

# Quantum Mechanics

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## 1. Schrödinger equation

### 1.1. Time evolution

In qm time is a parameter and not an observable/operator.  $H$  labels a state given at a time. This implies that we may envisage that, if a physical system is represented by  $|\alpha\rangle$  at the time  $t_0$ , at a later time, the state may differ. The new “evolved” state from  $t_0$  to  $t$  can be denoted by

$$|\alpha, t_0; t\rangle, \quad t \geq t_0$$

such that

$$|\alpha, t_0; t_0\rangle = |\alpha\rangle$$

The transition from  $|\alpha\rangle$  to  $|\alpha, t_0; t\rangle$  is given by the so-called time-evolution operator  $\hat{U}(t, t_0)$ , such that

$$|\alpha, t_0; t\rangle = \hat{U}(t, t_0)|\alpha\rangle \quad (1.1)$$

The time-evolution operator has the following properties:

- Unitarity. Suppose that at  $t_0$   $|\alpha\rangle$  is expanded in terms of the (orthonormal) eigenstates  $|a\rangle$  of some observable  $A$

$$|\alpha\rangle = \sum_a |a\rangle \langle a|\alpha\rangle, \quad |\alpha, t_0; t\rangle = \sum_a c_a(t_0) |a\rangle \quad (1.2)$$

Thus, at a later time  $t > t_0$  we have

$$|\alpha, t_0; t\rangle = \sum_a c_a(t) |a\rangle \quad (1.3)$$

where  $c_a(t)$  is the prob. amp. of finding  $|\alpha\rangle$  at  $|a\rangle$ . Although we expect that in general for any particular  $a$ , the prob. amplitudes differ,

$$c_a(t) \neq c_a(t_0)$$

the sum of all probabilities must be unity at all times, i.e.,

$$\sum_a |c_a(t)|^2 = \sum_a |c_a(t_0)|^2 = 1 \quad (1.4)$$

(as long as the states  $|\alpha\rangle$  and  $|a\rangle$  are correctly normalized). This implies that

$$\langle \alpha | \alpha \rangle = \sum_{a'} \sum_a c_{a'}^*(t_0) \langle a' | a \rangle c_a(t_0) = \sum_a |c_a(t_0)|^2 = 1 \quad (1.5)$$

and consequently

$$\langle \alpha, t_0; t | \alpha, t_0; t \rangle = \sum_a |c_a(t)|^2 = 1 \quad (1.6)$$

That is, the state  $|\alpha\rangle$  remain normalized at all times. For terms of 1.1, this is satisfied as long as

$$\hat{U}^\dagger(t, t_0) \hat{U}(t, t_0) = 1 \quad (1.7)$$

- Time composition:

$$\hat{U}(t_2, t_0) = \hat{U}(t_2, t_1) \hat{U}(t_1, t_0) \quad t_0 < t_1 < t_2 \quad (1.8)$$

- Infinitesimal form. We can also propose an infinitesimal form for  $\hat{U}(t, t_0)$ , such that

$$|\alpha, t_0; t_0 + dt\rangle = \hat{U}(t_0 + dt, t_0) |\alpha\rangle, \quad \lim_{dt \rightarrow 0} \hat{U}(t_0 + dt, t_0) = 1 \quad (1.9)$$

We can note that this together with the previous properties is satisfied if

$$\hat{U}(t_0 + dt, t_0) = 1 - i\hat{\Omega}dt \quad (1.10)$$

when  $\hat{\Omega}$  is a Hermitian operator. This can be shown as follows. Notice that

$$\begin{aligned} \hat{U}^\dagger(t_0 + dt, t_0) \hat{U}(t_0 + dt, t_0) &= (1 + i\hat{\Omega}dt)(1 - i\hat{\Omega}dt) \\ &= 1 + i(\hat{\Omega}^\dagger - \hat{\Omega})dt + O(dt^2) \end{aligned}$$

which equals the identity as long as  $\hat{\Omega} = \hat{\Omega}^\dagger$  (Hermitian) and  $dt^2$  is negligible. Furthermore, the time composition is also direct

$$\begin{aligned} \hat{U}(t_0 + dt_1 + dt_2, t_0 + dt_1) \hat{U}(t_0 + dt_1, t_0) &= (1 - i\hat{\Omega}dt_2)(1 - i\hat{\Omega}dt_1) \\ &= 1 - i\hat{\Omega}(dt_2 + dt_1) + O(dt_1 dt_2) \\ &= \hat{U}(t_0 + dt_1 + dt_2, t_0) \end{aligned}$$

Borrowing from classical mechanics the idea that the Hamiltonian  $H$  is the generator of time evolution, we can assert that

$$\hat{\Omega} = \frac{\hat{H}}{\hbar} \quad (1.11)$$

which has the right dimensions of frequency. Therefore, we obtain

$$\hat{U}(t_0 + dt, t_0) = 1 - \frac{i}{\hbar} \hat{H} dt \quad (1.12)$$

for (1.11) we have used  $\hbar$  to get the right dimensions, recalling the Planck radiation  $E = \hbar\omega$ .

These properties lead to an interesting differential equation. Using the time decomposition property, we see that

$$\hat{U}(t + dt, t_0) = \hat{U}(t + dt, t) \hat{U}(t, t_0) = (1 - \frac{i}{\hbar} \hat{H} dt) \hat{U}(t, t_0) \quad (1.13)$$

Note now that

$$\hat{U}(t + dt, t_0) - \hat{U}(t, t_0) = -\frac{i}{\hbar} \hat{H} dt \hat{U}(t, t_0) \quad (1.14)$$

which implies

$$i\hbar \frac{\hat{U}(t + dt, t_0) - \hat{U}(t, t_0)}{dt} = \hat{H} \hat{U}(t, t_0)$$

that can be written in the differential form

$$i\hbar \frac{\partial}{\partial t} \hat{U}(t, t_0) = \hat{H} \hat{U}(t, t_0) \quad (1.15)$$

This is the Schrödinger equation for the time-evolution operator and is the most fundamental equation of quantum mechanics. Multiplying (1.15) by the ket  $|\alpha\rangle = |\alpha, t_0; t_1\rangle$ , we find

$$\begin{aligned} i\hbar \frac{\partial}{\partial t} \hat{U}(t, t_0) |\alpha\rangle &= \hat{H} \hat{U}(t, t_0) |\alpha\rangle \\ i\hbar \frac{\partial}{\partial t} |\alpha, t_0; t\rangle &= \hat{H} |\alpha, t_0; t\rangle \end{aligned} \quad (1.16)$$

which is the Schrödinger eq. for a state.

The Hamiltonian  $\hat{H}$  has in principle an arbitrary form (it can be an abstract op., or a differential op., or have matrix form) and be time-dependent or time-independent. We can straightforwardly obtain the form of  $\hat{U}(t, t_0)$  for a time-independent Hamiltonian by solving (1.15):

$$\hat{U}(t, t_0) = \exp\left(-\frac{i}{\hbar} \hat{H}(t - t_0)\right) \quad (1.17)$$

which makes sense in general as the series

$$\hat{U}(t, t_0) = 1 - \frac{i}{\hbar} \hat{H}(t - t_0) + \frac{1}{2!} \left( \frac{-i}{\hbar} \hat{H}(t - t_0) \right)^2 \quad (1.18)$$

A time-dependent Hamiltonian is more complicated. We must distinguish two cases:

- when  $[\hat{H}(t_1), \hat{H}(t_2)] = 0$ , and
- when  $[\hat{H}(t_1), \hat{H}(t_2)] \neq 0$

In the former case, the solution of (1.15) reads

$$\hat{U}(t, t_0) = \exp\left(-\frac{i}{\hbar} \int_{t_0}^t dt' \hat{H}(t')\right) \quad (1.19)$$

To solve (1.15) when  $[\hat{H}(t_1), \hat{H}(t_2)] \neq 0$ , we first point out that this differential equation can be restated as

$$\hat{U}(t, t_0) \Big|_{t=t_0}^t = -\frac{i}{\hbar} \int_{t_0}^t dt' \hat{H}(t') \hat{U}(t', t_0)$$

recalling that  $\hat{U}(t_0, t_0) = 1$  we then have

$$\hat{U}(t, t_0) = 1 - \frac{i}{\hbar} \int_{t_0}^t dt' \hat{H}(t') \hat{U}(t', t_0) \quad (1.20)$$

By iteration, we can find the approximate solution

$$\begin{aligned}
\hat{U}(t, t_0) &= 1 - \frac{i}{\hbar} \int_{t_0}^t dt' \hat{H}(t') \left[ 1 - \frac{i}{\hbar} \int_{t_0}^{t'} dt'' \hat{H}(t'') \hat{U}(t'', t_0) \right] \\
&= 1 - \frac{i}{\hbar} \int_{t_0}^t dt' \hat{H}(t') + \left( \frac{-i}{\hbar} \right)^2 \int_{t_0}^t dt' \int_{t_0}^{t'} dt'' \hat{H}(t') \hat{H}(t'') \hat{U}(t'', t_0) \\
&= 1 + \sum_{n=1}^{\infty} \left( \frac{-i}{\hbar} \right)^n \int_{t_0}^t dt_1 \int_{t_0}^{t_1} dt_2 \cdots \int_{t_0}^{t_{n-1}} dt_n \hat{H}(t_1) \hat{H}(t_2) \cdots \hat{H}(t_n)
\end{aligned} \tag{1.21}$$

well-known as the Dyson series.

## 1.2. Time evolution of stationary energy eigenstates

Let us suppose that the Hamiltonian is time-independent. We would like to figure out the evolution of an arbitrary state  $|\alpha\rangle$ . This is particularly simple if we expand  $|\alpha\rangle$  in the basis of eigenstates  $|a\rangle$  of an operator  $\hat{A}$ , such that

$$[\hat{A}, \hat{H}] = 0 \tag{1.22}$$

i.e., compatible with  $\hat{H}$ . It follows that  $\{|a\rangle\}$  are energy eigenstates too, whose energy eigenvalues are

$$\hat{H}|a\rangle = E_a|a\rangle \tag{1.23}$$

Expanding the time-evolution operator in the states  $|a\rangle$ , we find

$$\begin{aligned}
\hat{U}(t, 0) &= \exp\left(-\frac{i}{\hbar} \hat{H}t\right) = \sum_{a'} |a'\rangle \langle a'| \exp(-i\hat{H}t/\hbar) \sum_a |a\rangle \langle a| \\
&= \sum_{a, a'} |a'\rangle \langle a| \langle a'| \exp(-i\hat{H}t/\hbar) |a\rangle \\
&= \sum_a \exp(-iE_a t/\hbar) |a\rangle \langle a|
\end{aligned} \tag{1.24}$$

Now expanding  $|\alpha\rangle$  in the basis of energy eigenstates, we see that

$$\begin{aligned}
|\alpha\rangle &= \sum_a \langle a|\alpha\rangle |a\rangle = \sum_a c_a(t=0) |a\rangle \\
\Rightarrow |\alpha; t\rangle &= U(t, 0)|\alpha\rangle = \sum_{a'} \exp\left(-iE_{a'}t/\hbar\right) |a'\rangle \langle a'| |\alpha\rangle c_a(t=0) \\
&= \sum_a c_a(t=0) \exp(-iE_a t/\hbar) |a\rangle \\
&\equiv \sum_a c_a(t) |a\rangle
\end{aligned}$$

where we realize the time evolution of the expansion coefficients.

Notice that for  $|\alpha\rangle = |a\rangle$ , we find that

$$|a; t\rangle = \exp(-iE_a t/\hbar) |a\rangle \tag{1.25}$$

so that any simultaneous eigenstate of  $H$  and  $A$  remains so at all time.

### 1.3. Expectation values

We are interested in the expectation value of an observable  $B$  in the state

$$|a; t\rangle = \hat{U}(t, 0)|a\rangle$$

as given in (1.26). The expectation value is computed as

$$\langle \hat{B} \rangle = \langle a | \hat{U}^\dagger(t, 0) \hat{B} \hat{U}(t, 0) | a \rangle = \langle a | \exp(iE_a t/\hbar) \hat{B} \exp(-iE_a t/\hbar) | a \rangle$$

Since  $U(t, 0)$  is a numeric phase, the result is

$$\langle \hat{B} \rangle = \langle a; t | \hat{B} | a; t \rangle = \langle a | \hat{B} | a \rangle \quad (1.26)$$

i.e., expectation values do not change in time for  $|\alpha\rangle = |a\rangle$ . This is why these states are called stationary.

The nonstationary states  $|\alpha\rangle = \sum_a c_a(0)|a\rangle$  are different. In this case, we have

$$\begin{aligned} \langle \hat{B} \rangle &= \langle \alpha; t | \hat{B} | \alpha; t \rangle = \langle \alpha | \hat{U}^\dagger(t, 0) \hat{B} \hat{U}(t, 0) | \alpha \rangle \\ &= \sum_a \sum_{a'} \langle a | c_a^*(0) \exp(iE_a t/\hbar) \hat{B} c_{a'}(0) \exp(-iE_{a'} t/\hbar) | a' \rangle \\ &= \sum_a \sum_{a'} c_a^*(0) c_{a'}(0) \exp(-i(E_{a'} - E_a)t/\hbar) \langle a | \hat{B} | a' \rangle \\ &= \sum_a \sum_{a'} c_a^*(0) c_{a'}(0) \exp(-i\omega_{a'a} t) \langle a | \hat{B} | a' \rangle \end{aligned}$$

where the Bohr's frequency  $\omega_{a'a} = \frac{1}{\hbar}(E_{a'} - E_a)$  is the oscillation frequency of the expectation value, and  $\langle a | \hat{B} | a' \rangle$  denote the matrix elements of the observable  $B$  in the basis of  $A$ .

### 1.4. Schrödinger's wave equation

#### 1.4.1. Schrödinger's equation in space representation

Let us define the position states  $|\vec{x}\rangle$  as the eigenstates of the position operator  $\hat{x}$ , such that

$$\hat{x}|\vec{x}'\rangle = \vec{x}'|\vec{x}'\rangle \quad (1.27)$$

Their orthonormality is given by

$$\langle \vec{x} | \vec{x}' \rangle = \delta^3(\vec{x} - \vec{x}') \quad (1.28)$$

Expanded in these terms, an arbitrary physical state  $|\alpha\rangle$  is given by

$$|\alpha\rangle = \int d^3x |\vec{x}\rangle \langle \vec{x} | \alpha \rangle \equiv \int d^3x \psi_\alpha(\vec{x}) |\vec{x}\rangle$$

where  $\psi_\alpha(\vec{x})$  is usually called the wave function of the state  $|\alpha\rangle$  in the representation of positions. Note that the interpretation of  $\psi_\alpha(\vec{x})$  is analogous to that of the expansion coefficient  $c_a = \langle a | \alpha \rangle$  of our previous discussion, i.e., probability amplitud, such that  $|\psi_\alpha(\vec{x})|^2 d^3x$  is the probability of finding in a narrow region  $d^3x$  around  $\vec{x}$ .

One may be interested in finding the element  $\langle \beta | \hat{A} | \alpha \rangle$  for some states  $|\alpha\rangle$ ,  $|\beta\rangle$ , in the position basis

$$\begin{aligned} \langle \beta | \hat{A} | \alpha \rangle &= \int d^3x \int d^3x' \langle \beta | \vec{x} \rangle \langle \vec{x} | \hat{A} | \vec{x}' \rangle \langle \vec{x}' | \alpha \rangle \\ &= \int d^3x \int d^3x' \psi_\beta^* \langle \vec{x} | \hat{A} | \vec{x}' \rangle \psi_\alpha(\vec{x}') \end{aligned} \quad (1.29)$$

which depends on the matrix element  $\langle \vec{x} | \hat{A} | \vec{x}' \rangle$  of  $\hat{A}$ . This expression is greatly simplified if  $\vec{A}$  is a polynomial of the operator  $\vec{x}$ , for example,  $\hat{A}$  can take the form<sup>1</sup> in one dimension

$$\hat{A} = \frac{1}{2} m \omega^2 \vec{x}^2 \quad (1.30)$$

In this case, the matrix element in (1.29) looks like

$$\begin{aligned} \langle \vec{x} | \hat{A} | \vec{x}' \rangle &\mapsto \langle x | \hat{A} | x' \rangle = \frac{1}{2} m \omega^2 \langle x | (\hat{x}^2) | x' \rangle \\ &= \frac{1}{2} m \omega^2 \langle x | (|x' \rangle \langle x'|^2) = \frac{1}{2} m \omega^2 x'^2 \langle x | x' \rangle \\ &= \frac{1}{2} m \omega^2 x'^2 \delta(x - x') \end{aligned}$$

where we have used (1.27) and (1.28) in their 1D versions. This expression can be generalized for  $\hat{A} = f(\hat{x})$  to

$$\langle \vec{x} | f(\hat{x}) | \vec{x}' \rangle = f(x') \delta^3(\vec{x} - \vec{x}') \quad (1.31)$$

and thus we obtain in this case

$$\langle \beta | \hat{A} | \alpha \rangle = \int d^3x \psi_\beta^* f(x) \psi_\alpha(x) \quad (1.32)$$

To arrive at the wave equation proposed by Schrödinger, we need further consider the representation of the momentum operator  $\hat{\vec{p}}$  in the position basis

$$\begin{aligned} \langle \vec{x} | \hat{\vec{p}} | \alpha \rangle &= \int d^3x' \langle \vec{x} | \vec{x}' \rangle \langle \vec{x}' | \hat{\vec{p}} | \alpha \rangle = \int d^3x' \delta^3(\vec{x} - \vec{x}') \int d^3x'' \langle \vec{x}' | \hat{\vec{p}} | \vec{x}'' \rangle \langle \vec{x}'' | \alpha \rangle \\ &= \int d^3x'' \langle \vec{x} | \hat{\vec{p}} | \vec{x}'' \rangle \psi_\alpha(\vec{x}'') \\ &= \int d^3x'' \delta^3(\vec{x} - \vec{x}'') (-i\hbar \vec{\nabla}'') \psi_\alpha(\vec{x}'') \\ &= -i\hbar \vec{\nabla} \psi_\alpha(\vec{x}) \end{aligned} \quad (1.33)$$

Let us now consider the Schrödinger equation, (1.16),

$$i\hbar \frac{\partial}{\partial t} |\alpha, t_0; t\rangle = \hat{H} |\alpha, t_0; t\rangle$$

in the position representation, with a Hamiltonian operator given by

$$\hat{H} = \frac{1}{2m} \hat{\vec{p}}^2 + \hat{V}(\hat{\vec{x}}) \quad (1.34)$$

where  $\hat{V}(\hat{\vec{x}})$  is a polynomial function of  $\hat{\vec{x}}$  and, thus, a Hamiltonian operator. From (1.33), we see that

$$\langle \vec{x} | \hat{V}(\hat{\vec{x}}) | \vec{x}' \rangle = V(\vec{x}'') \delta^3(\vec{x} - \vec{x}') \quad (1.35)$$

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<sup>1</sup>As we shall see, this form corresponds to the potential of the harmonic oscillator.

Multiplying the Schrödinger equation for states (1.16) by the bra  $\langle \vec{x}|$  we find

$$\begin{aligned} i\hbar \frac{\partial}{\partial t} \langle \vec{x}|\alpha, t_0; t \rangle &= \langle \vec{x}|\frac{1}{2m}\hat{p}^2|\alpha, t_0; t \rangle + \langle \vec{x}|\hat{V}(\hat{x})|\alpha, t_0; t \rangle \\ &= \frac{1}{2m}(-i\hbar\vec{\nabla})^2 \langle \vec{x}|\alpha, t_0; t \rangle + V(\vec{x}) \langle \vec{x}|\alpha, t_0; t \rangle \end{aligned} \quad (1.36)$$

where we have used (1.33) and (1.32) with  $\langle \beta| = \langle \vec{x}|$ . Defining

$$\psi_\alpha(\vec{x}, t) \equiv \langle \vec{x}|\alpha, t_0; t \rangle, \quad (1.37)$$

we obtain the time-dependent Schrödinger's wave equation

$$i\hbar \frac{\partial}{\partial t} \psi_\alpha(\vec{x}, t) = -\frac{\hbar^2}{2m} \vec{\nabla}^2 \psi_\alpha(\vec{x}, t) + V(\vec{x}) \psi_\alpha(\vec{x}, t) \quad (1.38)$$

The stationary wave equation is obtained when  $|\alpha \rangle = |a \rangle$  is an energy eigenstate and thus stationary. From (1.25) we have that

$$\begin{aligned} \langle \vec{x}|\alpha; t \rangle &= \langle \vec{x}|a; t \rangle = \exp(-iE_a t/\hbar) \langle \vec{x}|a \rangle = \exp(-iE_a t/\hbar) \langle \vec{x}|\alpha \rangle \\ &\Leftrightarrow \psi_\alpha(\vec{x}, t) = \exp(-iE_a t/\hbar) \psi_\alpha(\vec{x}) \end{aligned} \quad (1.39)$$

where we have defined  $\psi_\alpha(\vec{x}) \equiv \langle \vec{x}|\alpha \rangle$  in analogy with (1.37). It then follows that (1.38) becomes

$$E_\alpha \psi_\alpha(\vec{x}) = -\frac{\hbar^2}{2m} \vec{\nabla}^2 \psi_\alpha(\vec{x}) + V(\vec{x}) \psi_\alpha(\vec{x}) \quad (1.40)$$

when the time dependence carried by the time-evolution operator has been factorized and eliminated.

#### 1.4.2. Schrödinger's equation in momentum representation

As in the configuration space, we can now define the eigenstates of the momentum operator  $\hat{p}$  as

$$\hat{p}|\vec{p}' \rangle = \vec{p}'|\vec{p}' \rangle, \quad (1.41)$$

satisfying the orthonormality condition

$$\langle \vec{p}|\vec{p}' \rangle = \delta^3(\vec{p} - \vec{p}') \quad (1.42)$$

and the completeness relation

$$1 = \int d^3p |\vec{p} \rangle \langle \vec{p}| \quad (1.43)$$

An arbitrary state in the basis reads

$$|\alpha \rangle = \int d^3p |\vec{p} \rangle \langle \vec{p}|\alpha \rangle = \int d^3p |\phi_\alpha(\vec{p})| \vec{p} \rangle \quad (1.44)$$

where we have defined the momentum-space wave function

$$\phi_\alpha(\vec{p}) \equiv \langle \vec{p}|\alpha \rangle \quad (1.45)$$

normalized as

$$\int d^3p |\phi_\alpha(\vec{p})|^2 = \int d^3p \langle \alpha|\vec{p} \rangle \langle \vec{p}|\alpha \rangle = \langle \alpha|\alpha \rangle = 1 \quad (1.46)$$

if  $|\alpha\rangle$  is correctly normalized. The wave function  $\phi_\alpha(\vec{p})$  is related to  $\psi_\alpha(\vec{x})$  by

$$\phi_\alpha(\vec{p}) = \langle \vec{p} | \alpha \rangle = \int d^3x \langle \vec{p} | \vec{x} \rangle \langle \vec{x} | \alpha \rangle = \int d^3x \langle \vec{p} | \vec{x} \rangle \psi_\alpha(\vec{x}) \quad (1.47)$$

and conversely

$$\psi_\alpha(\vec{x}) = \int d^3p \langle \vec{x} | \vec{p} \rangle \phi_\alpha(\vec{p}) \quad (1.48)$$

where  $\langle \vec{p} | \vec{x} \rangle = \langle \vec{x} | \vec{p} \rangle^*$  is yet unknown. To figure out  $\langle \vec{p} | \vec{x} \rangle$ , let us compute

$$\begin{aligned} \langle \vec{x} | \hat{\vec{p}} | \vec{p} \rangle &= \int d^3x' \langle \vec{x} | \hat{\vec{p}} | \vec{x}' \rangle \langle \vec{x}' | \vec{p} \rangle \\ &= -i\hbar \vec{\nabla} \langle \vec{x} | \vec{p} \rangle \end{aligned} \quad (1.49)$$

which, using (1.41), implies

$$\vec{p} \langle \vec{x} | \vec{p} \rangle = -i\hbar \vec{\nabla} \langle \vec{x} | \vec{p} \rangle \quad (1.50)$$

whose solution reads

$$\langle \vec{x} | \vec{p} \rangle = N \exp(i\vec{p} \cdot \vec{x} / \hbar) \quad (1.51)$$

The normalization factor can easily be found by computing

$$\begin{aligned} \langle \vec{x} | \vec{x}' \rangle &= \int d^3p \langle \vec{x} | \vec{p} \rangle \langle \vec{p} | \vec{x}' \rangle \\ &= N^2 \int d^3p \exp\left(-i\vec{p} \cdot (\vec{x} - \vec{x}') / \hbar\right) \\ &= N^2 (2\pi\hbar)^3 \delta^3(\vec{x} - \vec{x}') \end{aligned} \quad (1.52)$$

where the last step is the result of realizing that the integral is just a Fourier transform. Notice that the l.h.s. is actually just  $\delta^3(\vec{x} - \vec{x}')$ , which finally lead to

$$N^2 = (2\pi\hbar)^{-3} \quad (1.53)$$

Replacing into (1.51), we find that

$$\langle \vec{x} | \vec{p} \rangle = \frac{1}{(2\pi\hbar)^{3/2}} \exp(i\vec{p} \cdot \vec{x} / \hbar) \quad (1.54)$$

It then follows from (1.47) and (1.48) that  $\phi_\alpha(\vec{p})$  is just the Fourier transform of  $\psi_\alpha(\vec{x})$ , and vice versa

$$\psi_\alpha(\vec{x}) = \int d^3p \frac{1}{2\pi\hbar}^{3/2} \exp\left(i\vec{p} \cdot \vec{x} / \hbar\right) \phi_\alpha(\vec{p}) \quad (1.55)$$

$$\phi_\alpha(\vec{p}) = \int d^3x \frac{1}{2\pi\hbar}^{3/2} \exp\left(-i\vec{p} \cdot \vec{x} / \hbar\right) \psi_\alpha(\vec{x}) \quad (1.56)$$

Let us use these results to find the Schrödinger's wave equation in momentum space for the Hamiltonian (1.34), defining additionally

$$\phi_\alpha(\vec{p}, t) \equiv \langle \vec{p} | \alpha, t_0; t \rangle \quad (1.57)$$

We have



$$\begin{aligned}
i\hbar \frac{\partial}{\partial t} \langle \vec{p} | \alpha, t_0; t \rangle &= \langle \vec{p} | \frac{1}{2m} \hat{p}^2 | \alpha, t_0; t \rangle \\
&+ \int d^3 p' \langle \vec{p} | \hat{V}(\hat{x}) | \vec{p}' \rangle \langle \vec{p}' | \alpha, t_0; t \rangle \\
&= \frac{1}{2m} \vec{p}^2 \phi_\alpha(\vec{p}, t) + \int d^3 p' \int d^3 x \int d^3 x' \langle \vec{p} | \vec{x} \rangle \langle \vec{x} | \hat{V}(\hat{x}) | \vec{x}' \rangle \langle \vec{x}' | \vec{p}' \rangle \phi_\alpha(\vec{p}', t) \\
&= \frac{\vec{p}^2}{2m} \phi_\alpha(\vec{p}, t) \\
&+ \frac{1}{(2\pi\hbar)^3} \int d^3 p' \int d^3 x \int d^3 x' \exp(-i\vec{p} \cdot \vec{x}/\hbar) V(\vec{x}) \delta^3(\vec{x} - \vec{x}') \exp(i\vec{p}' \cdot \vec{x}'/\hbar) \phi_\alpha(\vec{p}', t) \\
&= \frac{\vec{p}^2}{2m} \phi_\alpha(\vec{p}, t) \\
&+ \frac{1}{(2\pi\hbar)^{3/2}} \int d^3 p' \int \frac{d^3 x}{(2\pi\hbar)^{3/2}} \exp(-i(\vec{p} - \vec{p}') \cdot \vec{x}/\hbar) V(\vec{x}) \phi(\vec{p}', t)
\end{aligned}$$

which finally yields

$$i\hbar \frac{\partial}{\partial t} \phi_\alpha(\vec{p}, t) = \frac{\vec{p}^2}{2m} \phi_\alpha + \frac{1}{(2\pi\hbar)^{3/2}} \int d^3 p' V(\vec{p} - \vec{p}') \phi_\alpha(\vec{p}', t) \quad (1.58)$$

where

$$V(\vec{p} - \vec{p}') \equiv \frac{1}{(2\pi\hbar)^{3/2}} \int \exp(-i(\vec{p} - \vec{p}') \cdot \vec{x}/\hbar) V(\vec{x}) d^3 x \quad (1.59)$$

### 1.4.3. Continuity equation

If we let the probability density of a system to be defined as

$$\rho \equiv \psi_\alpha^*(\vec{x}, t) \psi_\alpha(\vec{x}, t) \quad (1.60)$$

then one easily finds that (omitting for simplicity the dependence on  $\vec{x}$  &  $t$ )

$$\begin{aligned}
\frac{\partial \rho}{\partial t} &= \psi_\alpha^* \frac{\partial}{\partial t} \psi_\alpha + \frac{\partial}{\partial t} \psi_\alpha^* \psi_\alpha \\
&= \psi_\alpha^* \frac{1}{i\hbar} \left( -\frac{\hbar^2}{2m} \vec{\nabla}^2 \psi_\alpha + V \psi_\alpha \right) - \frac{1}{i\hbar} \left( -\frac{\hbar^2}{2m} \vec{\nabla}^2 \psi_\alpha^* + V \psi_\alpha^* \right) \psi_\alpha \\
&= \frac{i\hbar}{2m} \left[ \psi_\alpha^* \vec{\nabla}^2 \psi_\alpha - (\vec{\nabla}^2 \psi_\alpha^*) \psi_\alpha \right] \\
&= \frac{i\hbar}{2m} \left[ \vec{\nabla} (\psi_\alpha^* \cdot \vec{\nabla} \psi_\alpha) - \vec{\nabla} \psi_\alpha^* \cdot \vec{\nabla} \psi_\alpha - \vec{\nabla} \cdot ((\vec{\nabla} \psi_\alpha^*) \psi_\alpha) + \vec{\nabla} \psi_\alpha^* \cdot \vec{\nabla} \psi_\alpha \right] \\
&= \frac{i\hbar}{2m} \vec{\nabla} \cdot \left( \psi_\alpha^* \vec{\nabla} \psi_\alpha - (\vec{\nabla} \psi_\alpha^*) \psi_\alpha \right) \\
&= -\vec{\nabla} \cdot \vec{j}
\end{aligned} \quad (1.61)$$

where

$$\vec{j} \equiv \left( \psi_\alpha^* \vec{\nabla} \psi_\alpha - (\vec{\nabla} \psi_\alpha^*) \psi_\alpha \right) \frac{i\hbar}{2m} \quad (1.62)$$

Rewriting (1.61) in these terms, we find

$$\frac{\partial \rho}{\partial t} + \vec{\nabla} \cdot \vec{j} = 0 \quad (1.63)$$

which coincides with the well-known continuity equation for a probability density  $\rho$  and a probability current  $\vec{j}$ , defined respectively by (1.60) & (1.62). It's interesting to realize that any continuity relation reveals a conserved quantity. In this case, the probability

$$P = \int d^3x \rho \quad (1.64)$$

is conserved (as expected):

$$\frac{dP}{dt} = \int_V d^3x \frac{\partial \rho}{\partial t} = - \int_{\partial V} d\vec{S} \cdot \vec{j} = 0 \quad (1.65)$$

which vanishes only if (1.62) vanishes at the boundary  $\partial V$  located infinitely far away. This condition is satisfied because  $\lim_{\vec{x} \rightarrow \infty} \psi_\alpha(\vec{x}, t) = 0$  is a condition of square integrability, i.e., for  $P$  in (1.64) to be finite always.

## 1.5. Elementary 1 quantum systems

### 1.5.1. Harmonic oscillator in operator method

The harmonic oscillator is one of the most important quantum systems because any physical system is some approximation (around its stability point) resembles an harmonic oscillator. The potential of this system is given, as in classical mechanics, by

$$\hat{V}(\hat{x}) = \frac{1}{2} m \omega^2 \hat{x}^2 \quad (1.66)$$

We would like to find the eigenstates of the Hamiltonian as a time-independent system. To do so, we shall follow the elegant procedure envisaged by Heisenberg based on the non-hermitian “ladder operators”

$$\hat{a} = \sqrt{\frac{m\omega}{2\hbar}} \left( \hat{x} + \frac{i\hat{p}}{m\omega} \right) \quad (1.67a)$$

$$\hat{a}^\dagger = \sqrt{\frac{m\omega}{2\hbar}} \left( \hat{x} - \frac{i\hat{p}}{m\omega} \right) \quad (1.67b)$$

which let one identify the operators  $\hat{x}$  and  $\hat{p}$  as

$$\hat{x} = \sqrt{\frac{\hbar}{2m\omega}} (\hat{a} + \hat{a}^\dagger) \quad (1.68a)$$

$$\hat{p} = i\sqrt{\frac{m\hbar\omega}{2}} (\hat{a}^\dagger - \hat{a}) \quad (1.68b)$$

The ladder operators do not commute:

$$\begin{aligned} [\hat{a}, \hat{a}^\dagger] &= \frac{m\omega}{2\hbar} \left( [\hat{x}, -\frac{i\hat{p}}{m\omega}] + [-\frac{i\hat{p}}{m\omega}, \hat{x}] \right) \\ &= \frac{i}{2\hbar} \left( -[\hat{x}, \hat{p}] + [\hat{p}, \hat{x}] \right) \end{aligned} \quad (1.69)$$

Defining the so-called “number operator”

$$\hat{N} \equiv \hat{a}\hat{a}^\dagger \quad (1.70)$$

it is easy to realize that in terms of  $\hat{x}$  and  $\hat{p}$

$$\begin{aligned} \hat{N} &= \frac{m\omega}{2\hbar} \left( \hat{x}^2 + \frac{\hat{p}^2}{m^2\omega^2} \right) + \frac{i}{2\hbar} [\hat{x}, \hat{p}] \\ &= \frac{1}{\hbar\omega} \left( \frac{\hat{p}^2}{2m} + \frac{1}{2}m\omega^2\hat{x}^2 \right) - \frac{1}{2} \end{aligned} \quad (1.71)$$

Therefore, it follows that the Hamiltonian can be expressed in terms of the number operator as

$$\hat{H} = \hbar\omega \left( \hat{N} + 1/2 \right) \quad (1.72)$$

Note that  $\hat{H}$  and  $\hat{N}$  are compatible (and hermitian) operators because

$$[\hat{H}, \hat{N}] = 0;$$

they can thus be simultaneously diagonalized, so that the eigenstates of  $\hat{N}$  are also energy eigenstates. The eigenstates of  $\hat{N}$  are defined by

$$\hat{N}|n\rangle = n|n\rangle \quad (1.73)$$

where  $n$  is an arbitrary real number at this point. It then follows that

$$\hat{H}|n\rangle = \hbar\omega(n + 1/2)|n\rangle \quad (1.74)$$

Inserting (1.74) in the time-independent Schrödinger’s equation, we identify

$$E_n = \hbar\omega(n + 1/2)$$

as the energy associated with the energy eigenstate  $|n\rangle$ . We can determine the values that  $n$  can take by exploring the action of the ladder operators on  $|n\rangle$ . Note first that

$$\begin{aligned} [\hat{N}, \hat{a}^\dagger] &= [\hat{a}^\dagger\hat{a}, \hat{a}^\dagger] = \hat{a}^\dagger[\hat{a}, \hat{a}^\dagger] = \hat{a}^\dagger \\ [\hat{N}, \hat{a}] &= [\hat{a}^\dagger\hat{a}, \hat{a}] = [\hat{a}^\dagger, \hat{a}]\hat{a} = -\hat{a} \end{aligned} \quad (1.75)$$

where (1.69) has been used. Applying these results, we readily see that

$$B' = -\nabla \times E, \quad (1.76a)$$

$$E' = \nabla \times B - 4\pi j, \quad \hat{N}(\hat{a}^\dagger|n\rangle) = ([\hat{N}, \hat{a}^\dagger] + \hat{a}^\dagger\hat{N})|n\rangle = (n+1)(\hat{a}^\dagger|n\rangle) \quad (1.76b)$$

$$\hat{N}(\hat{a}|n\rangle) = ([\hat{N}, \hat{a}] + \hat{a}\hat{N})|n\rangle = (n-1)(\hat{a}|n\rangle) \quad (1.76c)$$

Eq (1.76b) implies that  $\hat{a}^\dagger|n\rangle$  has the same eigenvalues as the state  $|n+1\rangle$  or, in other words, that the action of the operator  $\hat{a}^\dagger$  raises a state from  $|n\rangle$  to  $|n+1\rangle$ . Analogously,  $\hat{a}$  lowers the state from  $|n\rangle$  to  $|n-1\rangle$ . This is why the ladder operators  $\hat{a}$  and  $\hat{a}^\dagger$  are known as lowering and rising operators, respectively. It immediately follows that the following relations must hold

$$\begin{aligned} \hat{a}|n\rangle &= c|n-1\rangle \\ \hat{a}^\dagger|n\rangle &= d|n+1\rangle \end{aligned} \quad (1.77)$$

where  $c$  and  $d$  are numbers that may be determined by imposing the orthogonality (we shall impose orthonormality) of the eigenstates of the Hermitian operator  $\hat{N}$ . The norm of  $\hat{a}|n\rangle$  is computed as

$$|c|^2 \langle n-1|n-1\rangle = \langle n|\hat{a}^\dagger\hat{a}|n\rangle = n \langle n|n\rangle,$$

which implies

$$|c|^2 = n \quad (1.78)$$