MATHEMATICAL METHODS IN QUANTUM MECHANICS

Lecture notes by Răzvan Gelca

Contents

This course was taught for the first time in the Department of Mathematics and Statistics at Texas Tech University in Spring 2015, the second time in Fall 2017, and the third time in Fall 2020. Most of the materials were taken from the lecture notes of the Quantum Mechanics course that I took at University of Michigan with Alejandro Uribe, and from the books *Quantum Theory for Mathematicians* by Brian Hall, *Lectures on Quantum Mechanics for Mathematics Students* by L.D. Faddeev and O.A. Yakubovskii, and *Theta Functions and Knots* by myself.

Chapter 1

Classical Mechanics

This chapter is inspired by V.I. Arnol'd's *Mathematical Methods of Classical Mechanics*, and the students can consult this book for further details and examples. Because classical mechanics we can "see", it is phrased in the language of geometry. Originally geometry was an intuitive science of points, lines, and circles, but the coordinates introduced by Descartes were a powerful tool, and they led to differential geometry. It is in the language of differential geometry that we will tell the story of classical mechanics, and we do it from two points of view: the lagrangian and the hamiltonian formalism.

1.1 The Lagrangian formalism of classical mechanics

1. Lagrangian mechanics describes the motion of a system using the configuration space. A Lagrangian mechanical system is defined by a *finite dimensional manifold* (which parametrizes all possible configurations of the mechanical system) and a *function* on the tangent bundle to this manifold (the Lagrangian).

A smooth n-dimensional manifold M looks locally like $\mathbb{R}^n$. In formulas we will use a local coordinate $\mathbf{q} = (q_1, q_2, \ldots, q_n)$. By Whitney's embedding theorem, there is a smooth embedding $f : M \to \mathbb{R}^{2n}$. Thus we arrive at Poincarè's description of a manifold: a set in $\mathbb{R}^N$ that is locally the graph of a map $\phi : \mathbb{R}^r \to \mathbb{R}^s$.

Definition. The abstract definition of a smooth manifold (due to O. Veblen) is that it is a topological space M together with an open covering $(U_\alpha)_{\alpha \in A}$ and a collection of maps $\phi_\alpha : U_\alpha \to \mathbb{R}^n$ whose images are open sets and the maps ϕ_α are homeomorphisms onto these images, such that $\phi_\beta \circ \phi_\alpha^{-1}$ is smooth (whenever this map is defined).

The maps $\phi_\beta \circ \phi_\alpha^{-1}$ are the changes of coordinates from one local chart to another, if the two charts overlap.

The tangent space at a point $\mathbf{q}_0$ to the manifold M, $T_{\mathbf{q}_0}M$, can be defined abstractly as the set of equivalence classes of curves that have the same derivative at $\mathbf{q}_0$ in some system of local coordinates. Concretely, it is the vector space of tangent vectors at $f(\mathbf{q}_0)$ to $f(M)$ (here f is the embedding of M into some $\mathbb{R}^n$). A vector is of the form $\frac{d}{dt}(f \circ \mathbf{q})(t)|_{t=0}$, were $\mathbf{q}(t)$ is a curve, $t \in (-\epsilon, \epsilon)$, with $\mathbf{q}(0) = \mathbf{q}_0$. It is customary to denote the tangent vector to the curve $\mathbf{q}$ by $\dot{\mathbf{q}}$, or more precisely, the tangent vector to the curve $\mathbf{q}$ at point $\mathbf{q}(t)$ by $\dot{\mathbf{q}}(t)$.

The coordinates of a vector are computed in a local chart as $\frac{d}{dt}(\phi_\alpha \circ f \circ \mathbf{q})dt$. Then, if in the local chart U_α we use coordinates $\mathbf{x} = (x_1, x_2, \ldots, x_n)$, we identify $T_{\mathbf{q}_0}M$ with the vector space with basis $\partial/\partial x_1, \partial/\partial x_2, \ldots, \partial/\partial x_n$. Also, if in the local chart U_β we use coordinates $\mathbf{y} = (y_1, y_2, ..., y_n)$ and we identify $T_{\mathbf{q}_0}M$ with the vector space with basis $\partial/\partial y_1, \partial/\partial y_2, \ldots, \partial/\partial y_n$, then for $\mathbf{v} = (\alpha_1, \alpha_2, \ldots, \alpha_n)$ in the coordinate chart U_α then in the coordinate chart U_β, $\mathbf{v}$ has coordinates

$$\begin{pmatrix} \beta_1 \\ \beta_2 \\ \cdots \\ \beta_n \end{pmatrix} = \left(\frac{\partial(\phi_\beta \circ \phi_\alpha^{-1})_j}{\partial x_k} \right)_{jk} \begin{pmatrix} \alpha_1 \\ \alpha_2 \\ \cdots \\ \alpha_n \end{pmatrix}.$$

Formally we can think that the coordinates in the chart ϕ_α are $\mathbf{x}$ and in the chart ϕ_β are $\mathbf{y}$, and then the matrix for the change of coordinates of vectors is $\frac{\partial \mathbf{y}}{\partial \mathbf{x}}$.

The tangent bundle, TM, is the union of all tangent spaces, which is given a structure of a smooth manifold using the charts defined by ϕ_α and the coordinates of the tangent vector in the chart ϕ_α. The transition functions (from one coordinates system to another) are

$$\left(\phi_\beta \circ \phi_\alpha^{-1}, \left(\frac{\partial(\phi_\beta \circ \phi_\alpha^{-1})_j}{\partial x_k} \right)_{jk} \right)$$

Locally it is diffeomorphic to $\mathbb{R}^n \times \mathbb{R}^n$. Heuristically we think of it as being the "system of positions and velocities of a system of several particles with constraints".

Example 1.1.1. Suppose that you have a rod that has one end fixed in the space and the other moves freely (like a pendulum). If you have a particle attached to the other end, then this particle is constrained to a sphere. All possible velocities of the particle determine the tangent bundle to the sphere. The situation is described in Figure 1.1.

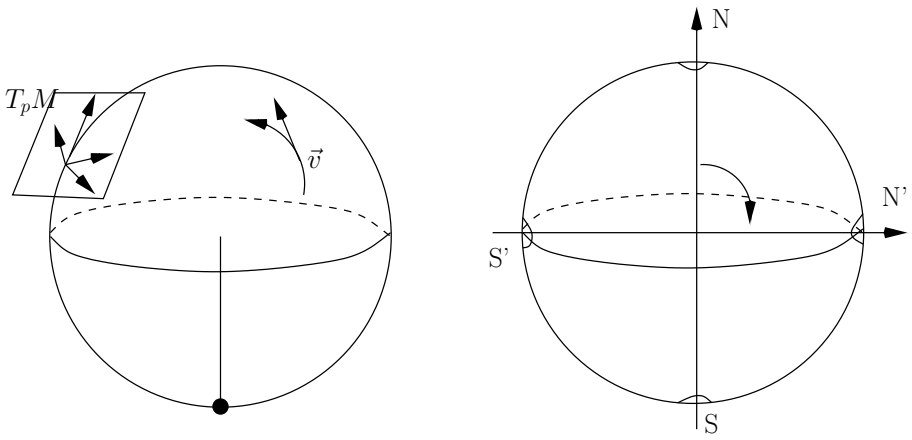

Figure 1.1: The sphere.

You can put Euler's coordinates on the sphere:

$$x = \sin\theta \cos\phi$$
$$y = \sin\theta \sin\phi$$
$$z = \cos\theta,$$

but these coordinates are ambiguous when $\theta = 0$ (north pole) and $\theta = \pi$ (south pole). The longitude ϕ is also ambiguous, as it is only defined modulo 2π. For disambiguation, remove the north pole and the south pole. Split into two charts, U_1 which does not contain the meridian 0, and U_2 which does not contain the meridian π. Now rotate the sphere by 90° so that the z-axis becomes the x-axis, and repeat to construct the charts U_3 and U_4. Now each chart is parametrized by a domain in the plane (two coordinates), with no ambiguity. The domain can be chosen as $(0, \pi) \times (0, 2\pi)$.

Example 1.1.2. If you have a charged particle and you want to keep it in a compact zone of the space you can use a magnetic field. Hopf's index theorem prevents you to keep it on a sphere (or inside a sphere), but you can force it to stay on the surface of a torus (Figure 1.2).

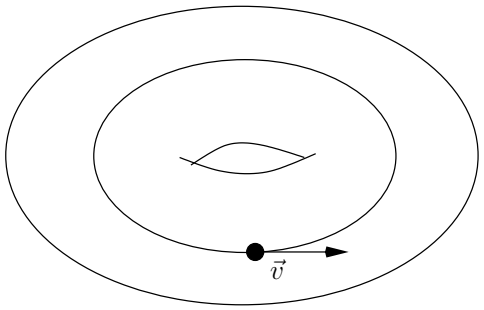

Figure 1.2: The torus.

You can identify the torus with the cartesian product of two circles (beware, the torus can be embedded in the 3-dimensional space but the cartesian product of two circles cannot be embedded isometrically in the space) and so it can be parametrized by $(e^{2\pi is}, e^{2\pi it})$. But again the parameters s, t are ambiguous, they are only defined modulo 2π. To turn this parametrization into an atlas, let $S = (0, 1) \times (0, 1)$ let also

$$U_{jk} = \{(e^{2\pi is}, e^{2\pi it}), s \neq j, t \neq k\}, \quad j, k = 0, \frac{1}{2}.$$
$$\phi_{jk}^{-1} : U_{jk} \to D, \phi_{jk}^{-1}(s, t) = (e^{2\pi is}, e^{2\pi it}).$$

The map $\phi_{jk} \circ \phi_{ml}^{-1}$ is just the translation by the vector $(m - j, l - k)$, i.e. $(x, y) \mapsto (x + m - j, y + l - k)$, which is smooth.

2. Let $L : TM \times \mathbb{R} \to \mathbb{R}$ be a smooth function, which plays the role of the Lagrangian of the system; it is particular to the given mechanical system and it is used to describe the evolution of the mechanical system as follows. Fix points $\mathbf{q}_0$ and $\mathbf{q}_1$ on M and consider the infinite dimensional space of smooth curves γ given by functions of the form $\mathbf{q} : [t_0, t_1] \to M$ such that $\mathbf{q}(t_0) = \mathbf{q}_0$ and $\mathbf{q}(t_1) = \mathbf{q}_1$. On this space define the functional

$$\Phi(\gamma) = \int_{t_0}^{t_1} L(\mathbf{q}, \dot{\mathbf{q}}, t) dt.$$

Here as the variable of L you plug in the point $\mathbf{q}(t) \in M$, the tangent vector $\dot{\mathbf{q}}(t)$ to the curve at this point, and the time t.

Hamilton's Minimal Action Principle: The motions of a mechanical system described by L coincide with the extrema of the functional, $\Phi(\gamma)$, called *action*.

A curve is an extremum of Φ if the "derivative" of Φ is zero at this curve. This is made precise within the framework of *variational calculus*.

Here we view M as a subset of some $\mathbb{R}^n$, and then consider paths γ in M and variations of these paths, written as $\gamma+\delta$, also in M. The differential of Φ is a linear functional F such that

$$\Phi(\gamma+\delta)-\Phi(\gamma)=F(\delta)+R$$

where $R(\gamma,\delta)=O(\delta^2)$ (if δ and its derivative are less than ϵ in absolute value, then $|R|<\epsilon^2$). We consider only variations in the space of paths in M that connect $\mathbf{q}_0$ and $\mathbf{q}_1$. So δ itself, which is a path in $\mathbb{R}^n$, is zero at t_0 and t_1. We have

$$\begin{aligned}\Phi(\gamma+\delta)-\Phi(\gamma)&=\int_{t_0}^{t_1}[L(\mathbf{q}+\mathbf{h},\dot{\mathbf{q}}+\dot{\mathbf{h}},t)-L(\mathbf{q},\dot{\mathbf{q}},t)]dt\\&=\int_{t_0}^{t_1}\left[\frac{\partial L}{\partial\mathbf{q}}\mathbf{h}+\frac{\partial L}{\partial\dot{\mathbf{q}}}\dot{\mathbf{h}}\right]dt+O(\mathbf{h}^2)=\int_{t_0}^{t_1}\left[\frac{\partial L}{\partial\mathbf{q}}-\frac{d}{dt}\left(\frac{\partial L}{\partial\dot{\mathbf{q}}}\right)\right]\mathbf{h}dt+O(\mathbf{h}^2).\end{aligned}$$

where for the last step we used integration by parts (and the fact that $\mathbf{h}$ is zero at the endpoints).

So the differential of Φ is this last integral, and the extrema of the action are those trajectories $t\mapsto\mathbf{q}(t)$ that satisfy the *Euler-Lagrange equations:*

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

3. Here is a reformulation of Newtonian mechanics (in $\mathbb{R}^n$). Let $L=T-V$ where $T=\sum_j m_j\dot{q_j}^2/2$ is the kinetic energy and $V=V(\mathbf{q})$ is the potential energy. Then

$$\frac{\partial L}{\partial\dot{q_j}}=\frac{\partial T}{\partial\dot{q_j}}=m_j\dot{q_j},\quad\frac{\partial L}{\partial q_j}=-\frac{\partial V}{\partial q_j}.$$

We obtain Newton's equations

$$m_j\ddot{q_j}+\frac{\partial V}{\partial q_j}=0,$$

which can be written in the more familiar form $ma_j=F_j$, were a_j is the jth component of the acceleration.

Remark 1.1.1. Newton's equations show that the *total energy* $E=T+V$ is conserved (i.e. does not change it time).

Example 1.1.3. A free particle moving in $\mathbb{R}$ has $L=T=\frac{m\dot{q}^2}{2}$. The Euler-Lagrange equation is

$$\frac{d}{dt}(m\dot{q})=0$$

which is equivalent to $\ddot{q}=0$. This means we are in the presence of uniform motion.

Example 1.1.4. The harmonic oscillator with no damping has a single force acting, and this is given by Hooke's law: $F = -kq$. We (forcefully) introduce the potential $V = \frac{kq^2}{2}$. Then

$$L = \frac{m\dot{q}^2}{2} - \frac{kq^2}{2}.$$

Then

$$\frac{\partial L}{\partial q} = -kq, \quad \frac{\partial L}{\partial \dot{q}} = m\dot{q},$$

so we obtain the equation of the harmonic oscillator

$$m\ddot{q} + kq = 0,$$

whose solutions are of the form $q(t) = c_1 \cos \omega t + c_2 \sin \omega t$ with $\omega = \sqrt{k/m}$.

1.2 The Hamiltonian formalism of classical mechanics

4. To pass from the Lagrangian to the Hamiltonian formulation of classical mechanics we use the Legendre transform. This transform is still part of what is traditionally called Lagrangian mechanics, but it establishes the transition from one formulation to the other and turns Lagrangian mechanics into a subtheory of Hamiltonian mechanics.

First we assume, as it is usually the case in real life applications, that the potential energy depends on $\mathbf{q}$ only and the kinetic energy is a positive definite quadratic function in $\dot{\mathbf{q}}$. Then $L(\mathbf{q}, \dot{\mathbf{q}}, t)$ is convex in $\dot{\mathbf{q}}$.

In general for a convex function $f(x)$ the Legendre transform is a function g defined as follows. For a number p, the function $F(x, p) = px - f(x)$ has a maximum at some point $x(p)$. Then $g(p) = F(p, x(p))$.

So in our case $F(\dot{\mathbf{q}}, \mathbf{p}) = \mathbf{p}\dot{\mathbf{q}} - L(\mathbf{q}, \dot{\mathbf{q}}, t)$. To maximize we set its derivative with respect to $\dot{\mathbf{q}}$) equal to 0, and get $\mathbf{p} - \frac{\partial L}{\partial \dot{\mathbf{q}}} = 0$, that is $\mathbf{p} = \frac{\partial L}{\partial \dot{\mathbf{q}}}$. So the Legendre transform of L (with respect to the variable $\dot{\mathbf{q}}$) is the Hamiltonian function

$$H(\mathbf{p}, \mathbf{q}, t) = \mathbf{p}\dot{\mathbf{q}} - L(\mathbf{q}, \dot{\mathbf{q}}, t),$$

with variables the positions $\mathbf{q}$ and the (conjugate) *momenta* $\mathbf{p} = \frac{\partial L}{\partial \dot{\mathbf{q}}}$. It is obvious that $\partial H/\partial \mathbf{q} = -\partial L/\partial \mathbf{q}$. Here is a slightly more complicated computation using the chain rule, in which we point out that $\dot{\mathbf{q}}(p)$ is obtained by solving the equation $\mathbf{p} - \frac{\partial L}{\partial \dot{\mathbf{q}}} = 0$:

$$\frac{\partial H}{\partial \mathbf{p}} = \dot{\mathbf{q}} + \mathbf{p}\frac{\partial \dot{\mathbf{q}}}{\partial \mathbf{p}} - \frac{\partial L}{\partial \dot{\mathbf{q}}}\frac{\partial \dot{\mathbf{q}}}{\partial \mathbf{p}} = \dot{\mathbf{q}} + \mathbf{p}\frac{\partial \dot{\mathbf{q}}}{\partial \mathbf{p}} - \mathbf{p}\frac{\partial \dot{\mathbf{q}}}{\partial \mathbf{p}} = \dot{\mathbf{q}}.$$

We can now see thta Euler-Lagrange equations are equivalent to Hamilton's equations

$$\dot{\mathbf{q}} = \frac{\partial H}{\partial \mathbf{p}}$$
$$\dot{\mathbf{p}} = -\frac{\partial H}{\partial \mathbf{q}}.$$

Indeed

$$-\frac{\partial H}{\partial \mathbf{q}} = \frac{\partial L}{\partial \mathbf{q}} = \frac{d}{dt}\frac{\partial L}{\partial \dot{\mathbf{q}}} = \frac{d\mathbf{p}}{dt} = \dot{\mathbf{p}}$$

by the Euler-Lagrange equations. Aha, so by introducing the momenta we turn a system of $n = \dim M$ second order differential equations into a system of $2n$ first order differential equations! So we have a change of variable and a reduction of order.

In the setting of Newtonian mechanics $\mathbf{p}\dot{\mathbf{q}}$ is twice the kinetic energy, so the Hamiltonian is $H = T + V$, which is the total energy of the system.

Example 1.2.1. For the harmonic oscillator,

$$p = \frac{\partial L}{\partial \dot{q}} = \frac{\partial}{\partial \dot{q}}\left(\frac{m\dot{q}^2}{2} - \frac{kq^2}{2}\right) = m\dot{q}.$$

Hence

$$H = \frac{p^2}{2m} + \frac{kq^2}{2}.$$

Hamilton's equations are

$$\begin{aligned} \dot{q} &= \frac{p}{m} \\ \dot{p} &= -kq. \end{aligned}$$

So the second order Euler-Lagrange equation now becomes Hamilton's system of first order differential equations. To solve this system you set

$$A = \begin{pmatrix} 0 & \frac{1}{m} \\ -q & 0 \end{pmatrix}.$$

and then write the system as $\frac{d\mathbf{x}}{dt} = A\mathbf{x}$. The solution is of the form $\mathbf{x}(t) = e^{tA}\mathbf{x}(0)$, and of course to find the exponential of the matrix you need to diagonalize it, and in the process of diagonalization you get the same characteristic equation as for the second order equation, etc...

5. Now we discuss a subtle point, namely the fact that one of the main differences between the Lagrangian and the Hamiltonian point of view is that one of them happens on the *tangent* bundle while the other happens on the *cotangent* bundle of the configuration space.

In general for a function $f(\mathbf{x})$, the gradient

$$\frac{\partial f}{\partial \mathbf{x}} = \left(\frac{\partial f}{\partial x_1}, \frac{\partial f}{\partial x_2}, \ldots, \frac{\partial f}{\partial x_n}\right)$$

is a covector. You can think of it as being the differential form $\frac{\partial f}{\partial \mathbf{x}} d\mathbf{x}$.

Because of the chain rule

$$\frac{\partial f}{\partial \mathbf{y}} = \frac{\partial f}{\partial \mathbf{x}}\frac{\partial \mathbf{x}}{\partial \mathbf{y}},$$

the matrix that changes the "vector" $\frac{\partial f}{\partial \mathbf{x}}$ to the "vector" $\frac{\partial f}{\partial \mathbf{y}}$ is

$$\frac{\partial \mathbf{x}}{\partial \mathbf{y}}^T = \left(\left(\frac{\partial \mathbf{y}}{\partial \mathbf{x}}\right)^{-1}\right)^T = \left(\left(\frac{\partial (\phi_\beta \circ \phi_\alpha^{-1})_j}{\partial x_k}\right)^{-1}_{jk}\right)^T.$$

This applies to Hamiltonian mechanics, which is obtained by replacing the variables $\dot{\mathbf{q}}$, which are tangent vectors, by the variables $\mathbf{p} = \partial L/\partial \dot{\mathbf{q}}$, which are covectors.

We can take the union of the cotangent space at every point of the configuration space M to produce the cotangent bundle T^*M. This can be endowed with the structure of a smooth manifold using the charts and the coordinates of forms. Locally T^*M looks like some $\mathbb{R}^{2n}$, where n is the number of degrees of freedom (parameters) of the classical system. The transition maps from one system of coordinates to another are

$$\left(\phi_\beta \circ \phi_\alpha^{-1}, \left(\left(\frac{\partial (\phi_\beta \circ \phi_\alpha^{-1})_j}{\partial x_k}\right)^{-1}_{jk}\right)^T\right).$$

So T^*M is a manifold, of *even dimension*. As we will see below, it is a manifold of a special type.

So to conclude, the Legendre transform maps functions that depend of the position and velocity, i.e. point and tangent vector, to functions that depend on the position and the momentum, i.e. point and cotangent vector.

You can think of this in terms of directional derivatives as follows. If V is a vector space and $f : V \to \mathbb{R}$ is a function, then for $v \in V$, $df_v \in V^*$. Indeed,

$$df_v(w) = \frac{d}{dt}(f(v + tw))|_{t=0}, \quad w \in V$$

which is a linear functional on V. We thus get a linear map $V \mapsto V^*$. For $f = L(\dot{\mathbf{q}})$, and $V = T_{\mathbf{q_0}}$ we obtain a map from $T_{\mathbf{q}_0}M \mapsto T^*_{\mathbf{q}_0}M$, which defines the Legendre transform $TM \to T^*M$.

So Hamiltonian mechanics happens on the *cotangent bundle* of the manifold of configurations. If M is the manifold of configurations, then its cotangent bundle T^*M is called the *phase space* of the classical mechanical system. The phase space parametrizes the positions and momenta of the mechanical system. Locally T^*M looks like some $\mathbb{R}^{2n}$, where n is the number of degrees of freedom (parameters) of the classical system.

Definition. A symplectic manifold is a manifold endowed with a globally defined symplectic form. A symplectic form is a nondegenerate closed 2-form.

We could think of a 2-form ω as an oriented area element, it tells us how to integrate functions on surfaces in M. The idea is that when computing the integral using Riemann sums, each term of the sum is of the form $f(t_j)\mu(D_j)$, where $\mu(D_j)$ is the signed area of a small domain D_j multiplied by some weight. In the limit we can choose D_j to be almost a parallelogram determined by two tangent vectors to the surface $\mathbf{v}_1$ and $\mathbf{v}_2$. The 2-form associates to the pair $(\mathbf{v}_1, \mathbf{v}_2)$ the number that is the signed area of the parallelogram formed by them multiplied by some weight that varies continuously with the point.

In other words 2-form is an object that, when restricted to a point on the manifold, it yields a bilinear antisymmetric map $T_{\mathbf{q}}M \times T_{\mathbf{q}}M \to \mathbb{R}$, and these bilinear maps vary smoothly with the point. A 2-form is called nondegenerate if for every vector $v \neq 0$ there is a vector w such that $\omega(v, w) \neq 0$. In local coordinates $(\mathbf{p}, \mathbf{q})$, the form can be written as

$$\omega = \sum_{j,k} f_{j,k}(\mathbf{p}, \mathbf{q}) dp_j \wedge dq_k = \sum_{j,k} f_{j,k}(\mathbf{p}, \mathbf{q}) d\mathbf{p} \wedge d\mathbf{q}.$$

Because the we work with signed areas, the sign changes when you flip a dp or dq over a dp or dq. Its differential is

$$d\omega = \sum_{l} \sum_{j,k} \left(\frac{\partial f_{j,k}}{\partial p_l} dp_l + \frac{\partial f_{j,k}}{\partial q_l} dq_l \right) \wedge dp_j \wedge dq_k.$$

The form is called closed if its differential is zero.

Theorem 1.2.1. The cotangent bundle T^*M is a symplectic manifold with the symplectic form

$$\omega = d\mathbf{p} \wedge d\mathbf{q} = \sum_{j=1}^{n} dp_j \wedge dq_j.$$

Proof. In local coordinates it is not hard to see that this form is nondegenerate and closed. It is harder to see that it is well defined. To prove this, we define a 1-form. Let v be a vector tangent to T^*M at a point $(\mathbf{q}, \mathbf{p})$. The differential of the projection $\pi : T^*M \to M$ maps v to $\pi_* v$. Let $\theta(v) = \mathbf{p}(\pi_* v)$. In local coordinates $\theta = \mathbf{p} d\mathbf{q}$, so $d\theta = \omega$. Note that $d\omega = d^2\theta = 0$. The form θ is called the canonical 1-form. □

6. Hamiltonian mechanics can be defined on a general symplectic manifold, not necessarily of the form T^*M, so we allow constraints at the level of momenta as well (not just at the level of positions). The two dimensional torus is a good example of a symplectic manifold that is not of the form T^*M.

Therefore, in Hamiltonian mechanics, the phase space is a pair (M, ω), where M is an even dimensional real manifold and ω is a symplectic form on M.

Theorem 1.2.2. (Darboux) Given a symplectic manifold (M, ω), every point has a neighborhood and a system of coordinates of that neighborhood in which the symplectic form is $d\mathbf{p} \wedge d\mathbf{q}$.

So for local computations we can assume that we are in $\mathbb{R}^{2n}$ with the standard symplectic form. For local computations we use the formulas

$$\omega\left(\frac{\partial}{\partial p_j}, \frac{\partial}{\partial q_k}\right) = \delta_{jk}, \omega\left(\frac{\partial}{\partial q_j}, \frac{\partial}{\partial p_k}\right) = -\delta_{jk}, \omega\left(\frac{\partial}{\partial p_j}, \frac{\partial}{\partial p_k}\right) = \omega\left(\frac{\partial}{\partial q_j}, \frac{\partial}{\partial q_k}\right) = 0,$$

The symplectic form defines an isomorphism

$$T_{\mathbf{q}_0}M \mapsto T^*_{\mathbf{q}_0}M$$

by $v \mapsto \omega(\cdot, v)$ where the latter maps a vector w to $\omega(w, v)$ (there is a fancy notation: $i_v\omega = \omega(v, \cdot)$).

Now we develop a formalism for **observable quantities**. This will bring us closer to the points of view of quantum mechanics. In classical mechanics we can see the particle, so we can determine its position and velocity precisely, or, in the present formalism, its position and momentum. Then we associate various quantities that are physically relevant and sometimes arise from conservation laws, such as angular momentum, total energy, kinetic energy, etc. But in quantum mechanics we cannot see particles. All we can do is set up some experiment which detects the existence of the particle and its "quantum" properties. Usually the outcome of the experiment is some data. We can reason the same way in the classical world, and think that what we measure are some numerical quantities that depend on the state of the particle. We thus are looking at functions of $(\mathbf{p}, \mathbf{q})$. For example, the kinetic energy is $T(p, q) = \frac{\mathbf{p}^2}{2m}$, the angular momentum is the vector valued function $\mathbf{L} = \mathbf{q} \times \mathbf{p}$, the total energy of the harmonic oscillator is $H(p, q) = \frac{p^2}{2m} + \frac{kq^2}{2}$ (here p and q are 1-dimensional variables as the oscillation happens on a line). The angular momentum has three components in the direction of the coordinate axes which are true functions $L_i = q_j p_k - q_k p_j$, where (i, j, k) is a cyclic permutation of $(1, 2, 3)$. We now develop the formalism for such functions, which we later translate to quantum mechanics via a procedure called quantization.

The symplectic form associates to each function f a Hamiltonian vector field $\mathbf{X}_f$ defined as the inverse through this isomorphism of df, that is

$$df = \omega(\cdot, \mathbf{X}_f).$$

In local coordinates

$$\mathbf{X}_f = \frac{\partial f}{\partial \mathbf{p}} \cdot \frac{\partial}{\partial \mathbf{q}} - \frac{\partial f}{\partial \mathbf{q}} \cdot \frac{\partial}{\partial \mathbf{p}}.$$

The convention is that two vectors (denoted by boldfaced letters) are multiplied using the dot product. Indeed,

$$\omega(\cdot, \mathbf{X}_f) = (d\mathbf{p} \wedge d\mathbf{q}) \left(\cdot, \frac{\partial f}{\partial \mathbf{p}} \frac{\partial}{\partial \mathbf{q}} - \frac{\partial f}{\partial \mathbf{q}} \frac{\partial}{\partial \mathbf{p}} \right) = \frac{\partial f}{\partial \mathbf{p}} d\mathbf{p} - \left(-\frac{\partial f}{\partial \mathbf{q}} d\mathbf{q} \right) = df.$$

There is a fancy way of writing this using the notation $i_{\mathbf{X}}\omega = \omega(\mathbf{X}, \cdot)$. Then $\mathbf{X}_f$ is defined by

$$df = -i_{X_f}\omega.$$

Example 1.2.2. For the coordinate functions q_j, p_j, we have

$$\mathbf{X}_{q_j} = -\frac{\partial}{\partial p_j}, \quad \mathbf{X}_{p_j} = \frac{\partial}{\partial q_j}.$$

7. There is a *Poisson bracket* for smooth functions on the symplectic manifold M which is defined by

$$\{f, g\} = -\omega(\mathbf{X}_f, \mathbf{X}_g).$$

In local coordinates

$$\begin{aligned}\{f,g\} &= -\omega(\mathbf{X}_f,\mathbf{X}_g) = -\omega\left(\sum_j \frac{\partial f}{\partial p_j}\frac{\partial}{\partial q_j} - \frac{\partial f}{\partial q_j}\frac{\partial}{\partial p_j}, \sum_k \frac{\partial g}{\partial p_k}\frac{\partial}{\partial q_k} - \frac{\partial g}{\partial q_k}\frac{\partial}{\partial p_k}\right)\\ &= -\sum_{jk}\frac{\partial f}{\partial p_j}\frac{\partial g}{\partial p_k}\omega\left(\frac{\partial}{\partial q_j},\frac{\partial}{\partial q_k}\right) + \sum_{jk}\frac{\partial f}{\partial p_j}\frac{\partial g}{\partial q_k}\omega\left(\frac{\partial}{\partial q_j},\frac{\partial}{\partial p_k}\right)\\ &\quad + \sum_{jk}\frac{\partial f}{\partial q_j}\frac{\partial g}{\partial p_k}\omega\left(\frac{\partial}{\partial p_j},\frac{\partial}{\partial q_k}\right) - \sum_{jk}\frac{\partial f}{\partial q_j}\frac{\partial g}{\partial q_k}\omega\left(\frac{\partial}{\partial p_j},\frac{\partial}{\partial p_k}\right)\\ &= \left(\frac{\partial f}{\partial \mathbf{q}}\right)^T\frac{\partial g}{\partial \mathbf{p}} - \left(\frac{\partial f}{\partial \mathbf{p}}\right)^T\frac{\partial g}{\partial \mathbf{q}}.\end{aligned}$$

Theorem 1.2.3. The Poisson bracket satisfies

1. $\{af+bg,h\} = a\{f,h\} + b\{g,h\}$ for a,b constants and f,g functions;
2. $\{f,g\} = -\{g,f\}$;
3. $\{fg,h\} = \{f,h\}g + f\{g,h\}$.
4. **(The Jacobi identity)** $\{f,\{g,h\}\} + \{g,\{h,f\}\} + \{h,\{f,g\}\} = 0$.

The conditions can be checked by working in local coordinates. The four conditions define in general what is called a Poisson structure on $C^\infty(T^*M)$. The first, second, and fourth conditions define a *Lie algebra* structure.

Definition. A manifold whose algebra of functions is endowed with a Poisson bracket is called a Poisson manifold.

Example 1.2.3. Not all Poisson manifolds are symplectic. For example $\mathbb{R}^3$ (which is not symplectic because it has odd dimension), has the a Poisson bracket defined by

$$\{f,g\} = \langle x, \nabla f_x \times \nabla g_x\rangle.$$

Recall that a Lie algebra is a vector space endowed with a *Lie bracket* (an antisymmetrical bilinear 2-form satisfying the Jacobi identity).

Example 1.2.4. The *Heisenberg Lie algebra* $\mathfrak{H}(\mathbb{R}^n)$ is the algebra generated by the coordinate and momentum functions $q_1,q_2,\ldots,q_n,p_1,p_2,\ldots,p_n$, satisfying

$$\{q_j,q_k\} = \{p_j,p_k\} = 0, \quad \{p_j,q_k\} = -\delta_{jk}.$$

These equations are the canonical commutation relations for the classical positions and momenta. This can be modeled using matrices as

$$q_j = \begin{pmatrix}0 & e_j^T & 0\\ 0 & 0 & 0\\ 0 & 0 & 0\end{pmatrix}, \quad p_k = \begin{pmatrix}0 & 0 & 0\\ 0 & 0 & e_k\\ 0 & 0 & 0\end{pmatrix}, \quad 1 = \begin{pmatrix}0 & 0 & 1\\ 0 & 0 & 0\\ 0 & 0 & 0\end{pmatrix},$$

We can also think of the Poisson bracket infinitesimally, as a Lie bracket of Hamiltonian vector fields:

$$[\mathbf{X}_f, \mathbf{X}_g] = \mathbf{X}_{\{f,g\}}.$$

Using the Poisson bracket we can write Hamilton's equations as

$$\dot{q_j} = \{q_j, H\}, \quad \dot{p_j} = \{p_j, H\}.$$

In general, if $f(\mathbf{p}, \mathbf{q})$ is a function of position and momentum, which we interpret as a observable (i.e. measurable) quantity then its value, as the particle moves along the trajectory $(\mathbf{q}, \mathbf{p})$, satisfies the system of first order partial differential equations

$$\frac{df}{dt} = \{f, H\}.$$

This is the general form of *Hamilton's equations*. This formula can be proved as follows

$$\frac{df}{dt} = \frac{df}{d\mathbf{p}}\dot{\mathbf{p}} + \frac{df}{d\mathbf{q}}\dot{\mathbf{q}} = -\frac{df}{d\mathbf{p}}\frac{\partial H}{\partial q} + \frac{df}{d\mathbf{q}}\frac{\partial H}{\partial p} = \{f, H\}.$$

where we have used Hamilton's equations. Note that this equation is of the form $f' = Af$, where A is a linear operator.

8. Diffeomorphisms of the symplectic manifold that preserve the symplectic form are called *symplectomorphisms*. So $\phi : M \to M$ is a symplectomorphism if $\phi^*\omega = \omega$. (Here $\phi^*\omega$ is the pull back of ω defined as $\phi^*\omega_x(X, Y) = \omega_x(\phi_*(X), \phi_*(Y))$, where $\phi_*(X)$ is the push forward of the vector X defined by $\phi_*(\gamma'(0)) = (\phi \circ \gamma)'(0)$.) Symplectomorphisms also preserve the Poisson bracket, so they preserve the Hamiltonian formalism. Symplectomorphisms are therefore the "isomorphisms" of classical mechanics.

Example 1.2.5. For a free 1-dimensional particle, the phase space is $T^*\mathbb{R} = \mathbb{R}^2$ with coordinates q the position and p the momentum. The symplectic form is $dp \wedge dq$. This is the same as the area element in the plane, taken with the sign that specifies orientation.

The linear transformations that preserve area and orientation form the special group

$$SL(2, \mathbb{R}) = \left\{ \begin{pmatrix} a & b \\ c & d \end{pmatrix} \mid ad - bc = 1 \right\},$$

which are therefore the linear syplectomorphisms.

Indeed, if ϕ is the symplectomorphism defined by the matrix $\left(\begin{smallmatrix} a & b \\ c & d \end{smallmatrix}\right)$, then the pullback of ω through ϕ is the symplectic form $\phi^*\omega$ defined by the condition

$$\begin{aligned} \phi^*\omega\left(\frac{\partial}{\partial p}, \frac{\partial}{\partial q}\right) &= \omega\left(a\frac{\partial}{\partial p} + b\frac{\partial}{\partial q}, c\frac{\partial}{\partial p} + d\frac{\partial}{\partial q}\right) = ad\omega\left(\frac{\partial}{\partial p}, \frac{\partial}{\partial q}\right) + bc\omega\left(\frac{\partial}{\partial q}, \frac{\partial}{\partial p}\right) \\ &= ad - bc = 1 = \omega\left(\frac{\partial}{\partial p}, \frac{\partial}{\partial q}\right). \end{aligned}$$

Example 1.2.6. For n free 1-dimensional particles, the phase space is $T^*\mathbb{R}^n = \mathbb{R}^{2n}$, with coordinates $\mathbf{q} = (q_1, q_2, \ldots, q_n)$ (positions) and $\mathbf{p} = (p_1, p_2, \ldots, p_n)$ (momenta). The symplectic form is $d\mathbf{p} \wedge d\mathbf{q} = \sum_j dp_j \wedge dq_j$.

The linear transformations that preserve the symplectic form are the elements of the symplectic group

$$Sp(2n, \mathbb{R}) = \left\{ h \in GL(2n, \mathbb{R}) \mid h^T \begin{pmatrix} 0 & I_n \\ -I_n & 0 \end{pmatrix} h = \begin{pmatrix} 0 & I_n \\ -I_n & 0 \end{pmatrix} \right\}.$$

Note that $Sp(2, \mathbb{R}) = SL(2, \mathbb{R})$, but for other n the symplectic group and the special linear group are different.

A particular example of symplectomorphisms, which play an important role in mechanics, are *Hamiltonian flows*. Let $H(\mathbf{p}, \mathbf{q})$ be a *time-independent Hamiltonian function* and let $\mathbf{X}_H$ be its associated vector field.

The Hamiltonian flow defined by H is a family of diffeomorphisms $\phi_t : M \to M$ indexed by some interval I containing 0 such that for every $x \in M$, $\phi_t(x) : I \to M$ is defined by the differential equation

$$\frac{d}{dt}\phi_t(x) = (\mathbf{X}_H)_{\phi_t(x)},$$

where the right term denotes the vector (of the vector field) at the point $\phi_t(x)$. Note that in local coordinates, $\phi_t(\mathbf{q}_0, \mathbf{p}_0) = (\mathbf{q}(t), \mathbf{p}(t))$, where $(\mathbf{q}(t), \mathbf{p}(t))$ are obtained by solving the initial value problem with Hamilton's equations and initial condition $(\mathbf{q}_0, \mathbf{p}_0)$. So the above equation is another way of phrasing Hamilton's equations. A Hamiltonian flow is shown schematically in Figure 1.3.

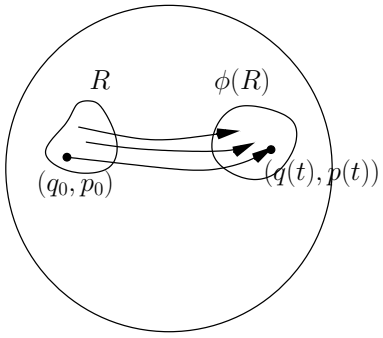

Figure 1.3: Hamiltonian flow.

Note that the uniqueness of the solution to the differential equation implies $\phi_t(\phi_s(x)) = \phi_{t+s}(x)$. Thus the Hamiltonian flow is a one parameter group of diffeomorphisms.

Theorem 1.2.4. The Hamiltonian flow defines a one parameter group of symplectomorphisms that conserve H.

Proof. As a consequence of the definition, each observable varies under the flow by the Hamilton equation $df/dt = \{f, H\}$. So the conservation of the energy follows from

$$\begin{aligned}\frac{dH}{dt} &= \left(\frac{\partial H}{\partial \mathbf{q}}\right)^T \frac{d\mathbf{q}}{dt} + \left(\frac{\partial H}{\partial \mathbf{p}}\right)^T \frac{d\mathbf{p}}{dt} = \left(\frac{\partial H}{\partial \mathbf{q}}\right)^T \dot{\mathbf{q}} + \left(\frac{\partial H}{\partial \mathbf{p}}\right)^T \dot{\mathbf{p}} \\ &= \left(\frac{\partial H}{\partial \mathbf{q}}\right)^T \frac{\partial H}{\partial \mathbf{p}} - \left(\frac{\partial H}{\partial \mathbf{p}}\right)^T \frac{\partial H}{\partial \mathbf{q}} = 0.\end{aligned}$$

This is the same as

$$\frac{dH}{dt} = \{H, H\} = 0.$$

To show that Hamiltonian flows are symplectomorphisms, let us prove that

$$\frac{d}{dt}\phi^*\omega = 0.$$

We have

$$\frac{d}{dt}\phi^*\omega = \phi_t^* \lim_{h\to 0} h^{-1}[\phi_h^*\omega - \omega] = \phi_t^* \mathcal{L}_{\mathbf{X}_H}\omega.$$

So let us check that

$$\mathcal{L}_{\mathbf{X}_H}\omega = \lim_{h\to 0} h^{-1}[\phi_h^*\omega - \omega] = 0.$$

We use the Cartan formula

$$\mathcal{L}_{\mathbf{X}} = i_{\mathbf{X}} \circ d + d \circ i_{\mathbf{X}},$$

to get

$$\mathcal{L}_{\mathbf{X}_H}\omega = i_{\mathbf{X}_H} \circ d\omega + d\omega(\mathbf{X}_H, \cdot) = 0 - d^2 H = 0.$$

Cartan's formula can be proved as follows. Both sides satisfy the "product rule" so they are derivations. So it suffices to check on functions, where it is trivial, and on 1-forms:

$$[i_X d + d i_{\mathbf{X}}]df = d i_{\mathbf{X}}(df) = d[i_{\mathbf{X}}(df)] = d(\mathbf{X}(f)) = d\mathcal{L}_{\mathbf{X}}(f) = \mathcal{L}_{\mathbf{X}} df.$$

Here is a different way to check that Hamiltonian flows are symplectomorphisms. Consider a surface Σ. It suffices to show that

$$\int_\Sigma \omega = \int_{\phi_t(\Sigma)} \omega$$

Using Stokes' Theorem we can write

$$\int_{\phi_t(\Sigma)} \omega - \int_\Sigma \omega = \int_{F\Sigma} d\omega = 0,$$

where $F\Sigma$ is the "flow" of Σ between time 0 and time t. So the "area element" is preserved. □

The equation $dH/dt = 0$ shows that total energy is preserved. The equation

$$\mathcal{L}_{\mathbf{X}_H}\omega = 0,$$

where the Lie derivative was defined above as $\mathcal{L}_{\mathbf{X}_H} = \lim_{h\to 0} h^{-1}[\phi_h^*\omega - \omega]$ implies Liouville's Theorem:

$$\mathcal{L}_{\mathbf{X}_H}(\omega^n) = 0,$$

which says that the volume form is preserved by the Hamiltonian flow.

Example 1.2.7. The torus $\mathbb{T}^2 = S^1 \times S^1 = \mathbb{R}^2/\mathbb{Z}^2$ is a symplectic manifold with symplectic form $dp \wedge dq$, where (p, q) mod 1 are the coordinates of the point on the torus.

Let us consider the Hamiltonian function $H(p, q) = \sin 2\pi p$. Then

$$X_H = 2\pi \cos 2\pi p \frac{\partial}{\partial q}.$$

The Hamiltonian flow is

$$\phi_t(p, q) = (p, 2\pi(tq \bmod 1) \cos 2\pi p).$$

On the other hand the family of translations

$$\phi_t(p, q) = (p + t, q)$$

which preserve the symplectic structure, are not a Hamiltonian flow. Indeed, we should have

$$\frac{\partial H}{\partial p} = 0 \text{ and } \frac{\partial H}{\partial q} = -1,$$

which would imply $H = -q + C$, C a constant. But this function is not well defined on the torus!

The map on the torus induced by the $(p, q) \mapsto (-q, p)$ is also a symplectomorphism, but it is not part of a one-parameter group, so it does not arise from a Hamiltonian flow.

Example 1.2.8. The rotation of the plane about the origin is not a Hamiltonian flow. Indeed, the map

$$(p, q) \mapsto ((\cos t)p - (\sin t)q, (\sin t)p + (\cos t)q)$$

yields, via Hamilton's equations,

$$-\frac{\partial H}{\partial q} = \dot{p} = -(\sin t)p - (\cos t)q$$
$$-\frac{\partial H}{\partial p} = \dot{q} = (\cos t)p - (\sin t)q.$$

These equations can be integrated, but the result is $H = -(\cos t)\frac{p^2+q^2}{2} - (\sin t)pq + C$, and this is time dependent (total energy varies with time). This might look strange, as rotation about the origin is a physical phenomenon, but remember, the physical rotation about the origin happens in the plane as the configuration space, here the plane is the phase space.

Chapter 2

"Matrix Mechanics" - A First Encounter with Quantum Mechanics

"I think I can safely say that nobody understands quantum mechanics." (Richard Feynman)

This chapter and the rest of the course are heavily inspired by Brian Hall's book *Quantum Theory for Mathematicians*, L.D. Faddeev and O.A. Yakubovskii *Lectures on Quantum Mechanics for Mathematics Students*, as well as on the lecture notes I took in the class of Alejandro Uribe at University of Michigan.

2.1 Problems that have led to the discovery of quantum mechanics

1. One should start with a quote by Lord Kelvin: "There is nothing new to be discovered in physics now, all that remains is more and more precise measurement". This was the universally accepted truth at the end of the 19th century. But then, as more and more precise measurements were made, problems started to appear.

There was the Michaelson-Morley experiment detecting the speed of light. It showed that the speed of light was always the same regardless of the orientation and the motion of the reference frame. To solve this inconsistency of physics, Einstein formulated a new set of postulates that fundamentally changed our understanding of physics at macroscopic level.

But then there were several experiments that raised questions about the microscopic world.

- The *black body radiation* is a first example where classical theory of physics produces nonsense. A black body is an object that emits radiation only when heated. An example is a stove element that turns red when heated. Experiments and classical physics predict that the black body emits electromagnetic radiation in all wavelengths, and the power of the radiation emitted by the black body is proportional to λ^{-4}, where λ is the wavelength. Thus as the frequency of the radiation increases, so does the power of the radiation, tending towards infinity (this is known as the utraviolet catastrophe).

  Experiments verify this law at low frequencies, but as the frequencies grow, the law fails badly. The problem was solved by Max Planck. In 1894, he has been comissioned

by the industry to increase the efficiency of light bulbs. This led him to the study of black body radiation. To predict correctly the radiation emitted by a black body, he had to postulate, in year 1900, what has since been known as Planck's law:

$$E = hn\nu,$$

namely that in a certain frequency, the energy emitted can only be an integer multiple of a certain quantity. Moreover, that quantity is proportional to the frequency. He called this assumption "an act of despair", but... it works.

- The *photoelectric effect* is a phenomenon in which a material emits electrons when light is shone upon it. A curious discovery was made at the end of 19th century, namely that the intensity of the electric current generated by the photoelectric effect does not depend on the intensity of light, but on the color of the light. Albert Einstein explained this in 1905 by making use of Planck's law: light can only be emitted in quanta, called photons, and only one photon with enough energy could make the material eject one electron, thus producing the observed effect.

- The distribution of *spectral lines of the hydrogen atom* is governed by a formula that was partially found by Johann Balmer in 1885, and then in its full extent by Johannes Rydberg in 1888. This formula gives the reciprocals of the wavelengths of the light emitted by the hydrogen atom as

$$\lambda^{-1} = R_H \left(\frac{1}{m^2} - \frac{1}{n^2} \right),$$

where R_H is a constant, and $m < n$ vary among the positive integers. This distribution of spectral lines was explained in 1913 by the model of the atom given by Niels Bohr.

2. Starting with the Rydberg formula, let us see how "matrix mechanics" appears. I will present you the explanation given by Werner Heisenberg in his book *The Physical Principles of the Quantum Theory*, as it is always advised to learn the great ideas directly from the great masters who have discovered them.

First, let us place ourselves in the classical setting from the end of 19th century. The view was that the hydrogen atom has a nucleus surrounded by electrons, and the motion of the electrical charges which are the electrons produce, according to Maxwell's equations, electromagentic waves. These electromagnetic waves is what we see in the spectrum. Now, the classical theory predicts that the motion of an electron should produce electromagnetic waves in one fundamental frequency together with harmonics in frequencies that are integer multiples of this fundamental frequency. Basically you have some periodic phenomenon which creates periodic waves, and these waves can be expanded into a Fourier series. The first sinusoidal wave of the expansion gives the fundamental frequency, and the others are the harmonics.

But the Balmer series contradicts this. We do not see a Fourier series expansion. Let us examine this closely.

In classical mechanics we observe functions on the phase space of the system. But what do we observe in the case of quantum theory? Well, for the hydrogen atom we observe its spectral lines, or more precisely, the frequencies given by the formula

$$\nu(m,n) = \frac{R_H}{m^2} - \frac{R_H}{n^2} = T_m - T_n.$$

These satisfy the Rydberg-Ritz combination principle:

$$\nu(m,k) + \nu(k,n) = \nu(m,n).$$

Planck's law tells us that energy is emitted in quanta, and that the energy of one quantum of light is proportional to the frequency: $E = h\nu$.

Thus what we observe are energies of quanta of light, and we have a sequence of energy levels $W_1, W_2, \ldots$ such that

$$\nu(m,n) = \frac{1}{h}(W_n - W_m).$$

Now we observe several waves with certain frequencies and certain amplitudes. These waves can be represented by exponential functions, and the amplitudes by real numbers. Using the above intuition we arrange these waves corresponding to the spectral lines as a matrix:

$$\left(q(m,n)e^{2\pi i\nu(m,n)t}\right), \quad m,n.$$

This departs from the classical situation where a wave is decomposed into elementary oscillations by the Fourier transform, and the elementary oscillations are summed with the inverse Fourier transform. So here we do not have a sum of elementary oscillations, but a matrix of elementary oscillations.

So, instead of the sum of elementary oscillations, here we have a matrix of elementary oscillations. Now, $\nu(m,m) = 0$, and moreover, we can allow $\nu(m,n) = -\nu(n,m)$ because of the differences, so that the entries of the matrix are defined for all pairs of integers, and the matrix is Hermitian.

In classical physics you can perform algebraic computations with Fourier series, and in this matrix model you can also perform algebraic computations with matrices, as the Rydberg-Ritz combination principle enables matrix multiplication. Note that amplitudes themselves are multiplied like matrices. But matrices do not commute!

The axioms of quantum mechanics are phrased based on these speculations. They define an abstract mathematical formalism, which tells us how to predict the results of particular experiments, and it turns out that these predictions match the results of the experiments. We should point out that the mathematical formalism of quantum mechanics is more amorphous than that of classical mechanics; once the general axioms are established there are many fixes and guesses that have to be applied to particular situations. Moreover, one should realize that the mathematics of quantum mechanics is not a rephrasing in rigorous terms of common language and intuition. While it comes from our intuition of the world, it introduces mathematical tools that we cannot identify with our senses.

2.2 The axioms of quantum mechanics

3. The states of a quantum system are in one-to-one correspondence with complex lines in a separable complex Hilbert space.

A Hilbert space is first a $\mathbb{C}$-vector space H endowed with a map $\langle , \rangle : H \times H \to \mathbb{C}$, called inner product, that is linear in the first variable, satisfies $\langle \phi, \phi \rangle \geq 0$ with equality if and only if $\psi = 0$, and $\langle \phi, \psi \rangle = \overline{\langle \psi, \phi \rangle}$. The inner product defines a norm $\|\phi\| = \sqrt{\langle \phi, \phi \rangle}$ and the norm induces the distance between ϕ and ψ as $\|\phi - \psi\|$. With this distance we can define a metric space structure on H, and for H to be a Hilbert space we require it to be complete as a metric space, namely that every Cauchy sequence has a convergent subsequence. A good example is the Hilbert space $L^2(\mathbb{R})$ with the inner product $\langle f, g \rangle = \int f(t)\overline{g(t)}dt$. Using the norm we can define the concept of continuity, either for functions from H to H or from H to $\mathbb{C}$. Linear functions from H to H will be called operators, linear functions from H to $\mathbb{C}$ will be called linear functionals.

The quantum phase space is

$$\mathbb{P}H = \{[\psi] \,|\, \psi \in H \backslash \{0\}\} \text{ where } [\psi] = \{\lambda\psi \,|\, \lambda \in \mathbb{C}^*\}.$$

4. *The observables are self-adjoint operators on the Hilbert space.*

To define the concept of a self-adjoint operator, the key ingredient is the Riesz lemma. This lemma states that every continuous linear functional $\phi : H \to \mathbb{C}$ is of the form $\phi \mapsto \langle \phi, \psi \rangle$.

Now we look at densely defined linear operators, namely at linear operators A that are defined on some dense subspace $\mathcal{D}(A) \subset H$ (here $\mathcal{D}(A)$ is called the domain of A). Whenever the linear functional $\phi \mapsto \langle A\phi, \psi \rangle$ is continuous on the dense subset $\mathcal{D}(A)$ it can be extended, by continuity to the whole H, and then by the Riesz representation theorem it is of the form $\phi \mapsto \langle \phi, \xi \rangle$. Then we can define $A^*\psi = \xi$. This is the adjoint of A and its domain is the set of all ψ for which this construction works.

Definition. $A : H \to H$, is self-adjoint if

- A is densely defined ($\overline{\mathcal{D}(A)} = H$),
- A is symmetric, that is $\langle A\psi, \phi \rangle = \langle \psi, A\phi \rangle$ for all $\phi, \psi \in \mathcal{D}(A)$,
- $\mathcal{D}(A^*) = \mathcal{D}(A)$.

The last two conditions can be summarized as $A = A^*$.

Example 2.2.1. Every Hermitian matrix is self-adjoint.

If A is a continuous linear operator that is densely defined, then it is defined everywhere. Continuity simply means that if $x_n \to x$ with $x_n, x \in \mathcal{D}(A), n \geq 1$, then $Ax_n \to Ax$. If $x_n \to x$ with $x_n \in \mathcal{D}(A)$ but $x \notin \mathcal{D}(A)$, then x_n is a Cauchy sequence, so $x_n - x_m \to 0$ when $n, m \to \infty$. Then Ax_n is a Cauchy sequence, and we can let Ax be its limit. This is how you extend A to the entire space. So for continuous operators, $\mathcal{D}(A)$ could, and should be assumed to be the entire Hilbert space. In that case a self-adjoint operator is defined by the

equality $\langle Ax, y\rangle = \langle x, Ay\rangle$. However, as we will see later, continuous operators do not suffice for modeling quantum mechanics.

Let us recall that a continuous linear operator $A : \mathcal{H} \to \mathcal{H}$ is also called a bounded operator, and for a good reason: The $\epsilon - \delta$ definition of continuity on metric spaces implies that a linear operator $A : \mathcal{H} \to \mathcal{H}$ is continuous if and only if there is $C > 0$ such that $\|Ax\| \leq C\|x\|$ for all x. The infimum of all C with this property is called the norm of A and is denoted by $\|A\|$. The norm satisfies $\|I\| = 1$, $\|A+B\| \leq \|A\|+\|B\|$ and $\|AB\| \leq \|A\|\|B\|$. Bounded operators are nice, for one thing because they are defined everywhere not just on a dense subset, but also because there is a rich theory of bounded linear operators. They form a C^* algebra, which by definition, is an algebra over $\mathbb{C}$ endowed with a conjugation (the operation of taking the adjoint) and a norm satisfying the above properties as well as the equality $\|A^*A\| = \|A\|^2$.

5. Recall that the spectrum of an operator is the complement of the set of complex numbers λ for which the operator $\lambda I - A$ has a continuous inverse (where I is the identity operator). If A is continuous, it is known that the spectrum is nonempty (this is a consequence of Liouville's Theorem in complex analysis applied to the holomorphic function $\lambda \mapsto \|(\lambda I - A)^{-1}\|$), moreover the spectrum is compact (i.e. closed and bounded in $\mathbb{C}$). But there are unbounded operators that have empty spectrum. But self-adjoints always have nonempty spectrum. Moreover, unlike the case of matrices, it might be that the space of a continuous operator is not discrete, or that the spectrum of an operator that is not continuous is unbounded.

We focus now on observables with discrete spectrum. *If A is an observable with discrete spectrum*, then A is diagonalizable. So

$$A\psi_j = \lambda_j \psi_j,$$

where ψ_j, $j \in J$, is an orthonormal basis.

In this situation,

- *An observation of A always results in some λ_j for some j.*
- *If $\psi = \sum_j c_j\psi_j$ and the system is in state $[\psi]$ and you observe A, the probability of obtaining the result λ_j is $|c_j^2|/||\psi||^2$.*

Note that $\|\psi\|^2 = \sum |c_j|^2$.

Consequently, if you measure A repeatedly with the system in state ψ, the average of the measurements is $\langle A\psi, \psi\rangle / \|\psi\|^2$ (the expected value of the observable):

$$Exp(A)_{[\psi]} = \frac{1}{\|\psi\|^2} \sum_j \lambda_j |c_j|^2.$$

6. *If A and B are quantum observables, then AB is the quantum observable: "observe B then observe A then multiply the resulting numbers".*

Then

$$\frac{1}{i\hbar}(AB - BA) = \frac{1}{i\hbar}[A, B]$$

is the self-adjoint measure of the lack of commutativity of observations.

7. Here is a new look at states. We determine that the system is in a certain state by performing experiments. In those experiments we measure certain observables. In classical mechanics, a given state yields always the same measurement for the observable. In quantum mechanics this is not the case. So each state ω of the system associates to a given quantum observable a probability distribution (and we have explained above what it is for observables that have discrete spectrum). So for an observable A the state ω defines a map $E \mapsto \omega_A(E)$ from the Borel sets in $\mathbb{R}$ to $[0,1]$ such that

$$\omega_A(\sqcup E_j) = \sum_j \omega_A(E_j), \quad \omega_A(\emptyset) = 0, \omega_A(\mathbb{R}) = 1.$$

We should also have a relationship of the form: if observable B is equal to $f(A)$ for some function f, then $\omega_B(E) = \omega_A(f^{-1}(E))$.

Note that a convex combination of probability measures is a probability measure, so if $\omega = t\omega_{1,A} + (1-t)\omega_{2,A}$ then ω should also correspond to a state, and we let this state be $t\omega_1 + (1-t)\omega_2$. A state that can be decomposed this way is called a *mixed state*, a state that cannot be decomposed this way is called a *pure state*, thus pure states are the extremal elements of the set of states.

The expected value of an observable A is

$$\omega(A) = \int_{-\infty}^{\infty} \lambda d\omega_A(\lambda).$$

Physical intuition suggests we should have the following properties

$$\begin{aligned}
&\omega(cI) = c \\
&\omega(\alpha A + \beta B) = \alpha\omega(A) + \beta\omega(B) \\
&\omega(A^2) = \omega(A^*A) \geq 0 \\
&\omega(A) \in \mathbb{R}.
\end{aligned}$$

Now we would like to have enough observables to be able to separate states, so let us assume that all self-adjoint operators are quantum observables. Extend by linearity ω to all operators, and in particular to the C^* algebra of bounded linear operators. Then we have a *positive linear functional* on the C^* algebra of bounded linear operators on H whose norm is equal to 1. Now let us imagine that instead we only have the observables and abstract states defined as positive linear functionals of observables (whose values are the results of experiments). The extremal points of the set of functionals are called pure states, the other states are the mixed states. Via the Gelfand-Naimark-Segal construction we can produce a Hilbert space such that the *pure states* are the linear functionals of the form

$$A \mapsto \langle A\psi, \psi \rangle,$$

for some $\psi \in H$ of norm 1.

Well, we are only concerned with pure states, we call these simply states, and hence the first axiom of quantum mechanics. Note that $\psi \mapsto e^{i\theta}\psi$ does not change the state, this is why we work with $\mathbb{P}H$.

In general, for a not necessarily diagonalizable operator A, the expected value of A in the state ψ ($\|\psi\| = 1$) is $\langle A\psi, \psi\rangle$.

8. A consequence of de Broglie's formula $p = \frac{h}{\lambda}$ is the Heisenberg uncertainty principle. Let us formulate the *Heisenberg uncertainty principle* in dimension 1. If q is the position and p is the momentum then the standard deviations σ_q and σ_p satisfy

$$\sigma_q \sigma_p \geq \frac{\hbar}{2}.$$

Recall from statistics that the standard deviation of a (finite) variable x is

$$s_x = \left(\frac{1}{n}\sum_{j=1}^{n}(x_j - \mu)^2\right)^{1/2},$$

where μ is the mean. In our case, for an operator A, it is

$$\sigma_A = \sqrt{\langle (A - A_{avg})^2\psi, \psi\rangle},$$

where A_{avg} is the expected value of A.

Theorem 2.2.1. Two operators that satisfy the *canonical commutation relation*

$$PQ - QP = -i\hbar I$$

satisfy the Heisenberg uncertainty principle

$$\sigma_Q \sigma_P \geq \frac{\hbar}{2}.$$

Proof. The standard deviation of the observable A in state ψ (here ψ is a unit vector) is

$$\sigma_A^2 = \left\langle (A - A_{avg})^2\psi, \psi\right\rangle = \left\langle A^2\psi, \psi\right\rangle - \left\langle A\psi, \psi\right\rangle^2,$$

where $A_{avg} = \langle A\psi, \psi\rangle$. Now we mimic the proof of the Cauchy-Schwarz inequality. We first assume $P_{avg} = Q_{avg} = 0$, by translating P and Q by some multiples of the identity operator. Then we start with

$$\langle (Q + itP)\psi, (Q + itP)\psi\rangle \geq 0, \quad \|\psi\| = 1, \alpha \in \mathbb{R},$$

and rewrite it as

$$\left\langle Q^2\psi, \psi\right\rangle + t^2\left\langle P^2\psi, \psi\right\rangle + it\left\langle (QP - PQ)\psi, \psi\right\rangle \geq 0.$$

This is the same as

$$t^2\left\langle P^2\psi, \psi\right\rangle - t\hbar\left\langle \psi, \psi\right\rangle + \left\langle Q^2\psi, \psi\right\rangle \geq 0.$$

This is a quadratic in t which is nonnegative, so its discriminant is nonpositive. We get

$$\left\langle P^2\psi, \psi\right\rangle\left\langle Q^2\psi, \psi\right\rangle \geq \frac{\hbar^2}{4}.$$

For the last step we used $\langle \psi, \psi\rangle = 1$. The conclusion follows by taking the square root. □

Remark 2.2.1. We have more generally, for two observables A and B measured in state ψ the inequality

$$\sigma_A \sigma_B \geq \frac{\hbar}{2} |\langle \{A, B\}_\hbar \psi, \psi \rangle|$$

where

$$\{A, B\}_\hbar = \frac{1}{i\hbar}[A, B],$$

is the *quantum Poisson bracket.*

9. We need to explain the name quantum Poisson bracket. If there is an association of operators P, Q to the classical observables (coordinate functions on the phase space) p and q, then the linear map from the Heisenberg Lie algebra to the space of linear operators on H, $L(H)$, defined by

$$p \mapsto P, \quad q \mapsto Q, \quad 1 \mapsto I$$

is a Lie algebra isomorphism onto the image, if we endow $L(H)$ with the Lie bracket $\{\cdot, \cdot\}_\hbar$.

2.3 In search of a nice quantization scheme

10. Quantum mechanics should be obtained based on the intuition provided by classical mechanics, while classical mechanics should be the large scale behavior of quantum mechanics. Thus there should exists an algorithm for passing from classical mechanics to quantum mechanics (aiding our understanding of quantum mechanics) and an algorithm for passing from quantum mechanics to classical mechanics (which should allow us to recover the mechanics we already know).

The quantization "algorithm" should associate

- to the classical phase space a Hilbert space;
- to classical observables self-adjoint operators on the Hilbert space.

This association should satisfy Dirac's quantization conditions, as outlined by Paul Dirac in 1926 in *Proc. Royal Soc. London A*:

(1) $\mathrm{op}(1) = I$, where I is the identity operator;

(2) if $f = c_1 f_1 + c_2 f_2$ then $\mathrm{op}(f) = c_1 \mathrm{op}(f_1) + c_2 \mathrm{op}(f_2)$;

(3) (the correspondence principle) $\mathrm{op}(\{f, g\}) = \frac{1}{i\hbar}[\mathrm{op}(f), \mathrm{op}(g)]$ (!);

(4) the representation of quantum observables on the Hilbert space is irreducible.

There is one model for which this "algorithm" is specified as an axiom, the Heisenberg quantization conditions:

- $C^\infty(\mathbb{R}^n \times \mathbb{R}^n)$ becomes $L^2(\mathbb{R}^n, d\mathbf{q})$;
- the coordinate function q_j becomes the operator $Q_j = M_{q_j}$ of multiplication by q_j;
- the momentum function p_j becomes the operator $P_j = -i\hbar\frac{\partial}{\partial q_j} = \frac{\hbar}{i}\frac{\partial}{\partial q_j}$.

The choice of these operators is quite intuitive, the position operator is the multiplication by the position function, and the momentum is the instantaneous translation (the exponential of the differentiation operator is translation). They also satisfy the cannonical commutation relations, as a consequence of the product rule:

$$P_jQ_j\psi = -i\hbar\frac{\partial}{\partial q_j}(q_j\psi) = -i\hbar\psi - i\hbar q_j\frac{\partial\psi}{\partial q_j} = -i\hbar\psi + Q_jP_j\psi.$$

Here $\psi \in L^2(\mathbb{R}^n)$ is a function of $\mathbf{q} = (q_1, q_2, ..., q_n)$. Unfortunately both operators are unbounded. This means that they are discontinuous maps on $L^2(\mathbb{R}^n)$ and in particular they are not defined everywhere. The domains of these operators are chosen so that they are self-adjoint (which is possible – I will post the description of the domain soon). One should point out that the "eigenvalues" of M_{q_j} are Dirac's delta functions and they are not in the Hilbert space. Can we do any better, by choosing nicer operators that satisfy the canonical commutation relations?

Clearly finite dimensional operators do not work, because if $[P_j, Q_j] = i\hbar I$ then

$$0 = \text{trace}([P_j, Q_j]) = \text{trace}(-i\hbar I) = -i\hbar\,\text{dim}H.$$

But the situation is even worse.

Theorem 2.3.1. (Wintner-Wielandt) The relation $PQ - QP = -i\hbar I$ cannot be satisfied by bounded linear operators P and Q.

Proof. Assume that this is possible and set $A = Q$, $B = (i/\hbar)P$. Then $AB - BA = I$. We prove by induction that

$$A^nB - BA^n = nA^{n-1}.$$

This is clearly true for $n = 1$. If we assume that this is true for n, then for $n + 1$ we have

$$A^{n+1}B - BA^{n+1} = A^n(AB - BA) + (A^nB - BA^n)A = A^n + nA^{n-1}A = (n + 1)A^n.$$

Then

$$n\|A^{n-1}\| = \|A^nB - BA^n\| \leq 2\|A^n\|\|B\| \leq 2\|A^{n-1}\|\|A\|\|B\|.$$

So $n \leq 2\|A\|\|B\|$ for all positive integers n. This is clearly impossible. □

So the definition of the position and momentum as multiplication and differentiation operators seems to be the best choice we have. In fact there is a "theorem" due to M. Stone saying that this is the "only" choice. We will make this theorem precise in the guise of the Stone-von Neumann theorem.

11. We therefore have the Lie algebra isomorphism onto the image

$$\mathfrak{H}(\mathbb{R}^n) \to L(L^2(\mathbb{R}^n)), \quad f \mapsto \mathrm{op}(f).$$

Can this be extended to a Lie algebra isomorphism onto the image from the Lie algebra $C^\infty(\mathbb{R}^n \times \mathbb{R}^n)$ or in general from the Lie algebra of smooth functions on the phase space of a classical system to the algebra of linear operators on the Hilbert space of the associated quantum system?

The question is how to extend the definition to functions of p_j, q_j.

Example 2.3.1. The operator associated to the kinetic energy of one particle is

$$\mathrm{op}\left(\sum_{j=1}^{3} \frac{1}{2m} p_j^2\right) = \sum_{j=1}^{3} \frac{1}{2m} \mathrm{op}(p_j^2) = \text{guess} = \sum_{j=1}^{3} \frac{1}{2m} \mathrm{op}(p_j)^2 = -\frac{\hbar^2}{2m}\Delta,$$

where Δ is the Laplacian, $\Delta = \sum_{j=1}^{3} \frac{\partial^2}{\partial q_j^2}$. This is the *quantum kinetic energy*.

It is a good guess to define the quantum Hamiltonian of a Newtonian system as

$$\mathrm{op}(H) = \mathrm{op}(T) + \mathrm{op}(V) = -\frac{\hbar^2}{2m}\Delta + \mathrm{op}(V) = -\frac{\hbar^2}{2m}\Delta + M_{V(\mathbf{q})}.$$

But we run into a big problem when trying to define $\mathrm{op}(p_j q_k)$ (this is necessary for example when you try to understand the operator associated to angular momentum). The problem is that the naive definition $\mathrm{op}(q_j p_k) = \mathrm{op}(q_j)\mathrm{op}(p_k)$ fails because the latter product is not a self-adjoint operator. And because then what should we do with $\mathrm{op}(p_k q_j)$? It is natural to define

$$\mathrm{op}(q_j p_k) = \frac{1}{2}\left(\mathrm{op}(q_j)\mathrm{op}(p_k) + \mathrm{op}(p_k)\mathrm{op}(q_j)\right).$$

As the following result shows, this works for quadratic polynomials.

Proposition 2.3.1. The quantization scheme from the Lie algebra of quadratic polynomials on $\mathbb{R}^{2n}$ in the variables $q_1, q_2, \ldots, q_n, p_1, p_2, \ldots, p_n$, defined by

$$\mathrm{op}(1) = I, \quad \mathrm{op}(q_j) = M_{q_j}, \quad \mathrm{op}(p_j) = -i\hbar\tfrac{\partial}{\partial q_j},$$
$$\mathrm{op}(q_j q_k) = M_{q_j q_k}, \quad \mathrm{op}(p_j p_k) = -\hbar^2 \tfrac{\partial^2}{\partial q_j \partial q_k}, \quad \mathrm{op}(q_j p_k) = -i\hbar\left(M_{q_j}\tfrac{\partial}{\partial q_k} + \tfrac{1}{2}\delta_{jk}\right)$$

acting on $L^2(\mathbb{R}^n)$, which are then extended linearly to the Lie algebra of quadratic polynomials, satisfies Dirac's quantization conditions (1), (2), (3).

Proof. Conditions (1) and (2) are straightforward by the linearity requirement. Condition (3) is straightforward for Poisson brackets of functions depending only on q_j's or only on p_j's. For the others the computations are tedious but involve just differentiation and multiplication. Here is an example:

$$\{q_j p_k, p_\ell\} = \delta_{j\ell} p_k.$$

and

$$\begin{aligned}\frac{1}{i\hbar}[\mathrm{op}(q_jp_k),\mathrm{op}(p_\ell)]\psi &= \frac{1}{i\hbar}(-i\hbar)^2\left[\left(M_{q_j}\frac{\partial}{\partial q_k}+\frac{\delta_{j\ell}}{2}\right)\frac{\partial\psi}{\partial q_\ell}+\frac{\partial}{\partial q_\ell}\left(M_{q_j}\frac{\partial\psi}{\partial q_k}+\frac{\delta_{j\ell}\psi}{2}\right)\right]\\ &= i\hbar\left[M_{q_j}\frac{\partial^2\psi}{\partial q_k\partial q_\ell}+\frac{\delta_{j\ell}}{2}\frac{\partial\psi}{\partial q_\ell}-M_{q_j}\frac{\partial^2\psi}{\partial q_j\partial q_\ell}-\delta_{j\ell}\frac{\partial\psi}{\partial q_k}-\frac{\delta_{j\ell}}{2}\frac{\partial\psi}{\partial q_\ell}\right]\\ &= -i\hbar\delta_{j\ell}\frac{\partial}{\partial q_k}=\delta_{j\ell}\mathrm{op}(p_k).\quad\square\end{aligned}$$

12. But then how should we quantize $q_j^mp_k^n$? Should we average over all possible ways to write the product? And is Dirac's quantization condition (3) still satisfied? The next result is bad news!

Theorem 2.3.2. (The Weak No-Go Groenewold Theorem) The above quantization scheme cannot be extended to polynomials of degree less than or equal to 4.

Proof. (following Mark J. Gotay) For all quadratic polynomials $R_1(x)$ and all cubic polynomials $R_2(x)$ we must have

$$\begin{aligned}&\mathrm{op}(R_2(q_j))=R_2(\mathrm{op}(q_j)),\quad \mathrm{op}(R_2(p_j))=R_2(\mathrm{op}(p_j))\\ &\mathrm{op}(R_1(q_j)p_k)=\frac{1}{2}[R_1(\mathrm{op}(q_j))\mathrm{op}(p_k)+\mathrm{op}(p_k)R_1(\mathrm{op}(q_j))]\\ &\mathrm{op}(q_jR_1(p_k))=\frac{1}{2}[\mathrm{op}(q_j)R_1(\mathrm{op}(p_k))+R_1(\mathrm{op}(p_k))\mathrm{op}(q_j)].\end{aligned}$$

Again all situations can be checked in a similar fashion and we only check the case $R(q_j)=q_j^3$. Write $\mathrm{op}(q_j^3)=\mathrm{op}(q_j)^3+T$. Then $\{q_j^3,q_k\}=0$ and $\{q_j^3,p_k\}=-3q_j^2\delta_{jk}$ implies that T commutes with all $\mathrm{op}(q_k)$, $\mathrm{op}(p_k)$, $k=1,2,\ldots,n$. But then T also commutes with $\mathrm{op}(q_j)\mathrm{op}(p_k)+\mathrm{op}(p_k)\mathrm{op}(q_j)$. Therefore

$$\begin{aligned}-\mathrm{op}(q_j^3)&=\frac{1}{3}\mathrm{op}(\{p_jq_j,q_j^3\})=\frac{1}{3i\hbar}[\mathrm{op}(p_jq_j),\mathrm{op}(q_j^3)]\\ &=\frac{1}{3i\hbar}\left[\frac{1}{2}(\mathrm{op}(q_j)\mathrm{op}(p_j)+\mathrm{op}(p_j)\mathrm{op}(q_j),\mathrm{op}(q_j)^3+T\right]\\ &=\frac{1}{6i\hbar}[\mathrm{op}(q_j)\mathrm{op}(p_j)+\mathrm{op}(p_j)\mathrm{op}(q_j),\mathrm{op}(q_j)^3]=-\mathrm{op}(q_j)^3.\end{aligned}$$

In other words, $T=0$.

With this fact at hand, consider the equality

$$\frac{1}{9}\{q_j^3,p_j^3\}=\frac{1}{3}\{q_j^2p_j,p_j^2q_j\}\quad\left(=q_j^2p_j^2\right).$$

Quantizing the equality we obtain

$$\frac{1}{9}[\mathrm{op}(q_j)^3,\mathrm{op}(p_j)^3]=\frac{1}{3}[\mathrm{op}(q_j^2p_j),\mathrm{op}(p_j^2q_j)].$$

The left-hand side is

$$-i\hbar^3\left(M_{q_j}^2\frac{\partial}{\partial q_j^2}+2M_{q_j}\frac{\partial}{\partial q_j}+\frac{2}{3}I\right)$$

while the right-hand side is

$$\frac{1}{3}\left(\mathrm{op}(q_j^2p_j)\mathrm{op}(p_j^2q_j)-\mathrm{op}(p_j^2q_j)\mathrm{op}(q_j^2p_j)\right).$$

This is a bit messy, so we compute separately

$$\begin{aligned}\mathrm{op}(q_j^2p_j)&=\frac{1}{2}\left(\mathrm{op}(q_j)^2\mathrm{op}(p_j)+\mathrm{op}(p_j)\mathrm{op}(q_j)^2\right)=-\frac{i\hbar}{2}\left(M_{q_j^2}\frac{\partial}{\partial q_j}+\frac{\partial}{\partial q_j}M_{q_j^2}\right)\\&=-i\hbar\left(M_{q_j}^2\frac{\partial}{\partial q_j}+M_{q_j}\right),\end{aligned}$$

and

$$\mathrm{op}(p_j^2q_j)=\frac{1}{2}\left(\mathrm{op}(p_j)^2\mathrm{op}(q_j)+\mathrm{op}(q_j)\mathrm{op}(p_j)^2\right)=-\hbar^2\left(\frac{\partial}{\partial q_j}+M_{q_j}\frac{\partial^2}{\partial q_j^2}\right).$$

So the right-hand side is

$$-i\hbar^3\left(M_{q_j}^2\frac{\partial}{\partial q_j^2}+2M_{q_j}\frac{\partial}{\partial q_j}+\frac{1}{3}I\right).$$

And they don't coincide. □

We would like to use this symmetric quantization scheme. To this end, let us note that the error is of the order of $\hbar^2$ (this after dividing by $\frac{1}{i\hbar}$). As $\hbar$ is already very small, $\hbar^2$ is even smaller. We are forced to allow an error in Dirac's third condition. It is nowadays standard to let the correspondence principle be

- (3) $\mathrm{op}(\{f,g\})=\frac{1}{i\hbar}[\mathrm{op}(f),\mathrm{op}(g)]+O(\hbar)$.

We say that an operator S is $O(\hbar)$ if there is a constant C such that $\|S\psi\|\leq C\hbar\|\psi\|$ for all vectors ψ.

This modification will allow the construction of at leaset one quantization scheme, as we will see below, but unfortunately it will also permit the coexistence of an entire menagerie of non-equivalent quantization schemes.

2.4 Weyl quantization

13. This quantization method was introduced by Hermann Weyl in 1931. It applies to systems of particles with no constraints, where the configuration space is some $\mathbb{R}^n$ and the phase space is $\mathbb{R}^{2n}$. It is a method that is versatile in terms of which smooth functions can

be quantized, but unfortunately it cannot be quantized, except to configuration spaces that are quotients of $\mathbb{R}^n$ by the action of some discrete group (such as tori).

For simplicity, let us consider first the 1-dimensional case, and let me explain the idea behind this quantization method. We have to find a rule for associating linear operators to smooth functions on $\mathbb{R}^2$. The quantization rules tell us how linear functions can be quantized, and we want to extend that method to a sufficiently large family of smooth functions. Linearity helps, if we can identify a smaller family of "nicer" functions that we can quantize, and any other function is a linear combination of these, then we are done. Maybe we can even force infinite sums, although this is not required by the axioms, or integrals of such functions.

This brings us to the idea of using the Fourier transform. Recall that for an L^1 function $f(p,q)$, which we also need to be smooth for our purpose, the Fourier transform is

$$\hat{f}(a,b) = \int f(p,q)e^{-2\pi i(ap+bq)}dpdq,$$

and the inverse Fourier transform yields

$$f(p,q) = \int \hat{f}(a,b)e^{2\pi i(ap+bq)}dadb.$$

So f is the "sum" of the functions $e^{i(ap+bq)}$, and if we know how to quantize these, then we know how to quantize f itself. We can use the power series expansion

$$e^{i(ap+bq)} = \sum_{j=0}^{\infty} \frac{(2\pi i)^j}{j!}(ap+bq)^j,$$

so we can define

$$\mathrm{op}(e^{i(ap+bq)}) = \sum_{j=0}^{\infty} \frac{(2\pi i)^j}{j!}(a\mathrm{op}(p)+b\mathrm{op}(q))^j,$$

but does this series converge? The answer is yes, this series converges. The operator $ap+bq$ is self-adjoint ($a,b \in \mathbb{R}$ and $\mathrm{op}(p), \mathrm{op}(q)$ are self-adjoint). And it is an established fact that if A is self-adjoint operator (bounded or not), then

$$U = e^{iA} = \sum_{j=0}^{\infty} \frac{(i)^j}{j!}A^j$$

can be extended to a *unitary* operator on the entire Hilbert space. Recall that a unitary operator is one that is continuous, one-to-one, onto, and preserves the Hilbert space norm (in fact it preserves the inner product). Said differently, U is unitary if it is bounded and $UU^* = U^*U = I$. That U preserves the inner product follows from

$$\left\langle e^{iA}\psi_1, e^{iA}\psi_2 \right\rangle = \left\langle \psi_1, e^{-iA}e^{iA}\psi_2 \right\rangle = \left\langle \psi_1, \psi_2 \right\rangle.$$

Here we used the fact that $(iA)^* = -iA$ and then applied it term by term to the series expansion of the exponential. Now extend U to the entire Hilbert space. Its inverse is the

continuous extension of e^{-iA}, so we have $U : H \to H$ unitary. Since U is unitary, $\|U\| = 1$, exactly how $|e^{i(ap+bq)}| = 1$, so if the "clasical" series converges, then so does the "quantized" series.

It is important to point out that the Fourier transform is not defined for nonconstant polynomials (the integral does not converge), but Weyl quantization requires that $\mathrm{op}((ap + bq)^j) = (a\mathrm{op}(p) + b\mathrm{op}(q))^j$, and this prompts us to make the following definition:

Definition. For monomials in the position q and momentum p the Weyl quantization is defined by

$$\mathrm{op}(q^m p^n) = \frac{1}{(m+n)!} \sum_{\sigma} \sigma(\mathrm{op}(q), \ldots, \mathrm{op}(q), \mathrm{op}(p), \ldots, \mathrm{op}(p))$$

where the sum is taken over all possible permutations of $m + n$ objects and σ of m copies $\mathrm{op}(q)$ and n copies of $\mathrm{op}(p)$ means the product of those $\mathrm{op}(q)$ and $\mathrm{op}(p)$ multiplied in the order specified by σ. Weyl quantization is then extended linearly to all polynomials.

Proposition 2.4.1. Weyl quantization satisfies

$$\mathrm{op}(aq + bp)^m) = (a\,\mathrm{op}(q) + b\,\mathrm{op}(p))^m.$$

Proof. This is just an easy consequence of the definition. Note for example that

$$\begin{aligned}(a\,\mathrm{op}(q) + b\,\mathrm{op}(p))^3 &= a^3\mathrm{op}(q)^3 + a^2 b\,\mathrm{op}(q)^2\mathrm{op}(p) + a^2 b\,\mathrm{op}(q)\mathrm{op}(p)\mathrm{op}(q) + a^2 b\,\mathrm{op}(p)\mathrm{op}(q)^2 \\ &\quad + ab^2\mathrm{op}(q)\mathrm{op}(p)^2 + ab^2\mathrm{op}(p)\mathrm{op}(q)\mathrm{op}(p) + ab^2\mathrm{op}(p)^2\mathrm{op}(2) + b^3\mathrm{op}(p)^3 \\ &= \mathrm{op}(a^3q^3) + 3\mathrm{op}(a^2bq^2p) + 3\mathrm{op}(ab^2qp^2) + \mathrm{op}(p^3) = \mathrm{op}((aq + bp)^3).\end{aligned}$$

□

This means that we can define the quantization of a power series in $aq+bp$ as well through a limiting process, if the partial sums of operators converges. In particular we can define

$$\mathrm{op}\left(e^{2\pi i(aq+bp)}\right) = \sum_{m=0}^{\infty} \frac{1}{m!}[2\pi i(a\,\mathrm{op}(q) + b\,\mathrm{op}(p))]^m = e^{2\pi i(a\,\mathrm{op}(q)+b\,\mathrm{op}(p))}.$$

This operator is defined by

$$\mathrm{op}\left(e^{2\pi i(aq+bp)}\right)\psi = \lim_{N\to\infty} \sum_{m=0}^{N} \frac{1}{m!}[2\pi i(a\,\mathrm{op}(q) + b\,\mathrm{op}(p))]^m\psi$$

for all smooth functions $\psi \in L^2(\mathbb{R})$, whenever the series converges in $L^2(\mathbb{R})$.

We emphasize that although $a\,\mathrm{op}(q) + b\,\mathrm{op}(p)$ is not a bounded operator and is only defined on a dense subset of the Hilbert space, $\mathrm{op}\left(e^{2\pi i(aq+bp)}\right)$ is a unitary operator that is defined everywhere.

14. Now we can define Weyl quantization in $\mathbb{R}^{2n}$. We first do it first for polynomials by generalizing the above definition, then for exponential functions:

$$\mathrm{op}\left(e^{2\pi i(\mathbf{a}^T\mathbf{q}+\mathbf{b}^T\mathbf{p})}\right) = e^{2\pi i(\mathbf{a}^T\mathrm{op}(\mathbf{q})+\mathbf{b}^T\mathrm{op}(\mathbf{p}))},$$

where $\mathrm{op}(\mathbf{q}) = (\mathrm{op}(q_1), \mathrm{op}(q_2), \ldots, \mathrm{op}(q_n))$ and $\mathrm{op}(\mathbf{p}) = (\mathrm{op}(p_1), \mathrm{op}(p_2), \ldots, \mathrm{op}(p_n))$.

And then, for $f \in C^\infty(\mathbb{R}^n \times \mathbb{R}^n, \mathbb{C})$ we expand f into elementary oscillations using the Fourier transform and replace each oscillation (exponential function) by its Weyl quantization. This means that we set

$$\hat{f}(\mathbf{a}, \mathbf{b}) = \int_{\mathbb{R}^{2n}} f(\mathbf{q}, \mathbf{p}) e^{-2\pi i(\mathbf{a}^T\mathbf{q} + \mathbf{b}^T\mathbf{p})} d\mathbf{q}d\mathbf{p},$$

so that

$$f(\mathbf{q}, \mathbf{p}) = \int_{\mathbb{R}^{2n}} \hat{f}(\mathbf{a}, \mathbf{b}) e^{2\pi i(\mathbf{a}^T\mathbf{q} + \mathbf{b}^T\mathbf{p})} d\mathbf{a}d\mathbf{b},$$

and then set

$$\mathrm{op}(f) = \int_{\mathbb{R}^{2n}} \hat{f}(\mathbf{a}, \mathbf{b}) e^{2\pi i(\mathbf{a}^T\mathrm{op}(\mathbf{q}) + \mathbf{b}^T\mathrm{op}(\mathbf{p}))} d\mathbf{a}d\mathbf{b}.$$

The fact that the exponential operators are unitary makes the convergence issues parallel those that appear in the real-valued case.

15. As we explained before we have the following result for Weyl quantization.

Proposition 2.4.2. If R_1 and R_2 are polynomials of total degree at most 2 in $q_1, q_2, \ldots, q_n$, $p_1, p_2, \ldots, p_n$, then

$$\mathrm{op}(\{R_1, R_2\}) = \frac{1}{i\hbar}[\mathrm{op}(R_1), \mathrm{op}(R_2)].$$

We also know that for Weyl quantization this condition fails for certain polynomials of degree 3, by the weak form of the Groenewold No-Go Theorem. The question that remains is whether we could construct a quantization model that satisfies the condition on the commutator exactly, but maybe it is defined differently for quadratic polynomials. The answer is negative.

Theorem 2.4.1. (Groenewold's No-Go Theorem) There exists no linear map from the Lie algebra of polynomials of degree at most 4 to the Lie algebra of differential operators (with bracket $\frac{1}{i\hbar}[\cdot, \cdot]$ that extends the quantization scheme for the Heisenberg Lie algebra.

Proof. Here is the idea of the proof. Let

$$D = \sum_k f_k(\mathbf{q}) \left(\frac{\partial}{\partial \mathbf{q}}\right)^{\mathbf{k}}.$$

If D commutes with all operators of the form M_{q_j} and $\frac{\partial}{\partial q_j}$, then it is a multiple of the identity operator.

One can show that every polynomial of degree at most 2 is a linear combination of Poisson brackets of polynomials of degree at most 2, and every polynomial of degree at most 3 is a linear combination of Poisson brackets of polynomial of degree at most 3. Then like in the first step in the proof of Theorem 2.2.2, we deduce that the quantization scheme coincides with Weyl quantization for polynomials of total degree less than or equal to 3. And then you redo the last step on the proof of Theorem 2.2.2. □

2.5 A menagerie of quantizations

First let us fix the notation: $\mathbf{a}^{\mathbf{j}} = a_1^{j_1} a_2^{j_2} \cdots a_n^{j_n}$. Here is a list of important quantization schemes:

1. Pseudodifferential operator quantization:

$$\mathrm{op}(\mathbf{q}^{\mathbf{j}}\mathbf{p}^{\mathbf{k}}) = \mathrm{op}(\mathbf{q})^{\mathbf{j}}\mathrm{op}(\mathbf{p})^{\mathbf{k}},$$

   with the caveat that the resulting operators are not self-adjoint, it is nevertheless a useful scheme because of the many facts that are known about differential and more generally pseudodifferential operators (see my functional analysis lecture notes).

2. Symmetrized pseudodifferential operator quantization:

$$\mathrm{op}(\mathbf{q}^{\mathbf{j}}\mathbf{p}^{\mathbf{k}}) = \frac{1}{2}(\mathrm{op}(\mathbf{q})^{\mathbf{j}}\mathrm{op}(\mathbf{p})^{\mathbf{k}} + \mathrm{op}(\mathbf{p})^{\mathbf{k}}\mathrm{op}(\mathbf{q})^{\mathbf{j}}).$$

3. Weyl quantization:

$$\mathrm{op}((\mathbf{a}^T\mathbf{q} + \mathbf{b}^T\mathbf{p})^m) = (\mathbf{a}^T\mathrm{op}(\mathbf{q}) + \mathbf{b}^T\mathrm{op}(\mathbf{p}))^m.$$

4. Wick quantization:

$$\mathrm{op}((\mathbf{q} + i\mathbf{p})^{\mathbf{j}}(\mathrm{op}(\mathbf{q} - i\mathbf{p})^{\mathbf{k}}) = (\mathrm{op}(\mathbf{q}) - i\mathrm{op}(\mathbf{p}))^k(\mathrm{op}(\mathbf{q}) + i\mathrm{op}(\mathbf{p}))^{\mathbf{j}}.$$

5. Anti-Wick quantization, also known as Toeplitz quantization:

$$\mathrm{op}((\mathbf{q} + i\mathbf{p})^{\mathbf{j}}(\mathrm{op}(\mathbf{q} - i\mathbf{p})^{\mathbf{k}}) = (\mathrm{op}(\mathbf{q}) + i\mathrm{op}(\mathbf{p}))^{\mathbf{j}}(\mathrm{op}(\mathbf{q}) - i\mathrm{op}(\mathbf{p}))^{\mathbf{k}}.$$

   The last two are important in the context of the Segal-Bargmann quantization model, with holomorphic functions, which is related to the model for quantum field theory constructed V.A. Fock.

Chapter 3

The Heisenberg Group

3.1 The Heisenberg group

16. It is now time to have our first encounter with Lie groups (named after Sophus Lie) and their representations. Lie groups lie at the intersection of algebra and geometry. They are both

• groups: they have a multiplication that is associative, has an identity element, and every element has an inverse

• manifolds: locally they look like $\mathbb{R}^n$ (well, the manifold can also be infinite dimensional but in this class we do not care about that case)

But what is important it that the algebraic and geometric structure are related. This is expressed by the fact that the multiplication: $(x, y) \mapsto xy$ and the operation of taking the inverse: $x \mapsto x^{-1}$ are continuous maps.

Since Lie groups are manifolds, they have tangent spaces at each point, and in particular at the origin. The multiplication of the Lie group induces an operation on the tangent space at the origin, called Lie bracket, which is the infinitesimal form of the multiplication of the Lie group. The Lie bracket $(X, Y) \mapsto [X, Y]$ satisfies the properties 1, 2 and 4 of the Poisson bracket as listed in Theorem 1.2.3. These are the defining properties of a Lie algebra structure on a vector space.

Example 3.1.1. The group of complex numbers of absolute value equal to 1

$$U(1) = \{e^{it} \,|\, t \in \mathbb{R}\}$$

is a Lie group because it is a manifold – the circle – and the multiplication and taking the inverse are continuous (because exponentiation is continuous). Its Lie algebra, denoted by $u(1)$, is the tangent line to the circle at 1. It is generated by $\partial/\partial t$, but the Lie bracket is trivial. The group and its Lie algebra are shown in Figure 3.1.

Returning to Lie groups, Cayley has noticed that every group is a group of transformations. But we do not like transformations (i.e. bijections of sets) at large, we prefer transformations of some geometric space, such as $\mathbb{R}^n$. So then the bigger question arises whether groups can be described as groups of linear transformations of a vector space. Well,

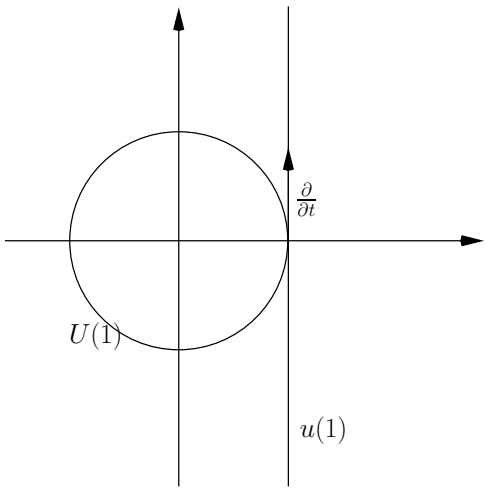

Figure 3.1: The group $U(1)$.

sometimes they can, sometimes they cannot. We make a compromise: study the homomorphisms from our group to a group of linear transformations. This is called representation theory; we say that our group is represented on a vector space, the homomorphism is called a representation. If the homomorphism is an isomorphism, the Lie group "is" a group of matrices. For example the group of rotations of the 3-dimensional is isomorphic to the group of orthgonal 3×3 matrices, and it is the convention to identify it with this group.

If the Lie group G is a closed subgroup of the group of $n \times n$ invertible matrices $GL(n, \mathbb{C})$, then its Lie algebra consists of the matrices X with the property that e^{tX} is in G for all $t \in \mathbb{R}$. In that case solving the equation

$$e^{tZ} = e^{tX} e^{tY}$$

yields, by the Baker-Campbell-Hausdorff formula,

$$Z = X + Y + \frac{t}{2}[X, Y] + O(t^2).$$

The Lie bracket is the commutator of matrices, and we notice that it is the first order term in the obstruction for $e^{tX}e^{tY}$ to equal $e^{t(X+Y)}$. More precisely (to get rid of the factor of $1/2$), it is the first order term in

$$e^{tX}e^{tY} - e^{tY}e^{tX}.$$

You can make sense of this in general as follows. Take a tangent vector at the identity element, X and extend it to a left invariant vector field by translating it all over the manifold (using the infinitesimal version of the maps $h \mapsto L_g h = gh$). Let $\bar{X}$ be its extension. This is a derivation on smooth functions. For two such left invariant vector fields, $\bar{X}, \bar{Y}$, that extend X, Y, define a left invariant vector field as the vector field that defines the derivation $[\bar{X}, \bar{Y}](f) = \bar{X}(\bar{Y}(f)) - \bar{Y}(\bar{X}(f))$. Restrict this vector field to the identity element to obtain $[X, Y]$.

Conversely, we can start with a Lie algebra of linear transformations of a vector space. Its elements are linear transformations and the Lie bracket is the commutator $[A, B] = AB - BA$. Then the exponentials

$$e^A = I + \frac{1}{1!}A + \frac{1}{2!}A^2 + \cdots + \frac{1}{n!}A^n + \cdots$$

form a group because by the Baker-Campbell-Hausdorff formula

$$e^A e^B = e^{A+B+c(A,B)}$$

where $c(A, B)$ is an expression obtained from A and B by applying only addition and commutators. So the exponent on the right also lies in the Lie algebra, and thus we are in th presence of a group. The Lie algebra is a vector space, so it has a manifold structure, and the exponential is a local diffeomorphism, inducing a manifold structure on the group of exponentials. Thus we obtain a Lie group.

While Lie groups arose in classical physics, quantum theory has a surprisingly deep connection to Lie groups and their representation theory.

17. If we start with the Heisenberg Lie algebra generated by the position and momentum operators, then the associated Lie group is the Heisenberg Lie group. More precisely, we start with the Lie algebra spanned as a real vector space by

$$Q_j = 2\pi i M_{q_j}, \quad P_j = 2\pi\hbar\frac{\partial}{\partial q_j}, \quad I,$$

acting on $L^2(\mathbb{R}^n)$. We exponentiate these as dictated by Weyl quantization. What we obtain is a Lie group. This is the Heisenberg group.

18. So let us study the Heisenberg group, which is determined by the Weyl quantizations of the exponential functions.

Let us denote

$$\exp(\mathbf{a}^T\mathbf{Q} + \mathbf{b}^T\mathbf{P} + tI) = e^{2\pi i t}\mathrm{op}\left(e^{2\pi i(\mathbf{a}^T\mathbf{q} + \mathbf{b}^T\mathbf{p})}\right).$$

Theorem 3.1.1. We have the following action on $L^2(\mathbb{R}^n)$:

$$\exp(\mathbf{a}^T\mathbf{Q} + \mathbf{b}^T\mathbf{P} + tI)\psi(\mathbf{q}) = e^{2\pi i\mathbf{a}^T\mathbf{q} + \pi i h\mathbf{a}^T\mathbf{b} + 2\pi i t}\psi(\mathbf{q} + h\mathbf{b})$$

One should stress out that in the statement of the theorem we have Planck's constant not the reduced Planck's constant.

Proof. The exponential of a differentiation operator is a translation and the exponential of the multiplication by a function is the multiplication by the exponential of that function. The second of these statements is easier to understand. Here is an explanation of the first. The expression

$$e^{t\frac{d}{dx}}\psi(x) = \sum_{n=0}^{\infty}\frac{t^n}{n!}\frac{d^n\psi}{dx^n}(x)$$

is just the MacLaurin expansion of ψ about x, so for all analytical functions it equals $f(x+t)$. But analytical functions are dense in $L^2(\mathbb{R})$ (because for examples functions of the form $e^{-x^2}P(x)$ where P is a polynomial are). So the exponential of differentiation equals translation for a dense set of functions. But both operators are continuous so they coincide everywhere.

So we have

$$\exp(b_j P_j)\psi(q_1,\dots,q_n) = \exp(2\pi i(-i\hbar b_j \frac{\partial}{\partial q_j}))\psi(q_1,\dots,q_n) = \psi(q_1,\dots,q_j+hb_j,\dots,q_n)$$
$$\exp(a_j Q_j)\psi(q_1,\dots,q_n) = e^{2\pi i a_j q_j}\psi(q_1,\dots,q_n)$$
$$\exp(2\pi i t I)\psi(q_1,\dots,q_n) = e^{2\pi i t}\psi(q_1,\dots,q_n).$$

Whenever two operators commute, their exponentials commute, so it is of no difficulty to prove the result for $\mathbf{a}=0$ or for $\mathbf{b}=0$. The general case follows from the Baker-Campbell-Hausdorff formula for operators whose second commutators are zero ($[A,[A,B]]=0$):

$$e^A e^B = e^{A+B+\frac{1}{2}[A,B]}. \quad \square$$

Corollary 3.1.1. We have the following multiplication rule for quantized exponentials

$$\exp(\mathbf{a}^T\mathbf{Q}+\mathbf{b}^T\mathbf{P}+tI)\exp(\mathbf{a}'^T\mathbf{Q}+\mathbf{b}'^T\mathbf{P}+t'I)$$
$$= \exp\left((\mathbf{a}+\mathbf{a}')^T\mathbf{Q}+(\mathbf{b}+\mathbf{b}')^T\mathbf{P}+t+t'-\frac{h}{2}(\mathbf{a}^T\mathbf{b}'-\mathbf{b}^T\mathbf{a}')\right).$$

(The minus sign is not a typo, note that $\mathbf{P}=2\pi i\mathrm{op}(\mathbf{p})$ and $\mathbf{Q}=2\pi i\mathrm{op}(\mathbf{q})$ so when taking the commutator, $i^2=-1$.)

19. This prompts us to make the following definition:

Definition. The *Heisenberg group* with real entries $\mathbf{H}(\mathbb{R}^n)$ is $\mathbb{R}^n\times\mathbb{R}^n\times U(1)$ with multiplication

$$(\mathbf{a},\mathbf{b},e^{2\pi it})\cdot(\mathbf{a}',\mathbf{b}',e^{2\pi it'}) = \left(\mathbf{a},\mathbf{a}',\mathbf{b}+\mathbf{b}',e^{2\pi i(t+t'-\frac{h}{2}(\mathbf{a}^T\mathbf{b}'-\mathbf{b}^T\mathbf{a}'))}\right).$$

For further use, we let

$$\omega((\mathbf{a},\mathbf{b}),(\mathbf{a}',\mathbf{b}')) = -h(\mathbf{a}^T\mathbf{b}'-\mathbf{b}^T\mathbf{a}')),$$

where h is Planck's constant.

The Heisenberg group is precisely the group of Weyl quantizations of exponential functions in the plane. One should emphasize that the Heisenberg group is a Lie group whose Lie algebra is the Heisenberg Lie algebra. Theorem 3.1.1 provides a faithful infinite dimensional representation

$$\rho_0 : \mathbf{H}(\mathbb{R}^n) \to U(L^2(\mathbb{R}^n))$$

of the Heisenberg group as a group of unitary operators on $L^2(\mathbb{R}^n)$. Here faithful means that it is injective. This is sometimes called the Schrödinger representation.

We have the following exponential form for the canonical commutation relations

$$\exp P_j \exp Q_k = e^{2\pi i h\delta_{jk}}\exp Q_k \exp P_j$$
$$\exp P_j \exp P_k = \exp P_k \exp P_j$$
$$\exp Q_j \exp Q_k = \exp Q_k \exp Q_j.$$

3.2 The Stone-von Neumann theorem

Theorem 3.2.1. (a) The representation ρ_0 of the Heisenberg group from Theorem 3.1.1 is irreducible.
(b) Any unitary irreducible representation ρ of $\mathbf{H}(\mathbb{R}^n)$ such that $\rho(\exp(tI)) = e^{2\pi it}I$ is unitary equivalent to ρ_0.

The proof (which we borrowed from Gerard Lion's book) is quite long and we divide it in steps.

20. To prove (a), note first that every element of

$$\mathbf{H}(\mathbb{R}^n) = \mathbb{R}^{2n} \times U(1) = B \times U(1),$$

is of the form $b\exp(tI)$, with $b \in B = \mathbb{R}^{2n}$. Identify B with $\mathbf{H}(\mathbb{R}^{2n})/U(1)$. Using this splitting we can turn a representation ρ of the Heisenberg group into a "representation" of the *additive group* $B = \mathbb{R}^{2n}$. Correctly stated, it is a projective representation.

Now let us consider compactly supported functions on the Heisenberg group with that are equivariant with respect to the circle action, meaning that

$$f(ue^{2\pi it}) = e^{-2\pi it}f(u).$$

We can identify these with functions on B, but we should always keep in mind the equivariance condition.

For a function f with compact support in $B = \mathbb{R}^{2n}$, which can be thought of as a compactly supported equivariant function on the Heisenberg group, we define the bounded operator $W_\rho(f) : H \to H$ by

$$\langle W_\rho(f)x, y\rangle = \int_B f(b)\,\langle \rho(b)x, y\rangle\, db.$$

This is called the Weyl transform of f. We can write it as

$$W_\rho(f) = \int_B f(b)\rho(b)db.$$

So we do not focus just on the operators defined by the representation of the Heisenberg group, we allow "sums" of such operators (in the sense of integration).

21. Let us now open a parenthesis in operator theory. A *Hilbert-Schmidt operator* on a Hilbert space H is an operator A satisfying

$$\|A\|_{HS}^2 = \sum_j \|Ae_j\|^2 < \infty$$

for some orthonormal basis $(e_j)_j$. This sum does not depend on the orthonormal basis, and defines the Hilbert-Schmidt norm $\|A\|_{HS}$. The Hilbert-Schmidt operators form a Hilbert space themselves (wow!), with the inner product

$$\langle A, B\rangle_{HS} = \sum_j \langle Ae_j, Be_j\rangle\,,$$

and with orthonormal basis the rank one operators

$$E_{j,k}(x) = \langle x, e_j \rangle \, e_k.$$

This means that every Hilbert-Schmidt operator is of the form

$$\sum_{jk} a_{jk} e_k \otimes e_j^* = \sum_j a_{jk} < \cdot, e_j > e_k,$$

with

$$\|A\|_{HS} = \left(\sum_{jk} |a_{jk}|^2 \right)^{1/2}.$$

We denote by $\mathcal{HS}(H)$ the Hilbert space of Hilbert-Schmidt operators on H. Note for example that $E_{j,j}$ are the projections onto the axes of coordinates. We should point out that every finite rank operator is Hilbert-Schmidt, and that finite rank operators form a dense set in the Hilbert space of Hilbert-Schmidt operators. So all Hilbert-Schmidt operators are compact, meaning that they map bounded sets to sets whose closure is compact (but not all compact operators are Hilbert-Schmidt).

We can define in general $E_{x,y} = \langle \cdot, y \rangle \, x$. Then

$$\langle E_{x,y}, E_{x',y'} \rangle_{HS} = \langle x, x' \rangle \, \overline{\langle y, y' \rangle}.$$

The map

$$(x, y) \mapsto E_{x,y}$$

defines a bilinear isometry from $H \otimes H^*$ to $\mathcal{HS}(H)$. As the Riesz Lemma implies that a Hilbert space is isometric with its dual, we have an isometry from $H \otimes H$ to the space of Hilbert-Schmidt operators that is linear in the first term and antilinear in the second. I know it is a bit hard to understand the tensor product $H \otimes H$, but if it is infinite dimensional and separable then it is isomorphic to $L^2(\mathbb{R})$ and then $H \otimes H$ is isomorphic to $L^2(\mathbb{R}^2)$, with $f \otimes g$ identified with $f(x)g(y)$.

22. We make the observation that $W_{\rho_0}(f)$ is a "sum" of operators of the form $\rho(u)$, $u \in \mathbf{H}(\mathbb{R}^n)$, so if ρ_0 has in invariant subspace, then so do all $W_{\rho_0}(f)$. Since we are proving (a), we already know that the Hilbert space is $H = L^2(\mathbb{R}^n)$. Thus it suffices to show that there is no subspace that is invariant under all operators $W_{\rho_0}(f)$, or under all operators in the closure of this set of operators (see Proposition 3.2.2 below). For that we need the following result:

Proposition 3.2.1. W_{ρ_0} extends to an isomorphism $L^2(\mathbb{R}^{2n}) \to \mathcal{HS}(L^2(\mathbb{R}^n))$ which is a constant multiple of an isometry.

Proof. Step 1. We show that the Weyl transform of a function is an integral operator.

Set $\ell = \mathbb{R}^n \oplus 0$, $\ell' = 0 \oplus \mathbb{R}^n$. We identify $L^2(\mathbb{R}^n)$ with $L^2(\ell')$. Then for $\psi \in L^2(\mathbb{R}^n)$, we have

$$(W_{\rho_0}(f)\psi)(\mathbf{q}) = \iint f(\mathbf{x},\mathbf{y}) e^{\pi i \mathbf{x}^T(2\mathbf{q}+h\mathbf{y})} \psi(\mathbf{q}+h\mathbf{y}) d\mathbf{x} d\mathbf{y},$$

where we used Theorem 3.1.1. Set $\xi = \mathbf{q} + h\mathbf{y}$ to turn this into

$$h^{-n} \iint e^{\pi i \mathbf{x}^T(\mathbf{q}+\xi)} f\left(\mathbf{x}, \frac{\xi - \mathbf{q}}{h}\right) e^{2\pi i \mathbf{x}^T \xi} \psi(\xi) d\mathbf{x} d\xi = h^{-n} \iint e^{\pi i \mathbf{x}^T(\mathbf{q}+\xi)} f\left(\mathbf{x}, \frac{\xi - \mathbf{q}}{h}\right) d\mathbf{x} \psi(\xi) d\xi.$$

Define

$$K_f(\mathbf{q},\xi) = h^{-n} \int e^{\pi i \mathbf{x}^T(\xi+\mathbf{q})} f\left(\mathbf{x}, \frac{\xi - \mathbf{q}}{h}\right) d\mathbf{x}.$$

This is what is called a *kernel*, which defines an *integral operator*

$$\psi(\mathbf{q}) \mapsto \int K(\mathbf{q},\mathbf{y}) \psi(\mathbf{y}) d\mathbf{y}.$$

So we recognize the Weyl transform of f to be an integral operator.

Step 2. We show that the Hilbert space of Hilbert-Schmidt operators on $L^2(\mathbb{R}^n)$ is the same as the space of integral operators with L^2 kernels, and moreover, the L^2 norm of the kernel is the Hilbert-Schmidt norm of the operator.

Indeed, if A is a Hilbert-Schmidt operator on $L^2(\mathbb{R}^n)$ then

$$\begin{aligned}(A\psi)(\mathbf{x}) &= \sum_{jk} a_{jk} < \psi, e_k > e_j = \sum_{jk} a_{jk} \left(\int_{\mathbb{R}^n} \psi(\mathbf{y}) \overline{e_k(\mathbf{y})} d\mathbf{y} \right) e_j(\mathbf{x}) \\ &= \int_{\mathbb{R}^n} \psi(\mathbf{y}) \left(\sum_{jk} a_{jk} e_j(\mathbf{x}) \overline{e_k(\mathbf{y})} \right) d\mathbf{y}.\end{aligned}$$

If we denote

$$K_A(\mathbf{x},\mathbf{y}) = \int_{\mathbb{R}^n} \sum_{jk} a_{jk} e_j(\mathbf{x}) \overline{e_k(\mathbf{y})} d\mathbf{y},$$

then it is not hard to see that

$$\|K_A(\mathbf{x},\mathbf{y})\|_2 = \|A\|_{HS}.$$

Therefore we wrote the operator A as an integral operator

$$(A\psi)(\mathbf{x}) = \int_{\mathbb{R}^n} \psi(\mathbf{y}) K_A(\mathbf{x},\mathbf{y}) d\mathbf{y}.$$

For the converse, note that $e_j(\mathbf{x})e_k(\mathbf{y})$, $j,k \geq 1$ is an orthonormal basis of $L^2(\mathbb{R}^{2n})$, so every function in $L^2(\mathbb{R}^{2n}$ can be written in the form $\sum_{jk} a_{jk} e_j(\mathbf{x}) \overline{e_k(\mathbf{y})}$ and so is in the image of the above map.

Step 3. We show that there is a constant such that the L^2 norm of a function is that constant times the L^2 norm of the kernel K_f of its Weyl transform.

The partial Fourier transform

$$(\mathcal{F}_{\mathbf{x}}f)(\mathbf{q},\mathbf{y}) = \int e^{-2\pi i \mathbf{x}^T \mathbf{q}} f(\mathbf{x},\mathbf{y})\, d\mathbf{x}$$

is a constant multiple of a unitary automorphism of $L^2(\mathbb{R}^{2n})$. And

$$K_f(\mathbf{q},\mathbf{y}) = h^{-n}(\mathcal{F}_{\mathbf{x}}f)\left(-\frac{\mathbf{q}+\mathbf{y}}{2}, \frac{\mathbf{q}-\mathbf{y}}{h}\right).$$

This change of coordinates is also a constant multiple of a unitary automorphism. So W_ρ is a constant multiple of an isometry, which therefore can be extended to the entire $L^2(\mathbb{R}^n)$. And because Fourier transform, the change of coordinates are isomorphisms, and the correspondence between kernels and Hilbert-Schmidt operators are onto, this extension is onto as well. Thus the proposition is proved. □

23. Now (a) is a direct corollary of the following result.

Proposition 3.2.2. If T is a bounded operator that commutes with all operators of the form $\rho(u)$, $u \in \mathbf{H}(\mathbb{R}^n)$, then T is a multiple of the identity operator.

Proof. If T commutes with every operator of the form $\rho(u)$, the T also commutes with operators of the form $W_\rho(f)$ with f compactly supported. But then by continuity it also commutes with all operators of the form $W_\rho(f)$, $f \in L^2(\mathbb{R}^{2n})$. In particular it commutes with the operators of the form $E_{jk} = e_j \otimes e_k^*$.

If T is not a multiple of the identity operator, then there are noncolline ar vectors ψ_1, ψ_2 such that $T\psi_1 = \psi_2$. Using Gram-Schmidt, and starting with ψ_1 and ψ_2, we can construct a basis such that $< Te_1, e_2 >= \lambda \neq 0$. But the fact that T commutes with $E_{1,2}$ means that

$$0 = (e_1 \otimes e_2^*)(e_1) = TE_{1,2}e_1 = E_{1,2}Te_1 = (e_1 \otimes e_2^*)T(e_1) = \lambda e_1 \neq 0.$$

This is a contradiction and the proposition is proved. □

Now if the Schrödinger representation were reducible, then we could write $L^2(\mathbb{R}^n) = H_1 \oplus H_2$, with H_1 and H_2 invariant subspaces (here it is important that the representation is unitary). But then the orthogonal projection onto H_1 would commute with $\rho(u)$ for all u, and this would contradict the above proposition.

24. We now start the proof of (b). Let ρ be a representation of $\mathbf{H}(\mathbb{R}^n)$ on a Hilbert space H such that $\rho(\exp tI) = e^{2\pi i t} I$, we do not even assume irreducibility. First, for two continuous compactly supported functions f, g on the Heisenberg group we define their convolution

$$f * g(u) = \int_{\mathbf{H}(\mathbb{R}^n)} f(v) g(v^{-1}u) dv,$$

where dv is the translation-invariant measure on $\mathbb{R}^n \times U(1)$.

Restrict the convolution to functions that satisfy the equivariance condition

$$f(ue^{2\pi it}) = e^{-2\pi it} f(u).$$

We can identify these with functions on B, but we should always keep in mind the equivariance condition. The convolution defined above becomes a convolution for functions on B (the minus sign appears when going from multiplicative to additive using the BCH formula):

$$(f * g)(b) = \int_B f(b')g(b - b')e^{\pi i\omega(b',b)}db.$$

Define also the left and right actions of the Heisenberg group on functions

$$(u * f)(b) = f(u^{-1}b), \quad (f * u)(b) = f(bu^{-1}).$$

We recall the definition of the Weyl transform, but beware, in this case we can no longer use the space $L^2(\mathbb{R}^n)$ so we can define this transform only for compactly supported functions. In fact, the Weyl transform can be extened to the algebra of rapidly decreasing functions $\mathcal{S}(\mathbb{R}^{2n})$ ($f \in \mathcal{S}(\mathbb{R}^{2n})$ if $\sup_{x\in R}^n |x^\alpha D^\beta f(x)| < \infty$ for all $\alpha, \beta \in \mathbb{Z}_+^n$) whose multiplication is the convolution $*$.

The Weyl transform is well behaved with respect to the convolution, in the sense that it satisfies the following properties, which are not hard to check:

(i) $W_\rho(f * g) = W_\rho(f)W_\rho(g)$,

(ii) $W_\rho(f^*) = W_\rho(f)^*$, where $f^*(v) = \overline{f(v^{-1})}$ (or when working with functions on B, $f^*(b) = \overline{f(-b)}$),

(iii) $W_\rho(u * f) = \rho(u)W_\rho(f)$,

(iv) $W_\rho(f * u) = W_\rho(f)\rho(u)$.

So the Weyl transform defines a homomorphism W_ρ from the algebra of rapidly decreasing functions endowed with the convolution $*$ into a subalgebra of operators $\mathcal{A}$ on H such that $W_\rho(f^*) = (W_\rho(f))^*$.

There is an isomorphism of $\mathcal{S}(\mathbb{R}^n)$ to the algebra $\mathcal{A}_0$ of integral operators on $L^2(\mathbb{R}^n)$ with rapidly decreasing kernel constructed using the Weyl transform W_{ρ_0} defined by the Schrödinger representation. Then $\Phi = W_\rho \circ W_{\rho_0}^{-1}$ is a homomorphism of operator algebras satsfying $\Phi(A^*) = (\Phi(A))^*$, for all A, and $\phi(I) = I$. This homomorphism goes from an algebra of integral operators that we understand, to an algebra of operators that we are trying to understand.

Fix $\psi_1 \in \mathcal{S}(\ell') = \mathcal{S}(0 \oplus \mathbb{R}^n)$ such that $\|\psi_1\|_2 = 1$. Consider the orthogonal projection operator $P_1 : L^2(\mathbb{R}^n) \to \mathbb{C}\psi_1$. Then P_1, being finite rank, is Hilbert-Schmidt, so it is defined by a kernel. In fact this kernel is $\psi_1(x)\overline{\psi_1(y)}$, so P_1 is of the form $W_{\rho_0}(g_1)$, for $g_1 \in \mathcal{S}(\mathbb{R}^{2n})$ (this is because g is mapped to $\psi_1(x)\overline{\psi_1(y)}$ by a partial Fourier transform and a change of coordinates, both of which preserve $\mathcal{S}(\mathbb{R}^{2n})$). Using the properties of projectors and property (iii) above we have

$$g_1 * g_1 = g_1, \quad g_1^* = g_1, \quad g_1 * u * g_1 = \langle \rho_0(u)\psi_1, \psi_1 \rangle \, g_1.$$

These relations and the multiplicative properties of the Weyl transform imply that $W_\rho(g_1)$ is a projector onto the space $H_1 = W_\rho(g_1)H$. Note that here we use the representation ρ, not the Schrödinger representation.

Lemma 3.2.1. The Hilbert space H is spanned by elements of the form $\rho(u)W_\rho(g_1)x$ for $u \in \mathbf{H}(\mathbb{R}^n)$ and $x \in H$.

Proof. Since ρ is a unitary representation, it suffices to show that the Hilbert space H is spanned by elements of the form $\rho(u)W_\rho(g_1)\rho(u^{-1})x$. Let $y \in H$ be orthogonal to all those elements. We have

$$\begin{aligned}
&\left\langle y, \rho(u)W_\rho(g_1)\rho(u^{-1})x\right\rangle = 0\\
&= \int_{\mathbb{R}^{2n}} \left\langle y, \rho(\exp(\mathbf{a}^T\mathbf{Q} + \mathbf{b}^T\mathbf{P})\rho(\exp(\mathbf{a}'^T\mathbf{Q} + \mathbf{b}'^T\mathbf{P}))\rho(\exp(-\mathbf{a}^T\mathbf{Q} - \mathbf{b}^T\mathbf{P}))x\right\rangle g_1(\mathbf{a}', \mathbf{b}')d\mathbf{a}'d\mathbf{b}'\\
&= \int_{\mathbb{R}^{2n}} \left\langle y, \rho(\exp(\mathbf{a}'^T\mathbf{Q} + \mathbf{b}'^T\mathbf{P})x\right\rangle g_1(\mathbf{a}', \mathbf{b}')e^{2\pi i\omega((\mathbf{a},\mathbf{b}),(\mathbf{a}',\mathbf{b}'))}d\mathbf{a}'d\mathbf{b}'.
\end{aligned}$$

This is a Fourier transform of $F(\mathbf{a}', \mathbf{b}') = \left\langle y, \rho(\exp(\mathbf{a}'^T\mathbf{Q} + \mathbf{b}'^T\mathbf{P})x\right\rangle g_1(\mathbf{a}', \mathbf{b}')$, which is identically equal to zero if and only if F is identically equal to zero. Since g_1 is not identically equal to zero, it follows that there is an element of the Heisenberg group, $u = \exp(\mathbf{a}'^T\mathbf{Q} + \mathbf{b}'^T\mathbf{P})$, such that $\langle y, \rho(u)x\rangle = 0$ for all $x \in H$. But $\rho(u)$ is unitary, so $\langle y, x\rangle = 0$ for all $x \in H$. Hence $y = 0$ and the lemma is proved. □

It is important to point out that if the representation ρ is irreducible, then in view of this result $W_\rho(g_1)H$ is one dimensional, because otherwise we have an invariant subspace spanned by $\rho(u)W_\rho(g_1)x_0$ for some x_0.

25. Before we proceed with the last step, we recall the notion of the tensor product of two vector spaces, and then of two Hilbert spaces.

For vector spaces V and W, construct the set $V \times W$. Now formally build a vector space whose basis consists of the elements of $V \times W$. Then factor this space by the relations
$(v_1 + v_2, w) = (v_1, w) + (v_2, w)$,
$(v, w_1 + w_2) = (v, w_1) + v(w_2)$,
$c(v, w) = (cv, w) = (v, cw)$,
for all $v, v_1, v_2 \in V$, $w, w_1, w_2 \in W$, $c \in C$. The resulting vector space is the tensor product $V \otimes W$. For example $\mathbb{C}^3 \otimes \mathbb{C}^2 = \mathbb{C}^6$.

If H_1, H_2 are Hilbert spaces, endow $H_1 \otimes H_2$ with the tensor product

$$\langle v_1 \otimes w_1, v_2 \otimes w_2\rangle = \langle v_1, v_2\rangle \langle w_1, w_2\rangle .$$

The completion of this space is the Hilbert space $H_1 \otimes H_2$. For example $L^2(\mathbb{R}^n) \otimes \mathbb{C}^2 = L^2(\mathbb{R}^n) \oplus L^2(\mathbb{R}^n)$ and $L^2(\mathbb{R}^n) \otimes L^2(\mathbb{R}^n) = L^2(\mathbb{R}^{2n})$.

26. Let us return to the space $H_1 = W_\rho(g_1)H$. We will prove that there exists a unitary isomorphism U between $H_0 \otimes H_1$ and H, where $H_0 = L^2(\mathbb{R}^n)$ is the Hilbert space onto which the Heisenberg group is represented via ρ_0 and H is the Hilbert space onto which the

Heisenberg group is represented via ρ, such that U interpolates between ρ_0 and ρ. Here ρ_0 acts on $H_0 \otimes H_1$ as $\rho_0 \otimes I_{H_1}$, that is $(\rho_0 \otimes I_{H_1})(u)(\psi \otimes x) = \rho_0(u)\psi \otimes x$. We define U by

$$U(\rho_0(u)\psi_1 \otimes w) = \rho(u)w, \quad u \in \mathbf{H}(\mathbb{R}^n), w \in H_1.$$

then extend this linearly.

It is important to point out that since ρ_0 is an irreducible representation, every nonzero element of $L^2(\mathbb{R}^n)$ is a cyclic vector for it, in particular so is ψ_1. Therefore the representation ρ_0 is irreducible, the set of linear combinations $\sum c_j \rho_0(u_j)\psi_1$ is dense in $L^2(\mathbb{R}^n)$. So if we check that U is an isometry, then we can extend it to the whole $H_0 \times H_1$. We check that for $w_1 = W_\rho(g_1)x_1$, $w_2 = W_\rho(g_1)x_2$, $u_1, u_2 \in \mathbf{H}(\mathbb{R}^n)$,

$$\langle \rho(u_1)w_1, \rho(u_2)w_2 \rangle_H = \langle \rho_0(u_1)\psi_1, \rho_0(u_2)\psi_1 \rangle_{H_0} \langle w_1, w_2 \rangle_{H_1}.$$

Indeed

$$\begin{aligned}
\langle \rho(u_1)W_\rho(g_1)x_1, \rho(u_2)W_\rho(g_1)x_2 \rangle_H &= \left\langle W_\rho(g_1)\rho(u_2)^{-1}\rho(u_1)W_\rho(g_1)x_1, x_2 \right\rangle_H \\
&= \left\langle W_\rho(g_1 * (u_2^{-1}u_1) * g_1)x_1, x_2 \right\rangle_H = \left\langle \rho_0(u_2^{-1}u_1)\psi_1, \psi_1 \right\rangle_{H_0} \langle W_\rho(g_1)x_1, x_2 \rangle_H \\
&= \left\langle \rho_0(u_2^{-1}u_1)\psi_1, \psi_1 \right\rangle_{H_0} \langle W_\rho(g_1)W_\rho(g_1)x_1, x_2 \rangle_H \\
&= \langle \rho_0(u_1)\psi_1, \rho_0(u_2)\psi_1 \rangle_{H_0} \langle W_\rho(g_1)x_1, W_\rho(g_1)x_2 \rangle_{H_1},
\end{aligned}$$

as claimed.

Now using Lemma 3.2.1, we deduce that U is surjective, so it is a unitary isomorphism between $H_0 \otimes H_1$ and H. The way U was defined shows that $U(\rho_0(u) \otimes I_{H_1})U^{-1} = \rho(u)$. So ρ is a multiple of the representation ρ_0. If ρ is irreducible, then H_1 is necessarily one-dimensional, and we have a unitary isomorphism between $L^2(\mathbb{R}^n) = H_0 = H_0 \otimes H_1 \cong H$. The theorem is proved.

3.3 The projective representation of the symplectic group

27. There is actually another irreducible representation of the Heisenberg group on $L^2(\mathbb{R}^n)$, given by

$$\exp(\mathbf{a}^T\mathbf{Q} + \mathbf{b}^T\mathbf{P} + tI)\psi(\mathbf{p}) = \exp(2\pi i(i\hbar\mathbf{a}^T\frac{\partial}{\partial \mathbf{p}} + \mathbf{b}^T\mathbf{p} + tI))\psi(\mathbf{p}),$$

where the position and momentum operators are

$$\mathrm{op}(q_j)\psi(\mathbf{p}) = i\hbar\frac{\partial \psi}{\partial p_j}, \quad \mathrm{op}(p_j)\psi(\mathbf{p}) = p_j\psi(\mathbf{p}).$$

This is the quantization in the *momentum representation*.

In this case

$$\exp(\mathbf{a}^T\mathbf{Q} + \mathbf{b}^T\mathbf{P} + tI)\psi(\mathbf{p})\psi(\mathbf{p}) = e^{2\pi i\mathbf{b}^T\mathbf{p} - \pi i h\mathbf{a}^T\mathbf{b} + 2\pi i t}\psi(\mathbf{p} - h\mathbf{a}),$$

which is the Schrödinger representation of the Heisenberg group in the momentum representation.

By the Stone-von Neumann Theorem, these two representations are unitary equivalent, in fact the unitary equivalence is the Fourier transform

$$(\mathcal{F}_h\psi)(\mathbf{p}) = h^{n/2}\int_{\mathbb{R}^n}\psi(\mathbf{q})e^{-2\pi ih\mathbf{p}^T\mathbf{q}}d\mathbf{q}.$$

28. Note that the above representation of the Heisenberg group, which arises from the quantization of the position and the momentum in the momentum representation can be obtained from the standard representation ρ_0 as follows: First apply to $(\mathbf{a}, \mathbf{b})$ the symplectomorphism with matrix

$$\begin{pmatrix} 0 & I \\ -I & 0 \end{pmatrix}$$

to change the element of the Heisenberg group, then apply ρ_0. We can do this for every element f of the symplectic group (the group of linear symplectomorphisms). By the Stone-von Neumann theorem, the resulting representation, which is irreducible, is unitary equivalent to ρ_0. We thus obtain a map $f \mapsto U_f$, from symplectomorphisms to unitary operators on $L^2(\mathbb{R}^n)$, which because of the Schur Lemma from representation theory satisfies

$$U_{f_1\circ f_2} = cU_{f_1}U_{f_2},$$

where c is a complex number of absolute value 1. This is the projective representation of the symplectic group. There is a way to get rid of the "projectivity" and turn this into a true group representation if we replace the symplectic group by its double cover, the *metaplectic group*. We obtain the Segal-Shale-Weil representation.

We should point out that this means that the symplectomorphisms which are elements of the symplectic group can be *quantized*. Note that the axioms of quantum mechanics make no provisions about the existence of a quantum version of the symplectomorphisms. In the next chapter we will encounter other symplectomorphisms that can be quantized.

Chapter 4

The Schrödinger Equation

4.1 The Schrödinger equation as the time evolution of a state

29. Up to this moment we have modeled the observables of quantum mechanical systems: what they are, how to measure them. We have modeled also the states of systems, which are the analogues of the locations of the particles in the phase space. Everything expressed in the language of functional analysis. We have not discussed time evolutions, how to make predictions of future measurements. What are the analogues of Hamilton's equations and Hamiltonian flows?

So let us assume that we know how to quantize every Hamiltonian, for example by using the method of Weyl quantization. We will use the short-hand notation $\hat{H}$ for $\mathrm{op}(H)$, the operator associated to H, and in general $\hat{f}$ for $\mathrm{op}(f)$.

In classical Hamiltonian mechanics, the state of a system is given by the position and momentum coordinates of the particle. The time evolution is described by Hamilton's equations.

In quantum mechanics, the quantum state of a system is given by a ray $[\psi]$ in a Hilbert space H (or by abuse of language by ψ itself). The *Schrödinger equation* describes the time evolution of the quantum state of a particle.

It has the form

$$i\hbar\frac{\partial\psi}{\partial t} = \hat{H}\psi.$$

"Where did we get that from? Nowhere. It is not possible to derive it from anything you know. It came out of the mind of Schrödinger." (R. Feynman)

There is a way of deriving this equation from physical intuition. Not a rigorous mathematical derivation, but an educated guess. It starts in simple relativity. This is modeled in a 4-dimensional space, the Minkowski space. The position is $\mathbf{x} = (\mathbf{q}, ict)$, where c is the speed of light and t is time, and the momentum is $\boldsymbol{\pi} = (\mathbf{p}, i\frac{E}{c})$ (because $E = mc^2$).

Now there are two ways in which we can quantize the momentum. The first is to use Dirac's rule for all four components and set

$$\hat{\pi} = -i\hbar\left(\frac{\partial}{\partial q_1}, \frac{\partial}{\partial q_2}, \frac{\partial}{\partial q_3}, \frac{\partial}{\partial ict}\right),$$

or to quantize the total energy E using a quantization rule and set

$$\hat{\pi} = -i\hbar \left(\frac{\partial}{\partial q_1}, \frac{\partial}{\partial q_2}, \frac{\partial}{\partial q_3}, \frac{i}{c}\hat{E} \right).$$

These "should" give the same answer, so $\hat{E} = i\hbar\frac{\partial}{\partial t}$. And since the total energy is equal to the Hamiltonian, for a time-dependent state $\psi(\mathbf{q}, t)$ we should have

$$i\hbar\frac{\partial \psi}{\partial t} = \hat{H}\psi.$$

Relativistic classical mechanics leads to Schrödinger's equation, but this equation is non-relativistic. Just to boost your intuition, we present an example, with more examples discussed in detail later on.

30. Schrödinger equation can be solved immediately. There is another way to look at this. Schrödinger's equation should remind us of the well known equation

$$\frac{dx}{dt} = ax.$$

This equation has the solution $x(t) = e^{at}x(0)$. For a system of first order equations

$$\frac{d\mathbf{x}}{dt} = A\mathbf{x},$$

we still have the solution $\mathbf{x}(t) = e^{At}\mathbf{x}(0)$. If we write the Schrödinger equation as

$$\frac{\partial \psi}{\partial t} = -\frac{i}{\hbar}\hat{H}\psi,$$

then

$$\psi(\mathbf{q}, t) = e^{-(i/\hbar)t\hat{H}}\psi(\mathbf{q}, 0).$$

And we are done. Well, not really... The operator $e^{-(i/\hbar)t\hat{H}}$ can be very hard to understand, and this is where the whole problem of solving the Schrödinger equation lies. What we can say is that, because $\hat{H}$ is a self-adjoint operator, and then so is $-(t/\hbar)\hat{H}$, then $e^{-(i/\hbar)t\hat{H}}$ is a unitary operator. Moreover, the additive property of exponentials induces a group structure on these operators

$$e^{-(i/\hbar)t\hat{H}}e^{-(i/\hbar)t'\hat{H}} = e^{-(i/\hbar)(t+t')\hat{H}}.$$

This group is the quantum analogue of the Hamiltonian flow discussed in the first chapter. So symplectomorphisms that come from Hamiltonian flows can be quantized, they are quantized by unitary operators.

31. We want to find a more profound connection between Schrödinger's equation and the Hamiltonian formalism of classical mechanics. Recall that the result of measuring an observable A at time t in state $\psi(t)$ should be

$$\langle A\psi(t), \psi(t)\rangle.$$

In the same vein as for Hamilton's equations, we want to understand the instantaneous rate of change of the result of the measurement. We compute

$$\langle A\psi(t), \psi(t)\rangle = \left\langle Ae^{-(i/\hbar)t\hat{H}}\psi(0), e^{-(i/\hbar)t\hat{H}}\psi(0)\right\rangle = \left\langle e^{(i/\hbar)t\hat{H}}Ae^{-(i/\hbar)t\hat{H}}\psi(0), \psi(0)\right\rangle.$$

Using the product rule for differentiation, we obtain

$$\frac{d}{dt}\langle A\psi(t), \psi(t)\rangle\,|_{t=0} = \left\langle [(i/\hbar)\hat{H}A - (i/\hbar)A\hat{H}]\psi(0), \psi(0)\right\rangle = \left\langle \frac{1}{i\hbar}[A, \hat{H}]\psi(0), \psi(0)\right\rangle.$$

We recognize the quantum Poisson bracket.

Now we can do this at every moment in time, not just $t = 0$, and write the Schrödinger equation for observables

$$\frac{d}{dt}\langle A\psi(t), \psi(t)\rangle = \left\langle \{A, \hat{H}\}_q\psi(t), \psi(t)\right\rangle.$$

This should remind us of Hamilton's equation

$$\frac{df}{dt} = \{f, H\},$$

where f is a classical observable. We arrive at the Heisenberg picture, with an equation that tells us how the result of the measurement of an observable changes in time as the state of the quantum system varies as specified by the Schrödinger equation.

We can interpret this differently assuming that the state (wave function) does not change, but the observable (operator) is the one that changes in time by the rule

$$A(t) = e^{(i/\hbar)t\hat{H}}Ae^{-(i/\hbar)t\hat{H}}.$$

Then we get the quantum version of Hamilton's equation, which is

$$\frac{dA}{dt} = \{A, \hat{H}\}_q.$$

Now we have a closer analogy between the Hamilton's equation for the time evolution of a classical observable f and the time evolution of its quantization:

$$\frac{df}{dt} = \{f, H\}, \quad \frac{d\mathrm{op}(f)}{dt} = \{\mathrm{op}(f), H\}_q,$$

It is also important to contrast the actions of the quantized Hamiltonian flow on states and on observables:

$$\psi \mapsto e^{-(i/\hbar)t\hat{H}}\psi, \quad A \mapsto e^{(i/\hbar)t\hat{H}}Ae^{-(i/\hbar)t\hat{H}}.$$

32. Let us return to the problem of solving the Schrödinger equation, and let us assume that H is time independent. Then so is $\hat{H}$. We try to separate variables by writing $\psi(t, q) = \phi(q)\tau(t)$. Then the equation becomes

$$i\hbar\phi\frac{d\tau}{dt} = (\hat{H}\phi)\tau.$$

So

$$\frac{i\hbar\frac{d\tau}{dt}}{\tau} = \frac{\hat{H}\phi}{\phi}.$$

Each of the two sides of the equality depends on a different set of variables (time on the left, space on the right), so both should be equal to a constant called E. Hence

$$\hat{H}\phi = E\phi.$$

This is the time-independent Schrödinger equation. It's solutions are also referred to as *stationary-state* solutions. We see that E is an eigenvalue of $\hat{H}$, so E is an energy level that can be measured by an experiment (according to the third axiom of quantum mechanics).

So Schrödinger's equation leads us naturally to considering the energy levels of the quantum Hamiltonian. There is another way to look at this, using the operator theoretical solution to Schrödinger's equation:

$$\psi(\mathbf{q}, t) = e^{(-i/\hbar)t\hat{H}}\psi(\mathbf{q}, 0).$$

What does this answer mean? If $\hat{H}$ is diagonalizable, then we can compute $e^{(-i/\hbar)t\hat{H}}$ the way we exponentiate a diagonalizable matrix. But for that we need to diagonalize $\hat{H}$, so we need to find the stationary state solutions. In other words, the trick of separation of variables is the method of diagonalizing operators.

Once we have diagonalized $\hat{H}$ in a basis ψ_j with eigenvalues E_j, we can write $\psi(\mathbf{q}, 0) = \sum_j c_j\psi_j$ and then the solution to Schrödinger's equation is

$$\psi(\mathbf{q}, t) = \sum_j c_j\psi_j e^{-\frac{i}{\hbar}E_j t}.$$

It is important to point out that not every self-adjoint operator can be diagonalized, but the intuition that diagonalization provides allows us to solve the problem even when the operator cannot be diagonalized. For those familiar with functional calculus, this is about passing from a projector valued measure on a discrete spectrum to a projector valued measure on a continuous spectrum, and from a sum to an integral.

We will now restrict ourselves to quantizations of systems that appear in Newtonian mechanics. If we take the quantization of the total energy of n particles in the three-dimensional space

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

then we have

$$\hat{H} = -\frac{\hbar^2}{2m}\sum_{j=1}^{3n}\frac{\partial^2}{\partial q_j^2} + M_{V(\mathbf{q})} = -\frac{hbar^2}{2m}\Delta + M_{V\mathbf{q}}.$$

So the Schrödinger equation is

$$i\hbar\frac{\partial\psi}{\partial t}(\mathbf{q}, t) = -\frac{\hbar^2}{2m}\Delta + V(\mathbf{q})\psi(\mathbf{q}, t).$$

This is a second order partial differential equation.

Example 4.1.1. A free particle in $\mathbb{R}^3$ has the Schrödinger equation

$$i\hbar\frac{\partial\psi}{\partial t} = -\frac{\hbar^2}{2m}\Delta\psi.$$

We try to apply the above considerations to this case. Now we are in trouble already! The self-adjoint operator that is the negative of the Laplacian, $-\Delta$, does not have discrete spectrum, in fact the spectrum of the Laplacian is $[0,\infty)$. So no eigenfunctions in $L^2(\mathbb{R}^3)$. We have to look for "eigenfunctions" outside of this space.

Now, we do not need to know what the spectrum of *Delta* is, all we need to know is that this operator is the square of a self-adjoint operator, so its spectrum is a subset of $[0,\infty)$ (this is for example the consequence of the fact that the spectrum of an operator of the form A^*A is a subset of $[0,\infty)$.

We have to solve the equation

$$-\Delta\psi = \lambda\psi,$$

where $\lambda = \frac{2mE}{\hbar^2}$. We have no (physical) constraints for the energy E, so we can let λ be any non-negative number. In fact we will look only at the cases where $\lambda > 0$. We do not even bother to find all stationary states ψ, only the ones where the variables can be separated:

$$\psi(q_1,q_2,q_3) = \psi_1(q_1)\psi_2(q_2)\psi_3(q_3).$$

The ψ_j satisfy the equation $\partial^2\psi_j/\partial q_j^2 = -\lambda\psi_j$, and this we know how to solve, and the solution is $c_1e^{k_1q_j} + c_2e^{-k_1'q_j}$. Putting these solutions together we obtain the stationary states

$$\phi(\mathbf{q}) = e^{i\mathbf{k}^T\mathbf{q}}, \quad \mathbf{k}\in\mathbb{R}^3, \|\mathbf{k}\|^2 = \lambda.$$

These are referred to as plane waves. Here $\mathbf{k}$ is the momentum of the plane wave, in agreement with

$$-i\hbar\frac{\partial}{\partial q_j}e^{i\mathbf{k}^T\mathbf{q}} = \hbar k_j e^{i\mathbf{k}^T\mathbf{q}}.$$

Now suppose the the initial condition is some $f(q_1,q_2,q_3) = \psi(q_1,q_2,q_3,0) \in L^2(\mathbb{R}^3)$. Then $f(\mathbf{k})e^{i\mathbf{k}^T\mathbf{q}}$ is a generalized solution to the Schrödinger equation. It is not a true solution because it is not square integrable. But by summing up such solutions we actually solve the initial value problem:

$$\psi(\mathbf{q},t) = \int_{\mathbb{R}^3} e^{-\frac{i\hbar}{m}t\|\mathbf{k}\|^2}e^{i\mathbf{k}^T\mathbf{q}}f(\mathbf{k})d\mathbf{k}.$$

The integral is the inverse Fourier transform of the function $e^{-\frac{i\hbar}{m}t\|\mathbf{k}\|^2}f(\mathbf{k})$ so it is a well defined L^2 function for every L^2 function f.

Notice that we no longer need to solve the case $\lambda = 0$, nor worry about solutions for which the variables are not separated, we do not need to worry even about knowing all the detailes of the spectral theory of the Laplacian. We imposed some friendly conditions on the way, and we solved the problem completely.

Example 4.1.2. Let us consider a free particle confined to a box $0 \leq q_1 \leq a, 0 \leq q_2 \leq b, 0 \leq q_3 \leq c$. To confine the particle to the box, we must consider a potential that is zero in the domain $0 < q_1 < a, 0 < q_2 < b, 0 < q_3 < c$ (the potential well) and is infinite outside of this domain.

Inside the box the Schrödinger equation is

$$i\hbar\frac{\partial\psi}{\partial t} = -\frac{\hbar^2}{2m}\Delta\psi.$$

We can repeat the above considerations and obtain the stationary states

$$\phi(\mathbf{q}) = e^{i\mathbf{k}^T\mathbf{q}}, \quad \mathbf{k} \in \mathbb{R}^3, \|\mathbf{k}\|^2 = \lambda.$$

Now let us examine what happens outside the potential well. There because V is infinite, the state must be zero. Oops! Exponentials cannot be zero. But let us return to the equation

$$-\Delta\psi = \lambda\psi$$

and notice that the solutions are all linear combinations (finite or infinite) of functions of the form $e^{i\mathbf{k}^T\mathbf{q}}, \quad \mathbf{k} \in \mathbb{R}^3$ where $\mathbf{k}$ is subject to the constraint $\|\mathbf{k}\|^2 = \lambda$. And some of these can be zero on the boundary of the well. Using Euler's formula $e^{ix} = \cos x + i \sin x$, we find that states of the form

$$\psi_{n_1,n_2,n_3}(q_1, q_2, q_3) = \sin\frac{n_1\pi q_1}{a}\sin\frac{n_2\pi q_2}{b}\sin\frac{n_3\pi q_3}{c}$$

are zero on the boundary. So the momentum and the energy must have discrete values

$$(k_1, k_2, k_3) = \left(\frac{n_1\pi}{a}, \frac{n_2\pi}{b}, \frac{n_3\pi}{c}\right), \quad E_{n_1,n_2,n_3} = \frac{\hbar^2}{2m}\left[\left(\frac{n_1\pi}{a}\right)^2 + \left(\frac{n_2\pi}{b}\right)^2 + \left(\frac{n_3\pi}{c}\right)^2\right].$$

In contrast with the previous situation ψ_{n_1,n_2,n_3} are L^2 functions (they are zero outside the well). And any solution to the Schrödinger equation is a finite or infinite linear combination of such solutions. Inside the well it is a Fourier series in sines only.

Before concluding this example, let us point out that the least value the energy can have is

$$E_{111} = \frac{\hbar^2}{2m}\left[\left(\frac{\pi}{a}\right)^2 + \left(\frac{\pi}{b}\right)^2 + \left(\frac{\pi}{c}\right)^2\right],$$

which is strictly positive. Moreover, if $a = b = c$ then $E_{211} = E_{121} = E_{122}$, so this level of energy has three linear independent states. In this case the three states are called degenerate because we cannot determine the state of the system from an experiment (the eigenvalue does not determine the state). So the spatial symmetry of the system makes it degenerate.

4.2 The harmonic oscillator

33. For the classical harmonic oscillator without damping, Hooke's law yields the following Hamiltonian

$$H = \frac{p^2}{2m} + \frac{m\omega^2}{2}q^2.$$

Recall that the solutions are sinusoidal with frequency ω. The quantum Hamiltonian is

$$\hat{H} = -\frac{\hbar^2}{2m}\frac{\partial^2}{\partial q^2} + \frac{m\omega^2}{2}M_{q^2} = \frac{1}{2m}\left(-\hbar^2\frac{\partial^2}{\partial q^2} + (m\omega)^2 M_{q^2}\right)$$

We have the Schrödinger equation

$$i\hbar\frac{\partial\psi}{\partial t} = \frac{1}{2m}\left(-\hbar^2\frac{\partial^2\psi}{\partial q^2} + (m\omega)^2 q^2\psi\right)$$

34. Let us solve the Schrödinger equation of the harmonic oscillator. We want to diagonalize the quantized Hamiltonian, namely we want to solve the stationary Schrödinger equation $\hat{H}\psi = E\psi$. With the substitutions $k^2 = \frac{2m}{\hbar^2}E$ and $\lambda = \frac{m\omega}{\hbar}$, we obtain the equation

$$\frac{d^2\psi}{dq^2} + (k^2 - \lambda^2 q^2)\psi = 0.$$

In the theory of differential equations this is known as Weber's differential equation. I have taken the approach to solving this equation from W. Greiner's *Quantum Mechanics'* book. We perform the change of variable $x = \lambda q^2$ to transform this equation into

$$x\frac{d^2\psi}{dx^2} + \frac{1}{2}\frac{d\psi}{dx} + \left(\frac{\kappa}{2} - \frac{1}{4}y\right)\psi = 0$$

where $\kappa = \frac{k^2}{2\lambda} = \frac{E}{\hbar\omega}$. Recall that we are trying to find states that are solutions to this equation, namely we are looking for square integrable functions. This means that for large x the terms that have a factor of x "dominate" and thus the asymptotic behaviour of the solution is dictated by the equation $\frac{d^2\psi}{dx^2} - \frac{1}{4}\psi = 0$. And square integrability tells us that the interesting solution to this equation is $e^{-x/2}$. So we should write

$$\psi(x) = e^{-x/2}\phi(x).$$

Differentiating we obtain

$$\frac{d\psi}{dx} = \left(-\frac{1}{2}\phi(x) + \frac{d\phi}{dx}\right)e^{-x/2}, \quad \frac{d^2\psi}{dx^2} = \left(\frac{1}{4}\phi(x) - \frac{d\phi}{dx} + \frac{d^2\phi}{dx^2}\right)e^{-x/2}.$$

Substituting in the Weber differential equation we obtain

$$x\frac{d^2\phi}{dx^2} + \left(\frac{1}{2} - x\right)\frac{d\phi}{dx} + \left(\frac{\kappa}{2} - \frac{1}{4}\right)\phi = 0.$$

This is a particular case of the Kummer differential equation (also known as the confluent hypergeometric equation):

$$x\frac{d^2\phi}{dx^2} + (c - x)\frac{d\phi}{dx} - a\phi = 0.$$

We have run out of tricks, so we switch to brute force. For this we expand ϕ into a power series around zero. Rewrite the equation as

$$\frac{d^2\phi}{dx^2} + \frac{c-x}{x}\frac{d\phi}{dx} - \frac{a}{x}\phi = 0 \quad \left(\frac{d^2\phi}{dx^2} + \frac{q(x)}{p(x)}\frac{d\phi}{dx} + \frac{r(x)}{p(x)} = 0\right).$$

Note that zero is a singular point, but it is a regular singular point ($\lim_{x\to 0} xq(x)/p(x)$ and $\lim_{x\to 0} x^2 r(x)/p(x)$ exist). So we can apply the Frobenius method, which means that we write

$$\phi(x) = x^r \sum_{m=0}^{\infty} a_m x^m = a_0 x^r + a_1 x^{r+1} + a_2 x^{r+2} + \cdots,$$

and we substitute in the Kummer differential equation. We obtain

$$\begin{aligned} r(r-1)a_0 x^{r-1} + (r+1)ra_1 x^r + (r+2)(r+1)a_2 x^{r+1} + \cdots + cra_0 x^{r-1} - ra_0 x^r + \\ c(r+1)a_1 x^r - (r+1)a_1 x^{r+1} + c(r+2)a_2 x^{r+1} - (r+2)a_2 x^{r+2} + \cdots \\ -aa_0 x^r - aa_1 x^{r+1} - aa_2 x^{r+2} - \cdots = 0. \end{aligned}$$

Setting the coefficient of x^{r-1} equal to zero yields $r(r-1+c) = 0$, so $r = 0$ and $r = 1-c$. For $r = 0$ we obtain the condition

$$\begin{aligned} 2\cdot 1a_2 x + 3\cdot 2a_3 x^2 + 4\cdot 3a_4 x^3 + 5\cdot 4x^4 + \cdots + ca_1 - a_1 x + 2ca_2 x - 2a_2 x^2 + ca_3 x^2 - 3a_3 x^3 + \cdots \\ -aa_0 - aa_1 x - aa_2 x^2 - aa_3 x^3 - \cdots = 0. \end{aligned}$$

We obtain the recursive relations

$$\begin{aligned} a_1 &= \frac{a}{c}a_0, \\ a_2 &= 2\frac{a+1}{2(c+1)}a_1 = \frac{a(a+1)}{c(c+1)}\cdot\frac{1}{2}, ..., \\ a_{m+1} &= \frac{m+a}{(m+1)(c+m)}a_m = \frac{a(a+1)\cdots(a+m)}{c(c+1)\cdots(c+m)}\cdot\frac{1}{m!}, ... \end{aligned}$$

We can set $a_0 = 0$. Then we define the Pochammer symbol $(a)_m = a(a+1)\cdots(a+m)$ and consider the confluent hypergeometric series

$${}_1F_1(a;c;x) = \sum_{m=0}^{\infty} \frac{(a)_m}{(c)_m}\frac{x^m}{m!}.$$

Then the solution to the Kummer differential equation is a constant multiple of this function. A similar computation shows that the solution correspoding to $r = 1-c$ is a constant multiple of $x_1^{1-c}F_1(a-c+1;2-c;x)$.

35. So the stationary state of the harmonic oscillator is of the form

$$\psi(x) = e^{-x/2}\left[A\ {}_1F_1(a;\tfrac{1}{2};x) + Bx^{1/2}\ {}_1F_1(a+\tfrac{1}{2};\tfrac{3}{2};x)\right], \text{ where } a = -\left(\frac{\kappa}{2} - \frac{1}{4}\right).$$

We want this to be a square integrable function. It is known that as $|x| \to \infty$,

$$ {}_1F_1(a;c;x) \sim \frac{\Gamma(c)}{\Gamma(a)} e^x x^{a-c}, \text{ where } \Gamma(z) = \int_0^\infty t^{z-1} e^{tz} dt. $$

We should point out that this formula defines Γ for positive real part of z, it is then extended analytically to a meromorphic function on the entire plane, and it has simple poles at the negative integers. For ψ to be square integrable, we should have $\Gamma(a) = 0$ and $B = 0$ or $\Gamma(a + \frac{1}{2}) = 0$ and $A = 0$. In each of these cases the power series is truncated. Thus we have two cases:

Case 1. $a = -n$ and $B = 0$ with $n = 0, -1, -2, ...$ in which case

$$ \frac{\kappa}{2} - \frac{1}{4} = n, $$

and we have the eigenfunction

$$ \psi_n(x) = e^{-(\lambda/2)q^2} \, {}_1F_1(-n; \tfrac{1}{2}; \lambda q^2), $$

corresponding to the energy level

$$ E_n = \hbar\omega(2n + \tfrac{1}{2}). $$

Case 2. $a = -n - \frac{1}{2}$ and $A = 0$, with

$$ \frac{\kappa}{2} - \frac{1}{4} = n + \frac{1}{2}, $$

having the eigenfunction

$$ \psi_n(x) = q e^{-(\lambda/2)q^2} \, {}_1F_1(-n; \tfrac{3}{2}; \lambda q^2), $$

corresponding to the energy level

$$ E_n = \hbar\omega(2n + 1 + \tfrac{1}{2}). $$

So the energy can only have the following values

$$ E_n = (n + \tfrac{1}{2})\hbar\omega, \quad n = 0, 1, 2, \cdots . $$

The energy is *quantized*! A peculiar fact is that the energy of the ground state is $\frac{1}{2}\hbar\omega$, which is positive?!

36. Here is an approach to solve the Schrödinger equation without actually solving the differential equation, using Dirac's "ladder operator" method. This trick is inspired by the representation theory of Lie groups, we will encounter it again in the next chapter. Based on the formula

$$ (a + ib)(a - ib) = a^2 + b^2 $$

we define the *annihilation operator* (a.k.a. *lowering operator*)

$$ a = \sqrt{\frac{m\omega}{2\hbar}} \left(\hat{q} + \frac{i}{m\omega} \hat{p} \right) = \sqrt{\frac{m\omega}{2\hbar}} \left(M_q + \frac{\hbar}{m\omega} \frac{\partial}{\partial q} \right) $$

and the *creation operator* (a.k.a. *raising operator*)

$$a^* = \sqrt{\frac{m\omega}{2\hbar}}\left(\hat{q} - \frac{i}{m\omega}\hat{p}\right) = \sqrt{\frac{m\omega}{2\hbar}}\left(M_q - \frac{\hbar}{m\omega}\frac{\partial}{\partial q}\right).$$

The $*$ is not accidental, one operator is the adjoint of the other. Moreover, using the canonical commutation relation we compute

$$a^*a = \frac{1}{\hbar\omega}\frac{1}{2m}\left(\hat{p}^2 + (m\omega)^2\hat{q}^2\right) - \frac{1}{2}I.$$

Therefore

$$\hat{H} = \hbar\omega\left(a^*a + \frac{1}{2}I\right).$$

Let us ignore the $\frac{1}{2}I$, known as the quantum correction, and proceed with understanding the spectral properties of a^*a. Right now the elimination of the term $\frac{1}{2}I$ seems ad hoc, but we point out that there is a rigorous procedure called the metaplectic correction for removing this term and adjusting the spectrum of the Hamiltonian to its correct values. Note that $a^*a \geq 0$, so its spectrum lies in $[0, \infty)$. We work in the assumption that a^*a has at least one eigenvalue (which is not necessarily true, but we assume it to be true for our considerations). First, we compute

$$[a, a^*] = I,$$

and then

$$[a, a^*a] = -a, \quad [a^*, a^*a] = a^*.$$

Lemma 4.2.1. Suppose ψ is an eigenvector of a^*a with eigenvalue λ. Then

$$a^*a(a\psi) = (\lambda - 1)(a\psi) \text{ and } a^*a(a^*\psi) = (\lambda + 1)(a^*\psi).$$

Thus $a\psi$ is either equal to zero, or it is an eigenvector of a^*a with eigenvalue $\lambda - 1$, and $a^*\psi$ is either equal to zero, or it is an eigenvector of a^*a with eigenvalue $\lambda + 1$.

Proof. We have

$$a^*a(a\psi) = (a(a^*a) - a)\psi = (\lambda - 1)a\psi$$

and the same for $a^*\psi$. □

As a corollary, since a^*a has only non-negative eigenvalues, the "lowering" of eigenvalues must end somewhere. So for an eigenvalue ψ there is $N \geq 0$ such that $a^N\psi \neq 0$ but $a^{N+1}\psi = 0$. Let

$$\psi_0 = a^N\psi/\|a^N\psi\|, \quad \psi_n = (a^*)^n\psi_0, \quad n \geq 0.$$

Theorem 4.2.1. The vectors ψ_n, $n \geq 0$ satisfy the following relations

$$\begin{aligned} a^*\psi_n &= \psi_{n+1} \\ a^*a\psi_n &= n\psi_n \\ \langle \psi_n, \psi_m \rangle &= n!\delta_{m,n} \\ a\psi_n &= n\psi_{n-1}. \end{aligned}$$

Proof. The only things to check are the third and fourth properties. For different indices the eigenvectors correspond to different eigenvalues, so they are orthogonal (a property of Hermitian operators). For equal indices have (inductively)

$$\begin{aligned} \langle \psi_n, \psi_n \rangle = \langle a^*\psi_{n-1}, a^*\psi_{n-1} \rangle = \langle \psi_{n-1}, aa^*\psi_{n-1} \rangle = \langle \psi_{n-1}, (a^*a + I)\psi_{n-1} \rangle \\ = ((n-1)+1) \langle \psi_{n-1}, \psi_{n-1} \rangle = n \cdot (n-1)! = n!. \end{aligned}$$

Also $a\psi_{n+1} = aa^*\psi_n = (a^*a + I)\psi_n = (n+1)\psi_n$. □

So if we have exactly one solution for the equation $a\psi_0 = 0$, then the vectors $\frac{1}{\sqrt{n!}}\psi_n$ form an orthonormal basis of the Hilbert space and they diagonalize $\hat{H}$ with diagonal entries $n+\frac{1}{2}$, $n \geq 0$. In this case we solved the time-independent Schrödinger equation

$$\hat{H}\psi = E\psi,$$

and moreover, we can make sense of the solution to the Schrödinger equation

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

since we know how to exponentiate diagonal operators.

37. Now we turn to analytical methods. The equation $a\psi_0 = 0$ reads

$$\frac{\hbar}{m\omega}\frac{\partial \psi_0}{\partial q} = -q\psi(q),$$

with solution

$$\psi_0(q) = Ce^{-\frac{m\omega}{2\hbar}q^2}.$$

Normalize this to have norm 1 by setting $C = \sqrt{\frac{\pi m\omega}{\hbar}}$. Thus

$$\psi_0(q) = \sqrt{\frac{\pi m\omega}{\hbar}}e^{-\frac{m\omega}{2\hbar}q^2}.$$

This shows that both assumptions in the ladder method work: that the Hamiltonian operator has eigenvalues, and that 0 is an eigenvalue with multiplicity one. The state ψ_0 is called the *ground state*, and the states ψ_n, $n > 0$, are called the *excited states*. For the *Hermite polynomials* defined by

$$H_n(q) = (-1)^n e^{q^2}\frac{d^n}{dq^n}e^{-q^2}, \quad n \geq 0$$

we have

$$\psi_n(q) = \sqrt[4]{\frac{m\omega}{\pi\hbar}}\frac{1}{2^n}e^{-\frac{m\omega}{2\hbar}q^2}H_n\left(\sqrt{\frac{m\omega}{2\hbar}}q\right).$$

Note that since the Hermite polynomials form an orthonormal basis of $L^2(\mathbb{R}^2, e^{-q^2/2}dq)$, we deduce that $\frac{1}{\sqrt{n!}}\psi_n$ is an orthonormal basis of $L^2(\mathbb{R}, dq)$, so indeed $\hat{H}$ is diagonalizable and the diagonal entries have one-dimensional eigenspaces. We arrived at the same Hermite polynomials that we encountered when working with hypergeometric functions.

There is another perspective in which we can think about this problem. The operators $a, a^*, a*a, I$ form a Lie algebra with the Lie bracket defined (formally) by $[a, a^*] = I, [a, a^*a] = -a, [a^*, a^*a] = a^*, [a, I] = [a^*, I] = [a^*a, I]$. We are trying to construct a representation of these operators on a Hilbert space. There is a way to construct this representation formally, through the process of "creation", but we can also construct this representation concretely, on $L^2(\mathbb{R})$ with Hermite polynomials. This is a pattern in representation theory and it is how representation theory is linked to differential equations and special functions.

38. We now describe a *holomorphic model* that produces the creation and annihilation operators. We consider the space of holomorphic functions for which

$$\frac{i}{2\pi}\int_{\mathbb{C}}|f(z)|^2e^{-|z|^2}dz\wedge d\bar{z} < \infty.$$

Note that the quantity on the left is positive. Endow this space with the inner product

$$\langle f, g\rangle = \frac{i}{2\pi}\int_{\mathbb{C}}f(z)\overline{g(z)}e^{-|z|^2}dz\wedge d\bar{z} < \infty.$$

so that that the previous expression is the square of the norm. Then this is a Hilbert space called the *Segal-Bargmann space*. An orthonormal basis is

$$f_n(z) = \frac{z^n}{\sqrt{n!}}, \quad n = 0, 1, \ldots.$$

We set

$$a^* = M_z, \quad a = \frac{d}{dz}.$$

These satisfy the same condition as the creation and annihilation operators:

$$[a, a^*] = I.$$

The inner product is chosen so that a and a^* are one the adjoint of the other! Note that we can reverse engineer $\hat{q}$ and $\hat{p}$ as

$$\hat{q} = \sqrt{\hbar}\frac{a + a^*}{\sqrt{2m\omega}} \quad \hat{p} = i\sqrt{2\hbar m\omega}(a^* - a),$$

obtaining another representation of the position and momentum operators, which of course is unitary equivalent to the Schrödinger representation because of the Stone-von Neumann Theorem.

It is important to point out that this construction has been adapted in 1932 by V.A. Fock to produce the Fock space as a model for quantum field theory, where the creation and annihilation operators actually create and annihilate particles. In loose terms the Fock space is the Segal-Bargmann space over $\mathbb{C}^\infty$ (seen as a limiting case of the Segal-Bargmann spaces over $\mathbb{C}^n$).

4.3 The WKB method

39. This method is named after Wentzel, Krammers, and Brillouin. The idea is to construct approximate solutions to the Schrödinger equation by trying to find solutions of the form

$$\psi(t, \mathbf{q}) = e^{(i/\hbar)S(t,\mathbf{q})}.$$

Substituting in the Schrödinger equation

$$i\hbar\frac{\partial\psi}{\partial t} = \left(-\frac{\hbar^2}{2m}\Delta + V(\mathbf{q})\right)\psi$$

we obtain

$$-\frac{\partial S}{\partial t} = -\frac{\hbar^2}{2m}\left(\frac{i}{\hbar}\Delta_{\mathbf{q}}S - \hbar^{-2}\|\nabla_{\mathbf{q}}S\|^2\right) + V.$$

We can view this as

$$-\frac{\partial S}{\partial t} = \frac{1}{2}\|\nabla_{\mathbf{q}}S\|^2 + V + O(\hbar).$$

Ignoring $O(\hbar)$ we obtain

$$-\frac{\partial S}{\partial t}(\mathbf{q}, t) = H(\mathbf{q}, \nabla_{\mathbf{q}}S).$$

This is the Hamilton-Jacobi equation from classical mechanics. Note that when S is the action ($S = \int L dt$), then this yields an equivalent formulation of classical mechanics.

40. Let us now turn to the time-independent Schödinger equation, and in order to simplify computations we work in the 1-dimensional case:

$$-\frac{\hbar^2}{2m}\frac{d^2\psi}{dq^2} + V(q)\psi = E\psi.$$

The substitution $\psi = e^{(i/\hbar)S(q)}$ yields the differential equation

$$\frac{1}{2m}\left(\frac{dS}{dq}\right)^2 + V(q) = E + \frac{i\hbar}{2m}\frac{d^2S}{dq^2}.$$

As above, we ignore the $\hbar$ term and turn this into the (time-independent version of the) Hamilton-Jacobi equation

$$H(q, S'(q)) = \frac{(S'(q))^2}{2m} + V(x) = E.$$

This gives

$$S'(x) = \pm\sqrt{2m(E - V(x))}.$$

Let us turn to the classical phase space $\mathbb{R}^2 = T^*\mathbb{R}$. Then the differential $dS = S'dq$ can be viewed as a map $dS : \mathbb{R} \to T^*\mathbb{R}$, $q \mapsto S'(q)dq$. Then S satisfies the Hamilton-Jacobi equation if and only if the image of dS lies in the level manifold $H^{-1}(E)$. This establishes a fundamental relationship between classical and quantum mechanics: "When the image of dS lies in a level manifold of the classical Hamiltonian, then $e^{(i/\hbar)S}$ can be taken as a first-order approximate solution to the Hamilton equation."

Here is another way to look at this. Using the expansion

$$S(q) = S(q, \hbar) = \sum_{n=0}^{\infty} (-i\hbar)^n S_n(q),$$

we obtain

$$\frac{1}{2m}(S_0')^2 = E - V(q), \quad S_0'S_1' = -\frac{1}{2}S_0''.$$

Let $p(q) = \sqrt{2m(E - V(x)}$ be the classical momentum of a particle with potential V and energy E. The solution to the first differential equation is

$$S_0 = \pm \int p(q)dq,$$

and from the second we get $S_1 = \frac{1}{2}\ln p$. So

$$\psi(q) = \frac{1}{\sqrt{|p(x)|}} \left(c_1 e^{(i/\hbar)\int p(q)dq} + c_2 e^{-(i/\hbar)\int p(q)dq} \right) (1 + O(\hbar))$$

41. We can go to a more general *Ansatz*, namely that

$$\psi(q) = A(q)e^{\pm(i/\hbar)S(q)},$$

where $S'(q) = p(q) = \sqrt{2m(E - V(q))}$. The amplitude function $A(q)$ is chosen to be independent of $\hbar$.

Proposition 4.3.1. For any two numbers E_1 and E_2 with $E_1 > \inf V(q)$ there exists a constant C and a nonzero compactly supported smooth function $A(q)$ with the following property: For every $E \in [E_1, E_2]$ the support of A is contained in the classically allowed region at energy E and the function ψ given by

$$\psi(q) = A(q) \exp\left(\pm \frac{i}{\hbar} \int p(q)dq \right)$$

satisfies

$$\|\hat{H}\psi - E\psi\| \leq C\hbar\|\psi\|.$$

Here by classically allowed region we mean the region where a particle with energy E can lie in the configuration space.

Proof. We will actually choose A to by any function with support in the classically allowed region for E_1 (which is included in the classically allowed region for E). We compute

$$\hat{H}\psi - E\psi = -\frac{\hbar^2}{2m}\left(A''(q) \pm 2\frac{i}{\hbar}A'(q)p(q) \pm \frac{i}{\hbar}p'(q)A(q)\right)e^{\pm(i/\hbar)\int p(q)dq}.$$

Thus

$$\|\hat{H}\psi - E\psi\| \leq \frac{\hbar^2}{2m}\|A''\| + \frac{\hbar}{2m}\|2A'p + Ap'\|.$$

Notice that the right-hand side is of order $\hbar\|\psi\|$ and $\|2A'p + Ap'\|$. □

We can even improve on the error by forcing $2A'p = -p'A$ so as to obtain

$$\hat{H}\psi - E\psi = -\frac{\hbar^2}{2m}A''\psi.$$

Chapter 5

The Hydrogen Atom. The Spin

This chapter follows closely the book of Faddeev and Yakubovskii.

5.1 The classical Kepler problem

42. We start with a short detour through classical mechanics. The classical Kepler problem asks to find the trajectories of planets orbiting the sun. The sun exterts the force

$$\mathbf{F} = -k\frac{\mathbf{q}}{\|\mathbf{q}\|^3},$$

which is given by the potential

$$V(\mathbf{q}) = -\frac{k}{\|\mathbf{q}\|}.$$

If m is the mass of the planet and M is the mass of the sun, then $k = GMm$.

Because the potential is invariant under rotations, the angular momentum $\mathbf{J} = \mathbf{q} \times \mathbf{p}$ is conserved (and we require it to be nonzero to avoid collision with the sun). This is a consequence of Noether's Theorem in mechanics: Every differentiable symmetry of the action of a physical system has a conserved observable. Note that the planet moves in a plane orthogonal to $\mathbf{J}$. We can compute explicitly

$$\mathbf{J} = \begin{vmatrix} \mathbf{i} & \mathbf{j} & \mathbf{k} \\ q_1 & q_2 & q_3 \\ p_1 & p_2 & p_3 \end{vmatrix} = (q_2p_3 - q_3p_2)\mathbf{i} + (q_3p_1 - q_1p_3)\mathbf{j} + (q_1p_2 - q_2p_1)\mathbf{k},$$

and the conservation of the momentum follows by an easy computation in which we use Hamilton's equations.

Definition. The *Runge-Lenz vector* is the vector-valued function on $\mathbb{R}^3\backslash\{0\} \times \mathbb{R}^3$ given by

$$\mathbf{A}(\mathbf{q}, \mathbf{p}) = \frac{1}{mk}\mathbf{p} \times \mathbf{J} - \frac{\mathbf{q}}{\|\mathbf{q}\|}.$$

Proposition 5.1.1. The Runge-Lenz vector is conserved.

Proof. Using Hamilton's equations: $\partial\mathbf{p}/\partial t = -\partial H/\partial\mathbf{q} = -\partial V/\partial\mathbf{q} = \mathbf{F}$, $\partial\mathbf{q}/\partial t = \mathbf{p}/m$, and the fact that $\mathbf{J}$ is constant, we compute

$$\begin{aligned}\dot{A}(t) &= \frac{1}{mk}\mathbf{F}\times\mathbf{J} - \frac{1}{\|\mathbf{q}\|}\frac{\mathbf{p}}{m} + \frac{\mathbf{q}}{\|\mathbf{q}\|^2}\cdot\sum_{j=1}^{3}\frac{\partial\|\mathbf{q}\|}{\partial q_j}\frac{dq_j}{dt}\\ &= -\frac{1}{m}\frac{1}{\|\mathbf{q}\|^3}\mathbf{q}\times(\mathbf{q}\times\mathbf{p}) - \frac{1}{\|\mathbf{q}\|}\frac{\mathbf{p}}{m} + \frac{\mathbf{q}}{\|\mathbf{q}\|^2}\cdot\sum_{j=1}^{3}\frac{\partial\|\mathbf{q}\|}{\partial q_j}\frac{dq_j}{dt}\\ &= \frac{1}{m}\left(-\frac{1}{\|\mathbf{q}\|^3}\mathbf{q}(\mathbf{q}\cdot\mathbf{p}) + \frac{1}{\|\mathbf{q}\|^3}\mathbf{p}(\mathbf{q}\cdot\mathbf{q}) - \frac{\mathbf{p}}{\|\mathbf{q}\|} + \frac{\mathbf{q}(\mathbf{q}\cdot\mathbf{p})}{\|\mathbf{q}\|^3}\right) = 0.\end{aligned}$$

Here we used $\mathbf{a}\times(\mathbf{b}\times\mathbf{c}) = \mathbf{b}(\mathbf{a}\cdot\mathbf{c}) - \mathbf{c}(\mathbf{a}\cdot\mathbf{b})$. □

43. Using the Runge-Lenz vector as an auxiliary tool we will now prove that the trajectory is either an *ellipse*, a *parabola*, or a *hyperbola*. Note that $\mathbf{A}$ lies in the plane of motion. Choose cartesian coordinates in this plane of motion such that $\mathbf{A}$ defines the x-axis, then switch to *polar coordinates*.

Theorem 5.1.1. In polar coordinates, the trajectory $(r(t), \theta(t))$ of the planet is given by

$$r(t) = \frac{\|\mathbf{J}\|^2}{mk}\frac{1}{1+\|\mathbf{A}\|\cos\theta(t)}.$$

Proof. Let $\mathbf{r}(t)$ be the trajectory. We substitute $\mathbf{r}(t)$ in the place of $\mathbf{q}$ in the Runge-Lenz vector and take the dot product with $\mathbf{r}$.

$$\mathbf{A}\cdot\mathbf{r} = \|A\|r\cos\theta = \frac{1}{mk}\mathbf{r}\cdot(\mathbf{p}\times\mathbf{J}) - r.$$

But

$$\mathbf{r}\cdot(\mathbf{p}\times\mathbf{J}) = (\mathbf{r}\times\mathbf{p})\cdot\mathbf{J} = \mathbf{J}\cdot\mathbf{J} = \|\mathbf{J}\|^2.$$

Now solve for r. □

Note that the equation from the statement is the equation of a conic

$$r = \frac{ep}{1+e\cos\theta}$$

where e is the excentricity and $|p|$ is the distance between the focus and the directrix.

5.2 Angular momentum in quantum mechanics

44. The angular momentum $\mathbf{J}$ is a vector valued function, so it is not an observable in the way we define observables. But its coordinates themselves are observables. Its coordinates

in classical mechanics are $q_2p_3 - q_3p_2$, $q_3p_1 - q_1p_3$, and $q_1p_2 - q_2p_1$. We quantize these as

$$L_1 = \hat{q}_2\hat{p}_3 - \hat{q}_3\hat{p}_2 = -i\hbar\left(M_{q_2}\frac{\partial}{\partial q_3} - M_{q_3}\frac{\partial}{\partial q_2}\right) = i\hbar\left(M_{q_3}\frac{\partial}{\partial q_2} - M_{q_2}\frac{\partial}{\partial q_3}\right)$$
$$L_2 = \hat{q}_3\hat{p}_1 - \hat{q}_1\hat{p}_3 = -i\hbar\left(M_{q_3}\frac{\partial}{\partial q_1} - M_{q_1}\frac{\partial}{\partial q_3}\right) = i\hbar\left(M_{q_1}\frac{\partial}{\partial q_3} - M_{q_3}\frac{\partial}{\partial q_1}\right)$$
$$L_3 = \hat{q}_1\hat{p}_2 - \hat{q}_2\hat{p}_1 = -i\hbar\left(M_{q_1}\frac{\partial}{\partial q_2} - M_{q_2}\frac{\partial}{\partial q_1}\right) = i\hbar\left(M_{q_2}\frac{\partial}{\partial q_1} - M_{q_1}\frac{\partial}{\partial q_2}\right).$$

Here it is important to notice that for $j \neq k$, $\hat{q}_j$ commutes with $\hat{p}_k$, so we do get self-adjoint operators. It is helpful to add to the bunch the operator

$$L^2 = L_1^2 + L_2^2 + L_3^2.$$

The following commutation relations are satisfied

$$[L_1, L_2] = i\hbar L_3, \quad [L_2, L_3] = i\hbar L_1, \quad [L_3, L_1] = i\hbar L_2, \quad [L_1, L^2] = [L_2, L^2] = [L_3, L^2] = 0.$$

The operator L^2 is positive, so $L = \sqrt{L^2}$ can be defined, in case it is needed.

Lemma 5.2.1. If in spherical coordinates ψ depends on $r = \sqrt{q_1^2 + q_2^2 + q_3^2}$ only, then $L_j\psi = 0$, $j = 1, 2, 3$.

Proof. We only check L_3, the others are exactly the same.

$$L_3\psi(r) = L_3\psi(q_1^2 + q_2^2 + q_3^2) = i\hbar\left(q_2\frac{\partial\psi}{\partial q_1} - q_1\frac{\partial\psi}{\partial q_2}\right) = i\hbar\psi'(q_1^2 + q_2^2 + q_3^2)(2q_2q_1 - 2q_1q_2) = 0.$$

□

45. Since we will need the Laplacian, which defines the quantum kinetic energy in the Schrödinger equation, we prove the following result.

Lemma 5.2.2. In spherical coordinates the Laplacian is given by the formula

$$\Delta = \frac{1}{r^2}\frac{\partial}{\partial r}\left(r^2\frac{\partial}{\partial r}\right) - \frac{1}{\hbar^2 r^2}L^2.$$

Proof. We compute

$$-\frac{1}{\hbar^2}L^2 = \left(M_{q_3}\frac{\partial}{\partial q_2} - M_{q_2}\frac{\partial}{\partial q_3}\right)^2 + \left(M_{q_1}\frac{\partial}{\partial q_3} - M_{q_3}\frac{\partial}{\partial q_1}\right)^2 + \left(M_{q_2}\frac{\partial}{\partial q_1} - M_{q_1}\frac{\partial}{\partial q_2}\right)^2$$
$$= \sum_j q_j^2\Delta - \sum_j q_j^2\frac{\partial^2}{\partial q_j^2} - 2\sum_{j<k} q_jq_k\frac{\partial^2}{\partial q_j\partial q_k} - 2\sum_j q_j\frac{\partial}{\partial q_j}.$$

The chain rule gives

$$\frac{\partial}{\partial q_j} = \frac{\partial}{\partial r}\frac{\partial r}{\partial q_j} = \frac{\partial}{\partial r}\frac{q_j}{r}.$$

Multiply this by q_j and add over $j = 1, 2, 3$ to obtain

$$r\frac{\partial}{\partial r} = \sum_j q_j \frac{\partial}{\partial q_j}.$$

Square to obtain

$$\left(r\frac{\partial}{\partial r}\right)^2 = \sum_j q_j^2 \frac{\partial^2}{\partial q_j^2} + 2\sum_{j<k} q_j q_k \frac{\partial^2}{\partial q_j \partial q_k} + \sum_j q_j \frac{\partial}{\partial q_j}.$$

Combine this with the previous computation to obtain the conclusion. □

5.3 The rotation group $SO(3)$ and its Lie algebra

46. Let us recall a few facts about the rotation group $SO(3)$ and its Lie algebra $so(3)$. The Lie group consists of orthogonal matrices of determinant 1. The Lie algebra $so(3)$ consists of the skew symmetric 3×3 matrices. It is generated by the matrices

$$A_1 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & -1 \\ 0 & 1 & 0 \end{pmatrix}, \quad A_2 = \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ -1 & 0 & 0 \end{pmatrix}, \quad A_3 = \begin{pmatrix} 0 & -1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}.$$

These are the infinitesimal rotations about the three coordinate axes. They satisfy

$$[A_j, A_{j+1}] = A_{j+2},$$

where indices are taken modulo 3. Every other skew symmetric matrix is a linear combination of these. The rotations by angle α about the coordinate axes are

$$e^{\alpha A_1}, \quad e^{\alpha A_2}, \quad e^{\alpha A_3}.$$

In general, if $\mathbf{n} = (n_1, n_2, n_3)$ is a unit vector, then the rotation by angle α about the axis of this vector is

$$g(\alpha \mathbf{n}) = e^{\alpha(n_1 A_1 + n_2 A_2 + n_3 A_3)}$$

Proposition 5.3.1. The Lie algebra $so(3)$ is isomorphic to the Lie subalgebra of quantum observables generated by L_1, L_2, L_3.

Proof. Recall that the Lie bracket of operators is $\frac{1}{i\hbar}[\cdot, \cdot]$. With this in mind, the isomorphism is $A_j \mapsto L_j$, $j = 1, 2, 3$. □

Theorem 5.3.1. The representation of $SO(3)$ on $L^2(\mathbb{R}^3)$ defined by

$$W(g)\psi(\mathbf{q}) = \psi(g^{-1}\mathbf{q})$$

satisfies

$$W(g(\alpha \mathbf{n})) = e^{-(i/\hbar)\alpha(n_1 L_1 + n_2 L_2 + n_3 L_3)}.$$

Proof. Let g_α be the rotation about $\mathbf{n}$ by α, so that g_α^{-1} is the rotation about $\mathbf{n}$ by $-\alpha$. Then

$$g_\alpha^{-1}\mathbf{q}\times\mathbf{q} = \|\mathbf{q}\|^2 \sin\alpha\ \mathbf{n}.$$

Multiply on the left by $\mathbf{q}$ and apply the bac-cab formula to obtain

$$-\|\mathbf{q}\|^2\sin\alpha\ \mathbf{q}\times\mathbf{n} = \mathbf{q}\times(g_\alpha^{-1}\mathbf{q}\times\mathbf{q}) = g_\alpha^{-1}\mathbf{q}(\mathbf{q}\cdot\mathbf{q}) - \mathbf{q}(g_\alpha^{-1}\mathbf{q}\cdot\mathbf{q})$$
$$= \|\mathbf{q}\|^2 g_\alpha^{-1}\mathbf{q} - \|\mathbf{q}\|^2\cos\alpha\ \mathbf{q}.$$

Hence

$$g_\alpha^{-1}\mathbf{q} = \sin\alpha\ \mathbf{q}\times\mathbf{n} + \cos\alpha\ \mathbf{q} = [(q_2n_3 - q_3n_2)\sin\alpha + q_1\cos\alpha]\mathbf{i}$$
$$+[(q_3n_1 - q_1n_3)\sin\alpha + q_2\cos\alpha]\mathbf{j} + [(q_1n_2 - q_2n_1)\sin\alpha + q_3\cos\alpha]\mathbf{k}.$$

Let us check this for the case where $\mathbf{n} = \mathbf{k}$ is unit vector defining the z-axis. In this case we have to prove that

$$W(g(\alpha\mathbf{k}))\psi(\mathbf{q}) = \psi(q_1\cos\alpha + q_2\sin\alpha, -q_1\sin\alpha + q_2\cos\alpha, q_3) = e^{-(i/\hbar)\alpha L_3}\psi(\mathbf{q}).$$

In other words

Consider the function

$$\Psi_{\mathbf{q}}(\alpha) = e^{-(i/\hbar)\alpha L_3}\psi(\mathbf{q}).$$

We have

$$\frac{d\Psi_{\mathbf{q}}(\alpha)}{d\alpha} = -(i/\hbar)L_3\psi_{\mathbf{q}}(\alpha) = \left(q_2\frac{\partial}{\partial q_1} - q_1\frac{\partial}{\partial q_2}\right)\Psi_{\mathbf{q}}(\alpha),$$

and $\Psi_{\mathbf{q}}(0) = \psi(\mathbf{q})$. It is not hard to check that $\psi(g(\alpha\mathbf{k})\mathbf{q})$ satisfies the same initial value problem, so they are equal.

The same is true for $\mathbf{n} = \mathbf{i}$ and $\mathbf{n} = \mathbf{j}$, and because the Lie algebra determined by the L_i's is isomorphic to the Lie algebra generated by A_i's, the equality holds in general (for example because of the Baker-Campbell-Hausdorff formula). The theorem is proved. □

47. Let us turn to spherical coordinates:

$$x = r\sin\phi\cos\theta, \quad y = r\sin\phi\sin\theta, \quad z = r\cos\phi.$$

The convention follows the Texas Tech calculus text book. Strangely there are two notational conventions, exchanging θ and ϕ. So be very careful!

We introduce a normalized version of the angular momentum operators, the so called *dimensionless momentum operators*, by

$$\tilde{L}_j = \frac{1}{\hbar}L_j.$$

Let $\tilde{L}^2 = \tilde{L}_1^2 + \tilde{L}_2^2 + \tilde{L}_3^2$. We focus on just two operators: $\tilde{L}^2$, which is part of the Laplacian and will be used for solving the Schrödinger equation of the hydrogen atom, and $\tilde{L}_3^2$, which is the angular momentum corresponding to the z-axis (which can be any line in the space by choosing the coordinates appropriately). The operators $\tilde{L}_1^2$ and $\tilde{L}_2^2$ are more like auxiliaries.

As Lemma 5.2.1 showed, the angular momentum operators depend only on the angle variables and not on r. Thus it is wise to consider the decomposition

$$L^2(\mathbb{R}^3) = L^2((0,\infty)) \otimes L^2(S^2)$$

where $L^2(0,\infty)$ is endowed with the integration measure $r^2 dr$ and $L^2(S^2)$ is endowed with the integration measure $\sin\phi d\phi d\theta$.

Switching to spherical coordinates, we have

$$\begin{aligned}
\tilde{L}_1 &= i\left(\sin\theta\frac{\partial}{\partial\phi} + \cot\phi\cos\theta\frac{\partial}{\partial\theta}\right)\\
\tilde{L}_2 &= -i\left(\cos\theta\frac{\partial}{\partial\phi} - \cot\phi\sin\theta\frac{\partial}{\partial\theta}\right)\\
\tilde{L}_3 &= -i\frac{\partial}{\partial\theta}.
\end{aligned}$$

We change our two auxiliary operators by replacing them by

$$L_\pm = \tilde{L}_1 \pm i\tilde{L}_2 = e^{\pm i\theta}\left(\pm\frac{\partial}{\partial\phi} + i\cot\phi\frac{\partial}{\partial\theta}\right).$$

Note that a representation of the Lie algebra generated by $\tilde{L}_1, \tilde{L}_2, \tilde{L}_3, \tilde{L}^2$ is the same a representation for the Lie algebra generated by $L_-, L_+, \tilde{L}_3, \tilde{L}^2$.

Let us understand what we want: to diagonalize $\tilde{L}_3$ and $\tilde{L}^2$ as operators acting on $L^2(S^2)$. This will yield, by the Spectral Mapping Theorem, the complete picture of the spectral theory of the operator $e^{i\alpha L_3}$, and so we will have the complete understanding of the spectral theory of the unitary operators that arise from the representation of $SO(3)$ on the space of states from Theorem 5.3.1.

First, note that if $\tilde{L}_3$ and $\tilde{L}^2$ are diagonalizable, then because they commute, they are simultaneously diagonalizable. We proceed as before, assuming that they are indeed diagonalizable, and then by understanding the properties of eigenvectors we compute them explicitly and thus prove that the operators are indeed diagonalizable. Let therefore $Y(\theta,\phi)$ be a common eigenvalue of $\tilde{L}_3$ and $\tilde{L}^2$:

$$\tilde{L}^2 Y = \lambda Y, \quad \tilde{L}_3 Y = mY.$$

Note that m is necessarily an integer, since Y is periodic in θ with period 2π. Also, note that $\tilde{L}^2 = \tilde{L}_1^2 + \tilde{L}_2^2 + \tilde{L}_3^2$ implies $\lambda \geq m^2$. Denote by $E_{\lambda,m}$ the common eigenspace of $\tilde{L}^2$ and $\tilde{L}_3$ with eigenvalues λ, m. We have

$$[\tilde{L}_3, L_+] = L_+, \quad [\tilde{L}_3, L_-] = -L_-, \quad [L_+, L_-] = 2\tilde{L}_3.$$

Consequently

$$L_+ : E_{\lambda,m} \to E_{\lambda,m+1}, \quad L_- : E_{\lambda,m} \to E_{\lambda,m-1},$$

so these are raising and lowering operators in the sense of representation theory. Because $\lambda \geq m^2$, there must be a moment when $L_+ = 0$, that is $L_+^{k+1} E_{\lambda,m} = 0$ where k is chosen to be the smallest with this property. So $L_+^{k+1} Y = 0$ and $L_+^k E_{\lambda,m} \subset E_{\lambda,m+k}$. Then since $L_- L_+ = L^2 - \tilde{L}_3^2 - \tilde{L}_3$,

$$0 = L_- L_+ L_+^k Y = (\lambda - (m+k)^2 - (m+k)) L_+^k Y = (\lambda - (m+k)(m+k+1)) L_+^k Y.$$

So $\lambda = \ell(\ell+1)$, where $\ell = m + k$. Let $Y_{\ell,\ell} \in E_{\ell(\ell+1),\ell}$. The equalities $\tilde{L}_3 Y_{\ell,\ell} = \ell Y_{\ell,\ell}$ and $L_+ Y_{\ell,\ell} = 0$ mean

$$\begin{aligned} -i \frac{\partial Y_{\ell,\ell}}{\partial \theta} &= \ell Y_{\ell,\ell} \\ \frac{\partial Y_{\ell,\ell}}{\partial \phi} + i \cot\phi \frac{\partial Y_{\ell,\ell}}{\partial \theta} &= 0. \end{aligned}$$

From the first equation we see that

$$Y_{\ell,\ell}(\theta, \phi) = e^{i\ell\theta} F_{\ell,\ell}(\phi).$$

Now substitute in the second equation to obtain

$$\frac{\partial F_{\ell,\ell}(\phi)}{\partial \phi} = l \cot\phi F_{\ell,\ell}(\phi).$$

So $F_{\ell,\ell}(\phi) = C \sin^\ell \phi$. Thus

$$Y_{\ell,\ell}(\theta, \phi) = C e^{i\ell\theta} \sin^\ell \phi.$$

We deduce that the spaces $E_{\lambda,m}$ are 1-dimensional. The lowering process gives, after a normalization,

$$\begin{aligned} Y_{\ell,m-1} &= -\frac{1}{\sqrt{(\ell+m)(\ell-m+1)}} L_- Y_{\ell,m} \\ &= -\frac{1}{\sqrt{(\ell+m)(\ell-m+1)}} e^{-i\theta} \left(-\frac{\partial}{\partial \phi} + i \cot\phi \frac{\partial}{\partial \theta} \right) Y_{\ell,m}. \end{aligned}$$

Some computations yield

$$Y_{\ell,m} = -\frac{1}{\sqrt{2\pi}} e^{im\theta} P_\ell^m(\cos\phi),$$

where

$$P_\ell^m(t) = \sqrt{\frac{(\ell+m)!}{(\ell-m)!}} \sqrt{\frac{2\ell+1}{2}} \frac{1}{2^\ell \ell!} (1-t^2)^{-\frac{m}{2}} \frac{d^{\ell-m}(t^2-1)^\ell}{dt^{\ell-m}}$$

are the (normalized) Legendre polynomials. The functions $Y_{\ell,m}$ are called *spherical functions*.

We conclude that a basis for an irreducible representation of the rotation group consists of

$$Y_{\ell,m}(\theta,\phi), \quad m=-\ell,-\ell+1,\dots,\ell-1,\ell,$$

for a fixed ℓ. Note that spherical functions are dense in $L^2(S^2)$, so the representation of $SO(3)$ on $L^2(S^2)$ can be decomposed into irreducible representations on spaces of spherical functions. Note that since we were able to solve the above differential equations, and since spherical functions span a dense subspace of $L^2(S^2)$, we deduce that $\tilde{L}^2$ and $\tilde{L}_3$ are indeed diagonalizable. What is the most amazing part of this investigation of the problem of diagonalizing $\tilde{L}_3$ is that it gives rise all finite dimensional irreducible representations of $SO(3)$.

5.4 The Schrödinger equation for the hydrogen atom

48. We have a proton (of mass m_1) and an electron (of mass m_2). So the quantum Hamiltonian of this system is

$$H=-\frac{\hbar}{2m_1}\Delta_1-\frac{\hbar}{2m_2}\Delta_2+V(\mathbf{q_1}-\mathbf{q_2}),$$

where $\mathbf{q_1}$ and $\mathbf{q}_2$ are the coordinates of the proton and the electron. We use the new variables

$$\mathbf{Q}=\frac{m_1\mathbf{q}_1+m_2\mathbf{q}_2}{m_1+m_2}, \quad \mathbf{q}=\mathbf{q}_1-\mathbf{q}_2,$$

where $\mathbf{Q}$ is the coordinate of the center of inertia, and $\mathbf{q}$ is the relative coordinate. In these coordinates

$$\hat{H}=-\frac{\hbar^2}{2M}\Delta_{\mathbf{Q}}-\frac{\hbar^2}{2\mu}\Delta_{\mathbf{q}}+V(\mathbf{q}),$$

with $M=m_1+m_2$, $\mu=m_1m_2/(m_1+m_2)$. The time-independent Schrödinger equation

$$\hat{H}\psi=E\psi$$

can be separated into the $\mathbf{Q}$ and $\mathbf{q}$ variables, and only the second has something interesting, as the first is just the Schrödinger equation for free motion.

49. So we focus on the Schrödinger equation

$$-\frac{\hbar}{2\mu}\Delta\psi+V(\mathbf{q})\psi=E\psi.$$

The hydrogen atom has the potential with spherical symmetry

$$V(\mathbf{q})=V(r)=-\frac{e^2}{r}.$$

Using Lemma 5.2.2 we can write the time-independent Schrödinger equation as

$$-\frac{\hbar^2}{2\mu r^2}\frac{\partial}{\partial r}\left(r^2\frac{\partial}{\partial r}\right)\psi + \frac{L^2}{2\mu r^2}\psi - \frac{e^2}{r}\psi = E\psi$$

We solve the equation in *atomic units* with $\hbar = 1$, $\mu = 1$, $e^2 = 1$. We look for solutions of the form

$$\psi(r, \theta, \phi) = R_\ell(r)Y_{\ell,m}(\theta, \psi).$$

They are eigenvectors of $L^2 = \tilde{L}^2$ and $L_3 = \tilde{L}_3$, so they describe states of the particle with definite values of the square of the angular momentum and its z-projection. We obtain the following equation for R_ℓ:

$$-\frac{1}{2r^2}\frac{d}{dr}\left(r^2\frac{dR_\ell}{dr}\right) + \frac{\ell(\ell+1)}{2r^2}R_\ell - \frac{1}{r}R_\ell = ER_\ell.$$

50. We introduce the function $f_\ell = rR_\ell$. This satisfies the radial Schrödinger equation

$$-\frac{1}{2}f_\ell''(r) + \frac{l(l+1)}{2r^2}f_\ell(r) - \frac{1}{r}f_\ell = Ef_\ell.$$

The radial equation coincides with the time-independent Schrödinger equation

$$-\frac{1}{2}\frac{d^2\psi}{dr^2} + V(r)\psi = E\psi$$

if we introduce the effective potential

$$V_{eff}(r) = \begin{cases} -\frac{1}{r} + \frac{\ell(\ell+1)}{2r^2} & \text{if } r > 0 \\ \infty & \text{if } r \leq 0. \end{cases}$$

Here we take into account that $f_\ell(r)$ is only defined for $r > 0$.

Let us examine the behavior of the solution when $r \to \infty$ and $r \to 0$. The first limit produces the equation

$$f_\ell'' + 2Ef_\ell = 0.$$

For $E > 0$ this has two linearly independent solutions e^{-ikr} and e^{ikr}, where $k^2 = 2E$. If $E < 0$, then this has two linearly independent solutions $e^{-\kappa r}$, $e^{\kappa r}$, $\kappa^2 = -2E$.

When $r \to 0$, note that the term with denominator r^2 dominates the term with denominator r and the term that contains E, so we should have something like

$$f_\ell'' - \frac{\ell(\ell+1)}{r^2}f_\ell = 0.$$

We look for solutions of the form r^α, which then yields the equation $\alpha^2 - \alpha - \ell(\ell+1) = 0$. This has the roots $\alpha = -\ell$ and $\alpha = \ell + 1$. So the differential equation has the linearly independent solutions $r^{-\ell}$ and $r^{\ell+1}$.

We are trying to find solutions to the Schrödinger equation of the form

$$\psi(\mathbf{q}) = \psi(r, \theta, \phi) = \frac{f_\ell(r)}{r} Y_{\ell m}(\theta, \phi)$$

that are continuous and are either square integrable, in which case they yield eigenfunctions, or are bounded, in which case they give a point in the spectrum that belongs to the continuous spectrum.

The continuity of ψ implies $f_\ell(0) = 0$, thus near zero the solution should be like $Cr^{\ell+1}$.

- For $E > 0$, the solution f_ℓ is bounded near ∞, all we need is to impose the right behavior at 0, so this gives the continuous spectrum.
- For $E < 0$, we should have a solution that behaves like $Ce^{-\kappa r}$ when $r \to \infty$. We will see that this situation gives rise to the eigenvalues.

51. We examine carefully the case $E < 0$, in which case the solution looks like $Cr^{\ell+1}$ near zero and $Ce^{-\kappa r}$ near infinity. It is therefore convenient to look for a solution of the form

$$f_\ell(r) = r^{l+1} e^{-\kappa r} \Lambda_\ell(r).$$

Substituting in the radial Schrödinger equation we obtain the second order equation for Λ_ℓ:

$$\Lambda_\ell'' + \left(\frac{2(\ell+1)}{r} - 2\kappa \right) \Lambda_\ell' + \left(\frac{2}{r} - \frac{2\kappa(\ell+1)}{r} \right) \Lambda_\ell = 0.$$

Here we used the fact that $E = -\kappa^2/2$.

Now we set

$$\Lambda_\ell(r) = \sum_{j=0}^{\infty} a_j r^j.$$

The second order differential equation yields a first order recursive relation in the coefficients (because of the way the powers of r appear). This recursive relation is

$$a_{j+1} = 2 \frac{\kappa(j+\ell+1) - 1}{(j+1)(j+2\ell+2)} a_j.$$

The ratio test shows that the series converges for all r. Let us examine the large r behavior of the resulting function. Note that

$$a_{j+1} \approx \frac{2\kappa}{j+1} a_j \text{ so } a_j \approx c \frac{(2\kappa)^j}{j!}.$$

Thus $\Lambda_\ell \approx ce^{2\kappa r}$. But then f_ℓ grows too fast. So the only hope is that the recursive relation hits a zero! This happens exactly when

$$\kappa = \frac{1}{j+\ell+1} \text{ for some } j.$$

We can start with ℓ being any nonnegative integer, and we can also choose j to be any nonnegative integer. So κ can be the reciprocal of any positive integer, and we we obtain the eigevalues of the Hamiltonian as the values E such that $-2E = \frac{1}{n^2}$. These are

$$E_n = -\frac{1}{2n^2},$$

with eigenfunctions of the form

$$\psi = r^\ell e^{-\frac{1}{j+\ell+1}r} \lambda_{j\ell}(r) Y_{lm}(\theta, \phi),$$

where $j + \ell + 1 = n$. And we obtain the formula for the frequencies of the spectral lines

$$h\nu_{mn} = E_n - E_m, \quad E_n > E_m,$$

where, when converting back from standard units, we have

$$E_n = -\frac{\mu e^4}{2n^2 h^2}.$$

5.5 The Lie group $SU(2)$ and the spin

52. We now introduce the 2-dimensional special unitary group

$$SU(2) = \left\{ \begin{pmatrix} a & b \\ -\bar{b} & \bar{a} \end{pmatrix} \mid |a|^2 + |b|^2 = 1 \right\}.$$

This is the group of 2×2 unitary matrices with determinant 1. We immediately notice that as a manifold $SU(2)$ is homeomorphic to the 3-dimensional sphere. Indeed, by writing $a = x_1 + iy_1$, $b = x_2 + iy_2$, we identify the elements of this group with the 4-tuples (x_1, y_1, x_2, y_2) satisfying $x_1^2 + y_1^2 + x_2^2 + y_2^2 = 1$. Building on this, we write

$$\begin{pmatrix} a & b \\ -\bar{b} & \bar{a} \end{pmatrix} = x_1 \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix} + y_1 \begin{pmatrix} i & 0 \\ 0 & -i \end{pmatrix} + x_2 \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix} + y_2 \begin{pmatrix} 0 & i \\ i & 0 \end{pmatrix}$$

This establishes an isomorphism between $SU(2)$ and the group U of unit quaternions, with the four matrices corresponding to $1, \mathbf{i}, \mathbf{j}, \mathbf{k}$. Recall that the quaterions are the elements of the form $a + b\mathbf{i} + c\mathbf{j} + d\mathbf{k}$ of an associative algebra with unit $\mathbb{H}$, subject to the rules $\mathbf{i}^2 = \mathbf{j}^2 = \mathbf{k}^2 = \mathbf{ijk} = -1$, $\mathbf{ij} = -\mathbf{ji}$, $\mathbf{jk} = -\mathbf{kj}$, $\mathbf{ik} = -\mathbf{ki}$. Here a, b, c, d range among the real numbers. The unit quaternions are those for which $a^2 + b^2 + c^2 + d^2 = 1$.

Now we construct a $2-1$ surjective group homomorphism

$$\pi : SU(2) \to SO(3),$$

which is also a covering map. For this, consider the Lie group homomorphism

$$U \to Aut(\mathbb{H}), \quad g \mapsto (x \mapsto gxg^{-1}),$$

where $Aut(\mathbb{H})$ are the automorphisms (invertible homomorphisms) of the quaternions. Note that image of U consists of isometries, and because it fixes the reals, it also fixes the orthogonal space to the reals, that is the imaginaries:

$$\text{Im } \mathbb{H} = \{bi + cj + dk \,|\, a, b, c \in \mathbb{R}\}.$$

We identify Im $\mathbb{H}$ with $\mathbb{R}^3$, thus we see that $SU(2)$ is represented by isometries of $\mathbb{R}^3$. Because $SU(2)$ is simply connected, these isometries must lie in $SO(3)$. Moreover, the kernel is $\{\pm 1\}$, and since $SU(2)$ and $SO(3)$ are compact and of the same dimension the map π that we constructed is onto (it is both open and closed).

53. As a corollary, the Lie algebras of $SU(2)$ and $SO(3)$ are isomorphic. But we can construct the Lie algebra $su(2)$ directly; it consists of the matrices that are trace zero and antihermitian:

$$su(2) = \left\{ \begin{pmatrix} ia & z \\ -\bar{z} & -ia \end{pmatrix} \,|\, a \in \mathbb{R}, z \in \mathbb{C} \right\}.$$

It is generated by the matrices

$$u_1 = \begin{pmatrix} 0 & i \\ i & 0 \end{pmatrix} \quad u_2 = \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix} \quad u_3 = \begin{pmatrix} i & 0 \\ 0 & -i \end{pmatrix}.$$

Note that $u_1 = i\sigma_1$, $u_2 = -i\sigma_2$, $u_3 = i\sigma_3$, where $\sigma_1, \sigma_2, \sigma_3$ are the Pauli matrices. We have the isomorphism

$$su(2) \to so(3), \quad u_j \mapsto 2A_j, \quad j = 1, 2, 3.$$

54. We now describe all irreducible representations of $SU(2)$. Let V^n be the space of homogeneous polynomials of degree n in two variables z_1 and z_2. The dimension of V^n is $n+1$, and we endow it with the inner product that makes $P_k(z_1, z_2) = z_1^k z_2^{n-k}$ an orthonormal basis. We let $SU(2)$ act by

$$(gP)(z_1, z_2) = P((z_1, z_2)g), \quad g \in SU(2).$$

Theorem 5.5.1. The representations V^n, $n \geq 0$ are irreducible and unitary and there are no other finite dimensional irreducible representations of $SU(2)$.

Proof. It is not hard to check that the representations are unitary, because they preserve lengths. To show that they are irreducible, we follow the book of Bröcker and tom Dieck. It suffices to show that every endomorphism A of V^n that is $SU(2)$ invariant is a multiple of identity. First let

$$g = \begin{pmatrix} e^i & 0 \\ 0 & e^{-i} \end{pmatrix} \in SU(2).$$

Then $gP_k = e^{(2k-n)i} P_k$, so $gAP_k = e^{(2k-n)i} AP_k$. Since the $e^{(2k-n)i}$-eigenspace of g acting on V^n is $\mathbb{C}P_k$, we must have $AP_k = c_k P_k$ for some $c_k \in \mathbb{C}$.

We now consider the real rotations

$$r_t = \begin{pmatrix} \cos t & -\sin t \\ \sin t & \cos t \end{pmatrix}.$$

We have

$$\begin{aligned} Ar_tP_n = A(z_1\cos t + z_2 \sin t)^n &= \sum_k \binom{n}{k} \cos^k t \cdot \sin^{n-k} t \cdot AP_k \\ &= \sum_k \binom{n}{k} \cos^k t \cdot \sin^{n-k} t \cdot c_k P_k. \end{aligned}$$

Also

$$r_tAP_n = \sum_k \binom{n}{k} \cos^k t \cdot \sin^{n-k} t \cdot c_n P_k.$$

This means that $c_k = c_n$, so $A = c_n I$.

For the converse, we use the fact that every representation of a Lie group into $GL(V)$ is smooth. With the following lemma we transfer the question about representations of the Lie group to that of representations of the Lie algebra.

Lemma 5.5.1. Let $SU(2) \to GL(V)$ be a smooth representation, and let $su(2) \to gl(V)$ be the induced representation of Lie algebras. Then V is an irreducible representation of $SU(2)$ if and only if it is an irreducible representation of $su(2)$.

Proof. Note that $SU(2)$ is obtained by exponentiating elements of $su(2)$. Let $v \in V\backslash\{0\}$. Then

$$\lim_{t\to 0} \frac{e^{Dt}v - v}{t} = Dv \text{ and } e^D v = \sum_{n=0}^{\infty} \frac{1}{n!} D^n v.$$

So every point in a minimal (closed) invariant subspace containing v of $SU(2)$ is in a minimal (closed) invariant subspace containing v of $su(2)$ and vice-versa. □

Let us study the invariant subspaces of $su(2)$. We complexify

$$sl(2,\mathbb{C}) = su(2) \otimes_{\mathbb{R}} \mathbb{C}$$

and work on the easier task of finding the irreducible representations of $sl(2,\mathbb{C})$. It is worth mentioning that this is the Lie algebra of the matrices with trace equal to zero.

We use the Pauli matrices $\sigma_1, \sigma_2, \sigma_3$ and define

$$L_3 = \sigma_3, \quad L_+ = \frac{1}{\sqrt{2}}(\sigma_1 + i\sigma_2), \quad L_- = \frac{1}{\sqrt{2}}(\sigma_1 - i\sigma_2),$$

so that $[L_3, L_\pm] = \pm L_\pm$, $[L_+, L_-] = L_3$.

Remark 5.5.1. We make this choice of operators because it is standard in representation theory. But we could have also worked with L_3, L_+, L_- defined as before (with no square root), same Lie algebra but with different generators. Now you might ask: are we in the same situation? Do we just cite what we have obtained there? In our previous situation we were not interested in the finite dimensional representations, we were interested in all representations. But we had an auxiliary operator, $\tilde{L}^2$, that we do not have here. So here we work in the additional hypothesis that the representation is finite, but we miss the auxiliary operator.The conclusion is almost the same, but with a "twist".

Let v be an eigenvector of L_3 in some irreducible representation of $sl(2,\mathbb{C})$, with eigenvalue λ. Then

$$L_3L_+v = (\lambda+1)L_+v \text{ and } L_3L_-v = (\lambda-1)L_-v.$$

So L_+ and L_- act as raising/lowering operators. Because the representation is finite dimensional, there is k such that $L_+^k v = 0$. We may actually assume $L_+v = 0$ (by changing notation). Then the irreducible representation is the span of $v, L_-v, L_-^2v, \ldots, L_-^nv$ for some n. But then the trace of L_3 is the sum of its eigenvalues, which are λ, $\lambda-1$, ...,$\lambda-n$. L_3, being a commutator, has trace zero, that is $(n+1)\lambda - \frac{n(n+1)}{2} = 0$, which implies $\lambda = \frac{n}{2}$. This means that there is only one irreducible representation of $sl(2,\mathbb{C})$ in each dimension, and we already know that there is one such representation that comes from $SU(2)$. So we have recovered the representations V^{n+1} of $SU(2)$, $n \geq 0$. □

55. We index the irreducible representations by $\ell = \frac{\dim\, V-1}{2} = \frac{n}{2}$, so we change the notation to $V_\ell = V^{2\ell+1}$. The number ℓ is called *spin*.

The eigenvector v is called the highest weight vector. It's eigenvalue is $\ell = \frac{n}{2}$. Note that $L_+(L_-^k v) = k(\ell - \frac{1}{2}(k-1))(L_-^{k-1}v)$. We define the inner product on V_ℓ by asking that the following vectors form an orthonormal basis:

$$|\ell, k\rangle = \sqrt{\frac{2^{\ell-k}(\ell-k)!}{(2\ell)!(\ell-k)!}} L_-^{\ell-k} v, \quad k = -\ell, -\ell+1, \ldots, \ell.$$

Example 5.5.1. If $\ell = 1/2$ we have the orthonormal basis

$$\left|\frac{1}{2}, \frac{1}{2}\right\rangle, \quad \left|\frac{1}{2}, -\frac{1}{2}\right\rangle.$$

with

$$L_+\left|\frac{1}{2}, \frac{1}{2}\right\rangle = 0, \quad L_+\left|\frac{1}{2}, -\frac{1}{2}\right\rangle = \frac{1}{\sqrt{2}}\left|\frac{1}{2}, \frac{1}{2}\right\rangle$$
$$L_-\left|\frac{1}{2}, -\frac{1}{2}\right\rangle = 0, \quad L_-\left|\frac{1}{2}, \frac{1}{2}\right\rangle = \frac{1}{\sqrt{2}}\left|\frac{1}{2}, -\frac{1}{2}\right\rangle.$$

The irreducible representations of $SO(3)$ are in bijective correspondence with the irreducible representations V^n of $SU(2)$ in which $-I$ acts as the identity. These are precisely the representations with integer spin, and we have encountered them before. The other

representations of $SU(2)$ in which $-I$ does not act as the identity only define projective representations of $SO(3)$. The representation (projective or exact) is defined by

$$gv = \pi^{-1}(g)v, \quad g \in SO(3).$$

56. The existence of the spin of an electron was proposed based on experiments. The first of them was the Stern-Gerlach experiment, where a beam of silver atoms was sent through an inhomogenous magnetic field. Classical physics would predict a broadening of the beam due to the perturbation of the trajectories by the inhomogeneity of the field. But in this experiment it was observed that the beam is split into two. Another experiment was the doublet splitting of sodium. It was observed that when the atom of sodium (which has one valence electron) goes from the first excited state to the ground state it produces two adjacent spectral lines. This lead Goudsmit and Uhlenbeck to propose that every electron has an intrinsic angular momentum (spin) equal to $\hbar/2$.

Now if we return to z component of the quantum angular momentum L_3, we notice that it is quantized, and its values are $\hbar$ multiplied by

$$-\ell, -\ell+1, \cdots 0, 1, 2, \cdots, \ell.$$

The intuition is that the electron also has an intrinsic rotation, and this can be one way or the other, this rotation is the spin and it is $\hbar$ multiplied by

$$-\frac{1}{2}, \frac{1}{2}$$

That would the correspond to $\ell = 1/2$. And this is how we arrive at the 2-dimensional irreducible representation of $SU(2)$. The Hilbert space of the electron should be

$$L^2(\mathbb{R}^3)\hat{\otimes}V_{1/2}.$$

But what does this mean?

Let us explain how the Schrödinger equation of an electron in an *electromagnetic field* should look like, following Greiner's *Quantum Mechanics*. If the electric field is $\mathbf{E}$ and the magnetic field is $\mathbf{B}$, then the Lorentz force (also known as the Heaviside-Lorentz force) is

$$\mathbf{F} = e\left(\mathbf{E} + \frac{\mathbf{v}}{c} \times \mathbf{B}\right).$$

Now the electric field and the magnetic field can be expressed in terms of the corresponding vector potential $\mathbf{A}(\mathbf{q}, t)$ and scalar potential $\phi(\mathbf{q}, t)$ as

$$\mathbf{E} = -\nabla\phi - \frac{1}{c}\frac{\partial \mathbf{A}}{\partial t}, \quad \mathbf{B} = \nabla \times \mathbf{A}.$$

The second equality can also be written as $\mathbf{B} = \operatorname{curl} \mathbf{A}$. This can be phrased in fancy language by combining $\mathbf{E}$ and $\mathbf{B}$ in a 2-form in the Minkowski space

$$F = (E_1 dx + E_2 dy + E_3 dz) \wedge dt + B_1 dy \wedge dz + B_2 dz \wedge dx + B_3 dx \wedge dy,$$

and the potential in a 1-form

$$A = A_1 dx + A_2 dy + A_3 dz + \phi dt,$$

and then stating that $F = dA$. Newton's equation

$$m\frac{d^2\mathbf{q}}{dt^2} = e\left(\mathbf{E} + \frac{1}{c}\mathbf{v}\times\mathbf{B}\right)$$

can be derived, via Hamilton's equation, from the Hamiltonian function

$$H = \frac{1}{2m}\left(\mathbf{p} - \frac{e}{c}\mathbf{A}\right)^2 + e\phi.$$

Indeed, when writing Hamilton's equations we obtain

$$\begin{aligned}\frac{d\mathbf{q}}{dt} &= \frac{\partial H}{\partial \mathbf{p}} = \frac{1}{m}\left(\mathbf{p} - \frac{e}{c}\mathbf{A}\right)\\ \frac{d\mathbf{p}}{dt} &= \frac{e}{mc}\left(\mathbf{p} - \frac{e}{c}\mathbf{A}\right)\frac{\partial \mathbf{A}}{\partial \mathbf{q}} - e\frac{\partial \phi}{\partial \mathbf{q}}.\end{aligned}$$

Here, be careful, here $\frac{\partial \mathbf{A}}{\partial \mathbf{q}}$ (indices of $\mathbf{A}$ on the rows, indices of $\mathbf{q}$ on the columns) is a 3×3 matrix multiplied on the right by a row vector. Differentiate the first equation with respect to t and substitute $d\mathbf{p}/dt$ from the second to obtain

$$m\frac{d^2\mathbf{q}}{dt^2} = -\frac{e}{c}\frac{\partial \mathbf{A}}{\partial t} - e\frac{\partial \phi}{\partial \mathbf{q}} + \frac{e}{c}\mathbf{v}\times \operatorname{curl}\mathbf{A},$$

and we are done.

After the quantization the Hamiltonian is

$$\hat{H} = \frac{1}{2m}\left(-i\hbar\hat{\mathbf{p}} - \frac{e}{c}\mathbf{A}\right)^2 + e\phi = -\frac{\hbar^2}{2m}\Delta + \frac{ie\hbar}{mc}\mathbf{A}\cdot\nabla + \frac{ie\hbar}{2mc}(\nabla\cdot\mathbf{A}) + \frac{e^2}{2mc^2}\mathbf{A}\cdot\mathbf{A} + e\phi.$$

Now we introduce the spin. The idea is that since the spin interacts with the magnetic field, the electron gains additional potential energy. This additional term is

$$\mu_B\hat{\sigma}\cdot\mathbf{B} = \frac{|e|\hbar}{2mc}(B_1\sigma_1 + B_2\sigma_2 + B_3\sigma_3),$$

where $\sigma_1, \sigma_2, \sigma_3$ are the Pauli matrices:

$$\sigma_1 = \begin{pmatrix} 0 & 1\\ 1 & 0\end{pmatrix}, \quad \sigma_2 = \begin{pmatrix} 0 & -i\\ i & 0\end{pmatrix}, \quad \sigma_3 = \begin{pmatrix} 1 & 0\\ 0 & -1\end{pmatrix}.$$

The wave function becomes vector valued! The Schrödinger equation becomes the *Pauli equation*

$$i\hbar\frac{\partial \Psi}{\partial t} = \left[\frac{1}{2m}\left(-i\hbar\hat{\mathbf{p}} - \frac{e}{c}\mathbf{A}\right)^2 + e\phi + \mu_B\hat{\sigma}\cdot\mathbf{B}\right]\Psi,$$

where $\Psi = (\psi_1, \psi_2)$ is the spinor wave function. The Pauli equation is a system of two differential equations for ψ_1 and ψ_2, describing electrons with the z-component of their spin up or down, respectively. Note that the values $1/2$ and $-1/2$ (or $\hbar/2$ and $-\hbar/2$ after introducing the Planck's constant in the components of the intrinsic angular momentum) are the eigenvalues of the z component of the angular momentum.

There are other values that the spin can have. By definition, a particle with integer spin is called a *boson* (after Bose) and a particle with half-integer spin is called a *fermion* (after Fermi). Electrons are fermions with spin $1/2$, photons are bosons with spin 1, the Higgs boson has spin 0. These are the only known values that the spin can have in nature.

To take into account the spin, the Hilbert space $L^2(\mathbb{R}^3)$ is modified so as to include internal degrees of freedom. The idea is that for each particle with "spin" ℓ, the Hilbert space should be

$$L^2(\mathbb{R}^3)\hat{\otimes}V_\ell,$$

where V_ℓ is an irreducible projective representation of $SO(3)$ of dimension $2\ell+1$ (which we dissambiguate as an irreducible representation of $SU(2)$). Here the hat denotes the Hilbert space tensor product.

57. Let us open a parenthesis and discuss the situation of *composite systems*. Based on the fact that for two 3-dimensional particles the Hilbert space is $L^2(\mathbb{R}^6) = L^2(\mathbb{R}^3)\hat{\otimes}L^2(\mathbb{R}^3)$, we introduce the following axiom of quantum mechanics.

The Hilbert space of a composite system made up of two subsystems is the Hilbert tensor product of the Hilbert spaces describing the two systems.

We recall that the Hilbert tensor product of $\mathcal{H}_1$ and $\mathcal{H}_2$, $\mathcal{H}_1\hat{\otimes}\mathcal{H}_2$ is obtained by endowing $\mathcal{H}_1 \otimes \mathcal{H}_2 = \mathrm{Span}\{x_1 \otimes x_2 \,|\, x_1 \in \mathcal{H}_1, x_2 \in \mathcal{H}_2\}$ with the inner product

$$\langle x_1 \otimes x_2, y_1 \otimes y_2\rangle = \langle x_1, y_1\rangle \langle x_2, y_2\rangle$$

and then completing it in the induced norm topology.

The quantum Hamiltonian for a non-interacting composite system is

$$\hat{H} = \hat{H}_1 \otimes I + I \otimes \hat{H}_2$$

where $\hat{H}_1, \hat{H}_2$ are the quantum Hamiltonian of the two systems. This is because of how the product rule for the differentiation of the tensor product interacts with the Schrödinger equation

$$\begin{aligned} i\hbar\frac{\partial}{\partial t}(\psi_1 \otimes \psi_2) &= i\hbar\frac{d\psi_1}{dt} \otimes \psi_2 + \psi_1 \otimes (i\hbar\frac{\partial\psi_2}{\partial t} \\ &= \hat{H}_1\psi_1 \otimes \psi_2 + \psi_1 \otimes \hat{H}_2\psi_2 = \hat{H}(\psi_1 \otimes \psi_2). \end{aligned}$$

Let us return to the case of a particle with spin ℓ. Its state is $\psi(\mathbf{x}) \otimes v \in L^2(\mathbb{R}^3)\hat{\otimes}V_\ell$. The system consisting of two such particles has the Hilbert space equal to

$$L^2(\mathbb{R}^6)\hat{\otimes}(V_\ell \otimes V_\ell).$$

Now one of the postulates of quantum mechanics is that *identical particles are indistinguishable*. This means that the state associated to the system formed by the first and the second particle should coincide, up to multiplication by a constant, with the system formed by the second and the first particle, that is

$$\psi(\mathbf{x}, \mathbf{y}) = c\psi(\mathbf{y}, \mathbf{x}).$$

Applying this reasoning twice, we see that $c^2 = 1$, so $c = \pm 1$.

Experimental considerations suggest that $c = (-1)^{2\ell}$. So $c = 1$ for bosons and $c = -1$ for fermions. So we have the following axiom of quantum physics:

Axiom: Consider a collection of N identical particles moving in $\mathbb{R}^3$ and having integer spin ℓ. Then the Hilbert space of such a collection is the subspace of $L^2(\mathbb{R}^{3N})\hat{\otimes}(V_\ell)^{\otimes N}$ consisting of those functions ψ that satisfy

$$\psi(\mathbf{x}_{\sigma(1)}, \mathbf{x}_{\sigma(2)}, \ldots, \mathbf{x}_{\sigma(N)}) = \psi(\mathbf{x}_1, \mathbf{x}_2, \ldots, \mathbf{x}_N)$$

for every permutation σ. Consider a collection of N identical particles moving in $\mathbb{R}^3$ and having half-integer spin ℓ. Then the Hilbert space of such a collection is the subspace of $L^2(\mathbb{R}^{3N})\hat{\otimes}(V_\ell)^{\otimes N}$ consisting of those functions ψ that satisfy

$$\psi(\mathbf{x}_{\sigma(1)}, \mathbf{x}_{\sigma(2)}, \ldots, \mathbf{x}_{\sigma(N)}) = \mathrm{sign}(\sigma)\psi(\mathbf{x}_1, \mathbf{x}_2, \ldots, \mathbf{x}_N)$$

for every permutation σ.

58. This now leads to the natural question of decomposing $V_\ell \otimes V_m$ into sums of irreducible representations.

Theorem 5.5.2. (Clebsch-Gordan Theorem) For any ℓ and m,

$$V_\ell \otimes V_m = \bigoplus_{j=|\ell-m|}^{\ell+m} V_j = V_{|\ell-m|} \oplus V_{|\ell-m|+1} \oplus \cdots \oplus V_{\ell+m-1} \oplus V_{\ell+m}.$$

Proof. Here are two sketches of proofs:
1. We argue on the particular example $V_1 \otimes V_{5/2}$, in which case we represent $e_j \otimes e_k = |1, j\rangle \otimes |5/2, k\rangle$ as the nodes of the diagram in the figure. Examining the eigenvectors of L_3 in $V_1 \otimes V_{5/2}$ we see that they are the subspaces that are spanned by the nodes of the diagonal lines drawn in the figure. It is now not hard to decompose $V_1 \otimes V_{5/2}$ as the direct sum of irreducibles: there is one that runs from the lower-left corner to the upper-right corner which is isomorphic to $V_{7/2}$ (note that it has dimension 8, since there are 7 steps to take you from the lower-left to the upper right). You can also compute its highest weight by hand and see that it is $7/2$. The next "diagonal" has dimension 2, so there will be a vector that is left out and that one will determine a 6-dimensional irreducible representation, with highest weight $5/2$, so we get a copy of $V_{5/2}$. Then we look at the next "diagonal", and we get the $V_{3/2}$. So

$$V_1 \otimes V_{5/2} = V_{3/2} \otimes V_{5/2} \otimes V_{7/2},$$

as claimed.

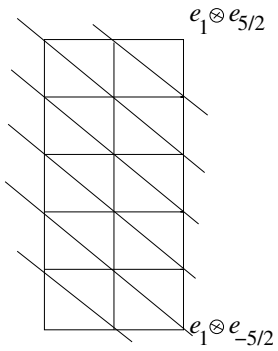

Figure 5.1:

2. Note that every element in $SU(2)$ is conjugate to an element of the form

$$e(t) = \begin{pmatrix} e^{it} & 0 \\ 0 & e^{-it} \end{pmatrix}.$$

Now we use the notion of a character of a representation, which is the trace of the representation. This is a function on the Lie group, or rather on the conjugacy classes of elements of the Lie group. In our case, each representation V will induce the character $\chi_V(t) = \text{trace}e(t)$. In the case of the representation V_ℓ, the value of this character is

$$\chi_\ell = \sum_{k=0}^{2\ell} e^{2i(\ell-k)t}.$$

Now characters are multiplicative with respect to the tensor product and additive with respect to sum. And they determine representations. Thus to check the Clebsch-Gordan theorem, it suffices to check the equality of characters, and this follows from

$$\left(\sum_{k=0}^{2\ell} e^{2i(\ell-k)}\right)\left(\sum_{n=0}^{2m} e^{2i(m-n)}\right) = \sum_{j=|\ell-m|}^{\ell+m}\left(\sum_{t=0}^{2j} e^{2i(j-t)}\right).$$

□

Chapter 6

Quantum Mechanics on Manifolds

Now let us assume that we have a classical mechanical system with constraints, and that we want to quantize it. As seen in the first chapter, the phase space is the cotangent bundle of the manifold which is the configuration space.

6.1 The prequantization line bundle and geometric quantization

59. So we want to apply the quantization scheme to the cotangent line bundle, or more generally to a symplectic manifold (M, ω). A good example to keep in mind is $M = S^1$, the circle, in which case $T^*(M) = S^1 \times \mathbb{R}$ is a cylinder, and this cylinder is endowed with the symplectic form $dp \wedge dq$, where $S^1 = \{e^{iq} \,|\, q \in \mathbb{R}\}$.

Thus we want

$$\begin{array}{ccc} (M, \omega) & \longrightarrow & \mathcal{H} \\ C^\infty(M) & \longrightarrow & L(\mathcal{H}). \end{array}$$

where $\mathcal{H}$ is a Hilbert space, and $L(\mathcal{H})$ are the linear operators on $\mathcal{H}$). Dirac's quantization conditions should be satisfied:

(1) $\hat{1} = I$

(2) $(af + bg)\hat{} = a\hat{f} + b\hat{g}$, $f, g \in C^\infty(M)$, $a, b \in \mathbb{C}$.

(3) The quantization should map the Poisson bracket to the quantum Poisson bracket:

$$\{f, g\}\hat{} = \frac{1}{i\hbar}[\hat{f}, \hat{g}] + O(\hbar).$$

(4) The representation of the algebra of quantum observables should be irreducible.

We could work with $L^2(M)$, square integrable functions on M. To make condition (3) work we can set $\hat{f} = -i\hbar \mathbf{X}_f$, because $[\mathbf{X}_f, \mathbf{X}_g] = -\mathbf{X}_{\{f,g\}}$. But then $\hat{1} = 0$, which conflicts with (1). Then we could modify the formula to $\hat{f} = -i\hbar \mathbf{X}_f + f$, but then (3) fails.

Now comes the geometric trick: If ω, the symplectic form, has what is called a symplectic potential θ, meaning that $d\theta = \omega$, then

$$\hat{f} = -i\hbar\left(\mathbf{X}_f - \frac{i}{\hbar}\theta(\mathbf{X}_f)\right) + f,$$

works! This formula is referred to as *geometric quantization* of observables, and it requires ω to be exact.

Example 6.1.1. In the case of a free 1-dimensional particle,

$$X_q = -\frac{\partial}{\partial p}, \quad X_p = \frac{\partial}{\partial q}.$$

Then $\theta = -qdp$, so

$$\mathrm{op}(q)\phi = -i\hbar\left(-\frac{\partial\phi}{\partial p} - \frac{i}{\hbar}q\phi\right) + q\psi = i\hbar\frac{\partial\phi}{\partial p}$$
$$\mathrm{op}(p)\phi = -i\hbar\frac{\partial\phi}{\partial q} + p\phi.$$

Wow, what is this? In fact you get the standard Schrödinger representation if you set $\phi(q,p) = e^{-(i/\hbar)qp}\psi(q)$. Then

$$\mathrm{op}(q)\psi = q\psi, \quad \mathrm{op}(p)\psi = -i\hbar\frac{\partial\psi}{\partial q}.$$

This seems to be an appropriate choice, because it yields an irreducible representation, as required by (4), while working with all functions does not.

60. Two problems arise in this example. The first is that the state ϕ depends on both p and q. In fact, for a general manifold we cannot separate the variables, and so we are forced to work with functions of both p and q. Well, not quite, as we will see below there is an elegant way to separate variables, and thus to restrict ourselves to only some functions, in this case to $\phi = \psi(q)e^{-(i/\hbar)pq}$, with $\psi \in L^2(\mathbb{R})$.

But there is a second problem, which is more subtle. This comes from the fact that while every symplectic form is locally of the form $\sum_j dp_j \wedge dq_j$, and so locally we can find the potential $\sum_j p_j dq_j$, or the potential $-\sum_j q_j dp_j$, this does not work globally. It does work globally, as we have seen, for cotangent bundles, but for a general symplectic manifold we only have a family of local potentials, and we have to patch those together. We do this by means of a *line bundle*.

Definition. A line bundle over a manifold M is a manifold L and a map $\pi : L \to M$ such that there exists a cover of M by open sets $(U_j)_j$, and homomorphisms $\phi_j : \pi^{-1}(U_j) \to U_j \times \mathbb{C}$ such that $\pi_j(\phi_j(x)) = \pi(x)$ for every $x \in L$ where $\pi_j : U_j \times \mathbb{C} \to U_j$ is $\pi_j(y,z) = y$. Moreover, the transition functions $\phi_j \circ \phi_k^{-1}$ should be of the form $(y,z) \mapsto (y, t_{jk}(x)z)$ with $t_{jk}(x)$ a linear map depending smoothly on x.

Here is another way to say this definition:

Definition. A (complex) *line bundle* $\mathcal{L}$ on a manifold M is defined by a cover $(U_j)_{j\in J}$ of M by contractible open sets such that $U_j \cap U_k$ is either empty or contractible for all j and k, and for every pair (j,k) with $U_j \cap U_k \neq \emptyset$ a smooth map

$$c_{jk} : U_j \cap U_k \to \mathbb{C}\backslash\{0\}.$$

The maps c_{jk} should satisfy the conditions

$$c_{kj} = (c_{jk})^{-1} \text{ and } c_{jk}c_{kl}c_{lj} = 1, \text{ for all } j,k,l \in J. \tag{6.1.1}$$

The line bundle itself is the quotient of the disjoint union of the sets $U_j \times \mathbb{C}$ by the equivalence relation which identifies $(x,z) \in U_j \times \mathbb{C}$ with $(x, c_{jk}z)$ for all $x \in U_j \cap U_k$, $z \in \mathbb{C}$.

The conditions (6.1.1) mean that c_{jk}, $j,k \in J$, is a *Čech cocycle*, more precisely a Čech 2-cocycle. If we consider a family of functions

$$d_j : U_j \to \mathbb{C}\backslash\{0\},$$

then the cocycle $d_j c_{jk}$ defines an equivalent line bundle.

Cocycles form an abelian group under the multiplication,

$$((c_{jk}),(c'_{jk})) \mapsto (c_{jk}c'_{jk})$$

which we denote by $\check{Z}^2((U_j),\mathbb{C})$. The tensor product of two line bundles over the same manifold is obtained by considering a collection of charts common to both and the associated cocycles. Then the cocycle of the tensor product line bundle is the product of the cocycles of the two line bundles.

Let $\check{B}^2((U_j),\mathbb{C})$ be the subgroup consisting of cocycles of the form $d_j d_k^{-1}$, $j,k \in J$, where the functions d_j are as above. If the quotient of two cocycles is in $\check{B}^2((U_j),\mathbb{C})$ then they define the same line bundle. In fact, it can be proved that this is a necessary and sufficient condition.

The quotient group

$$\check{H}^2((U_j),\mathbb{C}) = \check{Z}^2((U_j),\mathbb{C})/\check{B}^2((U_j),\mathbb{C})$$

is called the second *Čech cohomology group*. Each element of $\check{H}^2((U_j),\mathbb{C})$ defines a line bundle up to equivalence.

In this case we want to construct a line bundle with curvature $(1/\hbar)\omega$. We should point out that there is an obstruction to the existence of such a line bundle, known as the *Weil integrality condition*. This requires that

$$\frac{1}{2\pi}\omega \in H^2(M,\mathbb{Z}).$$

This means that when you integrate ω over every closed oriented surface the result is an integral multiple of 2π.

Here is the standard construction: The form ω being symplectic is closed, and so on each open set U_j it is exact, by Poincaré's Lemma. Hence there are real 1-forms θ_j on U_j called potentials, such that

$$d\theta_j = \omega.$$

We have $d(\theta_j - \theta_k) = 0$ on $U_j \cap U_k$, so by using again Poincaré's Lemma on the contractible domains $U_j \cap U_k$, we deduce that there are smooth functions

$$f_{jk} : U_j \cap U_k \to \mathbb{R}$$

such that $df_{jk} = \theta_j - \theta_k$.

Assume that

$$c_{jk} = e^{if_{jk}} \tag{6.1.2}$$

satisfies the cocycle condition (6.1.1). Then it defines a line bundle $\mathcal{L}$. The line bundle $\mathcal{L}$ is said to have curvature ω.

The 1-forms $-i\theta_j$ are the local expressions of a connection form on M. Indeed, by differentiating the relation (6.1.2) we obtain

$$dc_{jk} = ie^{if_{jk}} df_{jk} = ic_{jk}(\theta_j - \theta_k).$$

This can be rewritten as

$$-i\theta_k = c_{jk}^{-1} dc_{jk} + c_{jk}^{-1}(-i\theta_j)c_{jk},$$

and we recognize the formula expressing the change of the connection form under changes of coordinates.

The line bundle constructed above is called the *prequantization line bundle*. The Hilbert space will consist of sections of this line bundle.

Definition. A *section* of the line bundle L is a map $s : M \to L$ such that $\pi \circ s = 1_M$.

The operators defined by geometric quantization now act on smooth L^2 sections of L. There is however a problem with this, in that the Hilbert space is too big. For example for the case of a free 1-dimensional particle we get $L^2(\mathbb{R}^2)$ and not $L^2(\mathbb{R})$. We will learn in the next section how to pick the right sections.

6.2 Polarizations

Three examples should clarify the general situation. One is the standard example, where we work on $L^2(\mathbb{R}^n)$, with the variable $\mathbf{q}$, and have

$$\mathrm{op}(q_j) = M_{q_j}, \quad \mathrm{op}(p_j) = -i\hbar \frac{\partial}{\partial q_j}.$$

The second example corresponds to working in the momentum representation. This time $L^2(\mathbb{R}^n)$, with variable $\mathbf{p}$, and

$$\mathrm{op}(q_j) = i\hbar \frac{\partial}{\partial p_j}, \quad \mathrm{op}(p_j) = M_{p_j}.$$

The third situation was encountered when we studied the harmonic oscillator, where the Hilbert space was a space of harmonic functions. What these situations had in common was

that the functions that were states could be characterized by the conditions that they are annihilated by some differential operators: in the first case $\frac{\partial}{\partial p_j}\psi = 0$, in the second case $\frac{\partial}{\partial q_j}\psi = 0$, while in the third case $\frac{\partial}{\partial \bar{z}}\psi = 0$.

We now bring these conditions in a general unified framework, which works on manifolds. Let (M, ω) be a symplectic manifold of dimension $2n$. The tangent space $T_xM = \mathbb{R}^{2n}$ is itself endowed with the symplectic form ω. For a subspace W of T_xM, we define the orthogonal

$$W^\perp = \{v \in T_x \,|\, \omega(v, w) = 0 \text{ for all } w \in W\}.$$

The subspace W is called isotropic if $W \subset W^\perp$ and Lagrangian if it is maximal isotropic. Alternatively, Lagrangian subspaces are isotropic subspaces of dimension n. $\mathbf{L}$ is Lagrangian if and only if $\mathbf{L}^\perp = \mathbf{L}$.

We complexify the tangent space to $T_xM \otimes \mathbb{C} = \mathbb{C}^{2n}$, by placing complex coefficients in front of the basis vectors. The form ω extends to a symplectic form on $T_xM \otimes \mathbb{C}$. The notion of Lagrangian subspace extends to the complex situation as well, by considering complex subspaces.

Definition. A *polarization of a symplectic real vector space* (V, ω) is a Lagrangian subspace $\mathbf{L}$ of $V \otimes \mathbb{C}$.

A *complex distribution* $\mathbf{F}$ on M is a subbundle of the complexification of the tangent bundle of M. In that sense, at each point $x \in M$, the distribution associates a subspace of $T_xM \otimes \mathbb{C}$, and these subspaces vary smoothly with p and have all the same dimension. A polarization is called involutive if the Poisson bracket of two vector fields in $\mathbf{F}$ is also in $\mathbf{F}$.

A complex distribution $\mathbf{F}$ is called *Lagrangian* if

$$\dim_{\mathbb{C}} \mathbf{F} = \frac{1}{2} \dim_{\mathbb{R}} M \text{ and } \omega|\mathbf{F} \times \mathbf{F} = 0.$$

Definition. A *polarization of the symplectic manifold* (M, ω) is a complex Lagrangian involutive distribution $\mathbf{F}$ such that the function $p \mapsto \dim(\mathbf{F}_p \cap \overline{\mathbf{F}}_p)$ is constant on M, $\overline{\mathbf{F}}$ being the complex conjugate of $\mathbf{F}$.

Example 6.2.1. $M = \mathbb{R}^{2n}$ with the standard symplectic form. Then the distribution on M,

$$F = \text{Span}\left(\frac{\partial}{\partial p_1}, \frac{\partial}{\partial p_2}, \ldots, \frac{\partial}{\partial p_n}\right)$$

is a polarization.

Example 6.2.2. $M = \mathbb{R}^{2n}$ with the standard symplectic form. Then

$$F = \text{Span}\left(\frac{\partial}{\partial q_1}, \frac{\partial}{\partial q_2}, \ldots, \frac{\partial}{\partial q_n}\right)$$

is a polarization.

Example 6.2.3. $M = \mathbb{R}^2$ with the standard symplectic form. Then

$$F = \text{Span}\left(\frac{\partial}{\partial \bar{z}}\right)$$

is a polarization.

Now let us return to the prequantization line bundle L over the phase space M. This line bundle has a connection

$$\nabla = d - i\theta,$$

which is a differential operator that acts on sections. Locally sections are functions, and d is just the operator that turns a function into its differential. The operator d is not independent of the coordinates, and thus it is not globally defined. By adding the "correction" term $-i\theta$ we obtain a globally defined operator, meaning that the new differential operator behaves well under changes of coordinates.

Definition. The Hilbert space of the quantization consists of those sections of $\mathcal{L}$ that are *covariantly constant* in the direction of $\mathbf{F}$, meaning that

$$\nabla_{\mathbf{v}} s = 0, \text{ for all } \mathbf{v} \in \mathbf{F}.$$

Here $\nabla_{\mathbf{v}}$ is the directional derivative in the direction of $\mathbf{v}$.

Example 6.2.4. Let $M = \mathbb{R}^{2n}$ with the standard symplectic form and with polarization

$$F = \text{Span}\left(\frac{\partial}{\partial p_1}, \frac{\partial}{\partial p_2}, \ldots, \frac{\partial}{\partial p_n}\right).$$

Let

$$\theta = -(1/\hbar)\sum_{j=1}^{n} q_j dp_j$$

so that $d\theta = (1/\hbar)\omega$. Then the covariantly constant sections $s \in L^2(\mathbb{R}^{2n})$ are those that satisfy

$$\frac{\partial}{\partial p_j} s + (i/\hbar) q_j s = 0, \quad j = 1, 2, \ldots, n.$$

Solving the differential equation we obtain

$$s(\mathbf{q}, \mathbf{p}) = \psi(\mathbf{q}) e^{-(i/\hbar)\mathbf{p}^T \mathbf{q}}.$$

This should be compared with Example 6.1.1.

Example 6.2.5. Set $z_j = q_j - ip_j$. We consider the polarization

$$F = \text{Span}\left(\frac{\partial}{\partial \bar{z}_1}, \frac{\partial}{\partial \bar{z}_2}, \frac{\partial}{\partial \bar{z}_n}\right),$$

where

$$\frac{\partial}{\partial \bar{z}_j} = \frac{1}{2}\left(\frac{\partial}{\partial q_j} + i\frac{\partial}{\partial p_j}\right).$$

Take also

$$\theta = (1/\hbar)\sum_j p_j dq_j.$$

Then

$$\theta\left(\frac{\partial}{\partial \bar{z}_j}\right) = \frac{p_j}{2}.$$

So

$$\nabla_{\partial/\partial\bar{z}_j} = \frac{\partial}{\partial \bar{z}_j} - \frac{i}{2\hbar}p_j.$$

Thus the sections of the quantization line bundle (which are just functions over $\mathbb{R}^{2n} = \mathbb{C}^n$) should satisfy the system of differential equations

$$\frac{\partial \psi}{\partial \bar{z}_j} - \frac{i}{2\hbar}p_j\psi = 0, \quad j = 1, 2, \ldots, n.$$

If we set

$$\psi(\mathbf{z}, \bar{\mathbf{z}}) = \phi(\mathbf{z}, \bar{\mathbf{z}})e^{-\|Im\,\mathbf{z}\|^2/(2\hbar)},$$

then an easy computation shows that

$$\nabla_{\partial/\partial\bar{z}_j}\phi = e^{-\|Im\,\mathbf{z}\|^2/(2\hbar)}\frac{\partial \phi}{\partial \bar{z}_j}.$$

Thus $\nabla_{\partial/\partial\bar{z}_j} = 0$ for all j if and only if ϕ is a holomorphic function in the variables $z_j = q_j - ip_j$. If we require that the Hilbert space consists of functions that are square integrable over $\mathbb{R}^{2n}$ then ϕ should be in the Segal-Bargmann space of holomorphic square integrable functions on $\mathbb{C}^n$ with respect to the measure

$$e^{-\|Im\,\mathbf{z}\|^2/(2\hbar)}d\mathbf{p}d\mathbf{q}.$$

The argument for placing a Hilbert space structure on the space of sections that we used in the second example cannot be applied to the first example. So we need to be more "clever" when we choose the Hilbert space structure. Moreover, experiments show however that the line bundle constructed above is not quite right.

To solve both problems, there is a procedure called metalinear (or metaplectic) correction, which when applied to the prequantization line bundle yields a new line bundle, and that is the correct line bundle to work with. In short, the idea is the following.

First, for a point $x \in M$, we let $(v_1, v_2, \ldots, v_n)$ be a basis of F_x, called a frame at x. When changing from one coordinate system to another, specifically from U_j to U_k, the frame changes by multiplication by an element $C_{jk}(x) \in GL(n, \mathbb{C})$. The elements C_{jk} satisfy the cocycle condition that we have seen before, but they are matrices. Now consider the cocycle $c_{jk} = \det C_{jk}$. This cocycle defines the line bundle $\Lambda^n F$. Assume that we can find a line

bundle L' over M such that $L' \otimes L' = \Lambda^n F$. Then we replace L by $L \otimes L'$ and consider covariantly constant sections of this new line bundle.

Now what happens is that there is an inner product at each point such that

$$d\left\langle s, s'\right\rangle_x = \left\langle \nabla s, s'\right\rangle_x + \left\langle s, \nabla s'\right\rangle_x .$$

This inner product is constant on D where the complexification of D is $F \cap \overline{F}$. It is natural to define $\langle s, s'\rangle = \int M \left\langle s, s'\right\rangle_x \omega^n$, but this diverges. So instead we write $\left\langle s, s'\right\rangle_x \omega_x^n$ as $\left\langle s, s'\right\rangle_x \alpha \wedge \beta$ where α is the "volume form" of $\Lambda^n F$ and only compute $\int_{M/D} \langle s, s'\rangle\, x\beta$. This decomposition of the integrand is clearly possible locally, and it is precisely because of the alteration of the line bundle that we can decompose it globally.

Example 6.2.6. Let us consider the case of a free one-dimensional particle, whose phase space is $\mathbb{R}^2$. If we take $\theta = pdq$, and the polarization $\partial/\partial p$, then $\theta(\partial/\partial p) = 0$, so

$$\nabla_{\partial/\partial p}\psi = \frac{\partial \psi}{\partial p} = 0,$$

so ψ is a function of q only. Using the formula for geometric quantization we compute

$$\mathrm{op}(q)\psi = q\psi, \quad \mathrm{op}(p)\psi = -i\hbar \frac{\partial \psi}{\partial p}.$$

This gives exactly the formulas predicted by the Schrödinger representation of the Heisenberg Lie algebra. So we should be very happy. But...

$$\mathrm{op}(p^2) = -2i\hbar p \frac{\partial}{\partial q} - p^2.$$

And this introduces the variable p, thus it does not act on $L^2(\mathbb{R})$ with the variable q. This shows the limitations of geometric quantization. Even for the simplest example of a one-dimesional free particle you cannot quantize quadratic polynomials (which we did so easily in Chapter 2). Thus we cannot even quantize kinetic energy.

We conclude this chapter with the observation that we have on the one hand the Weyl quantization, which can only be performed when the phase space is $\mathbb{R}^{2n}$ with the standard symplectic form, but which can be applied to a large variety of observables. On the other hand we have a method of quantizing geometrically every phase space, but this method is very restrictive when it comes to decide which observables (Hamiltonians) can be quantized.