 + (\omega - \omega_{21})^2 T_2^2} \quad (11.4.41)$$

Note that in this analysis, the perturbing field need not be small. The approximation we have made is the two-state approximation. This is a decent approximation if the other eigenstates are far away from these two states in terms of energy levels. The other approximation is the rotating wave approximation which is good if the exciting frequency of the electric field is high and is much faster than the relaxation times  $T_1$  and  $T_2$ , and that it is close to  $\omega_{21}$ .

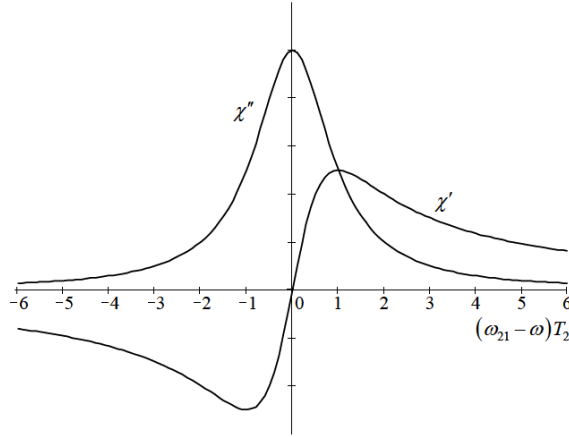


Figure 11.1: Plots of  $\chi'$  and  $\chi''$  according to (11.4.40) and (11.4.41) (from DAB Miller).

Also, this analysis is not a small perturbation analysis since the exciting field need not be weak. In fact,  $\Omega$  above, which is proportional to the exciting electric field, can be large. When it is large, we note from (11.4.29) and (11.4.33) that the fractional population difference approaches zero, or that  $\rho_{11} = \rho_{22}$ . In this case, the absorption transition is equal to the stimulated emission. From (11.4.39), when  $\Delta\omega = 0$ , or  $\omega = \omega_{21}$  we can express the denominator as

$$1 + 4\Omega^2 T_1 T_2 = 1 + \frac{I}{I_s} \quad (11.4.42)$$

where  $I$  is the field intensity, and  $I_s$  is the saturation intensity, since  $I$  is proportional to the square of the electric field. The saturation intensity is defined to be the intensity at which the absorption peak of  $\chi''$  will drop to half of its value compared to the weak excitation case where  $I$  is very small.

Equations (11.4.20) and (11.4.22) are also known as the optical Bloch equations. These equations were first used to analyze another two-state system, the nuclear magnetic resonance system. With a magnetic field pointing in the  $z$  direction, when the nuclear spins are pointing upward and aligned with the magnetic field, they are in the lower energy state. However, when they are pointing anti-parallel to the magnetic field, they are in a higher energy state. RF (radio frequency) field can be used to flip the spin states, and similar equations as above can be derived to describe the spin dynamics.



## Chapter 12

# Quantization of Classical Fields

### 12.1 Introduction

The quantum harmonic oscillator is extremely important for the quantization of classical fields, be it electromagnetic, acoustic, or elastic in nature. Classical fields can be thought of as due to a set of coupled classical harmonic oscillators. In the case of electromagnetic field, these classical harmonic oscillators are dipoles in space polarized by the electric field. In the case of acoustic or elastic field, they are due to coupled atomic or molecular vibration in a medium. Many experimental evidence suggest that there are photon particles associated with electromagnetic field, and phonon particles associated with vibrations of atoms or molecules in crystalline solids.

When the classical harmonic oscillators are replaced by quantum harmonic oscillators, the classical field becomes quantized as shall be seen. The energy that can propagate and associated with the field is quantized, as in the case of the quantum harmonic oscillator. In classical electromagnetics when we calculate energy, e.g. as stored energies

$$W_e = \frac{1}{2}\epsilon |\mathbf{E}|^2, \quad W_h = \frac{1}{2}\mu |\mathbf{H}|^2 \quad (12.1.1)$$

we think of  $W_e$  and  $W_h$  as a continuum, capable of assuming all values from zero to a large value. Also, the power flow by a plane wave

$$S = \frac{1}{2\eta_0} |\mathbf{E}|^2 \quad (12.1.2)$$

assumes continuous values in classical electromagnetics. As shall be shown, this cannot be the case if electromagnetic field consists is due to a set of coupled quantum harmonic oscillators. The energy carried by a plane wave is also quantized corresponding to packets of energy.

The quantized nature of electromagnetic radiation has been observed historically. This was first suggested for and used to deriving the Planck radiation law. Later, it was used by Einstein to explain experimental findings in the photoelectric effect. From the experimental data, it was deduced that the packet of energy associated with electromagnetic field is

$$E_{\text{photon}} = \hbar\omega \quad (12.1.3)$$

This quantized energy is very small when  $\omega$  is small but sizeable in the optical regime. The grainy nature of electromagnetic field is unimportant at radio frequency (RF), and microwave frequencies, but is important at optical frequencies. We shall show that this is the case when electromagnetic field is quantized when the classical harmonic oscillators are replaced by quantum harmonic oscillators.

To introduce the quantized nature of electromagnetic field, one way is to regard electromagnetic field to be in a resonant cavity or a box with periodic boundary condition. The electromagnetic field in such a system oscillates harmonically. We can replace the classical harmonic oscillators associated with such a system with the quantum harmonic oscillator; and hence, quantize the amplitude and the associated energy of the electromagnetic field.

A prevailing view is to think of a photon as a particle, but is associated with a wave field that satisfies Maxwell's equations. This is analogous to an electron being a particle and is associated with a wavefunction as its "halo". As photons are bosons, we can have a collection of photons described by the same wave field, with different amplitudes, since this wave field is now endowed with more energy. As shall be seen, the wave field associated with photons is quite different from the wave field associated with electrons, or other boson particles obeying Schrödinger equation.

We will first revisit the quantum harmonic oscillator due to its importance here. Next we study the wave on a linear atomic chain, first as a classical wave, and then as a set of coupled quantum harmonic oscillators. Then we study the quantization of electromagnetic wave field along similar spirit.

## 12.2 The Quantum Harmonic Oscillator Revisited

The governing equation for the quantum harmonic oscillator is

$$\hat{H}\psi = \left( -\frac{\hbar^2}{2m} \frac{d^2}{dz^2} + \frac{1}{2}m\omega^2 z^2 \right) \psi = E\psi \quad (12.2.1)$$

Letting  $\xi = \sqrt{\frac{m\omega}{\hbar}}z$ , we have

$$\frac{1}{2} \left( -\frac{d^2}{d\xi^2} + \xi^2 \right) \psi = \frac{E}{\hbar\omega} \psi \quad (12.2.2)$$

The above looks almost like  $A^2 - B^2$  with the exception that operators are involved instead of scalars. We can look at

$$\frac{1}{\sqrt{2}} \left( -\frac{d}{d\xi} + \xi \right) \frac{1}{\sqrt{2}} \left( \frac{d}{d\xi} + \xi \right) = \frac{1}{2} \left( -\frac{d^2}{d\xi^2} + \xi^2 \right) - \frac{1}{2} \left( \frac{d}{d\xi} \xi - \xi \frac{d}{d\xi} \right) \quad (12.2.3)$$

One can show easily that

$$\left( \frac{d}{d\xi} \xi - \xi \frac{d}{d\xi} \right) = \hat{I} \quad (12.2.4)$$

Hence,

$$\frac{1}{2} \left( -\frac{d^2}{d\xi^2} + \xi^2 \right) = \frac{1}{\sqrt{2}} \left( -\frac{d}{d\xi} + \xi \right) \frac{1}{\sqrt{2}} \left( \frac{d}{d\xi} + \xi \right) + \frac{1}{2} \quad (12.2.5)$$

We define

$$\hat{a}^\dagger = \frac{1}{\sqrt{2}} \left( -\frac{d}{d\xi} + \xi \right), \quad \text{creation or raising operator} \quad (12.2.6)$$

$$\hat{a} = \frac{1}{\sqrt{2}} \left( \frac{d}{d\xi} + \xi \right), \quad \text{annihilation or lowering operator} \quad (12.2.7)$$

It can easily be shown by integration by parts that  $\hat{a}^\dagger$  is in fact the conjugate transpose operator of  $\hat{a}$ .

Consequently, from (12.2.1) and (12.2.2), we have

$$\left( \hat{a}^\dagger \hat{a} + \frac{1}{2} \right) \psi = \frac{E}{\hbar\omega} \psi \quad (12.2.8)$$

and

$$\hat{H} = \hbar\omega \left( \hat{a}^\dagger \hat{a} + \frac{1}{2} \right) \quad (12.2.9)$$

Since

$$E_n = \left( n + \frac{1}{2} \right) \hbar\omega \quad (12.2.10)$$

we conclude that

$$\hat{a}^\dagger \hat{a} |\psi_n\rangle = n |\psi_n\rangle \quad (12.2.11)$$

In other words, if  $|\psi_n\rangle$  is an eigenstate of (12.2.1) or (12.2.8), the above must be true. We can define a number operator

$$\hat{n} = \hat{a}^\dagger \hat{a} \quad (12.2.12)$$

such that

$$\hat{n} |\psi_n\rangle = n |\psi_n\rangle \quad (12.2.13)$$

One can also show by direct substitution using (12.2.6) and (12.2.7) that

$$[\hat{a}, \hat{a}^\dagger] = \hat{a}\hat{a}^\dagger - \hat{a}^\dagger\hat{a} = \hat{I} \quad (12.2.14)$$

One can easily show, using the commutation relation, that

$$\hat{a}^\dagger \hat{a} \hat{a} |\psi_n\rangle = (\hat{a}\hat{a}^\dagger - 1) \hat{a} |\psi_n\rangle = (n-1) \hat{a} |\psi_n\rangle \quad (12.2.15)$$

Therefore

$$\hat{a} |\psi_n\rangle = A_n |\psi_{n-1}\rangle \quad (12.2.16)$$

Similarly, using the commutation relation,

$$\hat{a}^\dagger \hat{a} \hat{a}^\dagger |\psi_n\rangle = \hat{a}^\dagger (1 + \hat{a}^\dagger \hat{a}) |\psi_n\rangle = (n+1) \hat{a}^\dagger |\psi_n\rangle \quad (12.2.17)$$

Therefore

$$\hat{a}^\dagger |\psi_n\rangle = B_{n+1} |\psi_{n+1}\rangle \quad (12.2.18)$$

Testing (12.2.16) with  $\langle\psi_{n-1}|$ , and (12.2.18) with  $\langle\psi_{n+1}|$ , we have

$$\langle\psi_{n-1}|\hat{a}|\psi_n\rangle = A_n, \quad \langle\psi_{n+1}|\hat{a}^\dagger|\psi_n\rangle = B_{n+1} \quad (12.2.19)$$

By comparing the two equations in (12.2.19), and that

$$\langle\psi_{n-1}|\hat{a}|\psi_n\rangle = B_n^* \quad (12.2.20)$$

it is clear that

$$A_n = B_n^* \quad (12.2.21)$$

Since

$$\hat{a}^\dagger\hat{a}|\psi_n\rangle = A_n\hat{a}^\dagger|\psi_{n-1}\rangle = A_nB_n|\psi_n\rangle \quad (12.2.22)$$

$$= |A_n|^2|\psi_n\rangle = n|\psi_n\rangle \quad (12.2.23)$$

We conclude that

$$A_n = \sqrt{n} \quad (12.2.24)$$

For simplicity, we assume that  $A_n$  and  $B_n$  are both real. Therefore

$$\hat{a}|\psi_n\rangle = \sqrt{n}|\psi_{n-1}\rangle \quad (12.2.25)$$

$$\hat{a}^\dagger|\psi_n\rangle = \sqrt{n+1}|\psi_{n+1}\rangle \quad (12.2.26)$$

### 12.2.1 Eigenfunction by the Ladder Approach

With the raising operators, we can construct the requisite eigenfunctions if we only know the ground state eigenfunction. The ground state eigenfunction satisfies

$$\hat{a}|\psi_0\rangle = 0 \quad (12.2.27)$$

Written explicitly in coordinate space representation we have

$$\frac{1}{\sqrt{2}} \left( \frac{d}{d\xi} + \xi \right) \psi_0(\xi) = 0 \quad (12.2.28)$$

We can solve the above ordinary differential equation to get

$$\psi_0(\xi) = \frac{1}{\pi^{1/4}} e^{-\xi^2/2} \quad (12.2.29)$$

(We can also verify that  $\psi_0(\xi)$  satisfies (12.2.28) by direct substitution.) Consequently, one can show that

$$(\hat{a}^\dagger)^n |\psi_0\rangle = \sqrt{n!} |\psi_n\rangle \quad (12.2.30)$$

or that the  $n$ -th eigenstate is

$$|\psi_n\rangle = \frac{1}{\sqrt{n!}} (\hat{a}^\dagger)^n |\psi_0\rangle \quad (12.2.31)$$



## 12.3 Quantization of Waves on a Linear Atomic Chain—Phonons

When a particle is in simple harmonic motion, quantum mechanics allows us to quantize its motion using Schrödinger equation for the harmonic oscillator. When we have an atomic chain, there will be coupling forces between the atoms. When one of the atoms is set into simple harmonic oscillation, the whole chain will be in simple harmonic oscillation due to coupling between them. Since the energy of one single particle is quantized in the quantum realm, we expect that harmonic oscillation of the whole chain to be quantized in energy as well. This is a much easier notion to accept if we accept that the energy of one particle is quantized.

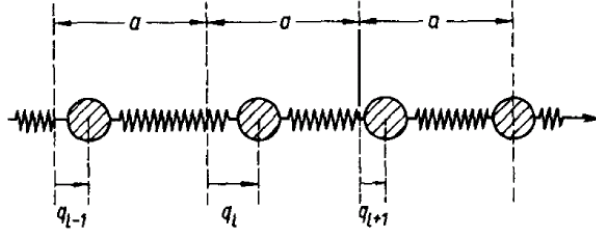


Figure 12.1: The linear atomic chain (from Haken).

For atoms on a linear atomic chain, assuming only nearest neighbor forces between the particles, the equation of motion is

$$\begin{aligned} m_0 \ddot{q}_\ell(t) &= f [q_{\ell+1}(t) - q_\ell(t)] - f [q_\ell(t) - q_{\ell-1}(t)] \\ &= f [q_{\ell+1}(t) - 2q_\ell(t) + q_{\ell-1}(t)] \end{aligned} \quad (12.3.1)$$

We impose the periodic boundary condition so that

$$q_\ell = q_{\ell+N} \quad (12.3.2)$$

Assuming that

$$q_\ell(t) = \frac{1}{\sqrt{N}} e^{ik_\nu \ell a} B_k(t) \quad (12.3.3)$$

where

$$k_\nu = \frac{2\nu\pi}{Na} \quad (12.3.4)$$

where  $\nu$  is an integer ranging from  $-\infty$  to  $\infty$ . Since  $k_\nu$  is countable, we will drop the subscript  $\nu$  and use  $k$  as an index. Using (12.3.3) in (12.3.1), we have

$$\ddot{B}_k(t) = g (e^{ika} + e^{-ika} - 2) B_k(t) \quad (12.3.5)$$

with

$$g = \frac{f}{m_0} \quad (12.3.6)$$

We let

$$B_k(t) = B_k(0)e^{\pm i\omega_k t} \quad (12.3.7)$$

Using (12.3.7) in (12.3.5) yields

$$-\omega_k^2 = 2g [\cos(ka) - 1] \quad (12.3.8)$$

or

$$\omega_k = \pm 2\sqrt{g} \sin\left(\frac{ka}{2}\right) \quad (12.3.9)$$

The kinetic energy of the system is

$$T = \sum_{\ell=1}^N \frac{m_0}{2} \dot{q}_\ell^2 \quad (12.3.10)$$

while the potential energy

$$V = \frac{1}{2}f \sum_{\ell=1}^N (q_\ell - q_{\ell+1})^2 \quad (12.3.11)$$

such that the force on the  $\ell$ -th atom is

$$F_\ell = -\frac{\partial}{\partial q_\ell} V \quad (12.3.12)$$

The classical Hamiltonian of the system is then

$$H = \sum_{\ell=1}^N \frac{p_\ell^2}{2m_0} + \frac{1}{2}f \sum_{\ell=1}^N (q_\ell - q_{\ell+1})^2 \quad (12.3.13)$$

The equation of motion (12.3.1) can be derived from the above using Hamiltonian equations of motion, yielding

$$\dot{q}_\ell(t) = \frac{1}{m_0} p_\ell(t), \quad \dot{p}_\ell(t) = f [q_{\ell+1}(t) - 2q_\ell(t) + q_{\ell-1}(t)]$$

$$q_\ell(t) = \frac{1}{\sqrt{N}} e^{ik\ell a} B_k(t) + \text{c.c.} \quad (12.3.14)$$

$$p_\ell(t) = \frac{1}{\sqrt{N}} e^{ik\ell a} A_k(t) + \text{c.c.} \quad (12.3.15)$$

so that

$$\dot{B}_k(t) = \frac{1}{m_0} A_k(t) \quad (12.3.16)$$

$$\dot{A}_k(t) = f [e^{ika} + e^{-ika} - 2] B_k(t) \quad (12.3.17)$$

In the above, “c.c.” stands for “complex conjugate”. Equation (12.3.17) can be written as

$$\dot{A}_k(t) = -\omega_k^2 m B_k(t) \quad (12.3.18)$$

where

$$\omega_k^2 = \frac{f}{m_0} [e^{ika} + e^{-ika} - 2] \quad (12.3.19)$$

Combining (12.3.16) and (12.3.18), we have

$$\ddot{A}_k(t) = -\omega_k^2 A_k(t) \quad (12.3.20)$$

$$\ddot{B}_k(t) = -\omega_k^2 B_k(t) \quad (12.3.21)$$

To have a travelling wave, we let

$$A_k(t) = A_k(0) e^{-i\omega_k t} \quad (12.3.22)$$

$$B_k(t) = B_k(0) e^{-i\omega_k t} \quad (12.3.23)$$

From the above, we see that

$$A_k(t) = -i\omega_k m_0 B_k(t) \quad (12.3.24)$$

We also find that

$$p_\ell^2(t) = \frac{2}{N} |A_k(t)|^2 + \left[ \frac{1}{N} e^{2ik\ell a} A_k^2(t) + \text{c.c.} \right] \quad (12.3.25)$$

where the second term is an oscillatory function of  $\ell$ . Also

$$\begin{aligned} (q_\ell - q_{\ell+1}) &= \frac{1}{\sqrt{N}} e^{ik\ell a} (1 - e^{ika}) B_k(t) + \text{c.c.} \\ &= -\frac{1}{\sqrt{N}} e^{ik\ell a} 2i \sin\left(\frac{ka}{2}\right) e^{ik\frac{\ell}{2} a} B_k(t) + \text{c.c.} \end{aligned} \quad (12.3.26)$$

Hence

$$(q_\ell - q_{\ell+1})^2 = \frac{8}{N} \sin^2\left(\frac{ka}{2}\right) |B_k(t)|^2 + \text{O.F.} \quad (12.3.27)$$

where “O.F.” stands for “oscillatory functions of  $\ell$ ”. The oscillatory terms sum to zero over  $\ell$  when we plug (12.3.25) and (12.3.27) into (12.3.13). Therefore,

$$\begin{aligned} H &= \frac{1}{m_0} |A_k(t)|^2 + 4f \sin^2\left(\frac{ka}{2}\right) |B_k(t)|^2 \\ &= \frac{1}{m_0} |A_k(t)|^2 + m_0 \omega_k^2 |B_k(t)|^2 \\ &= 2m_0 \omega_k^2 |B_k(t)|^2 \end{aligned} \quad (12.3.28)$$

after using (12.3.24).<sup>1</sup> If we let

$$4m_0\omega_k|B_k(t)|^2 = p_k^2(t) + q_k^2(t) \quad (12.3.29)$$

then

$$H = \frac{\omega_k}{2} [p_k^2(t) + q_k^2(t)] \quad (12.3.30)$$

which is identical to the Hamiltonian of a classical harmonic oscillator. The reason why the above simple form ensues is:

1. By imposing periodic boundary condition on the linear atomic chain, only countable, discrete modes with wavenumber  $k$  are allowed to propagate on the linear atomic chain;
2. Only a single mode is allowed to propagate on the linear atomic chain so as to make the solution and the corresponding Hamiltonian very simple.

The physical picture we shall have of the above is that because only one mode exists, all the atoms on the chain are oscillating at the same frequency and in unison except for a phase lag. So each of them is in a simple harmonic motion. Since we have learned how to convert a classical harmonic oscillator to a quantum harmonic oscillator, the same principles can be applied to turn the above into a chain of coupled quantum harmonic oscillators.

To convert the above into a quantum harmonic oscillator, we elevate the above amplitudes to operators, namely

$$\hat{H} = \frac{\omega_k}{2} [\hat{p}_k^2 + \hat{q}_k^2] \quad (12.3.31)$$

where

$$\hat{p}_k = -i\hbar \frac{d}{dq_k}, \quad \hat{q}_k^2 = q_k^2 \quad (12.3.32)$$

in coordinate space representation. We can further let

$$\xi_k = q_k/\sqrt{\hbar}, \quad \hat{\Pi}_k = \hat{p}_k/\sqrt{\hbar} = -i \frac{d}{d\xi_k} \quad (12.3.33)$$

so that

$$\hat{H} = \frac{\hbar\omega_k}{2} \left[ -\frac{d^2}{d\xi_k^2} + \xi_k^2 \right] = \frac{\hbar\omega_k}{2} (\hat{\Pi}_k^2 + \xi_k^2) \quad (12.3.34)$$

$$(12.3.35)$$

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<sup>1</sup>It is to be noted that because of the assumption of a right traveling wave in (12.3.22) and (12.3.23),  $|A_k(t)|^2$  and  $|B_k(t)|^2$  in the above are constants. What this means is that both the kinetic energy and the potential energy parts of the Hamiltonian are constants of motion. However, had we assumed that both right and left traveling waves exist in (12.3.22) and (12.3.23), this would not have been true: both  $|A_k(t)|^2$  and  $|B_k(t)|^2$  are time varying. However, it can be proved that their sum in the Hamiltonian is a constant.

The above is the quantum analogue of (12.3.30). Furthermore, we can show that

$$\begin{aligned}\hat{a}_k^\dagger &= \frac{1}{\sqrt{2}} \left( -i\hat{\Pi}_k + \xi_k \right) \\ \hat{a}_k &= \frac{1}{\sqrt{2}} \left( i\hat{\Pi}_k + \xi_k \right)\end{aligned}\quad (12.3.36)$$

Hence,

$$\hat{H} = \frac{\hbar\omega_k}{2} \left[ \hat{a}_k^\dagger \hat{a}_k + \hat{a}_k \hat{a}_k^\dagger \right] \quad (12.3.37)$$

We can rewrite (12.3.28) more suggestively as

$$H = m_0\omega_k^2 (B_k^*(t)B_k(t) + B_k(t)B_k^*(t)) \quad (12.3.38)$$

The quantum analogue of the above is

$$\hat{H} = m_0\omega_k^2 \left( \hat{B}_k^\dagger \hat{B}_k + \hat{B}_k \hat{B}_k^\dagger \right) \quad (12.3.39)$$

We can identify that

$$\begin{aligned}\hat{B}_k &= \sqrt{\frac{\hbar}{2m_0\omega_k}} \hat{a}_k \\ \hat{B}_k^\dagger &= \sqrt{\frac{\hbar}{2m_0\omega_k}} \hat{a}_k^\dagger\end{aligned}\quad (12.3.40)$$

Also, the above analogue is in the spirit of the Schrödinger picture where the time dependence is relegated to the state vector and the operators are time independent. An arbitrary phase factor can be added to the above but this result in a time shift for the field. With the above analogue, we can elevate the field to become operators, namely,

$$\hat{q}_\ell = \sqrt{\frac{\hbar}{2m_0\omega_k N}} e^{ik\ell a} \hat{a}_k + \text{c.c.} \quad (12.3.41)$$

$$\hat{p}_\ell = -i\sqrt{\frac{\hbar m_0\omega_k}{2N}} e^{ik\ell a} \hat{a}_k + \text{c.c.} \quad (12.3.42)$$

These field operators are time independent because they are in the Schrödinger picture. Notice that it is the Hamiltonian of the total system involving  $N$  atoms, or the total energy of the total system, that is quantized with with energy proportional to  $\hbar\omega_k$ . This quantized energy per atom is then proportional to  $\hbar\omega_k/N$ . This spread of the quantized energy per atom becomes smaller as  $N$  becomes larger, which becomes unphysical, unless we keep  $N$  finite. Hence, to be more realistic, one will have to consider a localized wave packet. We will return to this point later.

If we had assumed that the solutions are summation of different  $k$  modes, then (12.3.14) and (12.3.15) become

$$q_\ell(t) = \frac{1}{\sqrt{N}} \sum_k e^{ik\ell a} B_k(t) + \text{c.c.} \quad (12.3.43)$$

$$p_\ell(t) = \frac{1}{\sqrt{N}} \sum_k e^{ik\ell a} A_k(t) + \text{c.c.} \quad (12.3.44)$$

Consequently, (12.3.38), after using Parseval's theorem, becomes

$$H = m_0 \sum_k \omega_k^2 (B_k^*(t)B_k(t) + B_k(t)B_k^*(t)) \quad (12.3.45)$$

and (12.3.41) and (12.3.42) become

$$\hat{q}_\ell = \sum_k \sqrt{\frac{\hbar}{2m_0\omega_k N}} e^{ik\ell a} \hat{a}_k + \text{c.c.} \quad (12.3.46)$$

$$\hat{p}_\ell = \sum_k -i\sqrt{\frac{\hbar m_0\omega_k}{2N}} e^{ik\ell a} \hat{a}_k + \text{c.c.} \quad (12.3.47)$$

The above represents a multitude of propagating modes with different  $k$  values. They can come together to form a wave packet that more realistically represent a wave that propagates on an infinite linear atomic chain.

## 12.4 Schrödinger Picture versus Heisenberg Picture

At this juncture, it is prudent to discuss the difference between the Schrödinger picture versus the Heisenberg picture of quantum mechanics. In the Schrödinger picture, the time dependence is in the wavefunction or the state vector; whereas in the Heisenberg picture, the time dependence is in the operator that represents an observable.

The Schrödinger equation is given by

$$\hat{H}|\psi\rangle = i\hbar\partial_t|\psi\rangle \quad (12.4.1)$$

The formal solution to the above can be written as

$$|\psi\rangle = e^{-\frac{i}{\hbar}\hat{H}t}|\psi_0\rangle \quad (12.4.2)$$

where  $|\psi_0\rangle$  is the state vector at  $t = 0$  and is independent of time. The expectation value of an operator  $\hat{O}$  is

$$\langle\hat{O}\rangle = \langle\psi|\hat{O}|\psi\rangle = \langle\psi_0|e^{\frac{i}{\hbar}\hat{H}t}\hat{O}e^{-\frac{i}{\hbar}\hat{H}t}|\psi_0\rangle \quad (12.4.3)$$

Hence, in general, the relationship between an operator in the Schrödinger picture and in the Heisenberg picture is given by

$$\hat{O}_H = e^{\frac{i}{\hbar}\hat{H}t}\hat{O}_S e^{-\frac{i}{\hbar}\hat{H}t} \quad (12.4.4)$$

where the subscripts ‘‘H’’ and ‘‘S’’ mean the Heisenberg picture and the Schrödinger picture respectively. It can be shown easily that

$$\frac{d\hat{O}_H}{dt} = \frac{i}{\hbar} [\hat{H}, \hat{O}_H] \quad (12.4.5)$$

For the simple quantum harmonic oscillator,

$$\hat{H} = \hbar\omega \left( \hat{a}^\dagger \hat{a} + \frac{1}{2} \right) \quad (12.4.6)$$

It can be shown easily that

$$\left[ \hat{H}, \hat{a} \right] = -\hbar\omega \hat{a} \quad (12.4.7)$$

or that

$$\frac{d\hat{a}}{dt} = -i\omega \hat{a} \quad (12.4.8)$$

Consequently, solving give

$$\hat{a}(t) = \hat{a}(0)e^{-i\omega t} \quad (12.4.9)$$

Similarly, one can show that

$$\hat{a}^\dagger(t) = \hat{a}^\dagger(0)e^{i\omega t} \quad (12.4.10)$$

The above are the annihilation and creation operators in the Heisenberg picture. Using the above in (12.3.46) and (12.3.47), we have

$$\hat{q}_\ell(t) = \sum_k \sqrt{\frac{\hbar}{2m_0\omega_k N}} e^{ik\ell a - i\omega t} \hat{a}_k + \text{c.c.} \quad (12.4.11)$$

$$\hat{p}_\ell(t) = \sum_k -i\sqrt{\frac{\hbar m_0 \omega_k}{2N}} e^{ik\ell a - i\omega t} \hat{a}_k + \text{c.c.} \quad (12.4.12)$$

The above are the field operators in the Heisenberg picture. They are time dependent.

## 12.5 The Continuum Limit

In the continuum limit, we let

$$m_0 = \rho a \quad (12.5.1)$$

Then (12.3.1) becomes

$$\rho a \ddot{q}(x, t) = f a^2 \frac{\partial^2}{\partial x^2} q(x, t) \quad (12.5.2)$$

or

$$\rho \ddot{q}(t) = g \frac{\partial^2}{\partial x^2} q(x, t) \quad (12.5.3)$$

where  $g = fa$ . Assume a traveling wave

$$q(x, t) = \frac{1}{\sqrt{L}} e^{ikx} B_k(t) + \text{c.c.} \quad (12.5.4)$$

with periodic boundary condition; then

$$k_\nu = \frac{2\pi\nu}{L} \quad (12.5.5)$$

We define the momentum density to be

$$\pi(x) = \rho \dot{q}(x) \quad (12.5.6)$$

Then a Hamiltonian that will lead to (12.5.3) is

$$H = \frac{1}{2\rho} \int_0^L \pi^2(x) dx + \frac{g}{2} \int_0^L \left( \frac{\partial q(x)}{\partial x} \right)^2 dx \quad (12.5.7)$$

We let

$$q(x, t) = \frac{1}{\sqrt{L}} e^{ikx} B_k(t) + \text{c.c.} \quad (12.5.8)$$

$$\pi(x, t) = \frac{1}{\sqrt{L}} e^{ikx} A_k(t) + \text{c.c.} \quad (12.5.9)$$

Then

$$\dot{B}_k(t) = \frac{1}{\rho} A_k(t) \quad (12.5.10)$$

$$\dot{A}_k(t) = -gk^2 B_k(t) \quad (12.5.11)$$

From the above, we have

$$\ddot{B}_k(t) = -\frac{gk^2}{\rho} B_k(t) = -\omega_k^2 B_k(t) \quad (12.5.12)$$

where  $\omega_k^2 = gk^2/\rho$ , or

$$B_k(t) = B_k(0) e^{-i\omega_k t} \quad (12.5.13)$$

$$A_k(t) = -i\omega_k \rho B_k(t) \quad (12.5.14)$$

From (12.5.9), we obtain that

$$\pi^2(x, t) = \frac{2}{L} |A_k(t)|^2 + \text{O.T.} \quad (12.5.15)$$



$$\left(\frac{\partial q(x, t)}{\partial x}\right)^2 = \frac{2k^2}{L}|B_k(t)|^2 + \text{O.T.} \quad (12.5.16)$$

where O.T. stands for oscillatory term. When substituted into (12.5.7), the oscillatory terms integrate to zero, and

$$\begin{aligned} H &= \frac{1}{\rho}|A_k(t)|^2 + k^2 g|B_k(t)|^2 = 2\rho\omega_k^2|B_k(t)|^2 \\ &= \rho\omega_k^2 [B_k^*(t)B_k(t) + B_k(t)B_k^*(t)] \end{aligned} \quad (12.5.17)$$

Defining

$$4\rho\omega_k|B_k(t)|^2 = p_k^2(t) + q_k^2(t) \quad (12.5.18)$$

the above becomes

$$H = \frac{\omega_k}{2} [p_k^2(t) + q_k^2(t)] \quad (12.5.19)$$

which is the classical Hamiltonian of a simple harmonic oscillator. The above can be quantized to yield

$$\hat{H} = \frac{\hbar\omega_k}{2} (\hat{a}_k^\dagger \hat{a}_k + \hat{a}_k \hat{a}_k^\dagger) \quad (12.5.20)$$

and (12.5.17) can be elevated to quantized form giving

$$\hat{H} = \rho\omega_k^2 [\hat{B}_k^\dagger \hat{B}_k + \hat{B}_k \hat{B}_k^\dagger] \quad (12.5.21)$$

Hence, we conclude that

$$\hat{B}_k = \sqrt{\frac{\hbar}{2\rho\omega_k}} \hat{a}_k \quad (12.5.22)$$

up to an arbitrary phase. With this, we can elevate the fields to operators, namely,

$$\hat{q}(x, t) = \sqrt{\frac{\hbar}{2\rho\omega_k L}} e^{ikx - i\omega_k t} \hat{a}_k + \text{c.c.} \quad (12.5.23)$$

$$\hat{\pi}(x, t) = -i\sqrt{\frac{\hbar\omega_k\rho}{2L}} e^{ikx - i\omega_k t} \hat{a}_k + \text{c.c.} \quad (12.5.24)$$

The above are in Heisenberg picture where operators are functions of time. With multi-mode field, the above generalize to

$$\hat{q}(x, t) = \sum_k \sqrt{\frac{\hbar}{2\rho\omega_k L}} e^{ikx - i\omega_k t} \hat{a}_k + \text{c.c.} \quad (12.5.25)$$

$$\hat{\pi}(x, t) = -i \sum_k \sqrt{\frac{\hbar\omega_k\rho}{2L}} e^{ikx - i\omega_k t} \hat{a}_k + \text{c.c.} \quad (12.5.26)$$

## 12.6 Quantization of Electromagnetic Field

For electromagnetic field, the harmonic oscillators are dipoles in a medium that are polarized by an electric field. In general, the electric flux in a medium is given by

$$\mathbf{D} = \epsilon_0 \mathbf{E} + \mathbf{P} \quad (12.6.1)$$

The first term is the contribution to the electric flux  $\mathbf{D}$  due to vacuum, while the second term is the polarization density contributed from dipoles polarized in a material medium by an electric field. It is customary to write  $\mathbf{P} = \epsilon_0 \chi \mathbf{E}$  to indicate that the polarization density is proportional to the electric field.

The time variation of the electric flux,  $\partial_t \mathbf{D}$ , is instrumental in giving rise to displacement current, and hence, yielding a propagating wave. This electric flux exists even in vacuum. Hence, we can imagine that even vacuum is polarized by an electric field to produce dipole density.<sup>2</sup> In other words, the oscillating electric dipole will produce a magnetic field creating an inductive effect. Together, they can form a resonating system behaving like a harmonic oscillator. They resemble a tiny LC tank circuit.

As seen previously, the quantization of electromagnetic field is most easily described by the quantization of a modal field. A mode can be formed by a periodic boundary condition or a metallic perfect electric/magnetic conductor boundary condition. In this manner, the number of modes is countably infinite. In view of this, we can write

$$E_z = \alpha(t) D e^{ikx} + c.c. \quad (12.6.2)$$

By using periodic boundary condition, we get

$$k_\nu = \frac{2\nu\pi}{L} \quad (12.6.3)$$

making it countably infinite. (We will drop the subscript  $\nu$  subsequently.) The associated  $\mathbf{B}$  field is

$$B_y = \beta(t) \frac{D}{c} e^{ikx} + c.c. \quad (12.6.4)$$

where  $c$  is the velocity of light. The relation between  $\alpha(t)$  and  $\beta(t)$  can be derived from Maxwell's equations. From

$$\nabla \times \mathbf{E} = -\partial_t \mathbf{B} \quad (12.6.5)$$

we have

$$\partial_x E_z = \partial_t B_y \quad (12.6.6)$$

Using (12.6.2) and (12.6.4) in the above, we have

$$ik\alpha(t) D e^{ikx} + c.c. = \dot{\beta} \frac{D}{c} e^{ikx} + c.c. \quad (12.6.7)$$

---

<sup>2</sup>It is debatable as to what gives rise to these dipoles in vacuum. In my opinion, it could be electron-positron pairs that are embedded in vacuum. When vacuum is bombarded with an energetic photon, it is known to produce an electron-positron pair.

From the above, we gather that

$$\dot{\beta} = i\omega\alpha \quad (12.6.8)$$

using

$$\nabla \times \mathbf{B} = \mu_0\epsilon_0\partial_t\mathbf{E} \quad (12.6.9)$$

or more specifically

$$\partial_x B_y = \mu_0\epsilon_0\partial_t E_z \quad (12.6.10)$$

Finally, we find that

$$\dot{\alpha} = i\omega\beta \quad (12.6.11)$$

or that

$$\ddot{\alpha} = -\omega^2\alpha \quad (12.6.12)$$

$$\ddot{\beta} = -\omega^2\beta \quad (12.6.13)$$

A traveling wave solution is that  $\alpha = \alpha_0 e^{-i\omega t}$ ,  $\beta = \beta_0 e^{-i\omega t}$ . From (12.6.8) and (12.6.11), we deduce that  $\alpha_0 = \beta_0$

### 12.6.1 Hamiltonian

The Hamiltonian of the electromagnetic system is given by the total energy of the system which is

$$H = \int_0^L W(x) dx \quad (12.6.14)$$

where

$$W(x) = \frac{1}{2} \left( \epsilon_0 |\mathbf{E}|^2 + \frac{1}{\mu_0} |\mathbf{B}|^2 \right) \quad (12.6.15)$$

From (12.6.2), we derive that for a propagating mode,

$$E_z^2 = 2|\alpha(t)|^2 |D|^2 + \text{O.T.} \quad (12.6.16)$$

$$B_y^2 = 2|\beta(t)|^2 \frac{|D|^2}{c^2} + \text{O.T.} \quad (12.6.17)$$

where ‘‘O.T.’’ stands for ‘‘oscillatory terms’’. When substituted into (12.6.14) and integrated, the oscillatory terms disappear; we have

$$\begin{aligned} H &= L|D|^2 \epsilon_0 |\alpha(t)|^2 + \frac{L|D|^2}{\mu_0 c^2} |\beta(t)|^2 \\ &= L|D|^2 \epsilon_0 [|\alpha(t)|^2 + |\beta(t)|^2] \end{aligned} \quad (12.6.18)$$

Defining

$$D = \sqrt{\frac{\omega}{4L\epsilon_0}} \quad (12.6.19)$$

the above can be written as

$$H = \frac{\omega}{4} [|\alpha|^2 + |\beta|^2] = \frac{\omega}{2} |\alpha|^2 \quad (12.6.20)$$

By letting  $\alpha = p \pm iq$ , it becomes

$$H = \frac{\omega}{2} (p^2 + q^2) \quad (12.6.21)$$

The above resembles the Hamiltonian for a classical harmonic oscillator. Again, the reason for its simplicity is that we have assumed a single propagating mode. This single propagating mode is causing all the dipoles in the system to be oscillating time-harmonically in unison, but with a phase lag between them.<sup>3</sup>

But  $p$  and  $q$  are observables with Hermitian operators. Therefore, the quantum analogue is

$$\hat{H} = \frac{\omega}{2} (\hat{p}^2 + \hat{q}^2) \quad (12.6.22)$$

In coordinate space representation,

$$\hat{p} = -i\hbar \frac{d}{dq}, \quad \hat{q} = q \quad (12.6.23)$$

As in the linear atomic chain case, we can show that

$$\hat{H} = \frac{\hbar\omega}{2} (\hat{a}^\dagger \hat{a} + \hat{a} \hat{a}^\dagger) = \hbar\omega \left( \hat{a}_k^\dagger \hat{a}_k + \frac{1}{2} \right) \quad (12.6.24)$$

## 12.6.2 Field Operators

Alternatively, we can rewrite (12.6.20) as

$$H = \frac{\omega}{4} [\alpha(t)\alpha^*(t) + \alpha(t)\alpha^*(t)] \quad (12.6.25)$$

The above can be elevated to be quantum mechanical operators so that

$$\hat{H} = \frac{\omega}{4} [\hat{\alpha}\hat{\alpha}^\dagger + \hat{\alpha}\hat{\alpha}^\dagger] \quad (12.6.26)$$

Comparing the above with (12.6.24), we infer that

$$\hat{\alpha} = \sqrt{2\hbar}\hat{a} \quad (12.6.27)$$

---

<sup>3</sup>Notice that there is no radiation damping in these dipoles as the periodic boundary condition wraps the wave around, and energy is never lost from the system.

to within an arbitrary phase. Hence,  $\alpha$  in (12.6.2) can be elevated to be a field operator to yield, in the Heisenberg picture,

$$\hat{E}_z = \sqrt{\frac{\hbar\omega_k}{2L\epsilon_0}} \hat{a}_k e^{ikx - i\omega_k t} + \text{c.c.} \quad (12.6.28)$$

$$\hat{B}_y = \sqrt{\frac{\hbar\omega_k}{2L\epsilon_0}} \frac{1}{c} \hat{a}_k e^{ikx - i\omega_k t} + \text{c.c.} \quad (12.6.29)$$

The above can be generalized to a plane wave mode propagating in arbitrary direction, namely,

$$\hat{\mathbf{E}} = \sum_{\mathbf{k},s} \sqrt{\frac{\hbar\omega_k}{2V\epsilon_0}} \mathbf{e}_s \hat{a}_{\mathbf{k},s} e^{i\mathbf{k}\cdot\mathbf{r} - i\omega_k t} + \text{c.c.} \quad (12.6.30)$$

$$\hat{\mathbf{B}} = \sum_{\mathbf{k},s} \sqrt{\frac{\hbar\omega_k}{2V\epsilon_0}} \frac{1}{c} \hat{\mathbf{k}} \times \mathbf{e}_s \hat{a}_{\mathbf{k},s} e^{i\mathbf{k}\cdot\mathbf{r} - i\omega_k t} + \text{c.c.} \quad (12.6.31)$$

where  $\mathbf{e}_s$  is a unit vector denoting the polarization of the electric field. It is orthogonal to the  $\mathbf{k}$  vector. It can be either linearly polarized or circularly polarized. In the linearly polarized case,  $\mathbf{e}_s$  is either vertical or horizontal. In the circularly polarized case,  $\mathbf{e}_s$  is either right-hand or left-hand circularly polarized. In the above, we have defined the Hamiltonian as an integral over a volume  $V$ , and hence, the normalization above is with respect to  $V$ .

A photon is a packet of energy that propagates due to coupled quantum harmonic oscillators. This is similar to the phonon case which is a packet of energy due to coupled quantum harmonic oscillators of mechanical origin. A photon can also propagate in a dielectric medium, and it is quite clear then that in this case, the photon is due to coupled quantum harmonic oscillators of dipolar origin (or electronic origin).

The quantization of energy is with respect to the entire Hamiltonian which describes a set of coupled quantum harmonic oscillators. For the purely monochromatic photon, as is seen before, if the quantization energy is  $\hbar\omega_k$ , it will be shared among all the coupled harmonic oscillators. This is untenable physically as the number of coupled harmonic oscillators becomes large. Hence, it is more appropriate to think of a photon as quasi-monochromatic as a linear superposition of wavenumbers  $k$ , allowing a localized wave packet to be formed. A photon can be thought of as a ‘‘particle’’ possessing a packet of energy  $\hbar\omega_e$ , but it exists as a linear superposition of different modes with different  $k$ 's. Each mode can have a energy  $\hbar\omega_k$  associated with it, but as the particle is a linear superposition of different modes,  $\hbar\omega_e$  is the expectation value of the photon energy when the photon is in a superposition of different quantum states.<sup>4</sup>

### 12.6.3 Multimode Case and Fock State

When polarization is considered, we may write the Hamiltonian for a single mode photon as

$$\hat{H}_k = \sum_s \hat{H}_{k,s} = \sum_s \hbar\omega_k \left( \hat{a}_{k,s}^\dagger \hat{a}_{k,s} + \frac{1}{2} \right) \quad (12.6.32)$$

---

<sup>4</sup>This is my personal opinion.

where  $s$  stands for either horizontal or vertical polarization. The above can be easily derived by considering the electric field with both polarizations present.

In the above, we can denote eigenstate of the Hamiltonian as<sup>5</sup>

$$|\psi\rangle = |n_v\rangle|n_h\rangle \quad (12.6.33)$$

where  $v$  stands for vertical polarization and  $h$  stands for horizontal polarization. In the above, there are  $n_v$  photons in the vertical polarization state and there are  $n_h$  photons in the horizontal polarization state.

When many modes are considered simultaneously, it can be shown that the corresponding Hamiltonian is

$$\hat{H} = \sum_{\mathbf{k},s} \hat{H}_{\mathbf{k},s} = \sum_{\mathbf{k},s} \hbar\omega_{\mathbf{k}} \left( \hat{a}_{\mathbf{k},s}^\dagger \hat{a}_{\mathbf{k},s} + \frac{1}{2} \right) \quad (12.6.34)$$

The eigenstate of the above Hamiltonian can be written as

$$|\psi\rangle = \prod_{\mathbf{k},s} |n_{\mathbf{k},s}\rangle \quad (12.6.35)$$

The above can be written alternatively as

$$|\psi\rangle = \prod_s |n_{\mathbf{k}_1,s}\rangle |n_{\mathbf{k}_2,s}\rangle |n_{\mathbf{k}_3,s}\rangle \cdots = \prod_s |n_{\mathbf{k}_1,s}, n_{\mathbf{k}_2,s}, n_{\mathbf{k}_3,s}, n_{\mathbf{k}_4,s} \cdots\rangle = \prod_s |\{n_{\mathbf{k},s}\}\rangle \quad (12.6.36)$$

The above state is known as a Fock state or the occupational number state. It is customary to leave out the states where the occupational number is zero, listing only the states with non-zero number of photons. For example

$$|\psi\rangle = |3_{\mathbf{k}_1,v}, 2_{\mathbf{k}_2,h}, 1_{\mathbf{k}_3,v}\rangle \quad (12.6.37)$$

indicates a quantum state with three V (vertical) polarized photons in  $\mathbf{k}_1$  mode, two H (horizontal) polarized photons in  $\mathbf{k}_2$  mode, and one V polarized photon in  $\mathbf{k}_3$  mode.

#### 12.6.4 One-Photon State

The understanding of the one-photon state is important because of its usefulness in quantum information, quantum communication, and quantum cryptography. It is prudent to elaborate on it further. For example, when only one  $\mathbf{k}$  mode is present, to denote a photon in the vertical polarization state, and no photons in the horizontal polarization state, the state vector should be

$$|\psi\rangle = |1_v\rangle|0_h\rangle \quad (12.6.38)$$

Often, this is just written as

$$|\psi\rangle = |1_v\rangle \quad (12.6.39)$$

---

<sup>5</sup>When the Hamiltonian consists of sum of the Hamiltonians of two quantum systems, the states of the Hamiltonian operator can be represented by the product space of the states of the individual quantum system.

where the second part is understood. For a photon in arbitrary polarization, it is written as

$$|\psi\rangle = a_v|1_v\rangle + a_h|1_h\rangle \quad (12.6.40)$$

The above denotes a one-photon state where the photon is in the linear superposition of two orthonormal one-photon states. Hence,

$$|a_v|^2 + |a_h|^2 = 1 \quad (12.6.41)$$

Often, photons are generated by atomic transitions whereby the electron changes from a high energy state to a low energy state giving rise to one or more photons. Since this is a causal event, the wave field associated with a photon is a localized wave in space time.<sup>6</sup> The localized wave is formed by linear superposing wave field with different  $\mathbf{k}$  values. For instance, a localized one-photon, vertically polarized state is a linear superposition of a number of one-photon states with different  $\mathbf{k}$  values as follows:

$$|\psi\rangle = \sum_{\mathbf{k}} a_{\mathbf{k},v} |1_{\mathbf{k},v}\rangle \quad (12.6.42)$$

The wave field associated with the above state can be localized in space time, and the wave field satisfies causality.

It is to be noted that (12.6.34) is a quantum mechanical Hamiltonian involving the sum over individual Hamiltonians of different quantum systems. The quantum system arises from a set of coupled quantum harmonic oscillators. The choice of appropriate modes “diagonalizes” the system, giving rise to apparently uncoupled systems. When a photon is in a quantum state, it can be in a linear superposition of different quantum states of these different quantum systems. Therefore, one-photon states (12.6.40) and (12.6.42) should be thought of as a linear superposition of different quantum states, subject to the quantum interpretation of quantum mechanics. The particle is in a linear superposition of states before the measurement, and it collapses to one of the states after the measurement.

Even though a photon is associated with a packet of energy, when it is detected with a polarizer (that detects it either in the vertical or horizontal polarized state), it is found either in one state or the other. The packet of energy is never split between the two states indicated by (12.6.40). Hence, experiment evidence suggests that a photon is a quantum particle in the sense that an electron is a quantum particle at the quantum level. The same interpretation applies to (12.6.42). The subject of quantum interpretation will be discussed later.

### 12.6.5 Coherent State Revisited

The expectation value of the annihilation and creation operators with respect to the photon number states is always zero. Therefore, the expectation value of the field operators, which are proportional to the annihilation and creation operators, is always zero. Hence, the number states are non-classical. To arrive at an expectation value of field operators that are non-zero, and resemble a classical field, one has to work with the coherent states.

<sup>6</sup>See the appendix for a discussion on localized wave packets.

We have studied the coherent state previously. It is defined to be the eigenstate of the annihilation operator. Namely, if  $|\alpha\rangle$  represents a coherent state, then

$$\hat{a}|\alpha\rangle = \alpha|\alpha\rangle \quad (12.6.43)$$

where we have used the eigenvalue  $\alpha$  to index the coherent state. Since the number state  $|n\rangle$  is complete, we can expand the coherent state in terms of the number state, or

$$|\alpha\rangle = \sum_{n=0}^{\infty} C_n |n\rangle \quad (12.6.44)$$

When the annihilation operator is applied to both sides, we obtain

$$\hat{a}|\alpha\rangle = \sum_{n=0}^{\infty} C_n \hat{a}|n\rangle = \sum_{n=0}^{\infty} C_n \sqrt{n-1} |n-1\rangle = \sum_{n=0}^{\infty} C_{n+1} \sqrt{n} |n\rangle \quad (12.6.45)$$

Equating the above with  $\alpha|\alpha\rangle$ , we have

$$\sum_{n=0}^{\infty} C_{n+1} \sqrt{n} |n\rangle = \alpha \sum_{n=0}^{\infty} C_n |n\rangle \quad (12.6.46)$$

By the orthogonality of the number states, matching the coefficients, we have  $C_{n+1} = \alpha C_n / \sqrt{n+1}$ , or

$$C_n = C_{n-1} \alpha / \sqrt{n} = C_{n-2} \alpha^2 / \sqrt{(n-1)n} = C_0 \alpha^n / \sqrt{n!} \quad (12.6.47)$$

Consequently,

$$|\alpha\rangle = C_0 \sum_{n=0}^{\infty} \frac{\alpha^n}{\sqrt{n!}} |n\rangle \quad (12.6.48)$$

The above can be normalized to show that  $C_0 = \exp(-|\alpha|^2/2)$ .

The photon number operator is  $\hat{n} = \hat{a}^\dagger \hat{a}$ . It can be shown that the average number of photons associated with a coherent state is given by

$$\langle \hat{n} \rangle = \bar{n} = \langle \alpha | \hat{n} | \alpha \rangle = \langle \alpha | \hat{a}^\dagger \hat{a} | \alpha \rangle = |\alpha|^2 \quad (12.6.49)$$

Moreover

$$\langle \hat{n}^2 \rangle = \langle \alpha | \hat{n}^2 | \alpha \rangle = \langle \alpha | \hat{a}^\dagger \hat{a} \hat{a}^\dagger \hat{a} | \alpha \rangle = \langle \alpha | \hat{a}^\dagger \hat{a} \dagger \hat{a} \hat{a} + \hat{a}^\dagger \hat{a} | \alpha \rangle = |\alpha|^4 + |\alpha|^2 = \bar{n}^2 + \bar{n} \quad (12.6.50)$$

Then

$$\Delta n = \sqrt{\langle \hat{n}^2 \rangle - \langle \hat{n} \rangle^2} = \bar{n}^{1/2} \quad (12.6.51)$$

The above is characteristic of a Poisson process. The probability of detecting  $n$  photons in a coherent state is

$$P_n = |\langle n | \alpha \rangle|^2 = e^{-|\alpha|^2} \frac{|\alpha|^{2n}}{n!} = e^{-\bar{n}} \frac{\bar{n}^n}{n!} \quad (12.6.52)$$



typical of a Poisson distribution.<sup>7</sup> The fractional uncertainty in the photon number is

$$\frac{\Delta n}{\bar{n}} = \frac{1}{\sqrt{\bar{n}}} \quad (12.6.53)$$

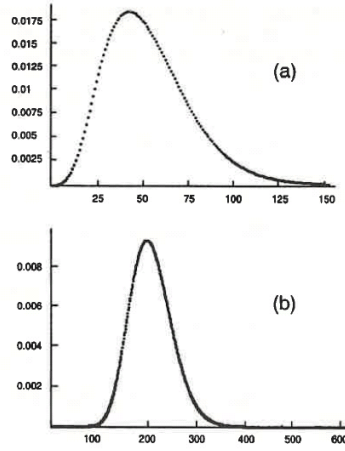


Figure 12.2: Two typical Poisson distributions. Case (a) has  $\bar{n} = 50$  while case (b) has  $\bar{n} = 200$  (from Haus).

### Time Evolution of the Coherent State

In the Schrödinger picture, the time dependence is with the eigenfunctions; hence, the time dependence of the coherent state is

$$|\alpha, t\rangle = e^{-\frac{1}{2}|\alpha|^2} \sum_{n=0}^{\infty} \frac{\alpha^n e^{i(n+1/2)\omega t}}{\sqrt{n!}} |n\rangle \quad (12.6.54)$$

The above can be written as

$$|\alpha, t\rangle = e^{i\omega t/2} e^{-\frac{1}{2}|\alpha|^2} \sum_{n=0}^{\infty} \frac{(\alpha e^{i\omega t})^n}{\sqrt{n!}} |n\rangle = e^{i\omega t/2} |\alpha e^{i\omega t}\rangle \quad (12.6.55)$$

In the Heisenberg picture, the coherent state is time independent, and time-dependent eigenstate can be obtained by applying the operator

$$\begin{aligned} |\alpha, t\rangle &= e^{-i\hat{H}t/\hbar} |\alpha\rangle = \sum_{n=0}^{\infty} \frac{\alpha^n}{\sqrt{n!}} e^{-i\hat{H}t/\hbar} |n\rangle \\ &= \sum_{n=0}^{\infty} \frac{\alpha^n}{\sqrt{n!}} e^{-i(n+1/2)\omega t} |n\rangle = \sum_{n=0}^{\infty} \frac{(\alpha e^{i\omega t})^n}{\sqrt{n!}} |n\rangle = e^{i\omega t/2} |\alpha e^{i\omega t}\rangle \end{aligned} \quad (12.6.56)$$

<sup>7</sup>An example of a Poisson distribution is a shot noise source, producing  $\bar{n}$  photons in interval  $\tau$ . The probability of find  $n$  photons in this time interval follows the Poisson distribution.

It can be further shown that the above can be summed in terms of a “Gaussian pulse”<sup>8</sup> indicating that this pulse evolves in time without shape distortion.

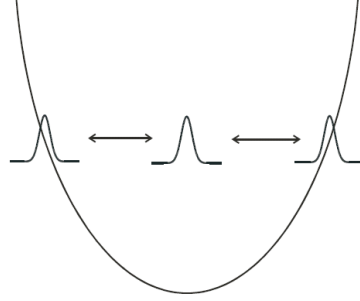


Figure 12.3: The time evolution of the coherent state. It follows the motion of a classical pendulum or harmonic oscillator (From Gerry and Knight).

The expectation value of the field operator in the Heisenberg picture with respect to the coherent state is then

$$\mathbf{E} = \langle \alpha | \hat{\mathbf{E}} | \alpha \rangle = \sum_{\mathbf{k}, s} \sqrt{\frac{\hbar \omega_{\mathbf{k}}}{2V \epsilon_0}} \mathbf{e}_s \langle \alpha | \hat{a}_{\mathbf{k}, s} | \alpha \rangle e^{i\mathbf{k} \cdot \mathbf{r} - i\omega_{\mathbf{k}} t} + \text{c.c.} = \sum_{\mathbf{k}, s} \sqrt{\frac{\hbar \omega_{\mathbf{k}}}{2V \epsilon_0}} \mathbf{e}_s \alpha_{\mathbf{k}, s} e^{i\mathbf{k} \cdot \mathbf{r} - i\omega_{\mathbf{k}} t} + \text{c.c.} \quad (12.6.57)$$

Similarly,

$$\mathbf{B} = \langle \alpha | \hat{\mathbf{B}} | \alpha \rangle = \sum_{\mathbf{k}, s} \sqrt{\frac{\hbar \omega_{\mathbf{k}}}{2V \epsilon_0}} \frac{1}{c} \hat{\mathbf{k}} \times \mathbf{e}_s \alpha_{\mathbf{k}, s} e^{i\mathbf{k} \cdot \mathbf{r} - i\omega_{\mathbf{k}} t} + \text{c.c.} \quad (12.6.58)$$

The above certainly look like a classical field where  $\alpha_{\mathbf{k}, s}$  are complex numbers.

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<sup>8</sup>It is to be noted that the Gaussian pulse we refer to is with respect to the displacement of the quantum harmonic oscillator, not with respect to the direction of propagation of the photon.

## Chapter 13

# Schrödinger Wave Fields

### 13.1 Introduction

We notice that electromagnetic field can be viewed as a collection of photons each with an field attached. These photons collectively form the coherent state that produces a wave field that is analogous to the classical electromagnetic wave field. Photons are bosons whose wave field satisfies Maxwell's equations.

Similarly, if a boson wavefunction satisfies Schrödinger equation, a collection of bosons gives rise to a wave field that satisfies Schrödinger equation. Similarly, a collection of fermions gives rise to a wave field that satisfies Schrödinger equation if each individual particle wavefunction satisfies the same equation. We will call such fields Schrödinger wave fields.

We have seen that the number of bosons in an electromagnetic wave field can be tracked by the annihilation and creation operators; similar operators are needed for tracking the particle of a Schrödinger wave field. The wave field can be a collection of bosons or a collection of fermions. Eventually, the bookkeeping of the many-particle system becomes simpler.

For Schrödinger equation, this is the second quantization. The first quantization was the discovery that electrons can be associated with a wavefunction, and the trapped modes in potential well are quantized. The second quantization refers to the fact that the Schrödinger wave field, which can be treated as a continuum, can be made granular or quantized. This is the advanced way of expressing the wave-particle duality concept inherent in a wave field.<sup>1</sup>

### 13.2 Fock Space for Fermions

For fermions, each eigenmode or eigenstate is either occupied or unoccupied. In general, there will be infinitely many modes that the fermions can occupy. We could denote a two-particle state as

$$|\psi_{2p}\rangle = |\cdots, 0, 1_m, 0, \cdots, 0, 1_v, 0, \cdots\rangle \quad (13.2.1)$$

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<sup>1</sup>The quantization of classical fields is also known as first quantization for them.

where only the  $m$  and  $v$  modes are occupied and the rest of the modes are unoccupied. The notation above is cumbersome, and it is common to denote the two-particle fermion Fock state as

$$|\psi_{2p}\rangle = |m, v\rangle \quad (13.2.2)$$

where for fermions, it is necessary that

$$|m, v\rangle = -|v, m\rangle \quad (13.2.3)$$

We define a fermion creation operator such that

$$\hat{b}_k^\dagger |m, n\rangle = |k, m, n\rangle \quad (13.2.4)$$

The newly created particle always occupies the first space in the Fock state vector. By definition then

$$\hat{b}_l^\dagger \hat{b}_k^\dagger |\dots\rangle = |l, k, \dots\rangle \quad (13.2.5)$$

$$\hat{b}_k^\dagger \hat{b}_l^\dagger |\dots\rangle = |k, l, \dots\rangle = -|l, k, \dots\rangle = -\hat{b}_l^\dagger \hat{b}_k^\dagger |\dots\rangle \quad (13.2.6)$$

Therefore

$$\hat{b}_k^\dagger \hat{b}_l^\dagger + \hat{b}_l^\dagger \hat{b}_k^\dagger = 0 \quad (13.2.7)$$

Similarly, we can show that

$$\hat{b}_k \hat{b}_l + \hat{b}_l \hat{b}_k = 0 \quad (13.2.8)$$

or that the above is the Hermitian conjugate of (13.2.7). To show (13.2.8) more rigorously, we start with a fermion Fock state

$$|\psi\rangle = |\dots, l, k, \dots\rangle \quad (13.2.9)$$

This state acquires a minus sign when we permute two of the particles. Finally, we can have a state such that

$$|\psi_F\rangle = |l, k, \dots\rangle = -|k, l, \dots\rangle \quad (13.2.10)$$

An annihilation operator can be defined such that

$$\hat{b}_l |l, k, \dots\rangle = |k, \dots\rangle \quad (13.2.11)$$

It can be shown that

$$\hat{b}_k \hat{b}_l |\psi_F\rangle = \hat{b}_k \hat{b}_l |l, k, \dots\rangle = \hat{b}_k |k, \dots\rangle = |\dots\rangle \quad (13.2.12)$$

Similarly,

$$-\hat{b}_l \hat{b}_k |\psi_F\rangle = \hat{b}_l \hat{b}_k |k, l, \dots\rangle = \hat{b}_l |l, \dots\rangle = |\dots\rangle \quad (13.2.13)$$

Therefore,

$$\hat{b}_k \hat{b}_l + \hat{b}_l \hat{b}_k = 0 \quad (13.2.14)$$

The above equations (13.2.7) and (13.2.14) imply that if  $k = l$ ,

$$\hat{b}_k^\dagger \hat{b}_k^\dagger = 0 \quad (13.2.15)$$

$$\hat{b}_k \hat{b}_k = 0 \quad (13.2.16)$$

It means that no two particles can be created in the same mode  $k$  or no two particles can be annihilated from the same mode  $k$ .

By the same token, we can show that for  $k \neq l$ ,

$$\hat{b}_k^\dagger \hat{b}_l + \hat{b}_l \hat{b}_k^\dagger = 0 \quad (13.2.17)$$

When  $k = l$ , we have a state  $|k, \dots\rangle$  and then

$$\begin{aligned} (\hat{b}_k^\dagger \hat{b}_k + \hat{b}_k \hat{b}_k^\dagger)|k, \dots\rangle &= \hat{b}_k^\dagger \hat{b}_k |k, \dots\rangle + \hat{b}_k \hat{b}_k^\dagger |k, \dots\rangle \\ &= \hat{b}_k^\dagger | \dots \rangle + 0 |k, \dots\rangle = \hat{b}_k^\dagger | \dots \rangle = |k, \dots\rangle \\ &= \hat{I} |k, \dots\rangle \end{aligned} \quad (13.2.18)$$

In conclusion,

$$\hat{b}_k^\dagger \hat{b}_l + \hat{b}_l \hat{b}_k^\dagger = \delta_{kl} \hat{I} \quad (13.2.19)$$

In summary,

$$\left[ \hat{b}_k^\dagger, \hat{b}_l \right]_+ = \hat{b}_k^\dagger \hat{b}_l + \hat{b}_l \hat{b}_k^\dagger = \delta_{kl} \hat{I} \quad (13.2.20)$$

$$\left[ \hat{b}_k^\dagger, \hat{b}_l^\dagger \right]_+ = \hat{b}_k^\dagger \hat{b}_l^\dagger + \hat{b}_l^\dagger \hat{b}_k^\dagger = 0 \quad (13.2.21)$$

$$\left[ \hat{b}_k, \hat{b}_l \right]_+ = \hat{b}_k \hat{b}_l + \hat{b}_l \hat{b}_k = 0 \quad (13.2.22)$$

where  $[A, B]_+ = AB + BA$ .

### 13.3 Field Operators

Just as in the case of electromagnetic field operators that represent many photons which are bosons, we can define field operators for fermions and bosons in general. A one-particle Schrödinger field operator is defined such that

$$\hat{\psi}(\mathbf{r}) = \sum_j \hat{b}_j \phi_j(\mathbf{r}) \quad (13.3.1)$$

where  $\phi_j$  is an eigenmode from a complete set of orthonormal functions. It may or may not be the  $j$ -th eigenstate of the quantum system governed by the one-particle Schrödinger equation. The above field operator acts on a state vector in the Fock space. For instance,

$$\begin{aligned}
\hat{\psi}(\mathbf{r})|m\rangle &= \hat{\psi}(\mathbf{r})\hat{b}_m^\dagger|0\rangle \\
&= \sum_j \phi_j(\mathbf{r})\hat{b}_j\hat{b}_m^\dagger|0\rangle \\
&= \sum_j \phi_j(\mathbf{r})(\delta_{jm} - \hat{b}_m^\dagger\hat{b}_j)|0\rangle \\
&= \phi_m(\mathbf{r})|0\rangle
\end{aligned} \tag{13.3.2}$$

Notice that the above field operator, when operating on a Fock vector with a  $m$ -th mode occupied, produces a vector tagged with the spatial dependence of the  $m$ -th eigenmode.

After using the completeness property of the orthonormal basis, plus the use of commutation relations for the annihilation and creation operators, the field operator can be shown to satisfy the following commutation relations:

$$\left[ \hat{\psi}^\dagger(\mathbf{r}), \hat{\psi}(\mathbf{r}') \right]_+ = \hat{\psi}^\dagger(\mathbf{r})\hat{\psi}(\mathbf{r}') + \hat{\psi}(\mathbf{r}')\hat{\psi}^\dagger(\mathbf{r}) = \delta(\mathbf{r} - \mathbf{r}') \tag{13.3.3}$$

$$\left[ \hat{\psi}^\dagger(\mathbf{r}), \hat{\psi}^\dagger(\mathbf{r}') \right]_+ = \hat{\psi}^\dagger(\mathbf{r})\hat{\psi}^\dagger(\mathbf{r}') + \hat{\psi}^\dagger(\mathbf{r}')\hat{\psi}^\dagger(\mathbf{r}) = 0 \tag{13.3.4}$$

$$\left[ \hat{\psi}(\mathbf{r}), \hat{\psi}(\mathbf{r}') \right]_+ = \hat{\psi}(\mathbf{r})\hat{\psi}(\mathbf{r}') + \hat{\psi}(\mathbf{r}')\hat{\psi}(\mathbf{r}) = 0 \tag{13.3.5}$$

As an extension, one can define two-particle field operator

$$\hat{\psi}(\mathbf{r}_1, \mathbf{r}_2) = \frac{1}{\sqrt{2}} \sum_{n_1, n_2} \hat{b}_{n_2}\hat{b}_{n_1}\phi_{n_1}(\mathbf{r}_1)\phi_{n_2}(\mathbf{r}_2) \tag{13.3.6}$$

where the factor  $1/\sqrt{2}$  is needed for normalization. Here,  $\phi_{n_1}(\mathbf{r}_1)$  and  $\phi_{n_2}(\mathbf{r}_2)$  may or may not be one-particle eigenstates of the Schrödinger equation. The requirements on them are that they are complete and orthonormal. In the above, the  $n_1 = n_2$  terms vanish because  $\hat{b}_n\hat{b}_n = 0$ . It can be shown easily that

$$\hat{\psi}(\mathbf{r}_1, \mathbf{r}_2)|l, m\rangle = \hat{\psi}(\mathbf{r}_1, \mathbf{r}_2)\hat{b}_l^\dagger\hat{b}_m^\dagger|0\rangle = \frac{1}{\sqrt{2}}[\phi_l(\mathbf{r}_1)\phi_m(\mathbf{r}_2) - \phi_l(\mathbf{r}_2)\phi_m(\mathbf{r}_1)]|0\rangle \tag{13.3.7}$$

It is to be noted that the above is zero when  $\mathbf{r}_1 = \mathbf{r}_2$ . It is a consequence of the field operator operating on the two-particle Fock state, and not the property of the field operator. A three-particle field operator is defined to be

$$\hat{\psi}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3) = \frac{1}{\sqrt{3!}} \sum_{n_1, n_2, n_3} \hat{b}_{n_3}\hat{b}_{n_2}\hat{b}_{n_1}\phi_{n_1}(\mathbf{r}_1)\phi_{n_2}(\mathbf{r}_2)\phi_{n_3}(\mathbf{r}_3) \tag{13.3.8}$$

In the above, none of the  $n_1, n_2, n_3$  can be repeated for the reasons of (13.2.15) and (13.2.16).

Then, it can be shown that

$$\begin{aligned}
 \hat{\psi}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3) |l, m, n\rangle &= \hat{\psi}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3) \hat{b}_l^\dagger \hat{b}_m^\dagger \hat{b}_n^\dagger |0\rangle \\
 &= \frac{1}{\sqrt{3!}} [\phi_l(\mathbf{r}_1) \phi_m(\mathbf{r}_2) \phi_n(\mathbf{r}_3) + \phi_l(\mathbf{r}_2) \phi_m(\mathbf{r}_3) \phi_n(\mathbf{r}_1) \\
 &\quad + \phi_l(\mathbf{r}_3) \phi_m(\mathbf{r}_1) \phi_n(\mathbf{r}_2) - \phi_l(\mathbf{r}_1) \phi_m(\mathbf{r}_3) \phi_n(\mathbf{r}_2) \\
 &\quad - \phi_l(\mathbf{r}_2) \phi_m(\mathbf{r}_1) \phi_n(\mathbf{r}_3) - \phi_l(\mathbf{r}_3) \phi_m(\mathbf{r}_2) \phi_n(\mathbf{r}_1)] |0\rangle \quad (13.3.9)
 \end{aligned}$$

In general, an  $N$ -particle field operator is

$$\hat{\psi}(\mathbf{r}_1, \mathbf{r}_2, \dots, \mathbf{r}_N) = \frac{1}{\sqrt{N!}} \sum_{n_1, n_2, \dots, n_N} \hat{b}_{n_N} \dots \hat{b}_{n_2} \hat{b}_{n_1} \phi_{n_1}(\mathbf{r}_1) \phi_{n_2}(\mathbf{r}_2) \dots \phi_{n_N}(\mathbf{r}_N) \quad (13.3.10)$$

In the above sum, only terms where  $n_1, n_2, \dots, n_N$  are distinct are contributing for the same reason given in (13.3.8). It is seen that

$$\hat{\psi}_{NP}(\mathbf{r}_1, \mathbf{r}_2, \dots, \mathbf{r}_N) = \frac{1}{\sqrt{N!}} \hat{\psi}_{1P}(\mathbf{r}_N) \dots \hat{\psi}_{1P}(\mathbf{r}_2) \hat{\psi}_{1P}(\mathbf{r}_1) \quad (13.3.11)$$

where  $\hat{\psi}_{1P}$  is a one-particle field operator while  $\hat{\psi}_{NP}$  is the  $N$ -particle field operator. However, the used of such factorized form has to be handled with caution, as each one-particle operator has to operate on distinct modes in order for the product of factors to be non-zero.

A general Fock state for  $N$  fermion particles can be written as

$$|\Phi_N\rangle = \hat{b}_{n_1}^\dagger \hat{b}_{n_2}^\dagger \hat{b}_{n_3}^\dagger \dots \hat{b}_{n_N}^\dagger |0\rangle = |n_1, n_2, n_3, \dots, n_N\rangle \quad (13.3.12)$$

In the above, the state vector changes sign when any two of the annihilation operators swap position, and that the state vector is zero if any two of the modes are identical.

When the Fock state is operated on by the  $N$ -particle annihilation operator, the coordinate space representation of the  $N$ -particle fermion state is obtained as shown by (13.3.7). This approach avoids the cumbersome use of Slater determinant in the book-keeping of the  $N$ -particle fermion states.

## 13.4 Similarity Transform

It is prudent to note from the above that one can change between coordinate space basis and Fock space basis via the algebra shown. This allows us to effect a change of basis for a complex quantum system. Some quantum systems, when expressed in the Fock space basis, is a lot simpler than the coordinate space basis. This was first shown for electromagnetic field. It can be shown for other quantum fields. Hence, it is worthwhile to review the mathematics of similarity transform next.

We can view the change of basis as a similarity transform when we change from eigenfunction space or coordinate space representation to Fock space representation. Given a matrix equation

$$\overline{\mathbf{A}} \cdot \mathbf{x} = \lambda \mathbf{x} \quad (13.4.1)$$

we can define a new representation for the unknowns as

$$\mathbf{x} = \bar{\mathbf{S}} \cdot \mathbf{y} \quad (13.4.2)$$

In the above,  $\bar{\mathbf{S}}$  is unitary since the length of the vectors does not change. The above equation becomes

$$\bar{\mathbf{A}} \cdot \bar{\mathbf{S}} \cdot \mathbf{y} = \lambda \bar{\mathbf{S}} \cdot \mathbf{y} \quad (13.4.3)$$

Multiplying the above by  $\bar{\mathbf{S}}^\dagger$ , we have

$$\bar{\mathbf{S}}^\dagger \cdot \bar{\mathbf{A}} \cdot \bar{\mathbf{S}} \cdot \mathbf{y} = \lambda \bar{\mathbf{S}}^\dagger \cdot \bar{\mathbf{S}} \cdot \mathbf{y} = \lambda \mathbf{y} \quad (13.4.4)$$

where the unitary property of the  $\bar{\mathbf{S}}$  operator has been used. Hence,

$$\bar{\mathbf{A}}_s \cdot \mathbf{y} = \lambda \mathbf{y} \quad (13.4.5)$$

where

$$\bar{\mathbf{A}}_s = \bar{\mathbf{S}}^\dagger \cdot \bar{\mathbf{A}} \cdot \bar{\mathbf{S}} \quad (13.4.6)$$

Equation (13.4.6) is in the form of a similarity transform.

## 13.5 Additive One-Particle Operator

We can consider a simple identical many particle Hamiltonian where the particles do not interact with each other, say via the Coulomb potential. For instance, for fermions, the only way they interact is via Pauli's exclusion principle. An example of such a Hamiltonian for  $N$  particles, in coordinate space representation, is

$$\hat{H}_r = \sum_{i=1}^N \hat{H}_{r_i} \quad (13.5.1)$$

where

$$\hat{H}_{r_i} = -\frac{\hbar}{2m} \nabla_i^2 + V(\mathbf{r}_i) \quad (13.5.2)$$

As an example, consider the three particle case. The three-particle wavefunction in coordinate space representation is

$$\begin{aligned} \psi_{3P}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3) = \frac{1}{\sqrt{3!}} & [\phi_l(\mathbf{r}_1) \phi_m(\mathbf{r}_2) \phi_n(\mathbf{r}_3) + \phi_l(\mathbf{r}_2) \phi_m(\mathbf{r}_3) \phi_n(\mathbf{r}_1) \\ & + \phi_l(\mathbf{r}_3) \phi_m(\mathbf{r}_1) \phi_n(\mathbf{r}_2) - \phi_l(\mathbf{r}_1) \phi_m(\mathbf{r}_3) \phi_n(\mathbf{r}_2) \\ & - \phi_l(\mathbf{r}_2) \phi_m(\mathbf{r}_1) \phi_n(\mathbf{r}_3) - \phi_l(\mathbf{r}_3) \phi_m(\mathbf{r}_2) \phi_n(\mathbf{r}_1)] \end{aligned} \quad (13.5.3)$$

If the eigenmodes above are the eigenmodes of the  $\hat{H}_{r_i}$  operator, then

$$\hat{H}_r \psi_{3P}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3) = \sum_{i=1}^N \hat{H}_{r_i} \psi_{3P}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3) = (E_l + E_m + E_n) \psi_{3P}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3) \quad (13.5.4)$$



The above approach gets unwieldy as the number of particles increases.

For a simpler approach, we can perform a change of basis or similarity transform on the original Hamiltonian. The new Hamiltonian is

$$\hat{H} = \int \sum_{i=1}^N \hat{\psi}_{NP}^\dagger(\mathbf{r}_1, \mathbf{r}_2, \dots, \mathbf{r}_N) \hat{H}_{ri} \hat{\psi}_{NP}(\mathbf{r}_1, \mathbf{r}_2, \dots, \mathbf{r}_N) d\mathbf{r}_1 d\mathbf{r}_2 \dots d\mathbf{r}_N \quad (13.5.5)$$

The operator above now acts on a vector in the Fock space, and transforms it to another vector in the same space.

Equation (13.5.5) can be written, after using (13.3.11), as

$$\begin{aligned} \hat{H} = \frac{1}{N!} \int \hat{\psi}_{1P}^\dagger(\mathbf{r}_N) \dots \sim \dots \hat{\psi}_{1P}^\dagger(\mathbf{r}_1) \left[ \sum_{i=1}^N \int d\mathbf{r}_i \hat{\psi}_{1P}^\dagger(\mathbf{r}_i) \hat{H}_{ri} \hat{\psi}_{1P}(\mathbf{r}_i) \right] \\ \hat{\psi}_{1P}(\mathbf{r}_1) \dots \sim \dots \hat{\psi}_{1P}(\mathbf{r}_N) d\mathbf{r}_1 \dots \sim \dots d\mathbf{r}_N \end{aligned} \quad (13.5.6)$$

where the  $\sim$  sign implies that  $\mathbf{r}_i$  is excluded from the sequence.

### 13.5.1 Three-Particle Case

It is quite complex to sort out the algebra of the above system. Much insight, however, can be gotten by studying a simpler three-particle case. In this case, a typical term of the transformed Hamiltonian is

$$\begin{aligned} \hat{H} = \frac{1}{3!} \int \hat{\psi}_{1P}^\dagger(\mathbf{r}_3) \hat{\psi}_{1P}^\dagger(\mathbf{r}_2) \left[ \int d\mathbf{r}_1 \hat{\psi}_{1P}^\dagger(\mathbf{r}_1) \hat{H}_{r1} \hat{\psi}_{1P}(\mathbf{r}_1) \right] \\ \hat{\psi}_{1P}(\mathbf{r}_2) \hat{\psi}_{1P}(\mathbf{r}_3) d\mathbf{r}_2 d\mathbf{r}_3 + \dots \end{aligned} \quad (13.5.7)$$

where the  $+\dots$  above refers to other two terms where  $\mathbf{r}_i = \mathbf{r}_2$  and  $\mathbf{r}_i = \mathbf{r}_3$ , and  $\mathbf{r}_i$  here refers to the coordinates for the term inside the square brackets.

The space-dependent parts involving  $\phi_{n_i}(\mathbf{r}_i)$  in the field operators can be grouped together and the spatial integrations can be performed first. These modes are orthonormal giving rise to  $\delta_{n_2 n_2'}$  and  $\delta_{n_3 n_3'}$  which can be used to reduce double summations into single summations. Finally, one obtains

$$\hat{H} = \frac{1}{6} \sum_{n_2, n_3} \hat{b}_{n_3}^\dagger \hat{b}_{n_2}^\dagger \left[ \sum_{n_1', n_1} \hat{b}_{n_1}^\dagger H_{n_1', n_1}^{(1)} \hat{b}_{n_1} \right] \hat{b}_{n_2} \hat{b}_{n_3} + \dots \quad (13.5.8)$$

The general case where  $\phi_{n_1}(\mathbf{r}_1)$  is not an eigenstate of the Hamiltonian  $\hat{H}_{r1}$  is assumed here. Hence,

$$H_{n_1', n_1}^{(1)} = \langle \phi_{n_1'} | \hat{H}_{r1} | \phi_{n_1} \rangle \quad (13.5.9)$$

In the event that  $\phi_{n_1}(\mathbf{r}_1)$  is an eigenstate, the above becomes

$$H_{n_1', n_1}^{(1)} = E_{n_1} \delta_{n_1' n_1} \quad (13.5.10)$$

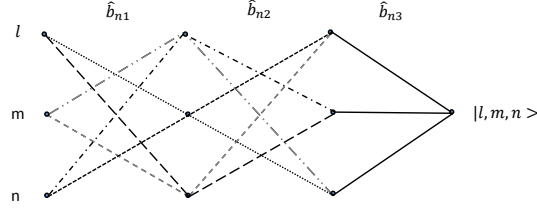


Figure 13.1: The three annihilation operators  $\hat{b}_{n_1}\hat{b}_{n_2}\hat{b}_{n_3}$  acting on the three-particle state  $|l, m, n\rangle$  produces six threads that eventually gives rise to six terms.

To evaluate the action of the above operator on the three-fermion-particle state  $|l, m, n\rangle = \hat{b}_l^\dagger \hat{b}_m^\dagger \hat{b}_n^\dagger |0\rangle$ , one needs to show

$$\begin{aligned} \hat{b}_{n_1}\hat{b}_{n_2}\hat{b}_{n_3}|l, m, n\rangle &= \hat{b}_{n_1}\hat{b}_{n_2}\hat{b}_{n_3}\hat{b}_l^\dagger\hat{b}_m^\dagger\hat{b}_n^\dagger|0\rangle \\ &= (\delta_{ln_3}\delta_{mn_2}\delta_{nm_1} - \delta_{ln_3}\delta_{mn_1}\delta_{nn_2} - \delta_{ln_2}\delta_{mn_3}\delta_{nn_1} \\ &\quad + \delta_{ln_1}\delta_{mn_3}\delta_{nn_2} + \delta_{ln_2}\delta_{mn_1}\delta_{nn_3} - \delta_{ln_1}\delta_{mn_2}\delta_{nn_3})|0\rangle \end{aligned} \quad (13.5.11)$$

A total of six terms is found. In the above  $\hat{b}_{n_3}$ , operating on the three-particle state, produces three terms each of which is a two-particle state as shown in Figure 13.1. Next,  $\hat{b}_{n_2}$  operating on them will produce two terms per each term of the two-particle state, producing a net of six terms of one-particle states. The final operation  $\hat{b}_{n_1}$  produces one term for each particle. In general, if we start with an  $N$ -particle state, the operation of  $\hat{b}_{n_1} \cdots \hat{b}_{n_N}$  on it will produce  $N!$  terms.

The  $\hat{b}^\dagger$  operators in (13.5.8) anti-commute with each other. Hence,  $\hat{b}_{n_3}^\dagger \hat{b}_{n_2}^\dagger$  can be moved to the right of the  $H_{n'_1, n_1}^{(1)}$  element, together with the summations. Therefore, we finally need to evaluate

$$\sum_{n_2, n_3} \hat{b}_{n_3}^\dagger \hat{b}_{n_2}^\dagger \hat{b}_{n_1} \hat{b}_{n_2} \hat{b}_{n_3} |l, m, n\rangle \quad (13.5.12)$$

The above can be shown to evaluate to

$$\begin{aligned} \sum_{n_2, n_3} \hat{b}_{n_3}^\dagger \hat{b}_{n_2}^\dagger \hat{b}_{n_1} \hat{b}_{n_2} \hat{b}_{n_3} |l, m, n\rangle &= 2 \left( \delta_{nn_1} \hat{b}_l^\dagger \hat{b}_m^\dagger + \delta_{ln_1} \hat{b}_m^\dagger \hat{b}_n^\dagger + \delta_{mn_1} \hat{b}_n^\dagger \hat{b}_l^\dagger \right) |0\rangle \\ &= 2\hat{b}_{n_1} \hat{b}_l^\dagger \hat{b}_m^\dagger \hat{b}_n^\dagger |0\rangle = 2\hat{b}_{n_1} |l, m, n\rangle \end{aligned} \quad (13.5.13)$$

The first equality above can be obtained by using (13.5.11). The second equality follows from working backward by expanding  $\hat{b}_{n_1} \hat{b}_l^\dagger \hat{b}_m^\dagger \hat{b}_n^\dagger |0\rangle$ . Consequently,

$$\hat{H}|l, m, n\rangle = \frac{1}{3} \sum_{n'_1, n_1} \hat{b}_{n'_1}^\dagger H_{n'_1, n_1}^{(1)} \hat{b}_{n_1} |l, m, n\rangle + \cdots \quad (13.5.14)$$

The other two terms indicated by the  $+\dots$  would contribute to exactly the same expression as they are from indistinguishable particles. Finally, we have

$$\hat{H}|l, m, n\rangle = \sum_{n'_1, n_1} \hat{b}_{n'_1}^\dagger H_{n'_1, n_1}^{(1)} \hat{b}_{n_1} |l, m, n\rangle \quad (13.5.15)$$

In general, the additive one-particle operator in Fock space is

$$\hat{H} = \sum_{n', n} \hat{b}_{n'}^\dagger H_{n', n}^{(1)} \hat{b}_n \quad (13.5.16)$$

The above can be used for the  $N$ -particle case as long as they are indistinguishable. For the case when the one-particle eigenstates are also the eigenstates of the  $\hat{H}_{r_i}$  operator, the above matrix becomes diagonal yielding,

$$\hat{H} = \sum_n E_n \hat{b}_n^\dagger \hat{b}_n \quad (13.5.17)$$

The above means that if we have an  $N$ -particle fermion field, it can be represented by the physics of one-particle Hamiltonian if the eigenstates chosen for the similarity transform are also the eigenstates of the one-particle Hamiltonian. Also, the above Hamiltonian is very similar to the Hamiltonian for photons derived in the previous chapter, except that photons are bosons and the above derivation is for fermions. A similar derivation for bosons shows that the Hamiltonian is similar to the above.

The above Hamiltonian does not distinguish between one particle or  $N$  particles. This difference of the quantum systems is expressed by the Fock states of the particles.

## 13.6 Additive Two-Particle Operator

In general, the  $N$  particles in a Schrödinger wave field will interact with each other. They may interact pair-wise, for instance, via Coulomb potential other than just the Pauli's exclusion principle. An example of an additive two particle operator in coordinate space is

$$V(\mathbf{r}_1, \mathbf{r}_2) = \frac{e^2}{4\pi\epsilon_0 |\mathbf{r}_1 - \mathbf{r}_2|} \quad (13.6.1)$$

This will be part of a two-particle Hamiltonian

$$\hat{H}_r(\mathbf{r}_1, \mathbf{r}_2) = -\frac{\hbar^2}{2m} (\nabla_1^2 + \nabla_2^2) + V(\mathbf{r}_1, \mathbf{r}_2) \quad (13.6.2)$$

When  $N$ -particle Hamiltonian is considered, the Coulomb interaction appears as

$$V = \sum_{i < j} V(\mathbf{r}_i, \mathbf{r}_j) = \frac{1}{2} \sum_{i \neq j} V(\mathbf{r}_i, \mathbf{r}_j) \quad (13.6.3)$$

For  $N$  particles, there are  $N(N-1)$  terms in the above summations. Again, we can transform the above from coordinate space representation to Fock-space representation as before arriving at

$$\hat{V} = \frac{1}{2} \sum_{i=1}^N \sum_{j=1, j \neq i}^N \frac{1}{N!} \int \hat{\psi}_{1P}^\dagger(\mathbf{r}_N) \cdots \sim \cdots \hat{\psi}_{1P}^\dagger(\mathbf{r}_1) \left[ \int d\mathbf{r}_i d\mathbf{r}_j \hat{\psi}_{1P}^\dagger(\mathbf{r}_i) \hat{\psi}_{1P}^\dagger(\mathbf{r}_j) V(\mathbf{r}_i, \mathbf{r}_j) \hat{\psi}_{1P}(\mathbf{r}_i) \hat{\psi}_{1P}(\mathbf{r}_j) \right] \hat{\psi}_{1P}(\mathbf{r}_1) \cdots \sim \cdots \hat{\psi}_{1P}(\mathbf{r}_N) d\mathbf{r}_1 \cdots \sim \cdots d\mathbf{r}_N \quad (13.6.4)$$

where the  $\sim$  indicates that  $\mathbf{r}_i$  and  $\mathbf{r}_j$  are excluded from the sequence. Without loss of generality, we can focus on the term where  $\mathbf{r}_i = \mathbf{r}_1$  and  $\mathbf{r}_j = \mathbf{r}_2$ . Then, the above becomes

$$\hat{V} = \frac{1}{2N!} \int \hat{\psi}_{1P}^\dagger(\mathbf{r}_N) \cdots \hat{\psi}_{1P}^\dagger(\mathbf{r}_3) \left[ \int d\mathbf{r}_1 d\mathbf{r}_2 \hat{\psi}_{1P}^\dagger(\mathbf{r}_2) \hat{\psi}_{1P}^\dagger(\mathbf{r}_1) V(\mathbf{r}_1, \mathbf{r}_2) \hat{\psi}_{1P}(\mathbf{r}_1) \hat{\psi}_{1P}(\mathbf{r}_2) \right] \hat{\psi}_{1P}(\mathbf{r}_3) \cdots \hat{\psi}_{1P}(\mathbf{r}_N) d\mathbf{r}_1 \cdots d\mathbf{r}_N + \cdots \quad (13.6.5)$$

where  $+\cdots$  above implies the rest of the terms; there are altogether  $N(N-1)$  terms. Performing the  $d\mathbf{r}_1 \cdots d\mathbf{r}_N$  integrations first, making use of mode orthonormality, we arrive at

$$\hat{V} = \frac{1}{2N!} \sum_{n_3, \dots, n_N} \hat{b}_{n_N}^\dagger \cdots \hat{b}_{n_3}^\dagger \left[ \sum_{n'_1, n_1, n'_2, n_2} \hat{b}_{n_2}^\dagger \hat{b}_{n_1}^\dagger V_{n'_1, n'_2, n_1, n_2}^{(1,2)} \hat{b}_{n_1} \right] \hat{b}_{n_2} \cdots \hat{b}_{n_N} + \cdots \quad (13.6.6)$$

where

$$V_{n'_1, n'_2, n_1, n_2}^{(1,2)} = \int d\mathbf{r}_1 d\mathbf{r}_2 \phi_{n'_1}^*(\mathbf{r}_1) \phi_{n'_2}^*(\mathbf{r}_2) V(\mathbf{r}_1, \mathbf{r}_2) \phi_{n_1}(\mathbf{r}_1) \phi_{n_2}(\mathbf{r}_2) \quad (13.6.7)$$

The above needs to operate on an  $N$ -particle Fock state, namely,

$$\hat{V}|l_1, \dots, l_N\rangle = \frac{1}{2N!} \sum_{n_3, \dots, n_N} \hat{b}_{n_N}^\dagger \cdots \hat{b}_{n_3}^\dagger \left[ \sum_{n'_1, n_1, n'_2, n_2} \hat{b}_{n_2}^\dagger \hat{b}_{n_1}^\dagger V_{n'_1, n'_2, n_1, n_2}^{(1,2)} \hat{b}_{n_1} \right] \hat{b}_{n_2} \cdots \hat{b}_{n_N} |l_1, \dots, l_N\rangle + \cdots \quad (13.6.8)$$

The right-most term that acts on the Fock state is of the form

$$\hat{b}_{n_1} \hat{b}_{n_2} \cdots \hat{b}_{n_N} |l_1, \dots, l_N\rangle \quad (13.6.9)$$

The above yields  $N!$  terms similar to (13.5.11). But  $\hat{b}_{n_N}^\dagger \cdots \hat{b}_{n_3}^\dagger$  commutes with  $\hat{b}_{n_2}^\dagger \hat{b}_{n_1}^\dagger$  in (13.6.6) above. We can then move  $\hat{b}_{n_N}^\dagger \cdots \hat{b}_{n_3}^\dagger$  to the right of  $V_{n'_1, n'_2, n_1, n_2}^{(1,2)}$ . Then we need to evaluate

$$\sum_{n_3, \dots, n_N} \hat{b}_{n_N}^\dagger \cdots \hat{b}_{n_3}^\dagger \hat{b}_{n_1} \hat{b}_{n_2} \cdots \hat{b}_{n_N} |l_1, \dots, l_N\rangle = (N-2)! \hat{b}_{n_1} \hat{b}_{n_2} |l_1, \dots, l_N\rangle \quad (13.6.10)$$

The above equality can be proved by induction from the three-particle case or Equation (13.5.13). As a result,

$$\hat{V} |l_1, \dots, l_N\rangle = \frac{1}{2N(N-1)} \sum_{n'_1, n_1, n'_2, n_2} \hat{b}_{n_2}^\dagger \hat{b}_{n_1}^\dagger V_{n'_1, n'_2, n_1, n_2}^{(1,2)} \hat{b}_{n_1} \hat{b}_{n_2} |l_1, \dots, l_N\rangle + \cdots \quad (13.6.11)$$

There are  $N(N-1)$  terms of the similar kind in the  $+\cdots$  above. Summing them up, we have

$$\hat{V} |l_1, \dots, l_N\rangle = \frac{1}{2} \sum_{n'_1, n_1, n'_2, n_2} \hat{b}_{n_2}^\dagger \hat{b}_{n_1}^\dagger V_{n'_1, n'_2, n_1, n_2}^{(1,2)} \hat{b}_{n_1} \hat{b}_{n_2} |l_1, \dots, l_N\rangle \quad (13.6.12)$$

In general, the additive two-particle operator in an  $N$ -particle fermion field is given by

$$\hat{V} = \frac{1}{2} \sum_{n'_1, n_1, n'_2, n_2} \hat{b}_{n_2}^\dagger \hat{b}_{n_1}^\dagger V_{n'_1, n'_2, n_1, n_2}^{(1,2)} \hat{b}_{n_1} \hat{b}_{n_2} \quad (13.6.13)$$

The above derivations for fermions can be repeated for the boson case. In this case, we will have a field of bosons. It is pleasing that by starting out with the  $N$ -particle coordinate space representation of the Schrödinger equation, one arrives at a much simpler Fock space representation of the same equation. It allows one to treat  $N$ -particle problems more succinctly using such a representation.

## 13.7 More on Field Operators

The above shows that the one-particle additive operator is independent of the number of particles. So what distinguishes one-particle quantum system from a two-particle one is in the Fock state: the first one will have a Fock state for one particle while the latter will have a Fock state for two particles.

A general one-particle Fock state is

$$|\psi\rangle = \sum_{\mu} c_{\mu} \hat{b}_{\mu}^{\dagger} |0\rangle \quad (13.7.1)$$

When the one-particle field operator acts on the above, it projects the coordinate space representation of the one-particle wavefunction. Namely,

$$\hat{\psi}(\mathbf{r}) |\psi\rangle = \sum_{\mu} c_{\mu} \phi_{\mu}(\mathbf{r}) |0\rangle = \psi(\mathbf{r}) |0\rangle \quad (13.7.2)$$

The above one-particle state is in a linear superposition of different eigenstates. Hence, it can constitute a wave packet.<sup>2</sup>

Also, from (13.3.1), we can easily show that

$$\hat{b}_\mu^\dagger = \int \phi_\mu(\mathbf{r}) \hat{\psi}^\dagger(\mathbf{r}) d\mathbf{r} \quad (13.7.3)$$

On combining with (13.7.3), we have

$$\begin{aligned} |\psi\rangle &= \sum_\mu c_\mu \int \phi_\mu(\mathbf{r}) \hat{\psi}^\dagger(\mathbf{r}) d\mathbf{r} |0\rangle \\ &= \int \left[ \sum_\mu c_\mu \phi_\mu(\mathbf{r}) \right] \hat{\psi}^\dagger(\mathbf{r}) d\mathbf{r} |0\rangle \\ &= \int f(\mathbf{r}) \hat{\psi}^\dagger(\mathbf{r}) d\mathbf{r} |0\rangle \end{aligned} \quad (13.7.4)$$

Therefore, a general one-particle state can be written as a linear superposition of the one-particle field operator  $\hat{\psi}^\dagger(\mathbf{r})$  weighted by  $f(\mathbf{r})$ . It is clear that  $f(\mathbf{r})$  satisfies the one-particle Schrödinger equation.

For a general two-particle state in Fock space,

$$|\psi\rangle = \sum_{\mu_1, \mu_2} c_{\mu_1, \mu_2} \hat{b}_{\mu_1}^\dagger \hat{b}_{\mu_2}^\dagger |0\rangle \quad (13.7.5)$$

upon substituting (13.7.3) into the above, we have

$$\begin{aligned} |\psi\rangle &= \int \left( \sum_{\mu_1, \mu_2} \phi_{\mu_1}(\mathbf{r}_1) \phi_{\mu_2}(\mathbf{r}_2) c_{\mu_1, \mu_2} \right) \hat{\psi}^\dagger(\mathbf{r}_1) \hat{\psi}^\dagger(\mathbf{r}_2) d\mathbf{r}_1 d\mathbf{r}_2 |0\rangle \\ &= \int f(\mathbf{r}_1, \mathbf{r}_2) \hat{\psi}^\dagger(\mathbf{r}_1) \hat{\psi}^\dagger(\mathbf{r}_2) d\mathbf{r}_1 d\mathbf{r}_2 |0\rangle \end{aligned} \quad (13.7.6)$$

where

$$f(\mathbf{r}_1, \mathbf{r}_2) = \sum_{\mu_1, \mu_2} \phi_{\mu_1}(\mathbf{r}_1) \phi_{\mu_2}(\mathbf{r}_2) c_{\mu_1, \mu_2}. \quad (13.7.7)$$

For fermions, it can be easily shown, using the commutation relations for the field operators, that

$$f(\mathbf{r}_1, \mathbf{r}_2) = -f(\mathbf{r}_2, \mathbf{r}_1) \quad (13.7.8)$$

The above implies that  $c_{\mu_1, \mu_2} = -c_{\mu_2, \mu_1}$ . The above satisfies the two-particle Schrödinger equation. It can be generalized to  $N$  particles, yielding

$$|\psi\rangle = \int f(\mathbf{r}_1, \mathbf{r}_2, \dots, \mathbf{r}_N) \hat{\psi}^\dagger(\mathbf{r}_1) \hat{\psi}^\dagger(\mathbf{r}_2) \dots \hat{\psi}^\dagger(\mathbf{r}_N) d\mathbf{r}_1 d\mathbf{r}_2 \dots d\mathbf{r}_N |0\rangle \quad (13.7.9)$$

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<sup>2</sup>See the appendix for a discussion of wave packets.

## 13.8 Boson Wave Field

The results for fermion wave field in the previous section can also be derived for boson wave field. The difference is that more than one boson can occupy a given state: The Pauli exclusion principle does not apply to bosons. A boson state involving two modes can be denoted as:

$$|\psi\rangle = |n_m, n_v\rangle \quad (13.8.1)$$

where there are  $n_m$  particles in mode  $m$ , and  $n_v$  particles in mode  $v$ . Since ordering is unimportant, it is necessary that

$$|n_m, n_v\rangle = |n_v, n_m\rangle \quad (13.8.2)$$

A boson creation operator can be defined such that

$$\hat{a}_k^\dagger |n_k, n_m, n_v\rangle = C_{n_k+1} |n_k + 1, n_m, n_v\rangle \quad (13.8.3)$$

It raises the number of bosons in state  $k$  by one. By definition,

$$\hat{a}_l^\dagger \hat{a}_k^\dagger |n_k, n_l, \dots\rangle = C_{n_k+1} C_{n_l+1} |n_k + 1, n_l + 1, \dots\rangle \quad (13.8.4)$$

$$\hat{a}_k^\dagger \hat{a}_l^\dagger |n_l, n_k, \dots\rangle = C_{n_l+1} C_{n_k+1} |n_l + 1, n_k + 1, \dots\rangle = \hat{a}_l^\dagger \hat{a}_k^\dagger |n_k, n_l, \dots\rangle \quad (13.8.5)$$

The last equality follows since ordering is unimportant. Consequently,

$$\hat{a}_k^\dagger \hat{a}_l^\dagger - \hat{a}_l^\dagger \hat{a}_k^\dagger = 0 \quad (13.8.6)$$

Similarly, one can define annihilation operators for bosons that have the opposite effect as the creation operators. Hence, going through a similar process, one also has

$$\hat{a}_k \hat{a}_l - \hat{a}_l \hat{a}_k = 0 \quad (13.8.7)$$

From definition,

$$\hat{a}^\dagger |n\rangle = C_{n+1} |n + 1\rangle \quad (13.8.8)$$

$$\hat{a} |n\rangle = B_n |n - 1\rangle \quad (13.8.9)$$

Furthermore, from (13.8.9), we have

$$\langle n + 1 | \hat{a}^\dagger |n\rangle = C_{n+1}, \quad C_{n+1}^* = \langle n | \hat{a} |n + 1\rangle = B_{n+1} \quad (13.8.10)$$

Then, we can show that

$$\hat{a}_l^\dagger \hat{a}_k |n_k, n_l, \dots\rangle = C_{n_k}^* C_{n_l+1} |n_k - 1, n_l + 1, \dots\rangle \quad (13.8.11)$$

$$\hat{a}_k \hat{a}_l^\dagger |n_l, n_k, \dots\rangle = C_{n_l+1}^* C_{n_k} |n_l + 1, n_k - 1, \dots\rangle \quad (13.8.12)$$

If  $C_n$  is a real number, then when  $l \neq k$ ,

$$\hat{a}_k \hat{a}_l^\dagger - \hat{a}_l^\dagger \hat{a}_k = 0 \quad (13.8.13)$$

For the case when  $l = k$ , we can drop the subscripts  $k$  and  $l$  and look at

$$\hat{a}^\dagger \hat{a} |n\rangle = C_n^2 |n\rangle \quad (13.8.14)$$

$$\hat{a} \hat{a}^\dagger |n\rangle = C_{n+1}^2 |n\rangle \quad (13.8.15)$$

If  $C_{n+1}^2 - C_n^2 = 1$ , then

$$\hat{a} \hat{a}^\dagger - \hat{a}^\dagger \hat{a} = \hat{I} \quad (13.8.16)$$

Then above implies that  $C_n^2 = n$ , since  $C_0 = 0$ . Also,  $C_n = \sqrt{n}$ . In summary, for bosons,

$$\hat{a}_k^\dagger \hat{a}_l^\dagger - \hat{a}_l^\dagger \hat{a}_k^\dagger = 0 \quad (13.8.17)$$

$$\hat{a}_k \hat{a}_l - \hat{a}_l \hat{a}_k = 0 \quad (13.8.18)$$

$$\hat{a}_k \hat{a}_l^\dagger - \hat{a}_l^\dagger \hat{a}_k = \hat{I} \delta_{kl} \quad (13.8.19)$$

The above is derived without resorting to the use of the quantum harmonic oscillator, but the assumption that  $C_n$  is real and that  $C_n^2 = n$ . One can further conclude that

$$\hat{a}^\dagger \hat{a} |n\rangle = n |n\rangle \quad (13.8.20)$$

which is the number operator, which has been derived using a different approach.

### 13.9 Boson Field Operators

Similar to fermions, we can define field operators for bosons, so that

$$\hat{\psi}(\mathbf{r}) = \sum_j \hat{a}_j \phi_j(\mathbf{r}) \quad (13.9.1)$$

In general, for  $N$  particles,

$$\hat{\psi}(\mathbf{r}_1, \mathbf{r}_2, \dots, \mathbf{r}_N) = \frac{1}{\sqrt{N!}} \sum_{n_1, n_2, \dots, n_N} \hat{a}_{n_N} \dots \hat{a}_{n_1} \phi_{n_1}(\mathbf{r}_1) \phi_{n_2}(\mathbf{r}_2) \dots \phi_{n_N}(\mathbf{r}_N) \quad (13.9.2)$$

We can illustrate with a three-particle field operator:

$$\hat{\psi}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3) = \frac{1}{\sqrt{3!}} \sum_{n_1, n_2, n_3} \hat{a}_{n_3} \hat{a}_{n_2} \hat{a}_{n_1} \hat{\phi}(\mathbf{r}_1) \phi(\mathbf{r}_1) \phi(\mathbf{r}_2) \phi(\mathbf{r}_3) \quad (13.9.3)$$

When this operates on  $|1_l, 1_m, 1_n\rangle$  where the three particles are in different modes, we get (13.3.9) except that  $-$  signs are now replaced by  $+$  signs. If we have a state denoted by

$$|2_l, 1_m\rangle = \frac{1}{\sqrt{2!}} \left( \hat{a}_l^\dagger \right)^2 a_m |0\rangle \quad (13.9.4)$$



then

$$\begin{aligned} & \hat{\psi}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3)|2_l, 1_m\rangle \\ &= \frac{1}{\sqrt{2!3!}} [2\phi_l(\mathbf{r}_1)\phi_m(\mathbf{r}_2)\phi_l(\mathbf{r}_3) + 2\phi_l(\mathbf{r}_2)\phi_m(\mathbf{r}_3)\phi_l(\mathbf{r}_1) + 2\phi_l(\mathbf{r}_3)\phi_m(\mathbf{r}_1)\phi_l(\mathbf{r}_2)]|0\rangle \end{aligned} \quad (13.9.5)$$

It can be shown that the above is the correctly normalized wavefunction. If the state is

$$|3_l\rangle = \frac{1}{\sqrt{3!}} (\hat{a}_l^\dagger)^3 |0\rangle \quad (13.9.6)$$

then

$$\hat{\psi}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3)|3_l\rangle = \frac{6}{3!} \phi_l(\mathbf{r}_1)\phi_l(\mathbf{r}_2)\phi_l(\mathbf{r}_3)|0\rangle \quad (13.9.7)$$

The above is in fact normalized.

## 13.10 Additive One-Particle Operator

In this case, we can illustrate with the three-particle case for bosons, arriving at an expression similar to (13.5.8)

$$\hat{H} = \frac{1}{6} \sum_{n_2, n_3} \hat{a}_{n_3}^\dagger \hat{a}_{n_2}^\dagger \left[ \sum_{n'_1, n_1} \hat{a}_{n'_1}^\dagger H_{n'_1, n_1}^{(1)} \hat{a}_{n_1} \right] \hat{a}_{n_2} \hat{a}_{n_3} + \dots \quad (13.10.1)$$

where

$$H_{n'_1, n_1}^{(1)} = \langle \phi_{n'_1} | \hat{H}_{r_1} | \phi_{n_1} \rangle \quad (13.10.2)$$

When the three particles are in the  $|1_l, 1_m, 1_n\rangle$  state, (13.5.11) follows for bosons except that we replace  $-$  signs with  $+$  signs. Equation (13.5.13) follows similarly for bosons, so does (13.5.14) and (13.5.15). When two of the particles are in the same state as (13.9.4), we have

$$\hat{a}_{n_1} \hat{a}_{n_2} \hat{a}_{n_3} |2_l, 1_m\rangle = \sqrt{2!} (\delta_{ln_3} \delta_{mn_2} \delta_{ln_1} + \delta_{ln_3} \delta_{mn_1} \delta_{ln_2} + \delta_{ln_2} \delta_{mn_3} \delta_{ln_1}) |0\rangle \quad (13.10.3)$$

Similar to (13.5.13), we need to evaluate

$$\begin{aligned} & \sum_{n_2, n_3} \hat{a}_{n_3}^\dagger \hat{a}_{n_2}^\dagger \hat{a}_{n_1} \hat{a}_{n_2} \hat{a}_{n_3} |2_l, 1_m\rangle \\ &= \sqrt{2!} \left( 2\delta_{ln_1} \hat{a}_l^\dagger \hat{a}_m^\dagger + \delta_{mn_1} \hat{a}_l^\dagger \hat{a}_l^\dagger \right) |0\rangle \\ &= \sqrt{2!} \hat{a}_{n_1} \left( \hat{a}_l^\dagger \right)^2 \hat{a}_m^\dagger |0\rangle = (2!) \hat{a}_{n_1} |2_l, 1_m\rangle \end{aligned} \quad (13.10.4)$$

The first equality is established using (13.10.3), while the second equality is obtained by working backward using commutation relations to expand  $\hat{a}_{n_1} \left( \hat{a}_l^\dagger \right)^2 \hat{a}_m^\dagger |0\rangle$ . When the three particles are in the same state as indicated by (13.9.6), we have

$$\hat{a}_{n_1} \hat{a}_{n_2} \hat{a}_{n_3} |3_l\rangle = (3!)^{3/2} \delta_{ln_3} \delta_{ln_2} \delta_{ln_1} |0\rangle \quad (13.10.5)$$

and

$$\begin{aligned} & \sum_{n_2 n_3} \hat{a}_{n_3}^\dagger \hat{a}_{n_2}^\dagger \hat{a}_{n_1} \hat{a}_{n_2} \hat{a}_{n_3} |3_l\rangle \\ &= (3!)^{3/2} \delta_{ln_1} \left(\hat{a}_l^\dagger\right)^2 |0\rangle = \frac{1}{3} (3!)^{3/2} \hat{a}_{n_1} \left(\hat{a}_l^\dagger\right)^3 |0\rangle = (2!) \hat{a}_{n_1} |3_l\rangle \end{aligned} \quad (13.10.6)$$

In general,

$$\sum_{n_2, n_3} \hat{a}_{n_3}^\dagger \hat{a}_{n_2}^\dagger \hat{a}_{n_1} \hat{a}_{n_2} \hat{a}_{n_3} |\psi_{3p}\rangle = (2!) \hat{a}_{n_1} |\psi_{3p}\rangle \quad (13.10.7)$$

where  $|\psi_{3p}\rangle$  is the three-particle Fock state of either

$$|1_l, 1_m, 1_n\rangle, |2_l, 1_m\rangle, |3_l\rangle \quad (13.10.8)$$

or other combinations. The proof for the rest follows that of fermion particles of (13.5.14) to (13.5.17). In general, the Hamiltonian for  $N$  bosons is given by

$$\hat{H} = \sum_{n', n} H_{n', n} \hat{a}_{n'}^\dagger \hat{a}_n \quad (13.10.9)$$

For the diagonal case, it becomes

$$\hat{H} = \sum_n E_n \hat{a}_n^\dagger \hat{a}_n \quad (13.10.10)$$

## 13.11 The Difference between Boson Field and Photon Field

The above is the Hamiltonian for a boson field or gas in Fock space. It looks strikingly similar to the Hamiltonian for photon field or electromagnetic field. However, it is to be noted that a boson field is quite different from the field of  $N$  photons.

1. A coordinate  $\mathbf{r}_i$  can be associated with every boson particle, but not with every photon;
2. All photons share the same electromagnetics field function in its field operator representation, but  $N$  bosons will have product of  $N$  one-particle field operator;
3. The Hamiltonian of a boson gas has no zero-point energy, but for photons, there exists the zero-point energy of  $\frac{1}{2} \hbar \omega$ .

A photon should be thought of as a quantized energy packet due to coupled quantum harmonic oscillation of “space” or “vacuum”, similar to the case of a phonon, which is coupled quantum harmonic oscillation of crystalline lattice.

## Chapter 14

# Interaction of Different Particles

### 14.1 Introduction

In the previous chapter, it was shown that the Hamiltonian of a many particle system can be written succinctly using Fock space representation. This opens up the study of complex systems, such as excitons that involve the interaction of electron-hole pair, as well as superconductivity that involves the interaction of electrons with phonons. In this chapter, we will study the interaction of the electron of an atom with photons of a cavity using the many-particle formalism.

### 14.2 Interaction of Particles

Say if we have two quantum systems that are initially non-interacting, their individual quantum systems can be described by eigenstates that satisfy

$$\hat{H}_1 |\psi_1\rangle = E_1 |\psi_1\rangle \quad (14.2.1)$$

$$\hat{H}_2 |\psi_2\rangle = E_2 |\psi_2\rangle \quad (14.2.2)$$

Even though the two systems are entirely non-interacting with each other, nevertheless, we can combine the two systems together and write

$$\hat{H}_0 |\psi\rangle = \left( \hat{H}_1 + \hat{H}_2 \right) |\psi_1\rangle |\psi_2\rangle = (E_1 + E_2) |\psi_1\rangle |\psi_2\rangle = E |\psi\rangle \quad (14.2.3)$$

where  $|\psi_1\rangle |\psi_2\rangle$  represents a state in the direct product space. It can be used to represent the eigenstates of the combined system. Equation (14.2.3) is entirely equivalent to (14.2.1) and (14.2.2). In the above,  $\hat{H}_1$  acts only on  $|\psi_1\rangle$  and  $\hat{H}_2$  acts on  $|\psi_2\rangle$ . The above allows us to add an interaction term between system 1 and system 2, and the same direct product space can be used to span the solution space. By using Fock space representation, the above Hamiltonians can be written in terms of annihilation and creation operators for a system

consisting of fermions and bosons,

$$\hat{H}_1 = \sum_j E_j \hat{b}_j^\dagger \hat{b}_j, \quad \hat{H}_2 = \sum_\lambda \hbar\omega_\lambda \hat{a}_\lambda^\dagger \hat{a}_\lambda \quad (14.2.4)$$

where  $\hat{H}_1$  can be the Hamiltonian for electrons, and  $\hat{H}_2$  can be the Hamiltonian for photons where the  $\frac{1}{2}\hbar\omega_\lambda$  term or the zero-point energy has been ignored. This term will just introduce a phase shift in the solution.

When electric dipole interaction exists between the electron of an atom and the electromagnetic field, the interaction Hamiltonian may be written as

$$\hat{H}_{ed,\mathbf{r}} = e\mathbf{E} \cdot \mathbf{r} \quad (14.2.5)$$

The above has been added as a perturbation to the unperturbed equation (14.2.3) where the electric field is treated classically in a previous chapter. We have enough knowledge now to treat both  $\mathbf{E}$  and  $\mathbf{r}$  quantum mechanically. For the electron in the  $i$ -th atom, the above becomes

$$\hat{H}_{ed,\mathbf{r}_i} = ie \sum_\lambda (\hat{a}_\lambda - \hat{a}_\lambda^\dagger) \sqrt{\frac{\hbar\omega}{2\epsilon_0}} \mathbf{u}_\lambda(\mathbf{r}_i) \cdot \mathbf{r}_i \quad (14.2.6)$$

We have used a quantized cavity mode electric field in (14.2.5) to arrive at the above. Also,  $\mathbf{r}$  in coordinate space representation remains unchanged. When  $N$  atoms are present,

$$\hat{H}_{ed,\mathbf{r}} = \sum_{i=1}^N ie \sum_\lambda (\hat{a}_\lambda - \hat{a}_\lambda^\dagger) \sqrt{\frac{\hbar\omega}{2\epsilon_0}} \mathbf{u}_\lambda(\mathbf{r}_i) \cdot \mathbf{r}_i \quad (14.2.7)$$

A similarity transform of the above Hamiltonian can be performed with the  $N$  particle field operator to yield

$$\hat{H}_{ed} = \int \hat{\psi}_{NP}^\dagger(\mathbf{r}_1, \dots, \mathbf{r}_N) \hat{H}_{ed,\mathbf{r}} \hat{\psi}_{NP}(\mathbf{r}_1, \dots, \mathbf{r}_N) d\mathbf{r}_1 \cdots d\mathbf{r}_N \quad (14.2.8)$$

As mentioned in the previous chapter, since only one-particle interaction is involved, the above can be transformed with the one-particle field operator yielding

$$\hat{H}_{ed} = \sum_{j,k,\lambda} H_{ed,\lambda,j,k} \hat{b}_j^\dagger \hat{b}_k (\hat{a}_\lambda - \hat{a}_\lambda^\dagger) \quad (14.2.9)$$

where

$$H_{ed,\lambda,j,k} = ie \sqrt{\frac{\hbar\omega}{2\epsilon_0}} \int d\mathbf{r} \phi_j^*(\mathbf{r}) \mathbf{u}_\lambda(\mathbf{r}) \cdot \mathbf{r} \phi_k(\mathbf{r}) \quad (14.2.10)$$

### 14.3 Time-Dependent Perturbation Theory

In the time-dependent perturbation theory, we first seek the solution of the unperturbed system. The eigenstate of the unperturbed system is defined as

$$\hat{H}_0 |N_{fm}; N_{bm}\rangle = E_m |N_{fm}; N_{bm}\rangle \quad (14.3.1)$$

where  $|N_{fm}; N_{bm}\rangle$  stands for a vector in the direct product space of fermions and bosons:  $N_{fm}$  stands for the  $m$ -th state of the fermions, while  $N_{bm}$  stands for the  $m$ -th state of the bosons.

When the system is perturbed, we denote the solution  $|\psi\rangle$  as a linear superposition of the eigenstates of the unperturbed solution. That is

$$|\psi\rangle = \sum_m c_m(t) e^{-i\omega_m t} |N_{fm}; N_{bm}\rangle \quad (14.3.2)$$

where  $\hbar\omega_m = E_m$ . The state  $|\psi\rangle$  evolves in time according to the equation

$$i\hbar \frac{\partial}{\partial t} |\psi\rangle = (\hat{H}_0 + \hat{H}_p) |\psi\rangle \quad (14.3.3)$$

Applying the time-dependent perturbation method as we have done previously, and testing the equation with  $\langle N_{fq}; N_{bq}|$  gives

$$\dot{c}_q^{(1)}(t) = \frac{1}{i\hbar} \sum_m c_m^{(0)} e^{-i(\omega_m - \omega_q)t} \langle N_{fq}; N_{bq} | \hat{H}_p | N_{fm}; N_{bm} \rangle \quad (14.3.4)$$

If the starting state is assumed such that  $c_s^{(0)} = 1$ , and all other states are zero, then the above becomes

$$\dot{c}_q^{(1)}(t) = \frac{1}{i\hbar} e^{i(\omega_q - \omega_s)t} \langle N_{fq}; N_{bq} | \hat{H}_p | N_{fs}; N_{bs} \rangle \quad (14.3.5)$$

with

$$\hat{H}_p = \hat{H}_{ed} = \sum_{j,k,\lambda} H_{ed,\lambda,j,k} \hat{b}_j^\dagger \hat{b}_k \left( \hat{a}_\lambda - \hat{a}_\lambda^\dagger \right) \quad (14.3.6)$$

### 14.3.1 Absorption

The electron in an atom is assumed to have two states: a lower energy state  $E_1$ , and an upper energy state  $E_2$ . We assume a photon in mode  $\lambda_1$  in the cavity. Then the starting state can be written as

$$|N_{fs}; N_{bs}\rangle = \hat{b}_1^\dagger \hat{a}_{\lambda_1}^\dagger |0\rangle \quad (14.3.7)$$

with energy

$$E_s = E_1 + \hbar\omega_{\lambda_1} \quad (14.3.8)$$

Consequently,

$$\hat{H}_p |N_{fs}; N_{bs}\rangle = \sum_{j,k,\lambda} H_{ed,\lambda,j,k} \hat{b}_j^\dagger \hat{b}_k \left( \hat{a}_\lambda - \hat{a}_\lambda^\dagger \right) \hat{b}_1^\dagger \hat{a}_{\lambda_1}^\dagger |0\rangle \quad (14.3.9)$$

It can be shown that

$$\hat{b}_j^\dagger \hat{b}_k \left( \hat{a}_\lambda - \hat{a}_\lambda^\dagger \right) \hat{b}_1^\dagger \hat{a}_{\lambda_1}^\dagger |0\rangle = \delta_{k1} \delta_{\lambda\lambda_1} \hat{b}_j^\dagger |0\rangle - \delta_{k1} \hat{b}_j^\dagger \hat{a}_\lambda^\dagger \hat{a}_{\lambda_1}^\dagger |0\rangle \quad (14.3.10)$$

In order for  $\langle N_{fq}; N_{bq} | \hat{H}_p | N_{fs}; N_{bs} \rangle$  to be nonzero, it is necessary that  $|N_{fq}; N_{bq}\rangle$  contains  $\hat{b}_j^\dagger |0\rangle$  or  $\hat{b}_j^\dagger \hat{a}_\lambda^\dagger \hat{a}_{\lambda_1}^\dagger |0\rangle$ . If

$$|N_{fq}; N_{bq}\rangle = \hat{b}_j^\dagger |0\rangle \quad (14.3.11)$$

then  $E_q = E_j$ , and

$$\dot{c}_q^{(1)}(t) = \frac{1}{i\hbar} e^{i(\omega_j - \omega_1 - \omega_{\lambda_1})t} \sum_{k,\lambda} H_{ed,\lambda,j,k} \delta_{k1} \delta_{\lambda\lambda_1} \langle 0 | \hat{b}_j \hat{b}_j^\dagger | 0 \rangle \quad (14.3.12)$$

$$= \frac{1}{i\hbar} e^{i(\omega_j - \omega_1 - \omega_{\lambda_1})t} H_{ed,\lambda_1,j,1} \quad (14.3.13)$$

where  $\hbar\omega_j = E_j$ . The above is an oscillatory function of  $t$  unless  $\omega_j - \omega_1 - \omega_{\lambda_1} = 0$ . When this happens,  $\dot{c}_q^{(1)}(t)$  will integrate to a large value, indicating the high likelihood of transition to his eigenstate. Therefore, for this to happen, we need

$$E_2 - E_1 = \hbar\omega_{\lambda_1} \quad (14.3.14)$$

In this process, we start with the electron in  $E_1$  and a photon in the cavity, and end up with the electron in  $E_2$  and no photon as indicated by (14.3.11).

If we integrate the above equation, a sinc function results, which can be approximated by a delta function for long integration time, as was done previously. One can show that

$$|c_q^{(1)}|^2 = \frac{2\pi}{\hbar} t_0 |H_{ed\lambda_1,j,1}|^2 \delta(E_2 - E_1 - \hbar\omega_{\lambda_1}) \quad (14.3.15)$$

One can derive the transition rate as before to arrive at

$$w_q = \frac{2\pi}{\hbar} |H_{ed\lambda_1,j,1}|^2 \delta(E_2 - E_1 - \hbar\omega_{\lambda_1}) \quad (14.3.16)$$

In the above, the other possibility is for the final state to be

$$|N_{fq}; N_{bq}\rangle = \hat{b}_j^\dagger \hat{a}_\lambda^\dagger \hat{a}_{\lambda_1}^\dagger |0\rangle \quad (14.3.17)$$

with energy

$$E_q = E_j + \hbar\omega_\lambda + \hbar\omega_{\lambda_1} \quad (14.3.18)$$

Hence

$$E_q - E_s = E_j - E_1 + \hbar\omega_\lambda \quad (14.3.19)$$

and  $E_q - E_s = 0$  in order for transition to occur or  $c_q^{(1)}$  to be large. However, (14.3.19) cannot be zero. If  $E_j = E_2$ ,  $E_q - E_s > 0$  always. If  $E_j = E_1$ , then  $\hbar\omega_\lambda > 0$ . What this means is that it is not possible to start with an electron in the ground state and a photon in the cavity to end up with the electron in the excited state with the emission of a photon.

## 14.4 Spontaneous Emission

In the case of spontaneous emission, we assume that the starting state of the electron is in the excited state with energy  $E_2$ . Then,  $E_s = E_2$  initially and

$$|N_{fs}; N_{bs}\rangle = \hat{b}_2^\dagger |0\rangle \quad (14.4.1)$$

Then

$$\hat{H}_p |N_{fs}; N_{bs}\rangle = \sum_{j,k,\lambda} H_{ed,\lambda,j,k} \hat{b}_j^\dagger \hat{b}_k (\hat{a}_\lambda - \hat{a}_\lambda^\dagger) \hat{b}_2^\dagger |0\rangle \quad (14.4.2)$$

The above can be reduced to

$$\hat{H}_p |N_{fs}; N_{bs}\rangle = - \sum_{j,k,\lambda} H_{ed,\lambda,j,k} \delta_{2k} \hat{b}_j^\dagger \hat{a}_\lambda^\dagger |0\rangle \quad (14.4.3)$$

Therefore, for  $\langle N_{fq}; N_{bq} | \hat{H}_p |N_{fs}; N_{bs}\rangle$  to be nonzero, we need

$$|N_{fq}; N_{bq}\rangle = \hat{b}_j^\dagger \hat{a}_\lambda^\dagger |0\rangle \quad (14.4.4)$$

with

$$E_q = E_j + \hbar\omega_\lambda \quad (14.4.5)$$

Therefore

$$E_q - E_s = E_j - E_2 + \hbar\omega_\lambda \quad (14.4.6)$$

In order for a sizeable  $c_q^{(1)}$ , we need  $E_j = E_1$ , and then

$$E_2 - E_1 = \hbar\omega_\lambda \quad (14.4.7)$$

Then electron starts with the excited state  $E_2$ , spontaneously emits a photon, and drops to lower state  $E_1$ . Then energy of the emitted photon satisfies (14.4.7) by energy conservation.

## 14.5 Stimulated Emission

In this case, the electron in an atom is in the excited state. The presence of a photon in the cavity stimulates the emission of another photon from the electron. The initial state is

$$|N_{fs}; N_{bs}\rangle = \hat{b}_2^\dagger \hat{a}_{\lambda_1}^\dagger |0\rangle \quad (14.5.1)$$

with

$$E_s = E_2 + \hbar\omega_{\lambda_1} \quad (14.5.2)$$

We can show that

$$\hat{H}_p |N_{fs}; N_{bs}\rangle = \sum_{j,k,\lambda} H_{ed,\lambda,j,k} \hat{b}_j^\dagger \hat{b}_k \left( \hat{a}_\lambda - \hat{a}_\lambda^\dagger \right) \hat{b}_2^\dagger \hat{a}_{\lambda_1}^\dagger |0\rangle \quad (14.5.3)$$

$$= \sum_{j,k,\lambda} H_{ed,\lambda,j,k} \delta_{k2} \left( \delta_{\lambda\lambda_1} \hat{b}_j^\dagger - \hat{b}_j^\dagger \hat{a}_\lambda^\dagger \hat{a}_{\lambda_1}^\dagger \right) |0\rangle \quad (14.5.4)$$

In order for the above to transition to the final state, or one requires a non-zero result for

$$\langle N_{fq}; N_{bq} | \hat{H}_p | N_{fs}; N_{bs} \rangle$$

it is necessary that

$$|N_{fq}; N_{bq}\rangle = \hat{b}_j^\dagger \hat{a}_\lambda^\dagger \hat{a}_{\lambda_1}^\dagger |0\rangle \quad (14.5.5)$$

with energy

$$E_q = E_j + \hbar\omega_\lambda + \hbar\omega_{\lambda_1} \quad (14.5.6)$$

or

$$E_q - E_s = E_j - E_2 + \hbar\omega_\lambda = 0 \quad (14.5.7)$$

In other words, the only possibility is for  $E_j = E_1$ , yielding

$$E_2 - E_1 = \hbar\omega_\lambda \quad (14.5.8)$$

The above is just the spontaneous emission of a photon with the above energy, regardless if we already have a photon with energy  $\hbar\omega_{\lambda_1}$  in the cavity.

Next we consider the case when  $\lambda = \lambda_1$ . Then

$$|N_{fq}; N_{bq}\rangle = \frac{1}{\sqrt{2!}} \hat{b}_j^\dagger \left( \hat{a}_{\lambda_1}^\dagger \right)^2 |0\rangle \quad (14.5.9)$$

with

$$E_q = E_j + 2\hbar\omega_{\lambda_1} \quad (14.5.10)$$

Consequently,

$$\begin{aligned} \langle N_{fq}; N_{bq} | \hat{H}_p | N_{fs}; N_{bs} \rangle &= H_{ed,\lambda_1,j,2} \langle 0 | \frac{1}{\sqrt{2!}} (\hat{a}_\lambda)^\dagger \hat{b}_j^\dagger \left( \hat{a}_{\lambda_1}^\dagger \right)^2 |0\rangle \\ &= \sqrt{2!} H_{ed,\lambda_1,j,2} \langle 0 | \frac{1}{\sqrt{2!}} (\hat{a}_\lambda)^\dagger \hat{b}_j^\dagger \frac{1}{\sqrt{2!}} \left( \hat{a}_{\lambda_1}^\dagger \right)^2 |0\rangle \\ &= \sqrt{2} H_{ed,\lambda_1,j,2} \end{aligned} \quad (14.5.11)$$

The  $\sqrt{2}$  factor is important implying that the transition is two times more likely to occur compared to the previous case. This is peculiar to stimulated emission where the emission of a photon is enhanced by the presence of a photon of the same frequency.



## 14.6 Multi-photon Case

In the multi-photon case, the transition rate for stimulated emission can be shown to be

$$w_q = \frac{2\pi}{\hbar} (n_{\lambda_1} + 1) |H_{ed,\lambda_1,1,2}|^2 \delta(E_1 - E_2 + \hbar\omega_{\lambda_1}) \quad (14.6.1)$$

implying that the presence of  $n_{\lambda_1}$  photons in the cavity enhances the emission by  $(n_{\lambda_1} + 1)$  times.

The spontaneous emission, however, is not affected by the presence of photons of other frequencies in the cavity. For the absorption case, it can be shown that the formula is

$$w_q = \frac{2\pi}{\hbar} n_{\lambda_1} |H_{ed,\lambda_1,1,2}|^2 \delta(E_2 - E_1 - \hbar\omega_{\lambda_1}) \quad (14.6.2)$$

## 14.7 Total Spontaneous Emission Rate

When an electron emits a photon into the cavity, there are many modes with the same frequency that the emission can occur. In general, the total spontaneous emission rate is

$$W_{\text{spont}} = \sum_q w_q = \frac{2\pi}{\hbar} \sum_{\lambda} |H_{ed,\lambda,1,2}|^2 \delta(E_1 - E_2 + \hbar\omega_{\lambda}) \quad (14.7.1)$$

In the above

$$\begin{aligned} H_{ed,\lambda,1,2} &= ie \sqrt{\frac{\hbar\omega_{\lambda}}{2\epsilon_0}} \int \phi_j^*(\mathbf{r}) [\mathbf{u}_{\lambda}(\mathbf{r}) \cdot \mathbf{r}] \phi_k(\mathbf{r}) d\mathbf{r} \\ &\simeq ie \sqrt{\frac{\hbar\omega_{\lambda}}{2\epsilon_0}} \mathbf{u}_{\lambda}(\mathbf{r}_0) \cdot \mathbf{r}_{jk} \end{aligned} \quad (14.7.2)$$

where

$$\mathbf{r}_{jk} = \int \phi_j^*(\mathbf{r}) \mathbf{r} \phi_k(\mathbf{r}) d\mathbf{r} \quad (14.7.3)$$

We have assumed here that  $\mathbf{u}_{\lambda}(\mathbf{r}_0)$  is slowly varying compared to  $\phi_l(\mathbf{r})$ ,  $l = j, k$ , and that  $\phi_l(\mathbf{r})$  is highly localized around an atom.

The modes in the cavity can be made countable by imposing a periodic boundary condition arriving at

$$\mathbf{u}_{\lambda}(\mathbf{r}) = \mathbf{e} \frac{1}{\sqrt{V_b}} e^{i\mathbf{k} \cdot \mathbf{r}} \quad (14.7.4)$$

The summation over different electromagnetic modes then becomes

$$\sum_{\lambda} \rightarrow \sum_{\text{pol}} \sum_{\mathbf{k}} \rightarrow \sum_{\text{pol}} \sum_{\mathbf{k}} \frac{V_b}{(2\pi)^3} \Delta\mathbf{k} = \sum_{\text{pol}} \int \frac{V_b}{(2\pi)^3} d\mathbf{k} \quad (14.7.5)$$

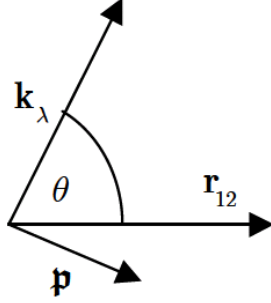


Figure 14.1: The vectors in the polarization of the emitted photon are aligned to simplify the calculation (from DAB Miller).

We can also pick the polarization such that one of them is always orthogonal to  $\mathbf{r}_{12}$ , so that

$$\mathbf{u}(\mathbf{r}_0) \cdot \mathbf{r}_{12} = u_{\lambda}(\mathbf{r}_0) r_{12} \sin \theta \quad (14.7.6)$$

Consequently,

$$W_{\text{spont}} = \frac{2\pi}{\hbar} \int \frac{V_b}{(2\pi)^3} \left| i e \sqrt{\frac{\hbar\omega_k}{2\epsilon_0}} \frac{1}{\sqrt{V_b}} e^{i\mathbf{k}\cdot\mathbf{r}_0} r_{12} \sin \theta \right|^2 \delta(E_1 - E_2 + \hbar\omega_k) d\mathbf{k} \quad (14.7.7)$$

or

$$\begin{aligned} W_{\text{spont}} &= \frac{e^2 |r_{12}|^2}{8\pi^2 \epsilon_0} \int \omega_k \sin^2 \theta \delta(E_1 - E_2 + \hbar\omega_k) d\mathbf{k} \\ &= \frac{e^2 |r_{12}|^2}{8\pi^2 \epsilon_0} \int_{k=0}^{\infty} \int_{\theta=0}^{\pi} \omega_k \delta(E_1 - E_2 + \hbar\omega_k) 2\pi \sin^3 \theta k^2 d\theta dk \end{aligned} \quad (14.7.8)$$

In the above,  $\omega_k = ck$ , or  $\hbar ck = \hbar\omega_k$ . Then

$$W_{\text{spont}} = \frac{e^2 |r_{12}|^2}{4\pi \epsilon_0 c^3 \hbar^4} \int_{k=0}^{\infty} \int_{\theta=0}^{\pi} (\hbar\omega_k) \delta(E_1 - E_2 + \hbar\omega_k) \sin^3 \theta d\theta (\hbar\omega_k)^2 d(\hbar\omega_k) \quad (14.7.9)$$

Since

$$\int_{\theta=0}^{\pi} \sin^3 \theta d\theta = - \int_{-1}^1 (1 - \cos^2 \theta) d \cos \theta = \frac{4}{3} \quad (14.7.10)$$

then,

$$W_{\text{spont}} = \frac{e^2 |r_{12}|^2 \omega_{12}^3}{3\pi \epsilon_0 \hbar c^3} \quad (14.7.11)$$

where  $\hbar\omega_{12} = E_2 - E_1$ . The life time of a state is then  $\tau = W_{\text{spont}}^{-1}$ .

## Chapter 15

# Quantum Information and Quantum Interpretation

### 15.1 Introduction

One important tenet of quantum mechanics is that one does not know what the state of the system is until one performs a measurement. After the measurement, the system collapses to the state that is “discovered” by the measurement. Before the measurement, the quantum system is described by a state that is in a linear superposition of different states. It is the existence as a linear superposition of states that greatly enriches the information content of a quantum system. The possibility of a system to be simultaneously in different states is peculiar to quantum mechanics. Objects in the classical world cannot be in such a state. It puts quantum systems in the realm of “ghosts” and “angels” where in fairy tales and ghost and angel stories of different cultures, they can be simultaneously in many places or in different states. It is because of these “ghost-angel” states, that we can have quantum cryptography, quantum communication, quantum circuits, and quantum computing.

### 15.2 Quantum Cryptography

Quantum cryptography can be used for secure quantum communication. It is secure because a quantum state cannot be replicated without disturbing the quantum state. Hence, whoever wants to replicate a quantum state to steal the data will be easily detected. The property of non-replication follows from the no-cloning theorem.

#### 15.2.1 No-cloning Theorem

First we assume that a quantum operator can be designed such that it can clone a quantum state without measuring it. But such a capability will violate the principle of linear superposition, as we shall show. Therefore, such an operator cannot be designed. This is a proof by contradiction.

A quantum system is described by a quantum state that evolves from an initial state to a final state following the laws of quantum mechanics. An example of such an evolutionary operator is

$$\hat{T} = e^{-i\frac{\hat{H}}{\hbar}t} \quad (15.2.1)$$

The above quantum operator is a unitary operator as well as a linear operator. First, assume that it has the capability of replicating a quantum state in system 2 to be identical to the state in system 1 after acting on such a quantum system. It does so without altering the quantum state of system 1, e.g., by a measurement. We denote this by

$$\hat{T}|\Psi_s\rangle_2|\Psi_a\rangle_1 = |\Psi_a\rangle_2|\Psi_a\rangle_1 \quad (15.2.2)$$

By the same token, it should replicate

$$\hat{T}|\Psi_s\rangle_2|\Psi_b\rangle_1 = |\Psi_b\rangle_2|\Psi_b\rangle_1 \quad (15.2.3)$$

Now if the state to be replicated is

$$|\Psi_c\rangle_1 = \frac{1}{\sqrt{2}} [|\Psi_a\rangle_1 + |\Psi_b\rangle_1] \quad (15.2.4)$$

Then

$$\begin{aligned} \hat{T}|\Psi_s\rangle_2|\Psi_c\rangle_1 &= \frac{1}{\sqrt{2}} \left[ \hat{T}|\Psi_s\rangle_2|\Psi_a\rangle_1 + \hat{T}|\Psi_s\rangle_2|\Psi_b\rangle_1 \right] \\ &= \frac{1}{\sqrt{2}} [|\Psi_a\rangle_2|\Psi_a\rangle_1 + |\Psi_b\rangle_2|\Psi_b\rangle_1] \neq |\Psi_c\rangle_2|\Psi_c\rangle_1 \end{aligned} \quad (15.2.5)$$

Clearly, the above violates the principle of linear superposition if the last equality is true. Hence, such a cloning operator is impossible due to the violation of the principle of linear superposition. The above proves the no-cloning theorem.

## 15.2.2 Entangled States

A multi-mode photon can be described by the following field operator

$$\mathbf{E}(\mathbf{r}) = \sum_{\mathbf{k},s} \sqrt{\frac{\hbar\omega_k}{2V\epsilon_0}} \mathbf{e}_s \hat{a}_{\mathbf{k},s} e^{i\mathbf{k}\cdot\mathbf{r} - i\omega t} + \text{c.c} \quad (15.2.6)$$

The single photon state can be denoted by

$$|\psi\rangle = |1_{\mathbf{k},v}\rangle \quad (15.2.7)$$

The above denotes a photon in a pure  $\mathbf{k}$  state with polarization  $v$ . It has a packet of energy  $E = \hbar\omega_k$ . However, a photon with a pure  $\mathbf{k}$  state is not localized. But a photon generated by a source like an atomic transition must be causal, and hence, localized. A localized wave packet describing this photon field can be constructed by using a linear superposition of wavenumber

$\mathbf{k}$  or frequencies. For high frequency photons, this localized state can have a center frequency with a small spread of frequencies around the center frequency. The single-photon Fock state can hence be written as

$$|\psi\rangle = \sum_{\mathbf{k}} c_{\mathbf{k}} |1_{\mathbf{k},v}\rangle \quad (15.2.8)$$

For quasi-monochromatic photons, the above will be dominated by one term and we can denote this photon approximately with the state vector (15.2.7).

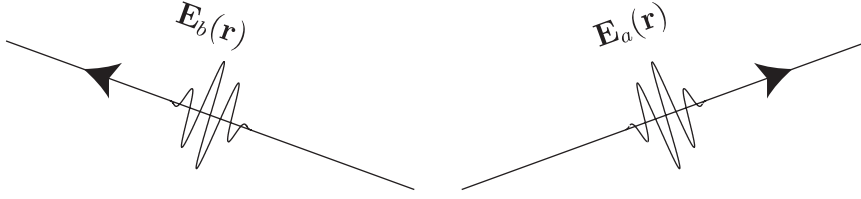


Figure 15.1: Two photons traveling in different directions.

With the above picture in mind, we can think of two localized photons traveling in different directions. A direct product space can be used to represent the state of these two photons:

$$|\psi\rangle_{ab} = |1_{\mathbf{k}_a,v}\rangle_a |1_{\mathbf{k}_b,v}\rangle_b \quad (15.2.9)$$

where we have assumed quasi-mono-chromatic photons. For simplicity, we denote a two-photon state as

$$|\psi\rangle_{12} = |V\rangle_1 |V\rangle_2 \quad (15.2.10)$$

If the two photons are generated from the same source, entangled photon states may result. Entangled two-particle states are those that cannot be written as a product (outer product or tensor product) of simpler states. An example is the EPR (Einstein, Podolsky, and Rosen) pair

$$|\Phi^+\rangle_{12} = \frac{1}{\sqrt{2}} (|H\rangle_1 |H\rangle_2 + |V\rangle_1 |V\rangle_2) \quad (15.2.11)$$

Other entangled states are

$$|\Phi^-\rangle_{12} = \frac{1}{\sqrt{2}} (|H\rangle_1 |H\rangle_2 - |V\rangle_1 |V\rangle_2) \quad (15.2.12)$$

$$|\Psi^+\rangle_{12} = \frac{1}{\sqrt{2}} (|H\rangle_1 |V\rangle_2 + |V\rangle_1 |H\rangle_2) \quad (15.2.13)$$

$$|\Psi^-\rangle_{12} = \frac{1}{\sqrt{2}} (|H\rangle_1 |V\rangle_2 - |V\rangle_1 |H\rangle_2) \quad (15.2.14)$$

The above four states are also called the Bell states. They are usually generated due to the particle pair needing to satisfy conservation of angular momentum. For instance, two photons are generated by atomic transitions where the initial angular momentum of the system is zero.

An angular momentum conserving state with two counter-propagating photon is

$$|\Phi\rangle = \frac{1}{\sqrt{2}}(|R\rangle_1|R\rangle_2 + |L\rangle_1|L\rangle_2) \quad (15.2.15)$$

where  $R$  and  $L$  stand for right-handed and left-handed circular polarizations, respectively. But a change of basis allows the above to be written as

$$|\Phi\rangle = -\frac{1}{\sqrt{2}}(|H\rangle_1|H\rangle_2 - |V\rangle_1|V\rangle_2) \quad (15.2.16)$$

one of the Bell states.

The entangled states are bewildering because it means that for two counter-propagating photons, if one measures one photon is in  $|H\rangle$  state, the other photon immediately collapses to an  $|H\rangle$  state as well, regardless of how far apart the two photons are. Whereas before the measurement, the photons are in a linear superposition of a  $|H\rangle$  and  $|V\rangle$  states.

### 15.2.3 A Simple Quantum Encryption Algorithm

Let us assume that Alice and Bob communicate by the use of simple photon polarizers. When Alice sends out a  $|V\rangle$  state photon, it represents a “1” and similarly, an  $|H\rangle$  state photon represents a “0”. If Bob receives with a similarly aligned polarizer, he receives the information correctly. (We call this the VH mode.) If Alice aligns her polarizer at  $45^\circ$  and Bob follows suit, he continues to receive the information correctly.

However, if Alice aligns her polarizer vertically, and Bob aligns his at  $45^\circ$ , the bit information received by him is only correct only 50% of the time. This is because the  $|+45\rangle$  state and  $|-45\rangle$  state are expressible as a linear superposition of the  $|H\rangle$  and  $|V\rangle$  states. Normally,

$$|+45\rangle = \frac{1}{\sqrt{2}}(|H\rangle + |V\rangle) \quad (15.2.17)$$

$$|-45\rangle = \frac{1}{\sqrt{2}}(|H\rangle - |V\rangle) \quad (15.2.18)$$

By the tenet of quantum measurement, these states are measured with equal likelihood of being  $|H\rangle$  or  $|V\rangle$ . Similarly, if Alice transmits with her polarizer aligned in the  $45^\circ$  angle, and Bob receives with polarizers with H and V polarization, he receives the polarization of  $|H\rangle$  and  $|V\rangle$  with equal likelihood (50% chance) regardless of what polarization Alice sends. This is because

$$|H\rangle = \frac{1}{\sqrt{2}}(|+45\rangle + |-45\rangle) \quad (15.2.19)$$

$$|V\rangle = \frac{1}{\sqrt{2}}(|+45\rangle - |-45\rangle) \quad (15.2.20)$$

Now if Alice decides to use her polarizer randomly, and Bob receives with the polarizer randomly so that they are equally likely to use VH mode or  $45^\circ$  mode. The probability that their polarizers are aligned is correctly 50%. During this time, they communicate with no error. The other 50% time, when their polarizers are misaligned, they communicate with 50% error. Hence, 25% of the data are wrong.

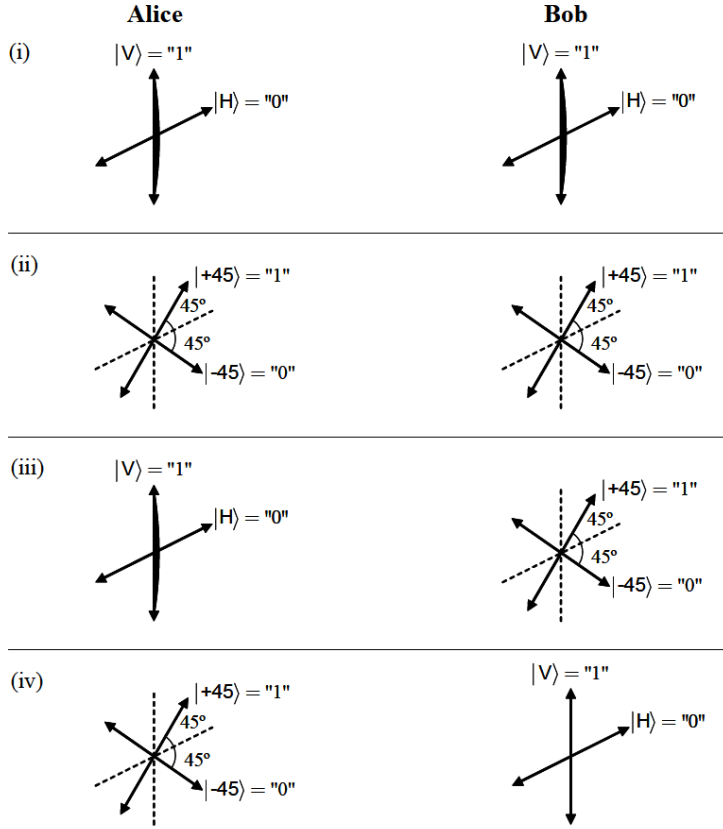


Figure 15.2: Communication between Alice and Bob using single-photon source and simplified polarizer measurement schemes (from DAB Miller).

After a preliminary quantum communication, Alice and Bob can communicate the information about the alignments of their polarizers, say, by a phone call. Bob will retain only the data when their polarizers are correctly aligned and discard the rest. For the preliminary communication, they can compare their data over the aligned case, and there should be error free in principle.

If an eavesdropper, Eve, attempts to steal the information, she does so by using a polarizer to intercept the photon. If Eve knows that Alice is using the VH mode, Eve aligns her polarizer in the VH mode. After Eve has received the data, she can retransmit the data to Bob, thus stealing the data. However, if the polarization used by Alice is random, Eve's polarizer is not aligned with Alice's half the time. Eve would have corrupted her data making the wrong transmission 50% of the time. This would increase the error in transmission of the data from Alice to Bob making it wrong 25% of the time. If Alice and Bob communicate by a phone call to check the security of their data transmission and found that it is wrong approximately

25% of the time, they would have suspected an eavesdropper.

Notice that Eve has to collapse the state of the photon sent out by Alice into one of the two states of Eve's polarizer before she can duplicate the photon and send it to Bob. Because of the no-cloning theorem, Eve cannot duplicate the state of the photon that Alice has sent without measuring it.

In a secure communication system, Alice and Bob do not send the real message in preliminary testing of the security of the channel. Alice will first send Bob the secret key and test if the channel is secure. If it is a secure channel, then she would send the rest of the information. The above is known as the BB84 protocol, attributed to Bennett and Brassard's work in 1984.

Also, notice that the above secure communication system does not work in a classical optical communication channel where a bunch of photons is sent. If a bunch of photons is sent by Alice to denote a V or an H polarization to send "1" and "0", when Eve eavesdrops with her misaligned polarizer by  $45^\circ$ , she would have noticed that equal number of photons are emerging from her two orthogonal polarizations. By checking the phase of the two streams of photons, she can easily duplicate a classical photon bunch and send it to Bob, meanwhile stealing the data off the communication channel. Hence, the security of the quantum communication channel comes from the interpretation of quantum mechanics: a particle is in the linear superposition of quantum states before the measurement. A measurement collapses the quantum state into one of the states "discovered" by the measurement.

The above discussion of a secure channel is based on ideal single-photon sources. In practice, non-ideality will give rise to communication errors. Quantum error correction schemes have been devised to minimize the errors in a quantum communication channel.

## 15.3 Quantum Computing

The distinguishing feature of quantum computing is quantum parallelism. Again, this follows from the tenet of quantum measurement. A quantum state can be in a linear superposition of many states before the measurement. After the measurement, the quantum state collapses to one of the quantum states. The prowess of quantum computing, as mentioned before, comes from the "ghost-angel" state of a quantum system.

### 15.3.1 Quantum Bits (Qubits)

A quantum bit or a qubit is a bit in a quantum state that is in the linear superposition of two states representing the  $|0\rangle$  bit and the  $|1\rangle$  bit. Namely,

$$|\psi\rangle = C_0 |0\rangle + C_1 |1\rangle \quad (15.3.1)$$

where  $|C_0|^2 + |C_1|^2 = 1$ . The two states  $|0\rangle$  and  $|1\rangle$  can be the vertical and horizontal polarization of a photon. Alternatively, it can be the up and down state of a spin, or any two energy levels of a multi-level system.

The richness of quantum information is already manifested in this very simple example. Unlike classical bits in classical computers, which can only have binary values, a qubit can have multitudes of possible values depending on the values of  $C_0$  and  $C_1$ . A two-state quantum



system, as has been shown in the spin case, can be represented by a Bloch sphere. Every point on the Bloch sphere represents a possible quantum state, depending on  $C_1$  and  $C_2$ , and there could be infinitely many states.

### 15.3.2 Quantum Gates

Analogous to classical logic gates, there are quantum gates that manipulate the  $|0\rangle$  and  $|1\rangle$  states of a qubit. A qubit as indicated by (15.3.1) can be represented by a column vector of length two. For example, the qubit in (15.3.1) can be represented by  $[C_0, C_1]^t$ . A quantum gate transforms the quantum state  $[C_0, C_1]^t$  to another state  $[C'_0, C'_1]^t$ . Such a matrix

$$\begin{bmatrix} C'_0 \\ C'_1 \end{bmatrix} = \begin{bmatrix} M_{11} & M_{12} \\ M_{21} & M_{22} \end{bmatrix} \begin{bmatrix} C_0 \\ C_1 \end{bmatrix} \tag{15.3.2}$$

has to be unitary since all quantum gates must operate by the time evolution according to

$$\hat{M} = e^{-i\frac{\hat{H}}{\hbar}t} \tag{15.3.3}$$

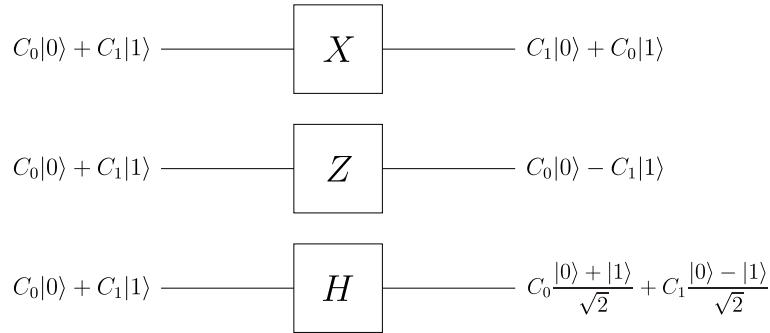


Figure 15.3: Single qubit gates showing their input and output states.

Examples of quantum gates, expressed in their matrix representations with a slight abuse of notation, are

$$\hat{X} = \begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix} \tag{15.3.4}$$

$$\hat{Z} = \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix} \tag{15.3.5}$$

$$\hat{H} = \frac{1}{\sqrt{2}} \begin{bmatrix} 1 & 1 \\ 1 & -1 \end{bmatrix} \tag{15.3.6}$$

$\hat{X}$  represents the NOT gate while  $\hat{Z}$  represents one that flips the sign of  $|1\rangle$  bit.  $\hat{H}$  is called the Hadamard gate that is almost like a “square root” gate.

When these quantum gates operate on the qubit denoted by (15.3.1), the results are as follows:

$$\hat{X}|\psi\rangle = C_1|0\rangle + C_0|1\rangle \quad (15.3.7)$$

$$\hat{Z}|\psi\rangle = C_0|0\rangle - C_1|1\rangle \quad (15.3.8)$$

$$\hat{H}|\psi\rangle = \frac{C_0}{\sqrt{2}}(|0\rangle + |1\rangle) + \frac{C_1}{\sqrt{2}}(|0\rangle - |1\rangle) \quad (15.3.9)$$

Moreover, one can show that

$$\hat{H}^2 = \hat{I} \quad (15.3.10)$$

When expressed in terms of matrix algebra,

$$\hat{X} \begin{bmatrix} C_0 \\ C_1 \end{bmatrix} = \begin{bmatrix} C_1 \\ C_0 \end{bmatrix} \quad (15.3.11)$$

$$\hat{Z} \begin{bmatrix} C_0 \\ C_1 \end{bmatrix} = \begin{bmatrix} C_0 \\ -C_1 \end{bmatrix} \quad (15.3.12)$$

$$\hat{H} \begin{bmatrix} C_0 \\ C_1 \end{bmatrix} = \begin{bmatrix} \frac{C_0+C_1}{\sqrt{2}} \\ \frac{C_0-C_1}{\sqrt{2}} \end{bmatrix} \quad (15.3.13)$$

In addition to one qubit gate, there are also two qubit gates. A very important one is the CNOT gate shown in Figure 15.4. Its transformation matrix is given by

$$\hat{U}_{\text{CNOT}} = \begin{bmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & 1 & 0 \end{bmatrix} \quad (15.3.14)$$

The above can also be implemented with a unitary transform.

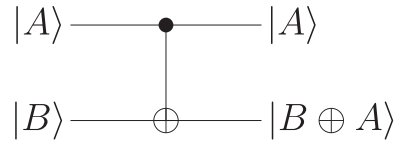


Figure 15.4: A quantum circuit of a two-qubit gate representing the CNOT gate. It takes two input streams, and has two output streams. The  $\oplus$  symbol represents an *exclusive or* operation, or addition modulus 2.

### 15.3.3 Quantum Computing Algorithms

As mentioned before, the most important aspect of quantum computing algorithm is quantum parallelism. We will illustrate this with the Deutsch algorithm. It can be implemented with the quantum circuit shown below:

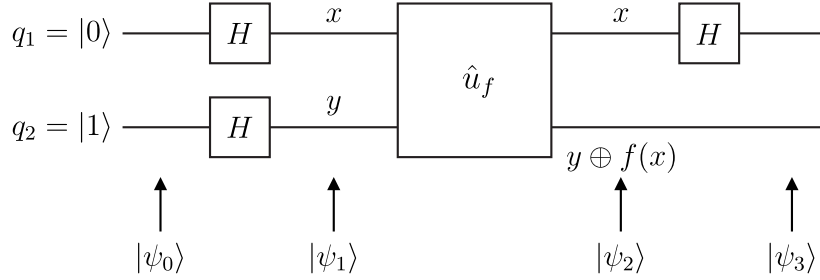


Figure 15.5: The implementation of the Deutsch algorithm with quantum gates and circuits.

We start with a state

$$|\psi_0\rangle = |0, 1\rangle \quad (15.3.15)$$

At stage  $|\psi_1\rangle$ , we have

$$x = \hat{H} \cdot q_1 = \frac{1}{\sqrt{2}}(|0\rangle + |1\rangle) \quad (15.3.16)$$

$$y = \hat{H} \cdot q_2 = \frac{1}{\sqrt{2}}(|0\rangle - |1\rangle) \quad (15.3.17)$$

Hence,

$$|\psi_1\rangle = \frac{1}{2} (|0\rangle + |1\rangle) (|0\rangle - |1\rangle) = \frac{1}{2} (|0, 0\rangle - |0, 1\rangle + |1, 0\rangle - |1, 1\rangle) \quad (15.3.18)$$

The unitary operator  $\hat{u}_f$  has no effect on the  $x$  qubit, but performs the operation  $y \oplus f(x)$  on the  $y$  qubit. The function  $f(x)$  takes input  $x$  which is either 0 or 1. It produces an output which is either constant or balanced, but the output is either 0 or 1.

Hence, according to the rules above,

$$|\psi_2\rangle = \frac{1}{2} (|0, f(0)\rangle - |0, 1 \oplus f(0)\rangle + |1, f(1)\rangle - |1, 1 \oplus f(1)\rangle) \quad (15.3.19)$$

If  $f(x)$  is a constant function, then  $f(0) = f(1)$ , and the above becomes

$$\begin{aligned} |\psi_2\rangle_{\text{const}} &= \frac{1}{2} (|0, f(0)\rangle - |0, 1 \oplus f(0)\rangle + |1, f(0)\rangle - |1, 1 \oplus f(0)\rangle) \\ &= \frac{1}{2} (|0\rangle + |1\rangle) (|f(0)\rangle - |1 \oplus f(0)\rangle) \end{aligned} \quad (15.3.20)$$

With the Hadamard operation on the upper qubit, we have

$$|\psi_3\rangle_{\text{const}} = |0\rangle \frac{1}{\sqrt{2}} (|f(0)\rangle - |1 \oplus f(0)\rangle) \quad (15.3.21)$$

If for the balanced case,  $f(0) \neq f(1)$ , hence  $f(1) = 1 \oplus f(0)$ . Then

$$\begin{aligned} |\psi_2\rangle_{\text{bal}} &= \frac{1}{2} (|0, f(0)\rangle - |0, 1 \oplus f(0)\rangle + |1, 1 \oplus f(0)\rangle - |1, f(0)\rangle) \\ &= \frac{1}{2} (|0\rangle - |1\rangle) (|f(0)\rangle - |1 \oplus f(0)\rangle) \end{aligned} \quad (15.3.22)$$

After the  $H$  gate operation, we have

$$|\psi_3\rangle_{\text{bal}} = |1\rangle \frac{1}{\sqrt{2}} (|f(0)\rangle - |1 \oplus f(0)\rangle) \quad (15.3.23)$$

The above shows the prowess of quantum parallelism with just one operation, one can determine if a function  $f(x)$  is balanced or constant. More sophisticated algorithms exploiting quantum parallelism, such as the Shor's algorithm and the Grover's algorithm, have been devised. The Shor's algorithm can perform a Fourier transform in  $(\log N)^2$  operations rather than the classical  $N \log N$  operations. The Grover's algorithm can search a data base with  $N$  data in  $\sqrt{N}$  operations rather than the classical  $N$  operations. Because of quantum parallelism, quantum computer can also perform quantum simulation of quantum system which is not possible on classical computers.

## 15.4 Quantum Teleportation

Quantum teleportation is the idea of Alice being able to send a photon of unknown state to Bob, without having to perform a measurement on this photon, nor disturb its state. We denote the photon, called photon 1, in the unknown state by

$$|\psi\rangle_1 = C_0 |0\rangle_1 + C_1 |1\rangle_1 \quad (15.4.1)$$

In the beginning, this photon is only accessible to Alice. Alice, however, is accessible to another photon, called photon 2, of an entangled Bell state. The other photon of the Bell state, photon 3, is not accessible to Alice but is accessible to Bob.

We can define the state of the three photons by the direct product state

$$|\Psi\rangle_{123} = \frac{1}{\sqrt{2}} (C_0 |0\rangle_1 + C_1 |1\rangle_1) (|0\rangle_2 |1\rangle_3 - |1\rangle_2 |0\rangle_3) \quad (15.4.2)$$

where the Bell state is assumed to be

$$|\Psi^-\rangle_{23} = \frac{1}{\sqrt{2}} (|0\rangle_2 |1\rangle_3 - |1\rangle_2 |0\rangle_3) \quad (15.4.3)$$

Expanding (15.4.2) gives rise to

$$\begin{aligned} |\Psi\rangle_{123} &= \frac{1}{\sqrt{2}} (C_0 |0\rangle_1 |0\rangle_2 |1\rangle_3 - C_0 |0\rangle_1 |1\rangle_2 |0\rangle_3 \\ &\quad + C_1 |1\rangle_1 |0\rangle_2 |1\rangle_3 - C_1 |1\rangle_1 |1\rangle_2 |0\rangle_3) \end{aligned} \quad (15.4.4)$$

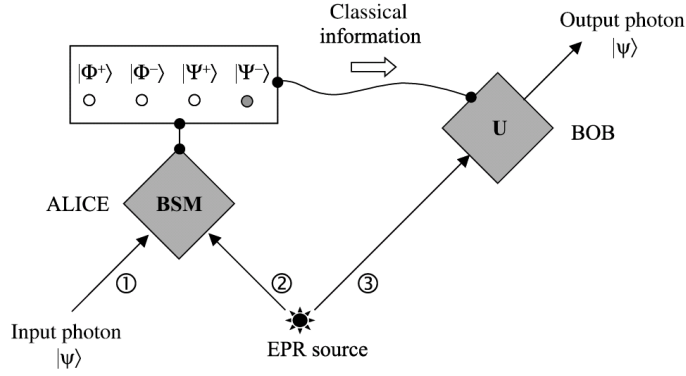


Figure 15.6: A teleportation experiment setup with photons. Alice makes measurement on the two photons accessible to her using the Bell state measurement device, and communicate the outcome to Bob via a classical channel. Bob then performs a unitary transformation on his photon to obtain the input photon state (from Fox).

The four Bell states are complete and orthogonal, and the states of photon 1 and photon 2 can be expanded in the four Bell states; namely

$$|\Phi^+\rangle_{12} = \frac{1}{\sqrt{2}} (|0\rangle_1 |0\rangle_2 + |1\rangle_1 |1\rangle_2) \quad (15.4.5)$$

$$|\Phi^-\rangle_{12} = \frac{1}{\sqrt{2}} (|0\rangle_1 |0\rangle_2 - |1\rangle_1 |1\rangle_2) \quad (15.4.6)$$

$$|\Psi^+\rangle_{12} = \frac{1}{\sqrt{2}} (|0\rangle_1 |1\rangle_2 + |1\rangle_1 |0\rangle_2) \quad (15.4.7)$$

$$|\Psi^-\rangle_{12} = \frac{1}{\sqrt{2}} (|0\rangle_1 |1\rangle_2 - |1\rangle_1 |0\rangle_2) \quad (15.4.8)$$

Projecting (15.4.4) onto the four Bell states, and subsequently expanding (15.4.4) in terms of them, we have

$$\begin{aligned} |\Psi\rangle_{123} = \frac{1}{2} [ & |\Phi^+\rangle_{12} (C_0 |1\rangle_3 - C_1 |0\rangle_3) \\ & + |\Phi^-\rangle_{12} (C_0 |1\rangle_3 + C_1 |0\rangle_3) \\ & + |\Psi^+\rangle_{12} (-C_0 |0\rangle_3 + C_1 |1\rangle_3) \\ & - |\Psi^-\rangle_{12} (C_0 |0\rangle_3 + C_1 |1\rangle_3) ] \end{aligned} \quad (15.4.9)$$

In the above, the first two photons, photon 1 and photon 2, are grouped into different Bell states. Moreover, the state of photon 3 resembles the state of the original photon 1. The quantum system now is in a linear superposition of different Bell states.

The Bell state measurement device projects the first two photons onto a Bell state. The Bell state measurement device collapses the quantum system into one of the four Bell states. For example, when Alice finds that the first two photons are in the Bell state  $|\Phi^+\rangle_{12}$ , then the third photon must be in the state

$$|\psi\rangle_3 = C_0 |1\rangle_3 - C_1 |0\rangle_3 \quad (15.4.10)$$

Bob can apply a unitary operator, similar to qubit gate operators described in the previous section, to obtain the original state of photon 1.

The above does not violate the no-cloning theorem, because the original state of the photon 1 is destroyed, and its semblance is reproduced in photon 3.

## 15.5 Interpretation of Quantum Mechanics

Quantum mechanics has the basic tenet that a quantum state is in a linear superposition of states before a measurement. A measurement projects a quantum state into one of the states. The ability of a quantum state to be in a linear superposition of states is surreal, and it has bothered a great many physicists. In the classical world, a system can only be in one state or the other, but not in a linear superposition of states. Only the world of ghosts and angels can we imagine that an object is in a linear superposition of states.

In the coordinate space, an electron, represented by its wavefunction, can be simultaneously at all locations where the wavefunction is non-zero. In the Young's double slit experiment, the electron, represented by its wavefunction, can go through both slits simultaneously like a wave.<sup>1</sup>

When the ghost-angel state concept is extended to classical objects, such as a cat, it gives rise to the ludicrous result: the story of the Schrödinger cat. The Schrödinger cat is a linear superposition of a dead cat and a live cat. To understand why the Schrödinger cat does not exist, we need to understand the concept of quantum coherence.

Two states are in quantum coherence if the phase relationships between them are deterministic and not random. When this coherence is lost, the phase relationship between them is lost. The quantum system has already collapsed into one of the two states. Hence, in practice, a measurement is not always necessary before the quantum system collapses into one or more of the quantum states. The interaction of quantum system with its environment can cause such a collapse.

From a statistical physics viewpoint, it is impossible for a quantum system to be completely isolated. Almost all systems are in a thermal bath of the universe with which they are seeking equilibrium. Macroscopic objects cannot be in a pure quantum state which has the characteristics of the ghost-angel state. It is impossible for the huge number of atoms in the Schrödinger cat to be coherent with respect to each other.

The density matrix is a nice way of representing a state of a quantum system where the physics of quantum coherence surfaces explicitly. This concept is expressed in the off-diagonal terms of the density matrix. If one allows time average or ensemble average<sup>2</sup> to the density matrix, when the system is expressed by quantum states that are not coherent, the

<sup>1</sup>Or the apparition of the ghost-angel state.

<sup>2</sup>Processes for which time average is the same as ensemble average are known as ergodic processes.

off-diagonal elements will average to zero. The system is in a mixed state rather than a pure quantum state. The system is similar to the local hidden variable theory: the state of the quantum system is already predetermined before the measurement.

Another uneasiness about the philosophical interpretation of quantum mechanics is that one does not know what state the quantum system is in before the measurement. This has prompted Einstein to ask, “Is the moon there if you don’t look at it?” The uncertainty of the state applied to quantum mechanics is only true for a linear superposition of coherent quantum states, which I term the ghost-angel state. This state has not been found to exist for macroscopic objects.

However, if one insists that, “One does not know if the moon is there before one looks at it.” as a true statement, it cannot be refuted nor confirmed by experiments. The mere act of an experiment already means that we have “looked” at the moon. The same claim goes that “If I saw a fallen tree in the forest, it did not necessary follow from the act of falling before I arrived there.” Alternatively, “If we found dinosaur bones, it did not necessary mean that dinosaurs roamed the earth over 200 million years ago.” We believe that the moon is there even if we do not look at it, the tree fell because it went through the act of falling, and that dinosaur bones were found because they roamed the earth 200 million years ago, because we believe in the realism of the world we live in. This realism cannot be proved but is generally accepted by those who live in this world. Hence, it is this surreal interpretation of quantum mechanics that causes the uneasiness among many physicists. But the interpretation of quantum mechanics is slightly better than the above: a quantum state is in a linear superposition of states, the precise one of which we are not sure of until a measurement is performed. However, this surrealism of this ghost-angel state exists in our minds in fairy tales and ghost stories of many cultures. Experimental effort has agreed with the surreal interpretation of quantum mechanics in terms of the Bell’s theorem, that will be discussed.

The ghost-angel state of a quantum system is what enriches the information in it. However, for a quantum system to be in such a state, the linear superposition of states must be coherent with each other. Quantum coherence is the largest stumbling block to the construction of quantum computers; however, rapid advances are being made, and one day, it can be a reality.

## 15.6 EPR Paradox

The interpretation of quantum mechanics went through difficult times. The fact that a particle can be in a superposition of many states before a quantum measurement, and the probabilistic interpretation of a quantum measurement behooves the challenge by many great physicists, especially Einstein. The most severe challenge of quantum mechanics and its interpretation comes from the EPR (Einstein, Podolsky and Rosen) paradox. To describe it simply, we imagine a pi meson that decays into an electron-positron pair:

$$\pi^0 \rightarrow e^- + e^+ \tag{15.6.1}$$

The pi meson originally has spin zero. So for conservation of angular momentum, the electron-positron pair will have opposite spins: spin up and spin down. Since the total angular

momentum is zero, they are in the singlet state which has total spin of zero, or

$$|\Psi\rangle = \frac{1}{\sqrt{2}} (|\uparrow\downarrow\rangle - |\downarrow\uparrow\rangle) \quad (15.6.2)$$

The electron-positron pair is in the linear superposition of two states, but the electron and positron are flying in opposite directions. According to the interpretation of quantum mechanics, one does not know the spin state of the electron nor the positron before the measurement. After one measures the spin state of, say the electron, irrespective of how far the positron is away from the electron, we immediately know the spin state of the positron according to the above equation. The spins of the two particles are always opposite to each other. This notion is unpalatable to many physicists, and hence, is called “spooky action at a distance” by Einstein. Information cannot travel faster than the speed of light. How could the state of one particle be immediately determined after a measurement is made at another particle far away? This paradox attempts to prove that quantum mechanics is incomplete by *reductio ad absurdum*.

## 15.7 Bell’s Theorem<sup>3</sup>

Quantum measurements are known to be random, and the data can only be interpreted probabilistically. If one were to measure the spin of one of the particle, it is equally likely to be in the spin up or spin down state randomly according to (15.6.2). In the hidden variable theory, it is suggested that the outcome of the experiment is already predetermined even before the measurement. The outcome is determined by a hidden random variable  $\lambda$ . It is the randomness of this variable that gives rise to the randomness of the outcome in quantum measurements.

Many hidden variable theories were proposed shortly after the EPR paradox was published. In 1964, J. S. Bell, in the spirit of proving the correctness of the hidden variable theory, came up with an inequality that showed the incompatibility of quantum mechanics and the hidden variable theory. If hidden variable theory is correct, the inequality will be satisfied, but if quantum mechanics is correct, the inequality is violated. This is known as Bell’s theorem.

We can discuss the derivation of the Bell’s theorem in the context of the two-photon experiment, since the experiment that verifies the theorem has been done using photons. The actual experiment is quite complex, but we will reduce it to a simplified case. The simplified experiment involves a photon source that produces an entangled photon pair, each of which is traveling in opposite directions. The photon pair is in one of the Bell state, say the EPR pair:

$$|\Psi\rangle = \frac{1}{\sqrt{2}} (|V_1V_2\rangle + |H_1H_2\rangle) \quad (15.7.1)$$

In the above state, if one of the photons is measured to be V (vertical) polarized, the other photon must be V polarized. However, if one photon is measured to be H (horizontal) polarized, the other photon must be H polarized. We will detect the photon state with a

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<sup>3</sup>This section is written with important input from Y. H. Lo and Q. Dai.



simple polarizer. In the above state, which is a linear superposition of two states, the photons are equally likely to be found in the first state,  $|V_1V_2\rangle$ , or the second state,  $|H_1H_2\rangle$ . The polarizer will detect an H or a V polarization with equal likelihood, but the moment that one photon is determined to be H polarized, the other photon is immediately known to be H polarized, and vice versa. This is the “spookiness” of quantum interpretation.

Imagine an atomic source that generates two photons propagating in opposite directions. The atom initially has zero angular momentum, so that the two photons are either both horizontally polarized or both vertically polarized. Hence, the photon can be described in one of the Bell states or an entangled state as shown in (15.7.1).

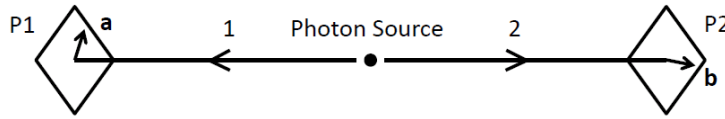


Figure 15.7: Experimental verification of Bell’s theorem

### 15.7.1 Prediction by Quantum Mechanics

Photon 1 is measured with a polarizer  $P_1$  with vertical polarization pointing in the  $\mathbf{a}$  direction. If  $P_1$  measures a V polarization, we set  $A(\mathbf{a}) = 1$ , and if  $P_1$  measures an H polarization, we set  $A(\mathbf{a}) = -1$ . Here,  $A(\mathbf{a})$  denotes the measurement outcome, and it is completely random according to quantum mechanics.

Similarly, polarizer  $P_2$  has its vertical polarization oriented in the  $\mathbf{b}$  direction. When it measures a V polarization, we set  $B(\mathbf{b}) = 1$ , and when it measures an H polarization, it sets  $B(\mathbf{b}) = -1$ . Again,  $B(\mathbf{b})$  is completely random. In the above, we set  $A$  and  $B$  to be functions of  $\mathbf{a}$  and  $\mathbf{b}$  respectively, as the experimental outcomes are expected to be functions of the orientation of the polarizers.

If  $\mathbf{a} = \mathbf{b}$ , we expect that

$$\langle A(\mathbf{a})B(\mathbf{b}) \rangle = E(\mathbf{a}, \mathbf{b}) = AB = 1 \tag{15.7.2}$$

If  $\mathbf{a} \perp \mathbf{b}$ , if  $P_1$  measures a V polarization,  $P_2$  will measure a H polarization, we expect that

$$\langle A(\mathbf{a})B(\mathbf{b}) \rangle = E(\mathbf{a}, \mathbf{b}) = AB = -1 \tag{15.7.3}$$

Even though  $A$  and  $B$  are random, their products are deterministic in the above two cases.

An interesting case ensues if  $\mathbf{a}$  and  $\mathbf{b}$  are at an incline with respect to each other. In this case, we know by quantum mechanics that

$$|V\rangle_a = \cos \theta |V\rangle_b - \sin \theta |H\rangle_b \tag{15.7.4}$$

$$|H\rangle_a = \sin \theta |V\rangle_b + \cos \theta |H\rangle_b \tag{15.7.5}$$

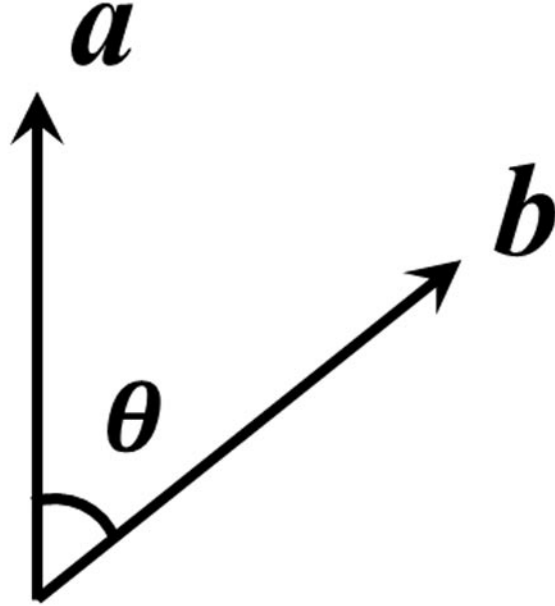


Figure 15.8: The case when the two polarizers are at an incline with respect to each other for proving the Bell's theorem.

If  $P_1$  measures an outcome with  $A(\mathbf{a}) = 1$ , then the photon that propagates to  $P_2$  must be polarized in the  $\mathbf{a}$  direction. However, according to (15.7.4), for  $P_2$ , such a photon has the probability of  $\cos^2 \theta$  being detected in the V polarization, and the probability of  $\sin^2 \theta$  being detected in the H polarization. Hence, the expectation value of  $B$ , or  $\langle B \rangle = \cos^2 \theta - \sin^2 \theta$ .

If  $P_1$  finds that  $A = -1$ , by similar arguments, the expectation value of  $B$  is  $\sin^2 \theta - \cos^2 \theta$ . Then

$$E(\mathbf{a}, \mathbf{b}) = \langle A(\mathbf{a})B(\mathbf{b}) \rangle = \cos^2 \theta - \sin^2 \theta = \cos(2\theta) \quad (15.7.6)$$

Notice that the above reduces to the special cases of: (i) when the polarizers  $P_1$  and  $P_2$  are aligned, namely when  $\theta = 0^\circ$ , as in (15.7.2), and (ii) when the polarizers are perpendicular to each other, with  $\theta = 90^\circ$ , as in (15.7.3).

### 15.7.2 Prediction by Hidden Variable Theory

In the hidden variable theory derivation, a particle is assumed to be already in a polarization state even before a measurement.  $A$  and  $B$  are random, but they are predetermined by a hidden random variable  $\lambda$ . We let

$$A(\mathbf{a}, \lambda) = \pm 1 \quad (15.7.7)$$

$$B(\mathbf{b}, \lambda) = \pm 1 \quad (15.7.8)$$

Here,  $A$  and  $B$  hence flip-flop between  $\pm 1$  entirely due to randomness of the variable  $\lambda$ . This theory predicts the randomness of quantum mechanics experiments nicely. The expectation value of  $AB$  then is

$$E(\mathbf{a}, \mathbf{b}) = \langle AB \rangle = \int \rho(\lambda) A(\mathbf{a}, \lambda) B(\mathbf{b}, \lambda) d\lambda \quad (15.7.9)$$

where  $\rho(\lambda)$  is the probability distribution of the random variable  $\lambda$ . The above is a very general representation of the hidden variable theory where we have not explicitly stated the functions  $\rho(\lambda)$ ,  $A(\mathbf{a}, \lambda)$  nor  $B(\mathbf{b}, \lambda)$ .

If  $\mathbf{a} = \mathbf{b}$ , then

$$A(\mathbf{a}, \lambda) = B(\mathbf{a}, \lambda) \quad (15.7.10)$$

and the above becomes

$$E(\mathbf{a}, \mathbf{a}) = \int \rho(\lambda) A^2(\mathbf{a}, \lambda) d\lambda = \int \rho(\lambda) d\lambda = 1 \quad (15.7.11)$$

Same as we would have found in (15.7.2).

If  $\mathbf{a} \perp \mathbf{b}$ , then

$$A(\mathbf{a}, \lambda) = -B(\mathbf{b}, \lambda) \quad (15.7.12)$$

and (15.7.9) becomes

$$E(\mathbf{a}, \mathbf{b}) = - \int \rho(\lambda) A^2(\mathbf{a}, \lambda) d\lambda = \int \rho(\lambda) d\lambda = -1 \quad (15.7.13)$$

as would have been found in (15.7.3). Hence, hidden variable theory is in good agreement with quantum mechanics interpretation of (15.7.2) and (15.7.3).

So far, everything is fine and dainty until when  $\mathbf{a}$  and  $\mathbf{b}$  do not belong to any of the above category, but in general are inclined with respect to each other. If  $\mathbf{a} \neq \mathbf{b}$  in general, then the hidden variable generates random  $A$  and  $B$  in such a manner that

$$E(\mathbf{a}, \mathbf{b}) \leq 1 \quad (15.7.14)$$

Consequently, we have

$$E(\mathbf{a}, \mathbf{b}) - E(\mathbf{a}, \mathbf{b}') = \int [A(\mathbf{a}, \lambda) B(\mathbf{b}, \lambda) - A(\mathbf{a}, \lambda) B(\mathbf{b}', \lambda)] \rho(\lambda) d\lambda \quad (15.7.15)$$

The above can be rewritten as

$$\begin{aligned} E(\mathbf{a}, \mathbf{b}) - E(\mathbf{a}, \mathbf{b}') &= \int A(\mathbf{a}, \lambda) B(\mathbf{b}, \lambda) [1 \pm A(\mathbf{a}', \lambda) B(\mathbf{b}', \lambda)] \rho(\lambda) d\lambda \\ &\quad - \int A(\mathbf{a}, \lambda) B(\mathbf{b}', \lambda) [1 \pm A(\mathbf{a}', \lambda) B(\mathbf{b}', \lambda)] \rho(\lambda) d\lambda \end{aligned} \quad (15.7.16)$$

We have just added and subtracted identical terms in the above. After using the triangular inequality

$$\begin{aligned} |E(\mathbf{a}, \mathbf{b}) - E(\mathbf{a}, \mathbf{b}')| &\leq \int |A(\mathbf{a}, \lambda)B(\mathbf{b}, \lambda)| |1 \pm A(\mathbf{a}', \lambda)B(\mathbf{b}', \lambda)| \rho(\lambda) d\lambda \\ &\quad + \int |A(\mathbf{a}, \lambda)B(\mathbf{b}', \lambda)| |1 \pm A(\mathbf{a}', \lambda)B(\mathbf{b}, \lambda)| \rho(\lambda) d\lambda \end{aligned} \quad (15.7.17)$$

Using the fact that

$$|AB| \leq 1, \quad 1 \pm AB \geq 0 \quad (15.7.18)$$

we have

$$\begin{aligned} |E(\mathbf{a}, \mathbf{b}) - E(\mathbf{a}, \mathbf{b}')| &\leq \int [1 \pm A(\mathbf{a}', \lambda)B(\mathbf{b}', \lambda)] \rho(\lambda) d\lambda \\ &\quad + \int [1 \pm A(\mathbf{a}', \lambda)B(\mathbf{b}, \lambda)] \rho(\lambda) d\lambda \end{aligned} \quad (15.7.19)$$

The above is the same as

$$|E(\mathbf{a}, \mathbf{b}) - E(\mathbf{a}, \mathbf{b}')| \leq 2 \pm [E(\mathbf{a}', \mathbf{b}') + E(\mathbf{a}', \mathbf{b})] \quad (15.7.20)$$

It is of the form

$$|X| \leq 2 \pm Y \quad (15.7.21)$$

which implies that

$$|X| \pm Y \leq 2 \quad (15.7.22)$$

or

$$|X + Y| \leq |X| + |Y| \leq 2 \quad (15.7.23)$$

Consequently, we have

$$|E(\mathbf{a}, \mathbf{b}) - E(\mathbf{a}, \mathbf{b}') + E(\mathbf{a}', \mathbf{b}') + E(\mathbf{a}', \mathbf{b})| \leq 2 \quad (15.7.24)$$

which is the Clauser-Horne-Shimony-Holt (CHSH) inequality.

In (15.7.20), if  $\mathbf{a}' = \mathbf{b}' = \mathbf{c}$ , then (15.7.20) becomes

$$|E(\mathbf{a}, \mathbf{b}) - E(\mathbf{a}, \mathbf{c})| \leq 2 \pm [E(\mathbf{c}, \mathbf{c}) + E(\mathbf{b}, \mathbf{c})] \quad (15.7.25)$$

From (15.7.11),  $E(\mathbf{c}, \mathbf{c}) = 1$ . We pick the smaller of the right-hand side of (15.7.25) and arrive at

$$|E(\mathbf{a}, \mathbf{b}) - E(\mathbf{a}, \mathbf{c})| \leq 1 - E(\mathbf{b}, \mathbf{c}) \quad (15.7.26)$$

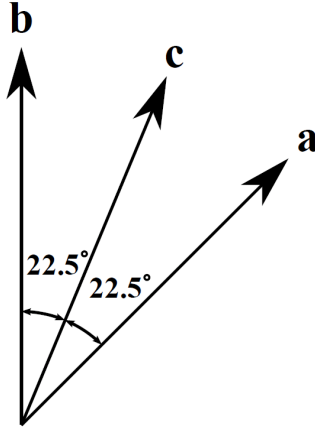


Figure 15.9: The case for **a**, **b**, and **c** where the quantum mechanics prediction violates the Bell's inequality.

The above is the Bell's inequality. It can be easily shown that it cannot be satisfied by quantum mechanics that predicts (15.7.6). Say if we pick **a**, **b**, and **c** as shown in the Figure 15.9, then according to quantum mechanics or (15.7.6),

$$E(\mathbf{a}, \mathbf{b}) = \cos(90^\circ) = 0 \tag{15.7.27}$$

$$E(\mathbf{a}, \mathbf{c}) = \cos(45^\circ) = \frac{1}{\sqrt{2}} \tag{15.7.28}$$

$$E(\mathbf{b}, \mathbf{c}) = \cos(45^\circ) = \frac{1}{\sqrt{2}} \tag{15.7.29}$$

Then in (15.7.26), we have

$$\left| 0 - \frac{1}{\sqrt{2}} \right| = 0.707 \not\leq 1 - \frac{1}{\sqrt{2}} = 0.293 \tag{15.7.30}$$

In experimental tests of Bell's theorem, quantum mechanics triumphs over hidden variable theory so far.<sup>4</sup> Hence, the spookiness of ghost-angel states will prevail in quantum mechanics.

## 15.8 A Final Word on Quantum Parallelism

Quantum interpretation gives rise to the spookiness of quantum mechanics, in a way giving it the capability that empowers ghosts and angels. Quantum parallelism is such an empowerment. It reminds me of a novel that I have read when I was young on the Monkey King. This

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<sup>4</sup>see a paper by A. Aspect, 1982.

Monkey King is of Indian origin, but has permeated Chinese culture to take on a different persona. According to the story in *The Journey to the West*, this Monkey King had unusual capabilities. One of them was that he could pull a strand of hair from his body, and with a puff of air, he could turn it into many duplicates of himself. He could then fight his enemies from all angles and all sides. He thus made himself invincible.

He was arrogant, mischievous, and wreaked havoc in the Heavenly Palace where the other gods lived. Finally, he could only be tamed and subdued by Lord Buddha, summoning him to accompany and protect the monk Xuan Zang in his treacherous journey to collect the Buddhist Sutra from the West (in this case India).

*Oh Lord, if we were to empower ourselves with the capabilities of ghosts and angels, who is there to curb our power!*



Figure 15.10: The Monkey King sending replicas of himself against his enemy, defeating everyone in his path.





# Appendix A

## Generators of Translator and Rotation

### A.1 Infinitesimal Translation

From Taylor series expansion, we have

$$\begin{aligned} f(x+a) &= f(x) + af(x) + \frac{a^2}{2!}f''(x) + \dots \\ &= \left[ 1 + a\frac{d}{dx} + \frac{a^2}{2!}\frac{d^2}{dx^2} + \dots \right] f(x) \\ &= e^{a\frac{d}{dx}} f(x) \end{aligned} \tag{A.1.1}$$

The exponential to an operator is interpreted as a Taylor series when it needs to be evaluated. The above can be written with a momentum operator (assuming that  $\hbar = 1$ )

$$f(x+a) = e^{ia\hat{p}} f(x) \tag{A.1.2}$$

or in Dirac notation

$$|f_a\rangle = e^{ia\hat{p}}|f\rangle \tag{A.1.3}$$

where  $|f_a\rangle$  is the state vector representation of the function  $f(x+a)$ . If a Hamiltonian is translational invariant, it will commute with the translation operator, namely

$$\hat{H}e^{ia\hat{p}}|f\rangle = e^{ia\hat{p}}\hat{H}|f\rangle \tag{A.1.4}$$

or

$$\left[ e^{ia\hat{p}}, \hat{H} \right] = 0 \tag{A.1.5}$$

For example, if the Hamiltonian is such that  $-\frac{1}{2}d^2/dx^2$ , then it is quite clear that the left-hand side of (A.1.4) is

$$-\frac{1}{2}\frac{d^2}{dx^2}e^{ia\hat{p}}f(x) = -\frac{1}{2}\frac{d^2}{dx^2}f(x+a) = -\frac{1}{2}f''(x+a) \tag{A.1.6}$$

The right-hand side of (A.1.4) is

$$e^{ia\hat{p}} \left[ -\frac{1}{2} \frac{d^2}{dx^2} f(x) \right] = e^{ia\hat{p}} \left[ -\frac{1}{2} f''(x) \right] = -\frac{1}{2} f''(x+a) \quad (\text{A.1.7})$$

which is the same as the left-hand side.

By assuming that  $e^{ia\hat{p}} \approx 1 + ia\hat{p} + \dots$ , the above also means that

$$\hat{H}\hat{p} = \hat{p}\hat{H}, \quad [\hat{H}, \hat{p}] = 0 \quad (\text{A.1.8})$$

This means that  $\hat{p}$  is a constant of motion as shown in Section 5.8, or

$$\frac{d\langle \hat{p} \rangle}{dt} = 0 \quad (\text{A.1.9})$$

In other words, momentum is conserved in a system where the Hamiltonian is translational invariant.

## A.2 Infinitesimal Rotation

Similarly, the  $L_z$  operator is  $-i\frac{\partial}{\partial\phi}$  which is a generator of rotation,<sup>1</sup>

$$f(\phi + \alpha) = e^{\alpha\frac{d}{d\phi}} f(\phi) = e^{i\alpha\hat{L}_z} f(\phi) \quad (\text{A.2.1})$$

Going through the derivation as we have before, for a Hamiltonian that is rotationally symmetric about the  $z$  axis, then it commutes with  $e^{i\alpha\hat{L}_z}$ . In this case

$$[e^{i\alpha\hat{L}_z}, \hat{H}] = 0 \quad (\text{A.2.2})$$

or Taylor series expanding the above, we derive that

$$[\hat{L}_z, \hat{H}] = 0 \quad (\text{A.2.3})$$

Hence, the  $\hat{z}$  component of the angular momentum is conserved for a system that is rotationally symmetric about the  $z$  axis.

The above has been motivated by  $\hat{L}_z$  that follows from the orbital angular momentum, whose form has been motivated by the classical angular momentum. What if the angular momentum has no classical analogue like the spin angular momentum? Or the state vector may not be written using wavefunctions at all. In this case, we can postulate a generalized generator of rotation of the form

$$e^{i\alpha\hat{J}_z} |j\rangle \quad (\text{A.2.4})$$

The above will take the angular momentum state  $|j\rangle$  and generate rotated state about the  $z$  axis. We postulate the above form for three reasons:

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<sup>1</sup>Assume again that  $\hbar = 1$ .

1. The generator of rotation for orbital angular momentum is already of this form. It must be the special case of the above form.
2. If the above generator commutes with the Hamiltonian with similar rotational symmetry, this component of angular momentum will be conserved.
3. The rotation will cause the expectation value of the angular momentum vector, an observable, to rotate according to a coordinate rotation about the  $z$  axis.

It is easy to see that the expectation value of the  $\hat{J}_z$  operator remains unchanged under this rotation. We can easily show that

$$\langle j | e^{-i\alpha\hat{J}_z} \hat{J}_z e^{i\alpha\hat{J}_z} | j \rangle = \langle j | \hat{J}_z | j \rangle \quad (\text{A.2.5})$$

In the above, functions of operator  $\hat{J}_z$  commute with each other, a fact that can be easily proved by expanding the functions as a power series. The conjugate transpose of  $e^{i\alpha\hat{J}_z}$  is  $e^{-i\alpha\hat{J}_z}$  since  $\hat{J}_z$  is Hermitian because it represents an observable. Hence, the last equality in (A.2.5) follows.

### A.3 Derivation of Commutation Relations

In general, we can define the angular momentum operator to be

$$\hat{\mathbf{J}} = \mathbf{i}\hat{J}_x + \mathbf{j}\hat{J}_y + \mathbf{k}\hat{J}_z \quad (\text{A.3.1})$$

The above is an observable so are the components of the operator in the  $x$ ,  $y$ ,  $z$  directions. Hence, the expectation value of the above with respect to the state  $|j\rangle$  gives rise to

$$\langle \hat{\mathbf{J}} \rangle = \mathbf{i}\langle \hat{J}_x \rangle + \mathbf{j}\langle \hat{J}_y \rangle + \mathbf{k}\langle \hat{J}_z \rangle = \mathbf{i}J_x + \mathbf{j}J_y + \mathbf{k}J_z = \mathbf{J} \quad (\text{A.3.2})$$

where we denote the expectation values of  $\hat{J}_i$  with scalar number  $J_i$ , and that of  $\hat{\mathbf{J}}$  with  $\mathbf{J}$ . We will test point 3 above with respect to this vector  $\mathbf{J}$ . This vector will have to rotate according to coordinate rotation as the state vector is rotated according to (A.2.4).

If  $e^{i\alpha\hat{J}_z}$  is a generator of rotation about the  $z$  axis, it will leave  $J_z$  unchanged as shown above. But it will not leave  $J_x$  and  $J_y$  components unchanged. Now, we can find the expectation value of  $\hat{J}_x$  under rotation about  $z$  axis, or that

$$J_x = \langle j | \hat{J}_x | j \rangle, \quad J'_x = \langle j | e^{-i\alpha\hat{J}_z} \hat{J}_x e^{i\alpha\hat{J}_z} | j \rangle \quad (\text{A.3.3})$$

Before rotation, this would have represented the expectation values of the  $x$  component of  $\mathbf{J}$ , the angular momentum. After rotation,

$$\mathbf{J} = \mathbf{i}'J'_x + \mathbf{j}'J'_y + \mathbf{k}J_z \quad (\text{A.3.4})$$

and  $J'_x$  will contain both the  $x$  and  $y$  component of  $\mathbf{J}$ . When we rotate the state vector by angle  $\alpha$ , the expectation value of the vector  $\mathbf{J}$  will rotate by  $\alpha$  in the same direction. But from the Figure A.1, it is clear that

$$J'_x = J_x \cos \alpha + J_y \sin \alpha \quad (\text{A.3.5})$$

The equivalence of rotations in (A.3.3) and (A.3.5) can be proved by lengthy algebra. To simplify the algebra, we can show the equivalence for infinitesimal rotation. To this end, we assume that  $\alpha$  is small so that we keep only the first two terms of the Taylor expansion for  $e^{i\alpha\hat{J}_z}$  or  $e^{i\alpha\hat{J}_z} \approx 1 + i\alpha\hat{J}_z$ . Then,

$$J'_x = \langle j | \hat{J}_x - i\alpha\hat{J}_z\hat{J}_x + i\alpha\hat{J}_x\hat{J}_z + \cdots | j \rangle \quad (\text{A.3.6})$$

For small  $\alpha$ , (A.3.5) becomes

$$J'_x \cong J_x + \alpha J_y \quad (\text{A.3.7})$$

But (A.3.6) can be written as

$$J'_x = \langle j | \hat{J}_x | j \rangle - i\alpha \langle j | \hat{J}_z \hat{J}_x - \hat{J}_x \hat{J}_z | j \rangle \quad (\text{A.3.8})$$

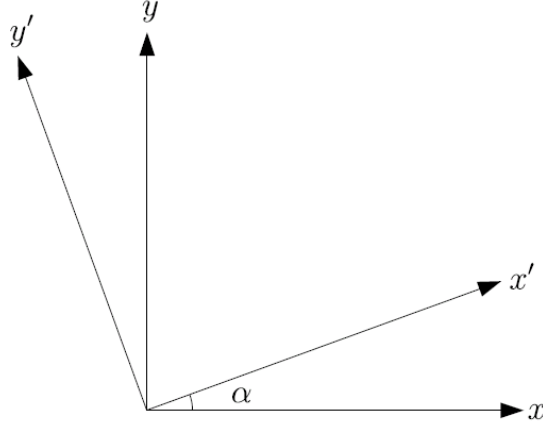


Figure A.1: Coordinate rotation of the  $xy$  plane about the  $z$  axis by angle  $\alpha$ .

Comparing (A.3.7) and (A.3.8), we have

$$\hat{J}_z \hat{J}_x - \hat{J}_x \hat{J}_z = i\hat{J}_y \quad (\text{A.3.9})$$

We can let  $\hbar \neq 1$  to arrive back at the previously postulated commutator relations. The above is arrived at using only rotational symmetry argument that the angular momentum operator is a generator of rotation. The other commutation relations for angular operators can be derived by similar arguments. In general, for all operators that represent angular momentum, we have

$$\left[ \hat{J}_x, \hat{J}_y \right] = i\hbar\hat{J}_z, \quad \left[ \hat{J}_y, \hat{J}_z \right] = i\hbar\hat{J}_x, \quad \left[ \hat{J}_z, \hat{J}_x \right] = i\hbar\hat{J}_y \quad (\text{A.3.10})$$

From the above, we can define raising and lowering operators and use the ladder approach to derive the properties of the eigenstates of angular momentum operators.

## Appendix B

# Quantum Statistical Mechanics

### B.1 Introduction

We would like to derive Maxwell-Boltzmann, Fermi-Dirac, and Bose-Einstein distributions. It is based on statistical mechanics that assumes that all identical energy levels have equal likelihood of being occupied. Only the counting is based on quantum mechanics.

We consider a quantum system with energy levels  $E_1, E_2, E_3, \dots$ . Each level  $E_i$  could have degeneracy of  $d_i, i = 1, 2, 3, \dots$ , and is filled by  $N_i, i = 1, 2, 3, \dots$  particles. When the system is coupled to its environment, energy will be exchanged with its environment via collision, vibration, and radiation. At thermal equilibrium, there is no net energy flowing into and out of the system. We expect that

$$E = \sum_{n=1}^{\infty} E_n N_n \quad (\text{B.1.1})$$

to be a constant due to energy conservation. Also we expect the number of particles to remain constant, or that

$$N = \sum_{n=1}^{\infty} N_n \quad (\text{B.1.2})$$

due to particle conservation.

The particles will acquire energy from their environment due to energy exchange, and distribute themselves across the energy levels subject to the above constraints, but also according to the availability of energy levels. Energy levels with higher degeneracies are more likely to be filled. We will consider three different cases: (i) when the particles are distinguishable, (ii) when they are identical fermions, and (iii) when they are identical bosons.

#### B.1.1 Distinguishable Particles

We will find the function  $Q(N_1, N_2, N_3, \dots)$  whose value equals the number of ways that there are  $N_1$  particles in  $E_1$ ,  $N_2$  particles in  $E_2$ ,  $N_3$  particles in  $E_3$  and so on. In order to fill

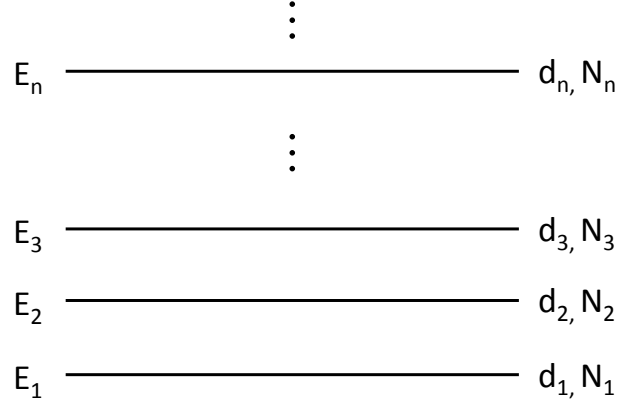


Figure B.1: Particles are distributed among different energy levels according to quantum statistics. An energy level is denoted with  $E_i$  with degeneracy  $d_i$  and  $N_i$  particles occupying it.

energy level  $E_1$  with  $N_1$  particles, the number of ways is

$$\binom{N}{N_1} = \frac{N!}{N_1!(N - N_1)!} \quad (\text{B.1.3})$$

If  $E_1$  has degeneracy  $d_1$ , all the degenerate levels are equally likely to be filled. Then there are  $d_1^{N_1}$  ways that the  $N_1$  particles can go into  $E_1$  level. Hence, the number of ways that  $E_1$  can be filled is

$$Q_{E_1} = \frac{N! d_1^{N_1}}{N_1!(N - N_1)!} \quad (\text{B.1.4})$$

The number of ways that  $E_2$  can be filled is

$$Q_{E_2} = \frac{(N - N_1)! d_2^{N_2}}{N_2!(N - N_1 - N_2)!} \quad (\text{B.1.5})$$

Then the total number of ways

$$\begin{aligned} Q(N_1, N_2, N_3, \dots) &= Q_{E_1} Q_{E_2} Q_{E_3} \dots \\ &= \frac{N! d_1^{N_1}}{N_1!(N - N_1)!} \frac{(N - N_1)! d_2^{N_2}}{N_2!(N - N_1 - N_2)!} \frac{(N - N_1 - N_2)! d_3^{N_3}}{N_3!(N - N_1 - N_2 - N_3)!} \dots \\ &= N! \frac{d_1^{N_1} d_2^{N_2} d_3^{N_3} \dots}{N_1! N_2! N_3! \dots} = N! \prod_{n=1}^{\infty} \frac{d_n^{N_n}}{N_n!} \end{aligned} \quad (\text{B.1.6})$$

Notice that the above is independent of the order in which the levels are filled, as expected. The above can also be interpreted in a different light: There are  $d_n^{N_n}$  ways that  $N_n$  particles can fit into the  $E_n$  energy level. But order is unimportant and a division by  $N_n!$  is necessary. But there are  $N!$  ways that these distinguishable particles can be selected in an ordered way, and hence, a prefactor of  $N!$  is needed in the above.

### B.1.2 Identical Fermions

Due to Pauli's exclusion principle, each energy level can admit only one particle. In this case, the number of ways energy  $E_n$  can be filled is

$$Q_{E_n} = \frac{d_n!}{N_n! (d_n - N_n)!}, \quad d_n \geq N_n \tag{B.1.7}$$

(The above has value 1 when  $d_n = 1$ .) The total number of ways is then

$$Q(N_1, N_2, N_3, \dots) = \prod_{n=1}^{\infty} \frac{d_n!}{N_n! (d_n - N_n)!} \tag{B.1.8}$$

### B.1.3 Identical Bosons

For bosons, repetition for filling a given state is allowed, but the particles are indistinguishable from each other. Then when  $N_n$  particles are picked to fill the  $d_n$  degenerate  $E_n$  levels, the first particle from  $N_n$  particles has  $d_n$  slots to fill in, while the second particle has  $d_n + 1$  slots to fill in, and the third particle has  $d_n + 2$  slots to fill in and so on, since repetition is allowed. That is the new particle can take the position of the old particle as well. So the number of ways that  $E_n$  with  $d_n$  degeneracy can be filled is

$$Q_{E_n} = \frac{d_n (d_n + 1) (d_n + 2) \dots (d_n + N_n - 1)}{N_n!} \tag{B.1.9}$$

$$= \frac{(N_n + d_n - 1)!}{N_n! (d_n - 1)!} \tag{B.1.10}$$

Therefore

$$Q(N_1, N_2, N_3, \dots) = \prod_{n=1}^{\infty} \frac{(N_n + d_n - 1)!}{N_n! (d_n - 1)!} \tag{B.1.11}$$

Unlike the distinguishable particle case, no prefactor of  $N!$  is needed for the identical particle case, since the order with which they are selected is unimportant.

## B.2 Most Probable Configuration

A way of filling the energy levels  $\{E_1, E_2, E_3, \dots\}$  with  $\{N_1, N_2, N_3, \dots\}$  is called a configuration. The most likely configuration to be filled is the one with the largest  $Q(N_1, N_2, N_3, \dots)$ . At statistical equilibrium, the system will gravitate toward this configuration. Hence, to find this configuration, we need to maximize  $Q$  with respect to different  $\{N_1, N_2, N_3, \dots\}$  subject to the energy conservation constraint (B.1.1) and particle conservation constraint (B.1.2).

We use the Lagrange multiplier technique to find the optimal  $Q$  subject to constraints (B.1.1) and (B.1.2). We define a function

$$G = \ln Q + \alpha \left[ N - \sum_{n=1}^{\infty} N_n \right] + \beta \left[ E - \sum_{n=1}^{\infty} N_n E_n \right] \quad (\text{B.2.1})$$

The optimal  $Q$  value is obtained by solving

$$\frac{\partial G}{\partial N_n} = 0, \quad n = 1, 2, 3, \dots, \quad \frac{\partial G}{\partial \alpha} = 0, \quad \frac{\partial G}{\partial \beta} = 0 \quad (\text{B.2.2})$$

### B.2.1 Distinguishable Particles

In this case,

$$G = \ln N! + \sum_{n=1}^{\infty} [N_n \ln d_n - \ln N!] \quad (\text{B.2.3})$$

$$+ \alpha \left[ N - \sum_{n=1}^{\infty} N_n \right] + \beta \left[ E - \sum_{n=1}^{\infty} N_n E_n \right] \quad (\text{B.2.4})$$

We use Stirling's formula that

$$\ln(z!) \approx z \ln z - z \quad (\text{B.2.5})$$

Consequently,

$$G = \sum_{n=1}^{\infty} [N_n \ln d_n - N_n \ln N_n + N_n - \alpha N_n - \beta E_n N_n] \quad (\text{B.2.6})$$

$$+ \ln N! + \alpha N + \beta E \quad (\text{B.2.7})$$

From the above

$$\frac{\partial G}{\partial N_n} = \ln d_n - \ln N_n - \alpha - \beta E_n = 0 \quad (\text{B.2.8})$$

As a result, we have

$$N_n = d_n e^{-(\alpha + \beta E_n)} \quad (\text{B.2.9})$$

The above is the precursor to the Maxwell-Boltzmann distribution from first principles.

### B.2.2 Identical Fermions

In this case

$$G = \sum_{n=1}^{\infty} [\ln(d'_n) - \ln(N'_n) - \ln((d_n - N_n)!)] \quad (\text{B.2.10})$$

$$+ \alpha \left[ N - \sum_{n=1}^{\infty} N_n \right] + \beta \left[ E - \sum_{n=1}^{\infty} N_n E_n \right] \quad (\text{B.2.11})$$



After applying Stirling's formula to  $N_n$  dependent terms, we have

$$G = \sum_{n=1}^{\infty} [\ln d_n! - N_n \ln N_n + N_n - (d_n - N_n) \ln(d_n - N_n) + (d_n - N_n) - \alpha N_n - \beta N_n E_n] \quad (\text{B.2.12})$$

$$+ \alpha N + \beta E \quad (\text{B.2.13})$$

From the above,

$$\frac{\partial G}{\partial N_n} = -\ln N_n + \ln(d_n - N_n) - \alpha - \beta E_n = 0 \quad (\text{B.2.14})$$

or

$$N_n = \frac{d_n}{1 + e^{\alpha + \beta E_n}} \quad (\text{B.2.15})$$

In the above derivation, we have assumed that  $d_n$  is large so that  $N_n$  is large since  $N_n \leq d_n$  in (B.1.7). The above is the precursor to the Fermi-Dirac distribution.

### B.2.3 Identical Bosons

In this case

$$G = \sum_{n=1}^{\infty} [\ln((N_n + d_n - 1)!) - \ln(N_n!) - \ln((d_n - 1)!)] \quad (\text{B.2.16})$$

$$+ \alpha \left[ N - \sum_{n=1}^{\infty} N_n \right] + \beta \left[ E - \sum_{n=1}^{\infty} N_n E_n \right] \quad (\text{B.2.17})$$

With Stirling's approximation to the  $N_n$  dependent terms,

$$G = \sum_{n=1}^{\infty} [(N_n + d_n - 1) \ln(N_n + d_n - 1) - (N_n + d_n - 1) - N_n \ln N_n + N_n] \quad (\text{B.2.18})$$

$$- \ln((d_n - 1)!) - \alpha N_n - \beta N_n E_n] + \alpha N + \beta E \quad (\text{B.2.19})$$

Then

$$\frac{\partial G}{\partial N_n} = \ln(N_n + d_n - 1) - \ln(N_n) - \alpha - \beta E_n = 0 \quad (\text{B.2.20})$$

yielding

$$N_n = \frac{d_n - 1}{e^{\alpha + \beta E_n} - 1} \simeq \frac{d_n}{e^{\alpha + \beta E_n} - 1} \quad (\text{B.2.21})$$

where we assume that  $d_n \gg 1$ . The above is the precursor to the Bose-Einstein distribution.

### B.3 The Meaning of $\alpha$ and $\beta$

We will apply the above to a simple quantum system in order to infer what  $\alpha$  and  $\beta$  should be. We consider  $N$  electrons inside a bulk material describable by a single electron in an effective mass approximation. In such a case, the kinetic energy of the electron is described by

$$E_k = \frac{\hbar^2 k^2}{2m} \quad (\text{B.3.1})$$

If we assume periodic boundary conditions on a box of lengths  $L_x$ ,  $L_y$ , and  $L_z$ , then

$$k^2 = \left(\frac{2n_x\pi}{L_x}\right)^2 + \left(\frac{2n_y\pi}{L_y}\right)^2 + \left(\frac{2n_z\pi}{L_z}\right)^2 \quad (\text{B.3.2})$$

In the  $k$  space, there is a state associated with a unit box of volume

$$\Delta V_k = \frac{(2\pi)^3}{L_x L_y L_z} = \frac{8\pi^3}{V} \quad (\text{B.3.3})$$

where  $V$  is the volume of the box of bulk material. In  $k$  space, in a spherical shell of thickness  $\Delta k$ , the number of states in the neighborhood of  $k$  is

$$d_k = \frac{4\pi k^2 \Delta k}{\Delta V_k} = \frac{V}{2\pi^2} k^2 \Delta k \quad (\text{B.3.4})$$

Then the total number of particles is, using (B.2.9) for Maxwell-Boltzmann,

$$\begin{aligned} N &= \sum_k d_k e^{-(\alpha + \beta E_k)} = \frac{V}{2\pi^2} e^{-\alpha} \sum_k k^2 \Delta k e^{-\beta E_k} \\ &= \frac{V}{2\pi^2} e^{-\alpha} \int_0^\infty e^{-\beta \hbar^2 k^2 / (2m)} k^2 dk \end{aligned} \quad (\text{B.3.5})$$

By using the fact that

$$I_1 = \int_0^\infty e^{-Ak^2} k^2 dk = -\frac{d}{dA} \int_0^\infty e^{-Ak^2} dk = -\frac{d}{dA} \frac{1}{2} \sqrt{\frac{\pi}{A}} = \frac{\sqrt{\pi}}{4} A^{-\frac{3}{2}} \quad (\text{B.3.6})$$

the above integrates to

$$N = V e^{-\alpha} \left( \frac{m}{2\pi\beta\hbar^2} \right)^{\frac{3}{2}} \quad (\text{B.3.7})$$

Similarly, the total  $E$  is given by

$$E = \sum_k E_k N_k = \sum_k E_k d_k e^{-(\alpha + \beta E_k)},$$

After using (B.3.1) and (B.3.4), we have

$$\begin{aligned} E &= \frac{V\hbar^2}{4\pi^2m} e^{-\alpha} \int_0^\infty e^{-\beta\hbar^2k^2/2m} k^4 dk \\ &= \frac{3V}{2\beta} e^{-\alpha} \left( \frac{m}{2\pi\beta\hbar^2} \right)^{\frac{3}{2}} = \frac{3N}{2\beta} \end{aligned} \quad (\text{B.3.8})$$

where we have made use of

$$I_2 = \int_0^\infty e^{-Ak^2} k^4 dk = -\frac{d}{dA} I_1 = \frac{3\sqrt{\pi}}{8} A^{-\frac{5}{2}} \quad (\text{B.3.9})$$

and (B.3.7) in getting the last equality.

From the equipartition theorem of energy from statistical mechanics we know that

$$\frac{E}{N} = \frac{3}{2} k_B T \quad (\text{B.3.10})$$

Hence, we conclude that

$$\beta = \frac{1}{k_B T} \quad (\text{B.3.11})$$

It is customary to write

$$\alpha = -\frac{\mu(T)}{k_B T} \quad (\text{B.3.12})$$

so that for Maxwell-Boltzmann distribution,

$$N_n = d_n e^{-(E_n - \mu)/(k_B T)} \quad (\text{B.3.13})$$

where  $\mu(T)$  is the chemical potential. When normalized with respect to degeneracy, we have

$$n_{mb} = e^{-(E_n - \mu)/(k_B T)}, \quad \text{Maxwell-Boltzmann} \quad (\text{B.3.14})$$

Fermi-Dirac and Bose-Einstein distribution become Maxwell-Boltzmann when  $(\beta E_n + \alpha) \gg 1$ . So they have the same  $\alpha$  and  $\beta$ . Therefore,

$$n_{fd} = \frac{1}{e^{(E_n - \mu)/(k_B T)} + 1}, \quad \text{Fermi-Dirac} \quad (\text{B.3.15})$$

$$n_{be} = \frac{1}{e^{(E_n - \mu)/(k_B T)} - 1}, \quad \text{Bose-Einstein} \quad (\text{B.3.16})$$

When  $(E_n - \mu) \gg k_B T$  the above distributions resemble the Maxwell-Boltzmann distribution. This is because  $n$  becomes small per energy level, and the indistinguishability of the particles plays a less important role. When  $(E_n - \mu) < 0$ , the Fermi-Dirac distribution “freezes” to 1 for all energy levels below  $\mu$ , since the Pauli’s exclusion principle only allows one fermion per level. When  $(E_n - \mu) \ll k_B T$ , the Bose-Einstein distribution diverges, implying that the particles condense to an energy level close to  $\mu$ . This condensation is more pronounced if  $k_B T$  is small.



# Appendix C

## Gaussian Wave Packet

### C.1 Introduction

In order to relate the wave nature to the particle nature of an object, it is necessary to obtain a wave packet picture of the object. Experimental observation has indicated that electrons are highly localized objects, so are photons. A wavefunction with a pure or single wavenumber  $k$  has equal amplitude everywhere. According to the probabilistic interpretation of quantum mechanics, it is untenable physically to think of a particle as all pervasive and can be found equally likely everywhere. A photon wave can have many  $k$  wavenumbers, and they can be linearly superposed to form a localized photon wave packet. Since the  $E$ - $k$  relationship of the photon wave is linear (or almost linear in the material-media case), the formation of such a wave packet is quite straightforward. It is less clear for an electron wave, since the  $E$ - $k$  relationship is nonlinear.

In this appendix, we will explain the particle nature of an electron in the classical limit when the momentum of the electron becomes large. One expects that the wavefunction resembles that of a localized particle in this limit so that the electron is found with high probability only in a confined location. This understanding can be achieved by studying the Gaussian wave packet solution of Schrödinger equation.

When the Schrödinger equation is solved in vacuum for electron, the solution is

$$e^{i\mathbf{k}\cdot\mathbf{r}}$$

This unlocalized solution means that the electron can be everywhere. A more realistic wavefunction for an electron is a wave packet which is localized. This is more akin to the motion that an electron is a localized particle. A wave packet can be constructed by linear superposing waves of different momenta  $\hbar\mathbf{k}$  or different energies  $\hbar\omega$ . We will derive the Gaussian wave packet solution to Schrödinger equation. This can be constructed by studying the solution to the wave equation.

## C.2 Derivation from the Wave Equation

It is well known that the wave equation

$$\left( \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} + \frac{\partial^2}{\partial z^2} + k^2 \right) \phi(x, y, z) = 0 \quad (\text{C.2.1})$$

admits solution of the form

$$\phi(x, y, z) = C \frac{e^{\pm ik\sqrt{x^2+y^2+(z-ib)^2}}}{\sqrt{x^2+y^2+(z-ib)^2}} \quad (\text{C.2.2})$$

for

$$\sqrt{x^2+y^2+(z-ib)^2} \neq 0 \quad (\text{C.2.3})$$

The above corresponds to a spherical wave where the source point is located in the complex coordinates  $x = 0$ ,  $y = 0$ , and  $z = ib$ . It was first suggested by G. Deschamps. From the above, it is quite clear that, after letting  $z = vt$ ,

$$\phi(x, y, vt) = A \frac{e^{\pm ik\sqrt{x^2+y^2+(vt-ib)^2}}}{\sqrt{x^2+y^2+(vt-ib)^2}} \quad (\text{C.2.4})$$

is a solution to

$$\left( \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} + \frac{1}{v^2} \frac{\partial^2}{\partial t^2} + k^2 \right) \phi(x, y, vt) = 0 \quad (\text{C.2.5})$$

The above equation can be factored

$$\left( v^{-1} \partial_t - i\sqrt{k^2 + \partial_x^2 + \partial_y^2} \right) \left( v^{-1} \partial_t + i\sqrt{k^2 + \partial_x^2 + \partial_y^2} \right) \phi(x, y, vt) = 0 \quad (\text{C.2.6})$$

where  $\partial_t = \partial/\partial t$ ,  $\partial_x^2 = \partial^2/\partial x^2$  and so on. In the above, function of an operator has meaning only when the function is Taylor expanded into an algebraic series. One can assume that  $k^2 \rightarrow \infty$ , while  $\partial_x^2$  and  $\partial_y^2$  are small.<sup>1</sup> Taylor expanding the above and keeping leading order terms only, we have

$$\left( v^{-1} \partial_t - ik - \frac{i}{2k} (\partial_x^2 + \partial_y^2) \right) \left( v^{-1} \partial_t + ik + \frac{i}{2k} (\partial_x^2 + \partial_y^2) \right) \phi(x, y, vt) \cong 0 \quad (\text{C.2.7})$$

In (C.2.4), we can let<sup>2</sup>  $|vt - ib|^2 \gg x^2 + y^2$  to arrive at the approximation

$$\phi(x, y, vt) \approx A \frac{e^{\pm \left( ik(vt-ib) + ik \frac{x^2+y^2}{vt-ib} \right)}}{vt-ib} \quad (\text{C.2.8})$$

<sup>1</sup>This is known as the paraxial wave approximation. It implies that the variation of the solution is mainly in the  $vt$  direction with slow variation in the  $x$  and  $y$  directions.

<sup>2</sup>It can be shown that when this is valid, the paraxial wave approximation used in (C.2.7) is good.

It can be shown that when we pick the plus sign above in the  $\pm$  sign, the above is the exact solution to

$$\left(v^{-1}\partial_t - ik - \frac{i}{2k}(\partial_x^2 + \partial_y^2)\right)\phi_+(x, y, vt) = 0 \quad (\text{C.2.9})$$

where

$$\phi_+(x, y, vt) = Ae^{ik(vt-ib)}\frac{e^{ik\frac{x^2+y^2}{2(vt-ib)}}}{2(vt-ib)} = e^{ik(vt-ib)}\psi(x, y, t) \quad (\text{C.2.10})$$

and

$$\psi(x, y, t) = -A_0\frac{ibe^{ik\frac{x^2+y^2}{2(vt-ib)}}}{(vt-ib)} \quad (\text{C.2.11})$$

Furthermore,  $\psi(x, y, t)$  is an exact solution to

$$\left(v^{-1}\partial_t - \frac{i}{2k}(\partial_x^2 + \partial_y^2)\right)\psi(x, y, t) = 0 \quad (\text{C.2.12})$$

If we multiply the above by  $i\hbar v$ , the above becomes

$$\left(i\hbar\partial_t + \frac{\hbar v}{2k}(\partial_x^2 + \partial_y^2)\right)\psi(x, y, t) = 0 \quad (\text{C.2.13})$$

By letting  $v = \hbar k/m$ , the above becomes

$$\left(i\hbar\partial_t + \frac{\hbar^2}{2m}(\partial_x^2 + \partial_y^2)\right)\psi(x, y, t) = 0 \quad (\text{C.2.14})$$

which is just Schrödinger equation. Equation (C.2.11) represents the Gaussian wave packet solution of Schrödinger equation. It can be studied further to elucidate its physical contents.

### C.3 Physical Interpretation

In (C.2.11), one can write

$$\frac{k(x^2 + y^2)}{2(vt - ib)} = \frac{k(x^2 + y^2)(vt + ib)}{2(v^2t^2 + b^2)} = \frac{k(x^2 + y^2)}{2R} + i\frac{x^2 + y^2}{W^2} \quad (\text{C.3.1})$$

where

$$R = \frac{v^2t^2 + b^2}{vt}, \quad W^2 = \frac{2b}{k}\left(1 + \frac{v^2t^2}{b^2}\right) \quad (\text{C.3.2})$$

Then Gaussian wave packet can be more suggestively written as

$$\psi(x, y, t) = \frac{A_0}{\sqrt{1 + v^2t^2/b^2}}e^{-\frac{x^2+y^2}{W^2}}e^{ik\frac{x^2+y^2}{2R}}e^{-i\varphi(t)} \quad (\text{C.3.3})$$

where

$$\varphi(t) = \tan^{-1}(vt/b) \quad (\text{C.3.4})$$

The above reveals a wave packet which is Gaussian tapered with width  $W$  and modulated by oscillatory function of space and time. However, the width of this packet is a function of time as indicated by (C.3.2).

At  $vt = 0$ , the width of the packet is given by

$$W_0 = \sqrt{\frac{2b}{k}}, \quad \text{or} \quad b = \frac{1}{2}kW_0^2 \quad (\text{C.3.5})$$

This width can be made independent of  $k$  if  $b$  is made proportional to  $k$ . Nevertheless, as time progresses with  $vt > 0$ , the width of the packet grows according to (C.3.2). However, to maintain a fixed-width  $W_0$ , it is necessary that  $b \rightarrow \infty$  as  $k \rightarrow \infty$ . Subsequently, the effect of  $v^2t^2/b^2$  becomes small in (C.3.2) as  $b \rightarrow \infty$ . This means that the width of the Gaussian wave packet remains almost constant for the time when  $vt \leq b$ , but  $b$  is a large number proportional to  $k$ . The duration over which the Gaussian wave packet's width remains unchange becomes increasingly longer as  $k$  becomes larger.

In the above,  $\hbar k$  represents the momentum of a particle. When the particle carries high momentum, it can be represented by a Gaussian wave packet that hardly changes in shape with respect to time. This is what is expected of a classical picture of a moving particle: the Gaussian wave packet does reproduce the classical picture of a high momentum particle.

It is quite easy to design a wave packet for a photon that does not spread with time by linear superposing waves with different frequencies. This is because the  $\omega$ - $k$  diagram (or  $E$ - $k$  diagram, since  $E = \hbar\omega$ ) for photons is a straight line. It implies that all waves with different  $k$ 's travel with the same phase velocity. These Fourier modes are locked in phase, and the pulse shape does not change as the wave packet travels.

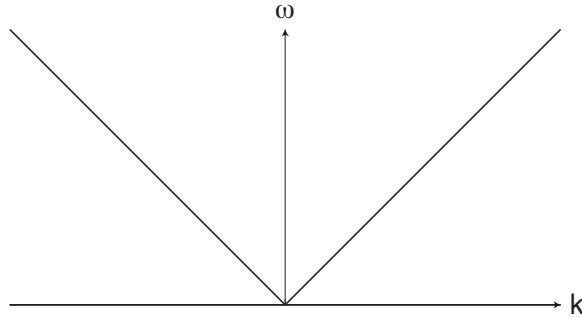


Figure C.1:  $\omega - k$  diagram for a photon wave which is a straight line.

But for electron wave in a vacuum, the  $\omega$ - $k$  diagram is quadratic. It implies that waves with different  $k$  numbers travel with different phase velocity, giving rise to distortion of the pulse shape. But if we have a narrow band pulse operating in the high  $k$  regime, the  $\omega$ - $k$  diagram is quasi-linear locally, and there should be little pulse spreading. Hence, one can construct a quasi-distortionless pulse in the high  $k$  regime.



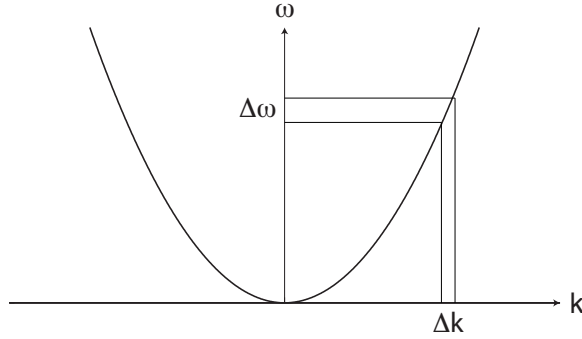


Figure C.2:  $\omega - k$  diagram of an electron wave in vacuum.

### C.4 Stability of the Plane Wave Solution

It is quite obvious from our understanding of quantum mechanics, wave packets, and coherence, that the plane wave solution  $e^{ikx}$  is not stable as  $k \rightarrow \infty$ . One can always express

$$e^{ikx} = \int_{-\infty}^{\infty} dx' e^{ikx'} \delta(x - x') \tag{C.4.1}$$

If one can think of  $\delta(x - x')$  as the limiting case of a wave packet, the above implies that a plane wave can be expanded as a linear superposition of wave packets at each location  $x$ , but bearing a phase  $\exp(ikx)$ . As  $k \rightarrow \infty$ , this phase is rapidly varying among the different wave packets. Hence, their coherence is extremely difficult to maintain, and upset easily by coupling to other quantum systems that exists in the environment. In other words, the particle cannot be stably maintained in the plane-wave state classically: it has to collapse to a wave-packet state. Consequently, classically, particles are localized when its momentum  $k$  becomes large.

