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On statistical mechanics:

what happens when there are too many particles?

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Theoretical Physics

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Introduction

In these lectures notes, we will cover the part of physics which studies systems composed by a large amount of constituents, like particles: thermodynamics and statistical mechanics. In the first part, we will recall some notions of thermodynamics, in the language of differential geometry: the laws of thermodynamics and thermodynamics potentials. In the second part, we will study classical statistical mechanics, starting from the basic notions of classical (hamiltonian) mechanics till the 3 ensembles: microcanonical, canonical and grandcanonical. In the third part, we will study quantum statistical mechanics, starting from the basic notions of quantum mechanics, in the language of canonical and second quantisation, till the properties of bosons and fermions. At the end of each of these last 2 parts, there will be some applications and exercises. In the last part, we will superficially introduce classical phase transitions and the classical Ising model.

Part I

Thermodynamics

Chapter 1

The 2 laws, which are 4

In this chapter, we will recall some notions of thermodynamics: states, equilibrium and the laws of thermodynamics.

1.1 Equilibrium

The topic of which thermodynamics studies is a class of systems composed by a large amount of particles, roughly speaking Avogadro number $N_A \simeq 6 \times 10^{23}$ constituents, once it reaches a macroscopical equilibrium configuration. To understand the notion of equilibrium, consider a system immersed in its surroundings which can interact with the latter by exchanging matter and/or energy (mechanical, electric, magnetic, chemical work) or it can be completely isolated. Once a sufficient amount of time has gone past, it reaches a configuration. Which particular configuration and its stability can be selected by the different boundary conditions the system finds itself in, i.e. the specification on how the system is in contact and how it interacts with its surroundings. In other words, there is only one and only one final equilibrium configuration towards to the system evolves, once boundary conditions have been given. However, the way the system reaches the equilibrium configuration is irreversible. Equilibrium therefore means that once the system has reached its final configuration, it will stay there forever.

1.2 States

A state is a macroscopic configuration. Mathematically speaking, it is a point in the manifold \mathcal{M} of thermodynamic states. To describe it, we need a chart, given by macroscopical physical quantities, called thermodynamical variables. They can be divided into two groups, one conjugate to the other, according to their behaviour when the physical system is rescaled, i.e. when volume and number of particles change: extensive variables do scale with it whereas intensive ones do not. Some of

Extensive	Intensive
energy E	-
entropy S	temperature T
volume V	pression p
number of particles N	chemical potential μ
polarization \mathbf{P}	electric field \mathbf{E}
magnetization \mathbf{M}	magnetic field \mathbf{B}

Table 1.1: Extensive and intensive thermodynamical variables.

them are written in Table 1.1. However, we have to be careful, since only volume is (by definition) extensive and all the others quantities can be considered extensive only if the surface terms are negligible when we take the thermodynamic limit.

Each physical system has an equation of state, i.e. a functional relation among thermodynamic quantities, which restrict the number of independent variables. Geometrically, it means that the only admissible states are a submanifold of the entire manifold of states, given by the constrain induced by the equation of state.

Example 1.1 (Perfect gas). Consider a perfect gas. A chart on its 3-dimensional manifold can be (p, V, T) and its equation of state is $PV = Nk_B T$. This means that the allowed states are in a 2-dimensional manifold embedded in \mathbb{R}^3 .

1.3 The laws of thermodynamics

Thermodynamics is governed by a set of laws which every system must obey. They are particular kind of laws, since they are limitation laws: they tell us only which processes cannot happen.

Law 1.1 (0th)

Let A and B be 2 thermodynamic systems in thermal contact. At equilibrium, only a subset of states $\mathcal{A} \subset \mathcal{M}_A \times \mathcal{M}_B$ is accessible and not the whole manifold. Mathematically, it means that there exists a functional relation of the kind

$$F_{AB}(a, b) = 0 \ , \quad (1.1)$$

with $a \in \mathcal{M}_A$ and $b \in \mathcal{M}_B$. Moreover, thermal equilibrium is an equivalence class, which can be proved that it means

$$F_{AB}(a, b) = f_A(a) - f_B(b) \ . \quad (1.2)$$

Putting together (1.1) and (1.2), we define the empirical temperature

$$t_A = f_A(a) = t_B = f_B(b) \ .$$

It is a limitation law because it limits which possible configuration a system can reach, when in thermal contact with a second one.

Law 1.2 (1st)

Let δQ be the infinitesimal heat and δL the infinitesimal work exchanged in a quasi-static process ($\delta Q > 0$ means absorbed by the system, $\delta L > 0$ means performed by the system). For any cyclic process, i.e. processes in which the initial and the final state coincide, we have

$$\oint (\delta Q - \delta L) = 0 .$$

This means that $\delta Q - \delta L$ is a 1-form, which vanishes when line-integrated along a closed curve in \mathcal{M} . This implies, by the Poincaré lemma, that it is an exact differential

$$dE = \delta Q - \delta L ,$$

called the internal energy. However, notice that heat and work are not exact differential, since $\oint \delta Q \neq 0$ and $\oint \delta L \neq 0$.

The generalisation to a system that can exchange matter is given by

$$\oint (\delta Q - \delta L + \mu dN) = \oint dE = 0 , \quad dE = \delta Q - \delta L + \mu dN , \quad (1.3)$$

where μ is the chemical potential (the necessary energy to add or remove a particle). Furthermore, we can express both δQ and δL as a linear combination of infinitesimal change of independent coordinates, e.g. $\delta L = p dV + B dM$. We assume that the internal energy is extensive and, therefore, the chemical potential is intensive.

It is a limitation law because it limits the configuration that a system can reach in isolation to those with $E = \text{const.}$

Law 1.3 (2nd)

For any cyclic process

$$\oint \frac{\delta Q}{T} \begin{cases} = 0 & \text{reversible process} \\ < 0 & \text{irreversible process} \end{cases} .$$

For reversible processes, $\frac{\delta Q}{T} = 0$ is an exact differential. This implies that we can define a function

$$S(a) - S(b) = \int_a^b \frac{\delta S}{T} ,$$

called the entropy. The integral is along any reversible path. Therefore, we have

$$dS \begin{cases} = 0 & \text{reversible process} \\ < 0 & \text{irreversible process} \end{cases} . \quad (1.4)$$

It is a limitation law because it limits the configuration that a system can reach in isolation to those in which entropy cannot decrease.

Law 1.4 (3rd)

Isothermal and adiabatic processes coincide when $T = 0$, or, equivalently, it is impossible to reach $T = 0$ with a finite number of processes. Mathematically,

$$\Delta S \rightarrow 0 \text{ as } T \rightarrow 0 .$$

Therefore, $T = 0$ is a singular point. Furthermore, if it were possible to reach $T = 0$, the second law $\delta Q \leq 0$ implies that it is impossible to raise the temperature. It is a thermodynamic feature, since it can be proved that it is impossible to realize an engine with efficiency $\eta = 1$.

It is a limitation law because it limits the configuration that a system can reach in isolation to those in which $T \neq 0$.

Chapter 2

Thermodynamic potentials

In this chapter, we will study thermodynamic potentials: energy, entropy, Helmholtz free energy, enthalpy, Gibbs free energy and grand potential. We will derive their definition, their differential and their equations of state.

2.1 Internal energy

Fundamental equation

Thermodynamic potentials are functions defined in the manifold, which are suited for a particular choice of the 3 coordinates and therefore they are useful if we find the system with the other coordinates constant. The first thermodynamic potential we are going to study is the internal energy E . It is defined by the first law of thermodynamic and its differential is

$$dE \leq TdS - pdV + \mu dN . \quad (2.1)$$

This relation is called the fundamental equation of thermodynamics.

Proof. In fact, we invert (1.3)

$$\delta Q = dE + \delta L - \mu dN ,$$

we use $\delta L = pdV$ and we put it into (1.4)

$$dS \leq \frac{\delta Q}{T} = \frac{dE + pdV - \mu dN}{T} . \quad (2.2)$$

Finally, we isolate dE

$$dE \leq TdS - pdV + \mu dN . \quad (2.3)$$

q.e.d.

Equation of state

Notice that the non-differential variables are intensive and the differential variables are extensive. This tells us that $E(S, V, N)$ is a function of the extensive variables S , V and N . The intensive variables T , p and μ can be derived from E by the following relations

$$T = \left. \frac{\partial E}{\partial S} \right|_{V,N}, \quad p = - \left. \frac{\partial E}{\partial V} \right|_{S,N}, \quad \mu = \left. \frac{\partial E}{\partial N} \right|_{S,V}. \quad (2.4)$$

This are called the equation of state of the system, since we can calculate one variable from it, e.g. $T = T(S, V, N)$, $p = p(S, V, N)$ or $\mu = \mu(S, V, N)$.

Proof. At constant V and N , (2.1) becomes

$$dE = TdS - p \underbrace{dV}_0 + \mu \underbrace{dN}_0 = TdS,$$

hence

$$T = \left. \frac{\partial E}{\partial S} \right|_{V,N}.$$

At constant S and N , (2.1) becomes

$$dE = T \underbrace{dS}_0 - pdV + \mu \underbrace{dN}_0 = -pdV,$$

hence

$$p = - \left. \frac{\partial E}{\partial V} \right|_{S,N}.$$

At constant S and V , (2.1) becomes

$$dE = T \underbrace{dS}_0 - p \underbrace{dV}_0 + \mu dN = \mu dN,$$

hence

$$\mu = \left. \frac{\partial E}{\partial N} \right|_{S,V}.$$

q.e.d.

Extensive and intensive

E is an extensive variable, i.e. an homogeneous function of degree one of the extensive variables

$$E(\lambda S, \lambda V, \lambda N) = \lambda E(S, V, N), \quad \forall \lambda > 0.$$

The physical meaning is that if we rescale the volume, the energy is rescaled by the same amount. Moreover, the intensive variables are homogeneous functions of degree zero of the extensive variables

$$T(S, V, N) = T\left(\frac{S}{N}, \frac{V}{N}\right), \quad p(S, V, N) = p\left(\frac{S}{N}, \frac{V}{N}\right), \quad \mu(S, V, N) = \mu\left(\frac{S}{N}, \frac{V}{N}\right). \quad (2.5)$$

By homogeneity properties

$$E = NE\left(\frac{S}{N}, \frac{V}{N}, 1\right) = Ne, \quad S = NS\left(\frac{E}{N}, \frac{V}{N}, 1\right) = Ns,$$

we can define specific energy and entropy

$$e = \frac{E}{N} = e(s, v), \quad s = \frac{S}{N} = s(e, v),$$

where $v = \frac{V}{N}$ is the specific volume. The Euler's theorem allows us to state that, if E is smooth, it can be written as

$$E = S \frac{\partial E}{\partial S} + V \frac{\partial E}{\partial V} + N \frac{\partial E}{\partial N},$$

or, using (2.1) and (2.4),

$$E = TS - pV + \mu N. \quad (2.6)$$

Integrability condition

In order to be an exact differential, the exterior derivative of the right hand side of (2.1) must have a null exterior derivative

$$-\frac{\partial T}{\partial V}\Big|_{S,N} = \frac{\partial p}{\partial S}\Big|_{V,N}, \quad \frac{\partial T}{\partial N}\Big|_{S,V} = \frac{\partial \mu}{\partial S}\Big|_{N,V}, \quad -\frac{\partial p}{\partial N}\Big|_{V,S} = \frac{\partial \mu}{\partial V}\Big|_{N,S}. \quad (2.7)$$

Proof. By means of the exterior derivative, we have

$$\begin{aligned} d(dE) &= d(TdS) - d(pdV) + d(\mu dN) \\ &= \frac{\partial T}{\partial S} \underbrace{dS \wedge dS}_0 + \frac{\partial T}{\partial V} dV \wedge dS + \frac{\partial T}{\partial N} dN \wedge dS - \frac{\partial p}{\partial S} dS \wedge dV - \frac{\partial p}{\partial V} \underbrace{dV \wedge dV}_0 \\ &\quad - \frac{\partial p}{\partial N} dN \wedge dV + \frac{\partial \mu}{\partial S} dS \wedge dN + \frac{\partial \mu}{\partial V} dV \wedge dN + \frac{\partial \mu}{\partial N} \underbrace{dN \wedge dN}_0 \\ &= \frac{\partial T}{\partial V} dV \wedge dS + \frac{\partial T}{\partial N} dN \wedge dS - \frac{\partial p}{\partial S} dS \wedge dV \\ &\quad - \frac{\partial p}{\partial N} dN \wedge dV + \frac{\partial \mu}{\partial S} dS \wedge dN + \frac{\partial \mu}{\partial V} dV \wedge dN. \end{aligned}$$

At constant N

$$\begin{aligned}
 0 = d^2 E &= \frac{\partial T}{\partial V} dV \wedge dS + \frac{\partial T}{\partial N} \underbrace{dN}_0 \wedge dS - \frac{\partial p}{\partial S} dS \wedge dV \\
 &\quad - \frac{\partial p}{\partial N} \underbrace{dN}_0 \wedge dV + \frac{\partial \mu}{\partial S} dS \wedge \underbrace{dN}_0 + \frac{\partial \mu}{\partial V} dV \wedge \underbrace{dN}_0 \\
 &= \frac{\partial T}{\partial V} dV \wedge dS - \frac{\partial p}{\partial S} dS \wedge dV = \frac{\partial T}{\partial V} dV \wedge dS + \frac{\partial p}{\partial S} dV \wedge dS ,
 \end{aligned}$$

hence, by the linear independence of V and S ,

$$-\frac{\partial T}{\partial V} \Big|_{S,N} = \frac{\partial p}{\partial S} \Big|_{V,N} .$$

At constant V

$$\begin{aligned}
 0 = d^2 E &= \frac{\partial T}{\partial V} \underbrace{dV}_0 \wedge dS + \frac{\partial T}{\partial N} dN \wedge dS - \frac{\partial p}{\partial S} dS \wedge \underbrace{dV}_0 \\
 &\quad - \frac{\partial p}{\partial N} dN \wedge \underbrace{dV}_0 + \frac{\partial \mu}{\partial S} dS \wedge dN + \frac{\partial \mu}{\partial V} \underbrace{dV}_0 \wedge dN \\
 &= \frac{\partial T}{\partial N} dN \wedge dS + \frac{\partial \mu}{\partial S} dS \wedge dN = \frac{\partial T}{\partial N} dN \wedge dS - \frac{\partial \mu}{\partial S} dN \wedge dS ,
 \end{aligned}$$

hence, by the linear independence of N and S ,

$$\frac{\partial T}{\partial N} \Big|_{S,V} = \frac{\partial \mu}{\partial S} \Big|_{N,V} .$$

At constant S

$$\begin{aligned}
 0 = d^2 E &= \frac{\partial T}{\partial V} dV \wedge \underbrace{dS}_0 + \frac{\partial T}{\partial N} dN \wedge \underbrace{dS}_0 - \frac{\partial p}{\partial S} \underbrace{dS}_0 \wedge dV \\
 &\quad - \frac{\partial p}{\partial N} dN \wedge dV + \frac{\partial \mu}{\partial S} \underbrace{dS}_0 \wedge dN + \frac{\partial \mu}{\partial V} dV \wedge dN \\
 &= -\frac{\partial p}{\partial N} dN \wedge dV + \frac{\partial \mu}{\partial V} dV \wedge dN = -\frac{\partial p}{\partial N} dN \wedge dV - \frac{\partial \mu}{\partial V} dN \wedge dV ,
 \end{aligned}$$

hence, by the linear independence of N and V ,

$$-\frac{\partial p}{\partial N} \Big|_{V,S} = \frac{\partial \mu}{\partial V} \Big|_{N,S} .$$

q.e.d.

2.2 Entropy

Gibbs-Duhem relation

The Gibbs-Duhem relation expresses the chemical potential μ in terms of the pressure p and the temperature T

$$SdT - Vdp + Nd\mu = 0, \quad d\mu = vdp - sdT. \quad (2.8)$$

Proof. Computing the differential of (2.6)

$$dE = TdS + SdT - pdV + \mu dN + Nd\mu$$

and comparing it with (2.1)

$$dE = TdS + SdT - pdV + -Vdp + \mu dN + Nd\mu = TdS - pdV + \mu dN,$$

we obtain

$$SdT - Vdp + Nd\mu = 0.$$

which can be written as

$$d\mu = \frac{V}{N}dp - \frac{S}{N}dT = vdp - sdT.$$

q.e.d.

Inverting (2.1), we obtained the entropy differential

$$dS = \frac{1}{T}dE + \frac{p}{T}dV - \frac{\mu}{T}dN. \quad (2.9)$$

Its equations of state are

$$\frac{1}{T} = \left. \frac{\partial S}{\partial E} \right|_{V,N}, \quad \frac{p}{T} = \left. \frac{\partial S}{\partial V} \right|_{E,N}, \quad -\frac{\mu}{T} = \left. \frac{\partial S}{\partial N} \right|_{E,V}. \quad (2.10)$$

2.3 Thermodynamic states as a manifold

S , V and N as independent (local) coordinates, i.e. a chart for \mathcal{M} . In our case, it can be thought as an open subset of \mathbb{R}^3 . A reversible process is a path. An irreversible process is an oriented path. Different thermodynamic systems are not connected by any process. Therefore, the manifold is path-connected and simply connected. Solving thermodynamics means find the fundamental equation

$$E = E(S, V, N).$$

However, we could have chosen as fundamental equation

$$S = S(E, V, N)$$

and a chart would have had E , V and N as coordinates. At least one of the local coordinates in any chart for \mathcal{M} must be extensive.

Proof. By the 0th law and (2.5), there must exist a functional relation between intensive variables. This means that one of the three is already fixed once the other two are given and they cannot be used all three as independent coordinates. q.e.d.

There are different thermodynamic potentials, which are functions of 3 independent variables that can be used to define a different chart for \mathcal{M} . Therefore, there are different approaches to thermodynamics.

Thermodynamic potentials can be obtained by various kind of Legendre transforms of (2.1), which exchange the role of an extensive variable to its conjugate intensive variable as independent variable. We require that the hypothesis of the inverse function theorem are satisfied, e.g.

$$\left. \frac{\partial^2 E}{\partial S^2} \right|_{V,N} \neq 0, \quad \left. \frac{\partial^2 E}{\partial V^2} \right|_{S,N} \neq 0, \quad \left. \frac{\partial^2 E}{\partial N^2} \right|_{S,V} \neq 0.$$

The thermodynamic potentials are the Helmholtz free energy F , the entalpy H , the Gibbs free energy G and the granpotential Ω .

2.4 Helmholtz free energy

The Helmholtz free energy is defined as

$$F = E - TS.$$

Its differential is

$$dF \leq -SdT - pdV + \mu dN. \quad (2.11)$$

Therefore

$$F = F(T, V, N).$$

Proof. By a Legendre transform, which means to complete a differential

$$dE \leq TdS - pdV + \mu dN = d(TS) - SdT - pdV + \mu dN,$$

hence

$$dF = d(E - TS) \leq -SdT - pdV + \mu dN.$$

q.e.d.

The equations of state are

$$S = -\frac{\partial F}{\partial T}\Big|_{V,N}, \quad p = -\frac{\partial F}{\partial V}\Big|_{T,N}, \quad \mu = \frac{\partial F}{\partial N}\Big|_{T,V}. \quad (2.12)$$

Proof. At constant V and N

$$dF = -SdT - p\underbrace{dV}_0 + \mu\underbrace{dN}_0 = -SdT,$$

hence

$$S = -\frac{\partial F}{\partial T}\Big|_{V,N}.$$

At constant T and N

$$dF = -S\underbrace{dT}_0 - pdV + \mu\underbrace{dN}_0 = -pdV,$$

hence

$$p = -\frac{\partial F}{\partial V}\Big|_{T,N}.$$

At constant T and V

$$dF = -S\underbrace{dT}_0 - p\underbrace{dV}_0 + \mu dN = \mu dN,$$

hence

$$\mu = \frac{\partial F}{\partial N}\Big|_{T,V}.$$

q.e.d.

The integrability conditions are

$$\frac{\partial S}{\partial V}\Big|_{T,N} = \frac{\partial p}{\partial T}\Big|_{V,N}, \quad -\frac{\partial S}{\partial N}\Big|_{T,V} = \frac{\partial \mu}{\partial T}\Big|_{N,V}, \quad -\frac{\partial p}{\partial N}\Big|_{V,T} = \frac{\partial \mu}{\partial V}\Big|_{N,T}. \quad (2.13)$$

Proof. By means of the exterior derivative, we have

$$\begin{aligned} d(dF) &= -d(SdT) - d(pdV) + d(\mu dN) \\ &= -\frac{\partial S}{\partial T}\underbrace{dT \wedge dT}_0 - \frac{\partial S}{\partial V}dV \wedge dT - \frac{\partial S}{\partial N}dN \wedge dT - \frac{\partial p}{\partial T}dT \wedge dV - \frac{\partial p}{\partial V}\underbrace{dV \wedge dV}_0 \\ &\quad - \frac{\partial p}{\partial N}dN \wedge dV + \frac{\partial \mu}{\partial T}dT \wedge dN + \frac{\partial \mu}{\partial V}dV \wedge dN + \frac{\partial \mu}{\partial N}\underbrace{dN \wedge dN}_0 \\ &= -\frac{\partial S}{\partial V}dV \wedge dT - \frac{\partial S}{\partial N}dN \wedge dT - \frac{\partial p}{\partial T}dT \wedge dV \\ &\quad - \frac{\partial p}{\partial N}dN \wedge dV + \frac{\partial \mu}{\partial T}dT \wedge dN + \frac{\partial \mu}{\partial V}dV \wedge dN. \end{aligned}$$

At constant N

$$\begin{aligned}
 0 = d^2 F &= -\frac{\partial S}{\partial V} dV \wedge dT - \frac{\partial S}{\partial N} \underbrace{dN}_0 \wedge dT - \frac{\partial p}{\partial T} dT \wedge dV \\
 &\quad - \frac{\partial p}{\partial N} \underbrace{dN}_0 \wedge dV + \frac{\partial \mu}{\partial T} dT \wedge \underbrace{dN}_0 + \frac{\partial \mu}{\partial V} dV \wedge \underbrace{dN}_0 \\
 &= -\frac{\partial S}{\partial V} dV \wedge dT - \frac{\partial p}{\partial T} dT \wedge dV = -\frac{\partial S}{\partial V} dV \wedge dT + \frac{\partial p}{\partial T} dV \wedge dT ,
 \end{aligned}$$

hence, by the linear independence of V and T ,

$$\left. \frac{\partial S}{\partial V} \right|_{T,N} = \left. \frac{\partial p}{\partial T} \right|_{V,N} .$$

At constant V

$$\begin{aligned}
 0 = d^2 F &= -\frac{\partial S}{\partial V} \underbrace{dV}_0 \wedge dT - \frac{\partial S}{\partial N} dN \wedge dT - \frac{\partial p}{\partial T} dT \wedge \underbrace{dV}_0 \\
 &\quad - \frac{\partial p}{\partial N} dN \wedge \underbrace{dV}_0 + \frac{\partial \mu}{\partial T} dT \wedge dN + \frac{\partial \mu}{\partial V} \underbrace{dV}_0 \wedge dN \\
 &= -\frac{\partial S}{\partial N} dN \wedge dT + \frac{\partial \mu}{\partial T} dT \wedge dN = -\frac{\partial S}{\partial N} dN \wedge dT - \frac{\partial \mu}{\partial T} dN \wedge dT ,
 \end{aligned}$$

hence, by the linear independence of N and T ,

$$-\left. \frac{\partial S}{\partial N} \right|_{T,V} = \left. \frac{\partial \mu}{\partial T} \right|_{N,V} .$$

At constant T

$$\begin{aligned}
 0 = d^2 F &= -\frac{\partial S}{\partial V} dV \wedge \underbrace{dT}_0 - \frac{\partial S}{\partial N} dN \wedge \underbrace{dT}_0 - \frac{\partial p}{\partial T} \underbrace{dT}_0 \wedge dV \\
 &\quad - \frac{\partial p}{\partial N} dN \wedge dV + \frac{\partial \mu}{\partial T} \underbrace{dT}_0 \wedge dN + \frac{\partial \mu}{\partial V} dV \wedge dN \\
 &= -\frac{\partial p}{\partial N} dN \wedge dV + \frac{\partial \mu}{\partial V} dV \wedge dN = -\frac{\partial p}{\partial N} dN \wedge dV - \frac{\partial \mu}{\partial V} dN \wedge dV ,
 \end{aligned}$$

hence, by the linear independence of N and V ,

$$-\left. \frac{\partial p}{\partial N} \right|_{V,T} = \left. \frac{\partial \mu}{\partial V} \right|_{N,T} .$$

q.e.d.

2.5 Enthalpy

The enthalpy is defined as

$$H = E + pV .$$

Its differential is

$$dH \leq TdS + Vdp + \mu dN . \quad (2.14)$$

Therefore

$$H = H(p, S, N) .$$

Proof. By a Legendre transform, which means to complete a differential

$$dE \leq TdS - pdV + \mu dN = TdS - d(pV) + Vdp + \mu dN ,$$

hence

$$dH = d(E + pV) \leq TdS + Vdp + \mu dN .$$

q.e.d.

The equations of state are

$$T = \left. \frac{\partial H}{\partial S} \right|_{p,N} , \quad V = - \left. \frac{\partial H}{\partial p} \right|_{S,N} , \quad \mu = \left. \frac{\partial H}{\partial N} \right|_{S,p} . \quad (2.15)$$

Proof. At constant p and N

$$dH = TdS + V \underbrace{dp}_0 + \mu \underbrace{dN}_0 ,$$

hence

$$T = \left. \frac{\partial H}{\partial S} \right|_{p,N} .$$

At constant S and N

$$dH = T \underbrace{dS}_0 + Vdp + \mu \underbrace{dN}_0 ,$$

hence

$$V = - \left. \frac{\partial H}{\partial p} \right|_{S,N} .$$

At constant S and p

$$dH = T \underbrace{dS}_0 + V \underbrace{dp}_0 + \mu dN ,$$

hence

$$\mu = \left. \frac{\partial H}{\partial N} \right|_{S,p} .$$

q.e.d.

The integrability conditions are

$$\left. \frac{\partial V}{\partial S} \right|_{p,N} = \left. \frac{\partial T}{\partial p} \right|_{S,N}, \quad \left. \frac{\partial V}{\partial N} \right|_{p,S} = \left. \frac{\partial \mu}{\partial p} \right|_{N,S}, \quad \left. \frac{\partial \mu}{\partial S} \right|_{N,p} = \left. \frac{\partial T}{\partial N} \right|_{S,p}. \quad (2.16)$$

Proof. By means of the exterior derivative, we have

$$\begin{aligned} d(dH) &= d(TdS) + d(Vdp) + d(\mu dN) \\ &= \frac{\partial T}{\partial S} \underbrace{dS \wedge dS}_0 + \frac{\partial T}{\partial p} dp \wedge dS + \frac{\partial T}{\partial N} dN \wedge dS + \frac{\partial V}{\partial S} dS \wedge dp + \frac{\partial V}{\partial p} \underbrace{dp \wedge dp}_0 \\ &\quad + \frac{\partial V}{\partial N} dN \wedge dp + \frac{\partial \mu}{\partial S} dS \wedge dN + \frac{\partial \mu}{\partial p} dp \wedge dN + \frac{\partial \mu}{\partial N} \underbrace{dN \wedge dN}_0 \\ &= \frac{\partial T}{\partial p} dp \wedge dS + \frac{\partial T}{\partial N} dN \wedge dS + \frac{\partial V}{\partial S} dS \wedge dp \\ &\quad + \frac{\partial V}{\partial N} dN \wedge dp + \frac{\partial \mu}{\partial S} dS \wedge dN + \frac{\partial \mu}{\partial p} dp \wedge dN. \end{aligned}$$

At constant N

$$\begin{aligned} 0 = d^2 H &= \frac{\partial T}{\partial p} dp \wedge dS + \frac{\partial T}{\partial N} \underbrace{dN \wedge dS}_0 + \frac{\partial V}{\partial S} dS \wedge dp \\ &\quad + \frac{\partial V}{\partial N} \underbrace{dN \wedge dp}_0 + \frac{\partial \mu}{\partial S} dS \wedge \underbrace{dN}_0 + \frac{\partial \mu}{\partial p} dp \wedge \underbrace{dN}_0 \\ &= \frac{\partial T}{\partial p} dp \wedge dS + \frac{\partial V}{\partial S} dS \wedge dp = \frac{\partial T}{\partial p} dp \wedge dS - \frac{\partial V}{\partial S} dS \wedge dp, \end{aligned}$$

hence, by the linear independence of S and p ,

$$\left. \frac{\partial V}{\partial S} \right|_{p,N} = \left. \frac{\partial T}{\partial p} \right|_{S,N}.$$

At constant S

$$\begin{aligned} 0 = d^2 H &= \frac{\partial T}{\partial p} dp \wedge \underbrace{dS}_0 + \frac{\partial T}{\partial N} dN \wedge \underbrace{dS}_0 + \frac{\partial V}{\partial S} \underbrace{dS \wedge dp}_0 \\ &\quad + \frac{\partial V}{\partial N} dN \wedge dp + \frac{\partial \mu}{\partial S} \underbrace{dS \wedge dN}_0 + \frac{\partial \mu}{\partial p} dp \wedge dN \\ &= \frac{\partial V}{\partial N} dN \wedge dp + \frac{\partial \mu}{\partial p} dp \wedge dN = \frac{\partial V}{\partial N} dN \wedge dp - \frac{\partial \mu}{\partial p} dN \wedge dp, \end{aligned}$$

hence, by the linear independence of N and p ,

$$\left. \frac{\partial V}{\partial N} \right|_{p,S} = \left. \frac{\partial \mu}{\partial p} \right|_{N,S}.$$

At constant p

$$\begin{aligned}
 0 = d^2 H &= \frac{\partial T}{\partial p} \underbrace{dp}_0 \wedge dS + \frac{\partial T}{\partial N} dN \wedge dS + \frac{\partial V}{\partial S} dS \wedge \underbrace{dp}_0 \\
 &+ \frac{\partial V}{\partial N} dN \wedge \underbrace{dp}_0 + \frac{\partial \mu}{\partial S} dS \wedge dN + \frac{\partial \mu}{\partial p} \underbrace{dp}_0 \wedge dN \\
 &= \frac{\partial T}{\partial N} dN \wedge dS + \frac{\partial \mu}{\partial S} dS \wedge dN = \frac{\partial T}{\partial N} dN \wedge dS - \frac{\partial \mu}{\partial S} dS \wedge dN ,
 \end{aligned}$$

hence, by the linear independence of S and N ,

$$\left. \frac{\partial \mu}{\partial S} \right|_{N,p} = \left. \frac{\partial T}{\partial N} \right|_{S,p} .$$

q.e.d.

2.6 Gibbs free energy

The Gibbs free energy is defined as

$$G = E - TS + pV = F + pV = H - TS .$$

Its differential is

$$dG \leq -SdT + Vdp + \mu dN . \quad (2.17)$$

Therefore

$$G = G(p, T, N) .$$

Proof. By a Legendre transform, which means to complete a differential

$$dE \leq TdS - pdV + \mu dN = d(TS) - SdT - d(pV) + Vdp + \mu dN ,$$

hence

$$dG = d(E - TS + pV) \leq -SdT + Vdp + \mu dN .$$

q.e.d.

The equations of state are

$$S = -\left. \frac{\partial G}{\partial T} \right|_{p,N} , \quad V = \left. \frac{\partial G}{\partial p} \right|_{T,N} , \quad \mu = \left. \frac{\partial G}{\partial N} \right|_{p,T} . \quad (2.18)$$

Proof. At constant p and N

$$dG = -SdT + V \underbrace{dp}_0 + \mu \underbrace{dN}_0 ,$$

hence

$$S = -\left. \frac{\partial G}{\partial T} \right|_{p,N} .$$

At constant T and N

$$dG = -S \underbrace{dT}_0 + V dp + \mu \underbrace{dN}_0 ,$$

hence

$$V = \left. \frac{\partial G}{\partial p} \right|_{T,N} .$$

At constant p and T

$$dG = -S \underbrace{dT}_0 + V \underbrace{dp}_0 + \mu dN ,$$

hence

$$\mu = \left. \frac{\partial G}{\partial N} \right|_{p,T} .$$

q.e.d.

The integrability conditions are

$$-\left. \frac{\partial V}{\partial T} \right|_{p,N} = \left. \frac{\partial S}{\partial p} \right|_{T,N} , \quad \left. \frac{\partial V}{\partial N} \right|_{p,T} = \left. \frac{\partial \mu}{\partial p} \right|_{N,T} , \quad -\left. \frac{\partial S}{\partial T} \right|_{T,p} = \left. \frac{\partial \mu}{\partial T} \right|_{N,p} . \quad (2.19)$$

Proof. By means of the exterior derivative, we have

$$\begin{aligned} d(dG) &= -d(SdT) + d(Vdp) + d(\mu dN) \\ &= -\frac{\partial S}{\partial T} \underbrace{dT \wedge dT}_0 - \frac{\partial S}{\partial p} dp \wedge dT - \frac{\partial S}{\partial N} dN \wedge dT + \frac{\partial V}{\partial T} dT \wedge dp + \frac{\partial V}{\partial p} \underbrace{dp \wedge dp}_0 \\ &\quad + \frac{\partial V}{\partial N} dN \wedge dp + \frac{\partial \mu}{\partial T} dT \wedge dN + \frac{\partial \mu}{\partial p} dp \wedge dN + \frac{\partial \mu}{\partial N} \underbrace{dN \wedge dN}_0 \\ &= -\frac{\partial S}{\partial p} dp \wedge dT - \frac{\partial S}{\partial N} dN \wedge dT + \frac{\partial V}{\partial T} dT \wedge dp \\ &\quad + \frac{\partial V}{\partial N} dN \wedge dp + \frac{\partial \mu}{\partial T} dT \wedge dN + \frac{\partial \mu}{\partial p} dp \wedge dN . \end{aligned}$$

At constant N

$$\begin{aligned}
 0 = d^2G &= -\frac{\partial S}{\partial p}dp \wedge dT - \frac{\partial S}{\partial N}\underbrace{dN}_0 \wedge dT + \frac{\partial V}{\partial T}dT \wedge dp \\
 &+ \frac{\partial V}{\partial N}\underbrace{dN}_0 \wedge dp + \frac{\partial \mu}{\partial T}dT \wedge \underbrace{dN}_0 + \frac{\partial \mu}{\partial p}dp \wedge \underbrace{dN}_0 \\
 &= -\frac{\partial S}{\partial p}dp \wedge dT + \frac{\partial V}{\partial T}dT \wedge dp = -\frac{\partial S}{\partial p}dp \wedge dT - \frac{\partial V}{\partial T}dp \wedge dT ,
 \end{aligned}$$

hence, by the linear independence of p and T ,

$$-\frac{\partial V}{\partial T}\Big|_{p,N} = \frac{\partial S}{\partial p}\Big|_{T,N} .$$

At constant T

$$\begin{aligned}
 0 = d^2G &= -\frac{\partial S}{\partial p}dp \wedge \underbrace{dT}_0 - \frac{\partial S}{\partial N}dN \wedge \underbrace{dT}_0 + \frac{\partial V}{\partial T}\underbrace{dT}_0 \wedge dp \\
 &+ \frac{\partial V}{\partial N}dN \wedge dp + \frac{\partial \mu}{\partial T}\underbrace{dT}_0 \wedge dN + \frac{\partial \mu}{\partial p}dp \wedge dN \\
 &= \frac{\partial V}{\partial N}dN \wedge dp + \frac{\partial \mu}{\partial p}dp \wedge dN = \frac{\partial V}{\partial N}dN \wedge dp - \frac{\partial \mu}{\partial p}dp \wedge dN ,
 \end{aligned}$$

hence, by the linear independence of p and N ,

$$\frac{\partial V}{\partial N}\Big|_{p,T} = \frac{\partial \mu}{\partial p}\Big|_{N,T} .$$

At constant p

$$\begin{aligned}
 0 = d^2G &= -\frac{\partial S}{\partial p}\underbrace{dp}_0 \wedge dT - \frac{\partial S}{\partial N}dN \wedge dT + \frac{\partial V}{\partial T}dT \wedge \underbrace{dp}_0 \\
 &+ \frac{\partial V}{\partial N}dN \wedge \underbrace{dp}_0 + \frac{\partial \mu}{\partial T}dT \wedge dN + \frac{\partial \mu}{\partial p}\underbrace{dp}_0 \wedge dN \\
 &= -\frac{\partial S}{\partial N}dN \wedge dT + \frac{\partial \mu}{\partial T}dT \wedge dN = -\frac{\partial S}{\partial N}dN \wedge dT - \frac{\partial \mu}{\partial T}dN \wedge dT ,
 \end{aligned}$$

hence, by the linear independence of N and T ,

$$-\frac{\partial S}{\partial N}\Big|_{T,p} = \frac{\partial \mu}{\partial T}\Big|_{N,p} .$$

q.e.d.

2.7 Grand potential

The grand potential is defined as

$$\Omega = E - TS - \mu N = F - \mu N . \quad (2.20)$$

Its differential is

$$d\Omega \leq -SdT - pdV - Nd\mu . \quad (2.21)$$

Therefore

$$\Omega = \Omega(T, V, \mu) .$$

Proof. By a Legendre transform, which means to complete a differential

$$dE \leq TdS - pdV + \mu dN = d(TS) - SdT - pdV + (\mu N) - Nd\mu ,$$

hence

$$d\Omega = d(E - TS - \mu N) \leq -SdT - pdV - Nd\mu .$$

q.e.d.

The equations of state are

$$S = -\frac{\partial \Omega}{\partial T} \Big|_{\mu, V} , \quad p = -\frac{\partial \Omega}{\partial V} \Big|_{T, \mu} , \quad \mu = -\frac{\partial \Omega}{\partial N} \Big|_{T, V} . \quad (2.22)$$

Proof. At constant μ and V

$$d\Omega = -SdT - p \underbrace{dV}_0 - N \underbrace{d\mu}_0 = -SdT ,$$

hence

$$S = -\frac{\partial \Omega}{\partial T} \Big|_{\mu, V} .$$

At constant T and μ

$$d\Omega = -S \underbrace{dT}_0 - p dV - N \underbrace{d\mu}_0 = -p dV ,$$

hence

$$p = -\frac{\partial \Omega}{\partial V} \Big|_{T, \mu} .$$

At constant T and V

$$d\Omega = -S \underbrace{dT}_0 - p \underbrace{dV}_0 - Nd\mu = -Nd\mu ,$$

hence

$$\mu = -\frac{\partial \Omega}{\partial N} \Big|_{T,V} .$$

q.e.d.

The integrability conditions are

$$\frac{\partial S}{\partial \mu} \Big|_{T,V} = \frac{\partial N}{\partial T} \Big|_{\mu,V} , \quad \frac{\partial S}{\partial V} \Big|_{T,\mu} = \frac{\partial p}{\partial T} \Big|_{V,\mu} , \quad \frac{\partial p}{\partial \mu} \Big|_{V,T} = \frac{\partial N}{\partial V} \Big|_{\mu,T} . \quad (2.23)$$

Proof. By means of the exterior derivative, we have

$$\begin{aligned} d(d\Omega) &= -d(SdT) - d(pdV) - d(Nd\mu) \\ &= -\frac{\partial S}{\partial T} \underbrace{dT \wedge dT}_0 - \frac{\partial S}{\partial V} dV \wedge dT - \frac{\partial S}{\partial \mu} d\mu \wedge dT - \frac{\partial p}{\partial T} dT \wedge dV - \frac{\partial p}{\partial V} \underbrace{dV \wedge dV}_0 \\ &\quad - \frac{\partial p}{\partial \mu} d\mu \wedge dV - \frac{\partial N}{\partial T} dT \wedge d\mu - \frac{\partial N}{\partial V} dV \wedge d\mu - \frac{\partial N}{\partial \mu} \underbrace{d\mu \wedge d\mu}_0 \\ &= -\frac{\partial S}{\partial V} dV \wedge dT - \frac{\partial S}{\partial \mu} d\mu \wedge dT - \frac{\partial p}{\partial T} dT \wedge dV \\ &\quad - \frac{\partial p}{\partial \mu} d\mu \wedge dV - \frac{\partial N}{\partial T} dT \wedge d\mu - \frac{\partial N}{\partial V} dV \wedge d\mu . \end{aligned}$$

At constant μ

$$\begin{aligned} 0 = d^2\Omega &= -\frac{\partial S}{\partial V} dV \wedge dT - \frac{\partial S}{\partial \mu} \underbrace{d\mu \wedge dT}_0 - \frac{\partial p}{\partial T} dT \wedge dV \\ &\quad - \frac{\partial p}{\partial \mu} \underbrace{d\mu \wedge dV}_0 - \frac{\partial N}{\partial T} dT \wedge \underbrace{d\mu}_0 - \frac{\partial N}{\partial V} dV \wedge \underbrace{d\mu}_0 \\ &= -\frac{\partial S}{\partial V} dV \wedge dT - \frac{\partial p}{\partial T} dT \wedge dV = -\frac{\partial S}{\partial V} dV \wedge dT + \frac{\partial p}{\partial T} dV \wedge dT , \end{aligned}$$

hence, by the linear independence of V and T ,

$$\frac{\partial S}{\partial V} \Big|_{T,\mu} = \frac{\partial p}{\partial T} \Big|_{V,\mu} .$$

At constant V

$$\begin{aligned} 0 = d^2\Omega &= -\frac{\partial S}{\partial V} \underbrace{dV \wedge dT}_0 - \frac{\partial S}{\partial \mu} d\mu \wedge dT - \frac{\partial p}{\partial T} dT \wedge \underbrace{dV}_0 \\ &\quad - \frac{\partial p}{\partial \mu} d\mu \wedge \underbrace{dV}_0 - \frac{\partial N}{\partial T} dT \wedge d\mu - \frac{\partial N}{\partial V} \underbrace{dV \wedge d\mu}_0 \\ &= -\frac{\partial S}{\partial \mu} d\mu \wedge dT - \frac{\partial N}{\partial T} dT \wedge d\mu = -\frac{\partial S}{\partial \mu} d\mu \wedge dT + \frac{\partial N}{\partial T} d\mu \wedge dT , \end{aligned}$$

hence, by the linear independence of μ and T ,

$$\left. \frac{\partial S}{\partial \mu} \right|_{T,V} = \left. \frac{\partial N}{\partial T} \right|_{\mu,V} .$$

At constant T

$$\begin{aligned} 0 = d^2\Omega &= -\frac{\partial S}{\partial V}dV \wedge \underbrace{\frac{dT}{0}} - \frac{\partial S}{\partial \mu}d\mu \wedge \underbrace{\frac{dT}{0}} - \frac{\partial p}{\partial T} \underbrace{\frac{dT}{0}} \wedge dV \\ &\quad - \frac{\partial p}{\partial \mu}d\mu \wedge dV - \frac{\partial N}{\partial T} \underbrace{\frac{dT}{0}} \wedge d\mu - \frac{\partial N}{\partial V}dV \wedge d\mu \\ &= -\frac{\partial p}{\partial \mu}d\mu \wedge dV - \frac{\partial N}{\partial V}dV \wedge d\mu = -\frac{\partial p}{\partial \mu}d\mu \wedge dV + \frac{\partial N}{\partial V}d\mu \wedge dV , \end{aligned}$$

hence, by the linear independence of N and V ,

$$\left. \frac{\partial p}{\partial \mu} \right|_{V,T} = \left. \frac{\partial N}{\partial V} \right|_{\mu,T} .$$

q.e.d.

Comments

The thermodynamic potential are not homogeneous functions since they depend on mixed extensive and intensive variables. However, they are extensive, i.e.

$$F = Nf(T, v) , \quad H = Nh(p, s) , \quad G = Ng(T, p) , \quad \Omega = Nf\omega(T, \mu) . \quad (2.24)$$

Notice that the chemical potential is also the Gibbs free energy per particle

$$g(T, p) = \mu(T, p) . \quad (2.25)$$

Proof. In fact, using (2.18) and (2.24)

$$\mu = \frac{\partial G}{\partial N} = \frac{\partial Ng}{\partial N} = g .$$

q.e.d.

Furthermore, notice that

$$\Omega = -pV .$$

Proof. Using (2.6) and (2.20)

$$\Omega = E - TS - \mu N = \mathcal{TS} - pV + \mu\mathcal{N} - \mathcal{TS} - \mu\mathcal{N} = -pV .$$

q.e.d.

2.8 Summary I

A summary of all charts and differentials is given by

$$E(S, V, N) , \quad dE = TdS - pdV + \mu dN ,$$

$$S(E, V, N) , \quad dS = dE/T + pdV/T - \mu dN/T ,$$

$$F(T, V, N) , \quad dF = -SdT - pdV + \mu dN ,$$

$$H(S, p, N) , \quad dH = TdS + Vdp + \mu dN ,$$

$$G(T, p, N) , \quad dG = -SdT + Vdp + \mu dN ,$$

$$\Omega(T, V, \mu) , \quad d\Omega = TdS - pdV + \mu dN .$$

A summary of all the equations of state is given by

$$T = \left. \frac{\partial E}{\partial S} \right|_{V,N} , \quad p = - \left. \frac{\partial E}{\partial V} \right|_{S,N} , \quad \mu = \left. \frac{\partial E}{\partial N} \right|_{S,V} ,$$

$$\frac{1}{T} = \left. \frac{\partial S}{\partial E} \right|_{V,N} , \quad \frac{p}{T} = \left. \frac{\partial S}{\partial V} \right|_{E,N} , \quad -\frac{\mu}{T} = \left. \frac{\partial S}{\partial N} \right|_{E,V} ,$$

$$S = - \left. \frac{\partial F}{\partial T} \right|_{V,N} , \quad p = - \left. \frac{\partial F}{\partial V} \right|_{T,N} , \quad \mu = \left. \frac{\partial F}{\partial N} \right|_{T,V} ,$$

$$T = \left. \frac{\partial H}{\partial S} \right|_{p,N} , \quad V = - \left. \frac{\partial H}{\partial p} \right|_{S,N} , \quad \mu = \left. \frac{\partial H}{\partial N} \right|_{S,p} ,$$

$$S = - \left. \frac{\partial G}{\partial T} \right|_{p,N} , \quad V = \left. \frac{\partial G}{\partial p} \right|_{T,N} , \quad \mu = \left. \frac{\partial G}{\partial N} \right|_{p,T} ,$$

$$S = - \left. \frac{\partial \Omega}{\partial T} \right|_{\mu,V} , \quad p = - \left. \frac{\partial \Omega}{\partial V} \right|_{T,\mu} , \quad \mu = - \left. \frac{\partial \Omega}{\partial N} \right|_{T,V} .$$

A summary of all integrability conditions is given by

$$- \left. \frac{\partial T}{\partial V} \right|_{S,N} = \left. \frac{\partial p}{\partial S} \right|_{V,N} , \quad \left. \frac{\partial T}{\partial N} \right|_{S,V} = \left. \frac{\partial \mu}{\partial S} \right|_{N,V} , \quad - \left. \frac{\partial p}{\partial N} \right|_{V,S} = \left. \frac{\partial \mu}{\partial V} \right|_{N,S} ,$$

$$\left. \frac{\partial S}{\partial V} \right|_{T,N} = \left. \frac{\partial p}{\partial T} \right|_{V,N} , \quad - \left. \frac{\partial S}{\partial N} \right|_{T,V} = \left. \frac{\partial \mu}{\partial T} \right|_{N,V} , \quad - \left. \frac{\partial p}{\partial N} \right|_{V,T} = \left. \frac{\partial \mu}{\partial V} \right|_{N,T} ,$$

$$\left. \frac{\partial V}{\partial S} \right|_{p,N} = \left. \frac{\partial T}{\partial p} \right|_{S,N} , \quad \left. \frac{\partial V}{\partial N} \right|_{p,S} = \left. \frac{\partial \mu}{\partial p} \right|_{N,S} , \quad \left. \frac{\partial \mu}{\partial S} \right|_{N,p} = \left. \frac{\partial T}{\partial N} \right|_{S,p} ,$$

$$- \left. \frac{\partial V}{\partial T} \right|_{p,N} = \left. \frac{\partial S}{\partial p} \right|_{T,N} , \quad \left. \frac{\partial V}{\partial N} \right|_{p,T} = \left. \frac{\partial \mu}{\partial p} \right|_{N,T} , \quad - \left. \frac{\partial S}{\partial N} \right|_{T,p} = \left. \frac{\partial \mu}{\partial T} \right|_{N,p} ,$$

$$\left. \frac{\partial S}{\partial \mu} \right|_{T,V} = \left. \frac{\partial N}{\partial T} \right|_{\mu,V} , \quad \left. \frac{\partial S}{\partial V} \right|_{T,\mu} = \left. \frac{\partial p}{\partial T} \right|_{V,\mu} , \quad \left. \frac{\partial p}{\partial \mu} \right|_{V,T} = \left. \frac{\partial N}{\partial V} \right|_{\mu,T} .$$

Notice that those which are independent are only 6

$$\begin{aligned}
 -\frac{\partial T}{\partial V}\Big|_{S,N} &= \frac{\partial p}{\partial S}\Big|_{V,N}, & \frac{\partial T}{\partial N}\Big|_{S,V} &= \frac{\partial \mu}{\partial S}\Big|_{N,V}, & -\frac{\partial p}{\partial N}\Big|_{V,S} &= \frac{\partial \mu}{\partial V}\Big|_{N,S}, \\
 \frac{\partial S}{\partial V}\Big|_{T,N} &= \frac{\partial p}{\partial T}\Big|_{V,N}, & -\frac{\partial S}{\partial N}\Big|_{T,V} &= \frac{\partial \mu}{\partial T}\Big|_{N,V}, \\
 \frac{\partial V}{\partial N}\Big|_{p,S} &= \frac{\partial \mu}{\partial p}\Big|_{N,S}, \\
 \frac{\partial p}{\partial \mu}\Big|_{V,T} &= \frac{\partial N}{\partial V}\Big|_{\mu,T}.
 \end{aligned}$$

Chapter 3

Maxwell's relations and stability conditions

In this

3.1 Maxwell's relations

Integrability condition can be written as jacobian determinant.

For the energy, they are

$$\frac{\partial(p, S, V)}{\partial(N, S, V)} = -\frac{\partial(\mu, S, N)}{\partial(V, S, N)} = \frac{\partial\mu, S, N}{\partial N, S, V} .$$

Proof. Using the first of (2.7)

$$-\frac{\partial T}{\partial V}\Big|_{S,N} = -\frac{\partial p}{\partial S}\Big|_{V,N} \rightarrow \frac{\partial(T, N, S)}{\partial(V, N, S)} = -\frac{\partial(p, N, V)}{\partial(S, N, V)} = \frac{\partial(p, N, V)}{\partial(V, N, S)} ,$$

hence, inverting the right-handed side

$$1 = \frac{\partial(T, N, S)}{\partial(V, N, S)} \frac{\partial(p, N, V)}{\partial(V, N, S)}^{-1} = \frac{\partial(T, N, S)}{\partial(V, N, S)} \frac{\partial(V, N, S)}{\partial(p, N, V)} = \frac{\partial(T, N, S)}{\partial(p, N, V)} .$$

Using the second of (2.7)

$$\frac{\partial T}{\partial N}\Big|_{S,V} = -\frac{\partial\mu}{\partial S}\Big|_{N,V} \rightarrow \frac{\partial(T, V, S)}{\partial(N, V, S)} = -\frac{\partial(\mu, V, N)}{\partial(S, V, N)} = \frac{\partial(\mu, V, N)}{\partial(N, V, S)} ,$$

hence, inverting the right-handed side

$$1 = \frac{\partial(T, V, S)}{\partial(N, V, S)} \frac{\partial(\mu, V, N)}{\partial(N, V, S)}^{-1} = \frac{\partial(T, V, S)}{\partial(N, V, S)} \frac{\partial(N, V, S)}{\partial(\mu, V, N)} = \frac{\partial(T, V, S)}{\partial(\mu, V, N)} .$$

Using the third of (2.7)

$$\left. \frac{\partial p}{\partial N} \right|_{V,S} = - \left. \frac{\partial \mu}{\partial V} \right|_{N,S} \rightarrow \frac{\partial(p, S, V)}{\partial(N, S, V)} = - \frac{\partial(\mu, S, N)}{\partial(V, S, N)} = \frac{\partial(\mu, S, N)}{\partial(N, S, V)} ,$$

hence, inverting the right-handed side

$$1 = \frac{\partial(p, S, V)}{\partial(N, S, V)} \frac{\partial(\mu, S, N)}{\partial(N, S, V)}^{-1} = \frac{\partial(p, S, V)}{\partial(N, S, V)} \frac{\partial(N, S, V)}{\partial(\mu, S, N)} = \frac{\partial(p, S, V)}{\partial(\mu, S, N)} .$$

q.e.d.

TO BE CONTINUED.

Not all the Maxwell's relations are independent, but only 6 of them

$$\begin{aligned} \frac{\partial(p, V, S)}{\partial(\mu, N, S)} &= 1 , & \frac{\partial(p, V, T)}{\partial(\mu, N, T)} &= 1 , & \frac{\partial(p, V, N)}{\partial(T, S, N)} &= 1 , \\ \frac{\partial(T, S, \mu)}{\partial(p, V, \mu)} &= 1 , & \frac{\partial(T, S, p)}{\partial(N, \mu, p)} &= 1 , & \frac{\partial(T, S, V)}{\partial(N, \mu, V)} &= 1 . \end{aligned}$$

The intergability conditions written in term of jacobian determinant give rise to the geometrical interpretation: the coordinate transformations, which mean that we changed into a different chart of independent variables, preserves the volume.

3.2 Stability conditions

Every thermodynamic potential has a natural chart. In fact, the configuration of stable equilibrium can be obtained by a set of variational principle, which can be derived by fixing to constants the natural independent variables. This variations principle derive from the second law of thermodynamics, since all systems evolve spontaneously to maximise the entropy. Therefore, minima of the thermodynamic potentials correspond to stable equilibrium under boundary condition which keep constant the natural variables

$$(T, V, N) = \text{const} \rightarrow \delta F = 0 , \delta^2 F > 0 ,$$

$$(S, p, N) = \text{const} \rightarrow \delta H = 0 , \delta^2 H > 0 ,$$

$$(T, p, N) = \text{const} \rightarrow \delta G = 0 , \delta^2 G > 0 ,$$

$$(T, V, \mu) = \text{const} \rightarrow \delta \Omega = 0 , \delta^2 \Omega > 0 .$$

Equilibrium of two subsystems requires that T , p and μ are equal.

Proof. Consider two subsystems A and B with extensive variables (E_A, V_A, N_A) and (E_B, V_B, N_B) . Therefore $E = E_A + E_B$, $V = V_A + V_B$ and $N = N_A + N_B$. The whole system is at fixed boundary conditions $E, V, S = \text{const}$. The entropy is additive

$$S = S_A + S_B = S_A(E_A, V_A, N_A) + S_B(E - E_A, V - V_A, N - N_A) .$$

Computing its derivative and imposing it to zero, using (2.10)

$$\begin{aligned} 0 = \delta S &= \frac{\partial S_A}{\partial E_A} \delta E_A + \frac{\partial S_A}{\partial E_A} \delta E_A + \frac{\partial S_A}{\partial V_A} \delta V_A + \frac{\partial S_A}{\partial N_A} \delta N_A \\ &\quad + \frac{\partial S_B}{\partial E_B} \underbrace{\delta(E - E_A)}_{-\delta E_A} + \frac{\partial S_B}{\partial V_B} \underbrace{\delta(V - V_A)}_{-\delta V_A} + \frac{\partial S_B}{\partial N_B} \underbrace{\delta(N - N_A)}_{-\delta N_A} \\ &= \frac{\partial S_A}{\partial E_A} \delta E_A + \frac{\partial S_A}{\partial V_A} \delta V_A + \frac{\partial S_A}{\partial N_A} \delta N_A - \frac{\partial S_B}{\partial E_B} \delta E_A - \frac{\partial S_B}{\partial V_B} \delta V_A - \frac{\partial S_B}{\partial N_B} \delta N_A \\ &= \delta E_A \left(\underbrace{\frac{\partial S_A}{\partial E_A}}_{\frac{1}{T_A}} - \underbrace{\frac{\partial S_B}{\partial E_B}}_{\frac{1}{T_B}} \right) + \delta V_A \left(\underbrace{\frac{\partial S_A}{\partial V_A}}_{\frac{p_A}{T_A}} - \underbrace{\frac{\partial S_B}{\partial V_B}}_{\frac{p_B}{T_B}} \right) + \delta N_A \left(\underbrace{\frac{\partial S_A}{\partial N_A}}_{-\frac{\mu_A}{T_A}} - \underbrace{\frac{\partial S_B}{\partial N_B}}_{-\frac{\mu_B}{T_B}} \right) \\ &= \delta E_A \left(\frac{1}{T_A} - \frac{1}{T_B} \right) + \delta V_A \left(\frac{p_A}{T_A} - \frac{p_B}{T_B} \right) + \delta N_A \left(-\frac{\mu_A}{T_A} + \frac{\mu_B}{T_B} \right) , \end{aligned}$$

hence, by the arbitrariness of δE_A , δV_A and δN_A ,

$$T_A = T_B , \quad p_A = p_B , \quad \mu_A = \mu_B .$$

q.e.d.

At $T, p, N = \text{const}$, the stability condition is

$$\begin{aligned} E_{SS} &= \left. \frac{\partial T}{\partial S} \right|_V > 0 , \quad E_{VV} = - \left. \frac{\partial p}{\partial V} \right|_S > 0 , \\ E_{SS} E_{VV} - E_{SV}^2 &= - \left. \frac{\partial T}{\partial S} \right|_V \left. \frac{\partial p}{\partial V} \right|_S - \left(\left. \frac{\partial p}{\partial S} \right|_V \right)^2 = - \left. \frac{\partial T}{\partial S} \right|_V \left. \frac{\partial p}{\partial V} \right|_S - \left(\left. \frac{\partial T}{\partial V} \right|_S \right)^2 > 0 , \end{aligned} \quad (3.1)$$

Proof. We know that $E = E(S, V, N)$. At constant N , its variation is

$$\begin{aligned} \delta E &= \underbrace{\left. \frac{\partial E}{\partial S} \right|_V}_T \delta S + \underbrace{\left. \frac{\partial E}{\partial V} \right|_S}_{-p} \delta V \\ &\quad + \frac{1}{2} \left(\underbrace{\left. \frac{\partial^2 E}{\partial S^2} \right|_V}_{E_{SS}} \delta S^2 + 2 \underbrace{\left. \frac{\partial^2 E}{\partial S \partial V} \right|_S}_{E_{SV}} \delta S \delta V + \underbrace{\left. \frac{\partial^2 E}{\partial V^2} \right|_S}_{E_{VV}} \delta V^2 \right) \\ &= T \delta S - p \delta V + \frac{1}{2} \left(E_{SS} \delta S^2 + 2 E_{SV} \delta S \delta V + E_{VV} \delta V^2 \right) . \end{aligned}$$

The first derivative terms vanishes, since

$$\begin{aligned}\delta G &= \delta E - T\delta S + p\delta V \\ &= T\delta S - p\delta V + \frac{1}{2}\left(E_{SS}\delta S^2 + 2E_{SV}\delta S\delta V + E_{VV}\delta V^2\right) - T\delta S + p\delta V \\ &= \frac{1}{2}\left(E_{SS}\delta S^2 + 2E_{SV}\delta S\delta V + E_{VV}\delta V^2\right) .\end{aligned}$$

The condition to be a minimum is that

$$E_{SS} > 0 , \quad E_{VV} > 0 , \quad E_{SS}E_{VV} - E_{SV}^2 > 0 .$$

Respectively, they become

$$E_{SS} = \frac{\partial}{\partial S} \underbrace{\frac{\partial E}{\partial S}}_T > 0 ,$$

$$E_{VV} = \frac{\partial}{\partial V} \underbrace{\frac{\partial E}{\partial V}}_{-p} > 0 ,$$

$$E_{SS}E_{VV} - E_{SV}^2 = -\frac{\partial T}{\partial S}\Big|_V \frac{\partial p}{\partial V}\Big|_S - \left(\frac{\partial p}{\partial S}\Big|_V\right)^2 = -\frac{\partial T}{\partial S}\Big|_V \frac{\partial p}{\partial V}\Big|_S - \left(\frac{\partial T}{\partial V}\Big|_S\right)^2 > 0 .$$

q.e.d.

We define the stability conditions in terms of the specific heat

$$C_V = T \frac{\partial S}{\partial T}\Big|_V > 0 ,$$

the adiabatic compressibility

$$\chi_S = -\frac{1}{V} \frac{\partial V}{\partial p}\Big|_S > 0$$

and the isothermal compressibility

$$\chi_T = -\frac{1}{V} \frac{\partial V}{\partial p}\Big|_T > 0 .$$

Proof. For the first, using (3.1) and $T > 0$

$$C_V = T \frac{\partial S}{\partial T}\Big|_V > 0 .$$

For the second, using (3.1) and $V > 0$

$$\chi_S = -\frac{1}{V} \frac{\partial V}{\partial p} \Big|_S > 0 .$$

For the third, using (3.1) and (??)

$$\begin{aligned} 0 &< \frac{\partial T}{\partial V} \Big|_S \frac{\partial T}{\partial V} \Big|_S + \frac{\partial T}{\partial S} \Big|_V \frac{\partial p}{\partial V} \Big|_S \\ &\quad - \frac{\partial T}{\partial V} \Big|_S \frac{\partial p}{\partial S} \Big|_V + \frac{\partial T}{\partial S} \Big|_V \frac{\partial p}{\partial V} \Big|_S \\ &= \frac{\partial(T, p)}{\partial(S, V)} \\ &= \frac{\partial(T, p)}{\partial(S, V)} = \frac{\partial(T, p)}{\partial(T, V)} \frac{\partial(T, V)}{\partial(S, V)} \\ &= \frac{\partial p}{\partial V} \Big|_T \frac{\partial T}{\partial S} \Big|_V \\ &= \frac{T}{C_V} \frac{\partial p}{\partial V} \Big|_T , \end{aligned}$$

hence, by $T > 0$, $C_V > 0$ and $V > 0$,

$$\chi_T = -\frac{1}{V} \frac{\partial V}{\partial p} \Big|_T > 0 .$$

q.e.d.

Consequently to stability, F is a concave of T and covex of V , whereas G is concave of both T and p .

Proof. For the concavity of F of T

$$C_V = T \frac{\partial S}{\partial T} \Big|_V = -T \frac{\partial^2 F}{\partial T^2} \Big|_V > 0 ,$$

hence

$$\frac{\partial^2 F}{\partial T^2} \Big|_V < 0 .$$

For the convexity of F of V

$$\chi_T = -\frac{1}{V} \frac{\partial V}{\partial p} \Big|_T = \left(V \frac{\partial^2 F}{\partial V^2} \Big|_T \right)^{-1} > 0 ,$$

hence

$$\frac{\partial^2 F}{\partial V^2} \Big|_T > 0 .$$

For the concavity of G of T

$$C_P = T \frac{\partial S}{\partial T} \Big|_P = -T \frac{\partial^2 G}{\partial T^2} \Big|_p > 0 ,$$

hence

$$\frac{\partial^2 G}{\partial T^2} \Big|_p < 0 .$$

For the concavity of G of p

$$\chi_T = -\frac{1}{V} \frac{\partial V}{\partial p} \Big|_T = -\frac{1}{V} \frac{\partial^2 G}{\partial p^2} \Big|_T > 0 ,$$

hence

$$\frac{\partial^2 G}{\partial p^2} \Big|_T < 0 .$$

q.e.d.

Furthermore, the second law of thermodynamics can be expressed, in order to maximise the entropy, by imposing that first derivatives vanish and the hessian, i.e. the matrix with its second derivatives, must be negative defined. Therefore, it must be (locally) concave in E , V and N .

When cease to work at constant N , the stability condition is

$$\frac{\partial N}{\partial \mu} \Big|_{V,T} = \frac{N^2}{V} \chi_T > 0 .$$

Proof. In fact

$$\begin{aligned} \frac{\partial N}{\partial \mu} \Big|_{V,T} &= \frac{\partial(N, V, T)}{\partial(\mu, V, T)} \\ &= \frac{\partial(N, V, T)}{\partial(N, p, T)} \frac{\partial(N, p, T)}{\partial(p, V, T)} \frac{\partial(p, V, T)}{\partial(\mu, V, T)} \\ &= \frac{\partial(N, V, T)}{\partial(N, p, T)} \frac{\partial(N, p, T)}{\partial(p, V, T)} \frac{\partial(p, V, T)}{\partial(\mu, N, T)} \frac{\partial(p, V, T)}{\partial(\mu, V, T)} \\ &= \frac{\partial(N, V, T)}{\partial(N, p, T)} \frac{\partial(N, p, T)}{\partial(\mu, N, T)} \frac{\partial(p, V, T)}{\partial(\mu, V, T)} \\ &= -\frac{\partial V}{\partial p} \Big|_{N,T} \frac{\partial p}{\partial \mu} \Big|_{V,T} \frac{\partial p}{\partial \mu} \Big|_{N,T} \end{aligned}$$

Now we use (2.8)

$$\frac{\partial p}{\partial \mu} \Big|_{V,T} = \frac{\partial p}{\partial \mu} \Big|_{N,T} = \left(\frac{\partial \mu}{\partial p} \Big|_T \right) = \frac{N}{V} ,$$

hence

$$\frac{\partial N}{\partial \mu} \Big|_{V,T} = -\frac{N^2}{V^2} \frac{\partial V}{\partial p} \Big|_{N,T} = \frac{N^2}{V} \chi_T > 0 .$$

q.e.d.

Moreover, we have the relation

$$\chi_T(C_P - C_V) = TV\alpha_p^2 ,$$

which implies that

$$C_P > C_V \iff \chi_T > \chi_S .$$

Proof. We start from

$$\begin{aligned} C_V &= T \frac{\partial S}{\partial T} \Big|_V = \frac{\partial E}{\partial T} \Big|_V , \\ C_p &= T \frac{\partial S}{\partial T} \Big|_p = \frac{\partial E}{\partial T} \Big|_p + p \frac{\partial V}{\partial T} \Big|_p , \end{aligned}$$

which imply that

$$\begin{aligned} TdS &= C_V dT + \left(\frac{\partial E}{\partial V} \Big|_T + p \right) dV = C_V dT + T \frac{\partial p}{\partial T} \Big|_V dV , \\ TdS &= C_p dT + \left(\frac{\partial E}{\partial p} \Big|_T + \frac{\partial V}{\partial p} \Big|_T \right) dp = C_p dT - T \frac{\partial V}{\partial T} \Big|_p dp . \end{aligned}$$

Comparing them

$$\begin{aligned} (C_p - C_V) dT &= T \left(\frac{\partial V}{\partial T} \Big|_p dp + \frac{\partial p}{\partial T} \Big|_V dV \right) , \\ (C_p - C_V) &= T \frac{\partial V}{\partial T} \Big|_p \frac{\partial p}{\partial T} \Big|_V . \end{aligned}$$

We use

$$\frac{\partial p}{\partial T} \Big|_V = \frac{\partial(p, V)}{\partial T, V} = \frac{\partial(p, V)}{\partial(p, T)} \frac{\partial(p, T)}{\partial(T, V)} = -\frac{\partial V}{\partial T} \Big|_p \frac{\partial p}{\partial V} \Big|_T ,$$

hence

$$C_p - C_V = \frac{T}{V\chi_T} \left(\frac{\partial V}{\partial T} \Big|_p \right)^2 ,$$

or, defining the thermal expansion coefficient

$$\alpha_p = \frac{1}{V} \frac{\partial V}{\partial T} \Big|_p , \tag{3.2}$$

we have

$$\chi_T(C_P - C_V) = TV\alpha_p^2 ,$$

Finally, we obtain

$$\frac{C_p}{C_V} = \frac{\chi_T}{\chi_S} .$$

q.e.d.

3.3 Statistical mechanics

It is important to say that this is all thermodynamics can tell us, thus in order to find the explicit expression of E , we must go into statistical mechanics.

Part II

Classical statistical mechanics

Chapter 4

Classical mechanics

A state constituted by a system of N particles is described by a point in a $2N$ -dimensional manifold \mathcal{M}^N , called the phase space, which is the Cartesian product of N single particle manifolds

$$\{(q^i, p_i)\} \in \mathcal{M}^N$$

where $i = 1, \dots, N$.

An observable is a smooth real function

$$f : \mathcal{M}^N \rightarrow \mathbb{R}$$

and its measurement in a fixed point $(\tilde{q}^i, \tilde{p}_i)$ is its value in it

$$f = f(\tilde{q}^i, \tilde{p}_i)$$

The time evolution is governed by a real function, called hamiltonian $H(q^i, p_i, t)$, which is the solution of the equations of motion, called the Hamilton's equations

$$\dot{q}^i = \frac{\partial H}{\partial p_i} \quad \dot{p}_i = -\frac{\partial H}{\partial q^i}$$

Theorem 4.1 (Conservation of energy)

If the hamiltonian does not depend explicitly on time, it can be interpreted physically as the energy of the system, which is constants

$$H(q^i(t), p_i(t)) = H(q^i(0), p_i(0)) = E = \text{const}$$

Since they are deterministic, once the initial conditions are given, the trajectory in phase space is completely determined

4.1 Probability density distribution

A macrostate is defined by setting the macroscopic thermodynamical quantities. A microstate is the knowledge of the phase space behaviour (q^i, p_i) .

In general, there are more microstates associated to the same macrostates, raising the concept of ensemble: fixing a macrostate, it is created a large number of copies of the same physical system but with different microstates. It can be studied with the introduction of a probability density distribution

$$\rho(q_i(t), p_i(t), t)$$

such that it satisfies the following properties

1. positivity, i.e.

$$\rho(q_i, p_i, t) \geq 0$$

2. normalisation, i.e.

$$\int_{\mathcal{M}^n} \underbrace{\prod_{i=1}^N d^d q^i d^d p^i}_{d\Gamma} \rho(q_i, p_i, t) = \int_{\mathcal{M}^n} d\Gamma \rho(q_i, p_i, t) = 1$$

To solve the dimensional problem of the volume element $d\Gamma$, which must be adimensional but it has the dimension of an action to the power of d , it can be introduced the adimensional volume element

$$d\Omega = \frac{d\Gamma}{h^d} = \frac{\prod_{i=1}^N d^d q^i d^d p^i}{h^d}$$

where the scale factor h has the dimension of an action.

The probability to find the system in a finite portion of the phase space $\mathcal{U} \subset \mathcal{M}^N$ is

$$\int_{\mathcal{U}} d\Gamma \rho(q_i, p_i, t)$$

4.2 Liouville's theorem

The flow of a system of particles keeps trasf of all their motions. See Figure.

Theorem 4.2 (Liouville)

The volume through the flow generated by the hamilton's equations is constant. See Figure. Mathematically

$$vol\Omega(t=0) = vol\Omega(t) \Rightarrow \frac{d\rho}{dt} = \frac{\partial\rho}{\partial t} + [\rho, H] = 0$$

Proof. Maybe in the future.

q.e.d.

The physical interpretation of this theorem is that particles do not appear nor disappear due to conservation of charge, mass, etc...

For stationary systems, i.e. when $\frac{\partial \rho}{\partial t} = 0$, the necessary condition for equilibrium is $[\rho, H] = 0$, which is satisfied only if

$$\rho = \text{const}$$

like in the microcanonical ensemble, and

$$\rho = \rho(H)$$

like in the canonical or the grandcanonical ensembles.

Proof. Maybe in the future.

q.e.d.

The average value of an observable is weighted by the probability density distribution

$$\langle f \rangle = \int_{\mathcal{M}^N} d\Gamma \rho(q^i, p_i) f(q^i, p_i) \quad (4.1)$$

and the standard deviation is

$$(\Delta f)^2 = \langle f^2 \rangle - \langle f \rangle^2$$

4.3 Time-independent Hamiltonian

Consider a time-independent hamiltonian. Since the energy is constant for the theorem 4.1.

$$\int_{\mathcal{M}^N} \rho = 1 \quad (4.2)$$

$$\frac{\partial S}{\partial E} = \frac{1}{T} \quad (4.3)$$

$$\frac{\partial F}{\partial T} = -S \quad (4.4)$$

$$\Omega = -pV = E - TS - \mu N \quad (4.5)$$

Chapter 5

Microcanonical ensemble

A microcanonical ensemble is a system which is isolated from the environment, i.e. it cannot exchange neither energy nor matter, so E , N and V are fixed. Since energy is conserved and the hamiltonian is time-independent, the trajectory of motion is restricted on the surface S_E and not on all the phase space.

Assume an a-priory uniform probability

$$\rho_{mc}(q^i, p_i) = C\delta(\mathcal{H}(q^i, p_i) - E)$$

where C is a normalisation constant, which can be evaluated by (4.2)

$$1 = \int_{\mathcal{M}^N} d\Omega \rho_{mc} = \int_{\mathcal{M}^N} d\Omega C\delta(\mathcal{H} - E) = C \int_{\mathcal{M}^N} d\Omega \delta(\mathcal{H} - E) = C\omega(E)$$

Hence

$$\rho_{mc}(q^i, p_i) = \frac{1}{\omega(E)}\delta(\mathcal{H}(q^i, p_i) - E)$$

Consider a displacement on an infinitesimal displacement of energy $\Delta E \ll 1$, then

$$\Gamma(E) = \int_E^{E+\Delta E} dE' \omega(E') \simeq \omega(E)\Delta E$$

and the distribution is

$$\rho_{mc}(q^i, p_i) = \begin{cases} \frac{1}{\Gamma(E)} & \mathcal{H} \in [E, E + \Delta E] \\ 0 & otherwise \end{cases}$$

Let $f(q^i, p_i)$ be an observable, then its microcanonical average is

$$\langle f(q^i, p_i) \rangle_{mc} = \int_{\mathcal{M}} d\Omega \rho_{mc} f = \int_{\mathcal{M}} d\Omega \frac{1}{\omega(E)}\delta(\mathcal{H} - E)f = \frac{1}{\omega(E)} \int_{S_E} dS_E f = \langle f \rangle_E \quad (5.1)$$

5.1 Thermodynamics potentials

The microcanonical entropy S_{mc} is defined by

$$S_{mc}(E, V, N) = k_B \ln \omega(E) \quad (5.2)$$

The logarithm is justified by the fact that the volume of a N -particle phase space is $(W_1)^N$, where W_1 is the volume of a single particle phase space. According to the properties of the logarithm, entropy becomes extensive.

In the thermodynamic limit, the following equations hold

$$s_{mc} = \lim_{td} \frac{S_{mc}}{N} = k_B \lim_{td} \frac{\log \omega(E)}{N} = k_B \underbrace{\lim_{td} \frac{\log \Sigma(E)}{N}}_{\mathcal{H} \in [0, E]} = k_B \underbrace{\lim_{td} \frac{\log \Gamma(E)}{N}}_{\mathcal{H} \in [E, E + \Delta E]}$$

Entropy is additive, so given two systems 1 and 2

$$s_{mc}^{tot} = s_{mc}^{(1)} + s_{mc}^{(2)}$$

Proof. Consider two isolated systems in contact at equilibrium with the same temperature $T = T_1 = T_2$. The total energy is $E = E_1 + E_2 + E_{surface}$ but, in the thermodynamic limit, the energy exchanged by the surface is a subleading term (E_1 and E_2 go as L^3 whereas $E_{surface}$ goes as L^2) and can be neglected. The energy density is

$$\begin{aligned} \omega(E) &= \int_{\mathcal{M}^N} d\Gamma_1 d\Gamma_2 \delta(\mathcal{H} - E) \\ &= \int dE_1 \int dS_{E_1} \int dE_2 \int dS_{E_2} \delta(E - E_1 - E_2) \\ &= \int dE_1 \int dE_2 \omega_1(E_1) \omega_2(E_2) \delta(E - E_1 - E_2) \\ &= \int_0^E dE_1 \omega_1(E_1) \omega_2(E_2 = E - E_1) \end{aligned}$$

Since the integrand is a positive function with a maximum in $E_1 \in [0, E]$

$$\begin{aligned} \int_0^E dE_1 \omega_1(E_1) \omega_2(E_2 = E - E_1) &\leq \omega_1(E_1^*) \omega_2(E_2^* = E - E_1^*) \int_0^E dE_1 \\ &= \omega_1(E_1^*) \omega_2(E_2^* = E - E_1^*) E \end{aligned} \quad (5.3)$$

On the other hand, it is always possible to find a value for ΔE in order to have

$$\Delta E \omega_1(E_1^*) \omega_2(E_2^*) \leq \omega(E) \quad (5.4)$$

Putting together (5.3) and (5.4)

$$\Delta E \omega_1(E_1^*) \omega_2(E_2^*) \leq \omega(E) \leq \omega_1(E_1^*) \omega_2(E_2^*) E$$

$$\omega_1(E_1^*)\Delta E\omega_2(E_2^*)\Delta E \leq \omega(E)\Delta E \leq \frac{E}{\Delta E}\omega_1(E_1^*)\Delta E\omega_2(E_2^*)\Delta E$$

$$\Gamma_1(E_1^*)\Gamma(E_2^*) \leq \Gamma(E) \leq \frac{E}{\Delta E}\Gamma(E_1^*)\Gamma(E_2^*)$$

Since the logarithm is a monotonic function

$$\log\left(\Gamma_1(E_1^*)\Gamma(E_2^*)\right) \leq \log\Gamma(E) \leq \log\left(\frac{E}{\Delta E}\Gamma(E_1^*)\Gamma(E_2^*)\right)$$

$$k_B \log\left(\Gamma_1(E_1^*)\Gamma(E_2^*)\right) \leq k_B \log\Gamma(E) \leq k_B \log\left(\frac{E}{\Delta E}\Gamma(E_1^*)\Gamma(E_2^*)\right)$$

$$\begin{aligned} k_B \log\Gamma_1(E_1^*) + k_B \log\Gamma(E_2^*) &\leq k_B \log\Gamma(E) \\ &\leq k_B \log\frac{E}{\Delta E} + k_B \log\Gamma(E_1^*) + k_B \log\Gamma(E_2^*) \end{aligned}$$

$$\begin{aligned} \frac{k_B \log\Gamma_1(E_1^*) + k_B \log\Gamma(E_2^*)}{N} &\leq \frac{k_B \log\Gamma(E)}{N} \\ &\leq \frac{k_B \log\frac{E}{\Delta E} + k_B \log\Gamma(E_1^*) + k_B \log\Gamma(E_2^*)}{N} \end{aligned}$$

In the thermodynamic limit, the last term vanishes, since $\lim_{td} \frac{1}{N} \log \frac{N}{\Delta N} = 0$. Hence

$$s_{mc}(E) = s_{mc}^{(1)} + s_{mc}^{(2)}$$

q.e.d.

The last result tells also that at equilibrium entropy is maximum.

In the thermodynamic limit, microcanonical entropy coincides with the thermodynamical one

$$s_{mc} = s_{td}$$

Proof. Since entropy is maximum at equilibrium, also $\Gamma_1(E_1)\Gamma_2(E_2)$ is so and

$$\begin{aligned} 0 &= \delta(\Gamma_1(E_1^*)\Gamma_2(E_2^* = E - E_1^*)) \\ &= \delta\Gamma_1(E_1^*)\Gamma_2(E_2^*) + \Gamma_1(E_1^*)\delta\Gamma_2(E_2^*) \\ &= \frac{\partial\Gamma_1}{\partial E_1}\Big|_{E_1^*} \delta E_1 \Gamma_2(E_2^*) + \Gamma_1(E_1^*) \frac{\partial\Gamma_2}{\partial E_2}\Big|_{E_2^*} \delta E_2 \end{aligned}$$

Since $E = \text{const}$, $0 = \delta E = \delta E_1 + \delta E_2$, $\delta E_2 = -\delta E_1$ and

$$\begin{aligned} 0 &= \frac{\partial\Gamma_1}{\partial E_1}\Big|_{E_1^*} \delta E_1 \Gamma_2(E_2^*) - \Gamma_1(E_1^*) \frac{\partial\Gamma_2}{\partial E_2}\Big|_{E_2^*} \delta E_1 \\ 0 &= \frac{\partial\Gamma_1}{\partial E_1}\Big|_{E_1^*} \Gamma_2(E_2^*) - \Gamma_1(E_1^*) \frac{\partial\Gamma_2}{\partial E_2}\Big|_{E_2^*} \end{aligned}$$

$$\begin{aligned}
\frac{\partial \Gamma_1}{\partial E_1} \Big|_{E_1^*} \Gamma_2(E_2^*) &= \Gamma_1(E_1^*) \frac{\partial \Gamma_2}{\partial E_2} \Big|_{E_2^*} \\
\frac{1}{\Gamma_1(E_1^*)} \frac{\partial \Gamma_1}{\partial E_1} \Big|_{E_1^*} &= \frac{1}{\Gamma_2(E_2^*)} \frac{\partial \Gamma_2}{\partial E_2} \Big|_{E_2^*} \\
\frac{\partial \log \Gamma_1}{\partial E_1} \Big|_{E_1^*} &= \frac{\partial \log \Gamma_2}{\partial E_2} \Big|_{E_2^*}
\end{aligned}$$

Using the thermodynamical relation (4.3)

$$S_{mc}(E) = S_{td}(E) \times \text{const}$$

where the constant can be chosen in order to have k_B in the same unit. q.e.d.

The universal Boltzmann's formula is

$$s_{mc} = s_{td} = k_B \log \omega(E) = -k_B \langle \log \rho_{mc} \rangle_{mc} . \quad (5.5)$$

Proof. Using (5.1),

$$\begin{aligned}
\langle \log \rho_{mc} \rangle_{mc} &= \int d\Gamma \rho_{mc} \log \rho_{mc} \\
&= \int d\Gamma \frac{1}{\omega(E)} \delta(\mathcal{H} - E) \log \left(\frac{1}{\omega(E)} \delta(\mathcal{H} - E) \right) \\
&= \int dS_E \frac{1}{\omega(E)} \log \frac{1}{\omega(E)} \\
&= -\frac{1}{\omega(E)} \log \omega(E) \int dS_E \\
&= -\log \omega(E)
\end{aligned}$$

q.e.d.

Chapter 6

Canonical ensemble

A canonical ensemble is a system which is immersed in a bigger environment or reservoir, which can exchange energy but not matter, so T , N and V are fixed. Globally, energy is conserved, since the universe composed by the union of the system and the environment can be considered as a microcanonical ensemble.

The canonical probability density distribution is

$$\rho_c(q^i, p_i) = \frac{1}{Z_N} \exp(-\beta \mathcal{H}(q^i, p_i)) \quad (6.1)$$

where β is

$$\beta = \frac{1}{k_B T}$$

and Z_N is the partition function

$$Z_N[V, T] = \int_{\mathcal{M}^N} d\Omega \exp(-\beta \mathcal{H}(q^i, p_i)) \quad (6.2)$$

which depends on the temperature through β and volume and temperature due to the integration domain $\mathcal{M}^N = V \otimes \mathbb{R}^d$.

Notice that the probability is a function of the hamiltonian, like Liouville's theorem said.

Proof. Consider the universe as a microcanonical ensemble. Its probability density distribution is

$$\rho_{mc}(q_i^{(1)}, p_i^{(1)}, q_i^{(2)}, p_i^{(2)}) = \frac{1}{\omega(E)} \delta(\mathcal{H}(q_i^{(1)}, p_i^{(1)}, q_i^{(2)}, p_i^{(2)}) - E)$$

where the total hamiltonian is

$$\mathcal{H}(q_i^{(1)}, p_i^{(1)}, q_i^{(2)}, p_i^{(2)}) = \mathcal{H}_1(q_i^{(1)}, p_i^{(1)}) + \mathcal{H}_2(q_i^{(2)}, p_i^{(2)})$$

Integrating it to all the possible state in the environment

$$\rho^{(1)} = \int d\Omega_2 \rho_{mc} = \int d\Omega_2 \frac{1}{\omega(E)} \delta(\mathcal{H} - E) = \frac{1}{\omega(E)} \int dS_{E_2} = \frac{1}{\omega(E)} \omega(E_2 = E - E_1)$$

and the corresponding entropy is

$$S_2(E_2) = k_B \ln \omega_2(E_2)$$

Applying small variation δE_1 to E_1 to preserve equilibrium, the entropy transforms, using (4.3)

$$k_B \ln \omega_2(E_2) = S_{mc}(E) - E_1 \left. \frac{\partial S_{mc}}{\partial E} \right|_{E_2} = S_{mc}(E) - E_1 \frac{1}{T}$$

$$\ln \omega_2(E_2) = \frac{S_{mc}(E)}{k_B} - E_1 \frac{1}{k_B T}$$

$$\omega_2(E_2) = \exp\left(\frac{S_{mc}(E)}{k_B} - E_1 \frac{1}{k_B T}\right) = \exp\left(\frac{S_{mc}(E)}{k_B}\right) \exp\left(-\frac{E_1}{k_B T}\right)$$

Putting together, dropping the indices

$$\rho_c = \frac{\omega(2)(E_2)}{\omega(E)} = \frac{1}{\omega(E)} \exp\left(\frac{S_{mc}(E)}{k_B}\right) \exp\left(-\frac{E_1}{k_B T}\right) = C \exp\left(-\frac{E_1}{k_B T}\right) \quad (6.3)$$

where C is a normalisation constant, which can be evaluated by (4.2)

$$1 = \int_{\mathcal{M}^N} d\Omega \rho = \int_{\mathcal{M}^N} d\Omega C \exp\left(-\frac{E_1}{k_B T}\right) = C \int_{\mathcal{M}^N} d\Omega \exp\left(-\frac{E_1}{k_B T}\right)$$

q.e.d.

The partition function can also be written as

$$Z_N[T, V] = \int_0^\infty dE \omega(E) \exp(-\beta E)$$

Proof. Foliating the phase space in energy hyper-surfaces

$$Z_N = \int_{\mathcal{M}^N} d\Omega \exp(-\beta \mathcal{H}) = \int_0^\infty dE \int dS_E \exp(-\beta \mathcal{H}) = \int_0^\infty dE \omega(E) \exp(-\beta E)$$

q.e.d.

Taking also in consideration indistinguishable particles, the partition function

$$Z_N = \int \frac{\prod_{i=1}^N d^d q^i d^d p^i}{h^{dN} \zeta_N} \exp(-\beta \mathcal{H})$$

where ζ_N is

$$\zeta_N = \begin{cases} 1 & \text{distinguishable} \\ N! & \text{indistinguishable} \end{cases}$$

The partition function of two systems is the multiplication of the single system ones

$$Z_N = Z_{N_1} Z_{N_2} \quad (6.4)$$

Proof. Since $\mathcal{H} = \mathcal{H}_1 + \mathcal{H}_2$,

q.e.d.

If the hamiltonian is the sum of N identical ones, like N non-interacting particles

$$\mathcal{H} = \sum_{i=1}^N \mathcal{H}_i$$

the partition function becomes

$$Z_N = \frac{(Z_1)^N}{\zeta_N}$$

Proof. Denominating Z_1 the single-particle partition function

$$\begin{aligned} Z_N &= \int_{\mathcal{M}^N = \mathcal{M}^{(1)} \otimes \dots \otimes \mathcal{M}^{(1)}} \prod_{i=1}^N \frac{d^d q^i d^d p^i}{h^{dN} \zeta_N} \exp(-\beta \mathcal{H}) \\ &= \int_{\mathcal{M}^N = \mathcal{M}^{(1)} \otimes \dots \otimes \mathcal{M}^{(1)}} \prod_{i=1}^N \frac{d^d q^i d^d p^i}{h^{dN} \zeta_N} \exp(-\beta \sum_{i=1}^N \mathcal{H}_i) \\ &= \int_{\mathcal{M}^N = \mathcal{M}^{(1)} \otimes \dots \otimes \mathcal{M}^{(1)}} \prod_{i=1}^N \frac{d^d q^i d^d p^i}{h^{dN} \zeta_N} \prod_{i=1}^N \exp(-\beta \mathcal{H}_i) \\ &= \int_{\mathcal{M}^N = \mathcal{M}^{(1)} \otimes \dots \otimes \mathcal{M}^{(1)}} \prod_{i=1}^N \frac{d^d q^i d^d p^i}{h^{dN} \zeta_N} \exp(-\beta \mathcal{H}_i) \\ &= \frac{Z_1 Z_1 \dots Z_1}{\zeta_N} = \frac{(Z_1)^N}{\zeta_N} \end{aligned}$$

q.e.d.

Let $f(q^i, p_i)$ be an observable, then its canonical average is

$$\langle f(q^i, p_i) \rangle_c = \int_{\mathcal{M}} d\Omega \rho_c f = \int_{\mathcal{M}} d\Omega \frac{\exp(-\beta \mathcal{H})}{Z_N} f$$

6.1 Thermodynamics variable

The canonical Helmotz free energy F is defined by

$$Z[V, T] = \exp(-\beta F[N, V, T]) \quad (6.5)$$

or, equivalently,

$$F[V, N, T] = -\frac{1}{\beta} \ln Z_N \quad (6.6)$$

Furthermore, the canonical internal energy is

$$E = \langle \mathcal{H} \rangle_c = \int d\Omega \frac{\exp(-\beta(\mathcal{H}))}{Z_N} \mathcal{H} \quad (6.7)$$

Proof. By normalisation condition

$$1 = \int d\Omega \frac{\exp(-\beta\mathcal{H})}{Z_N} = \int d\Omega \frac{\exp(-\beta\mathcal{H})}{\exp(-\beta F)} = \int d\Omega \exp(-\beta(\mathcal{H} - F))$$

Since F depends on the temperature, it is possible to derive with respect to β

$$\begin{aligned} 0 &= \frac{\partial}{\partial \beta} \left(\int d\Omega \exp(-\beta(\mathcal{H} - F)) \right) \\ &= \int d\Omega \exp(-\beta(\mathcal{H} - F)) \left(-(\mathcal{H} - F) + \beta \frac{\partial F}{\partial \beta} \right) \\ &= - \underbrace{\int d\Omega \frac{\exp(-\beta\mathcal{H})}{Z_N} \mathcal{H}}_E + F \underbrace{\int d\Omega \frac{\exp(-\beta\mathcal{H})}{Z_N}}_1 + \beta \frac{\partial F}{\partial \beta} \underbrace{\int d\Omega \frac{\exp(-\beta\mathcal{H})}{Z_N}}_1 \\ &= -E + F + \beta \frac{\partial F}{\partial \beta} \end{aligned}$$

Hence, using (4.4)

$$F = E + \beta \frac{\partial F}{\partial \beta} = E + T \frac{\partial F}{\partial T} = E - TS$$

showing that is indeed the Helmotz free energy.

q.e.d.

Notice that in the last result, the entropy can be also written as

$$S_c = \frac{E - F}{T} \quad (6.8)$$

The internal energy can also be written as

$$E = -\frac{\partial}{\partial \beta} \ln Z_N \quad (6.9)$$

Proof. Using (6.7),

$$-\frac{\partial}{\partial \beta} \ln Z_N = -\frac{1}{Z_N} \frac{\partial Z_N}{\partial \beta} = -\frac{1}{Z_N} \frac{\partial}{\partial \beta} \int d\Omega \exp(-\beta \mathcal{H}) = \int d\Omega \frac{\exp(-\beta \mathcal{H})}{Z_N} \mathcal{H} = \langle \mathcal{H} \rangle_c = E$$

q.e.d.

The universal Boltzmann's formula is still valid

$$S_c = -k_B \langle \ln \rho_c \rangle_c$$

Proof. Using (6.7) and (6.6)

$$\begin{aligned} -k_B \langle \ln \rho_c \rangle_c &= -k_B \int d\Omega \rho_c \ln \rho_c \\ &= -k_B \int d\Omega \rho_c \ln \frac{\exp(-\beta \mathcal{H})}{Z_N} \\ &= -k_B \int d\Omega \rho_c \ln \exp(-\beta \mathcal{H}) - k_B \int d\Omega \rho_c \ln Z_N \\ &= k_B \int d\Omega \beta \mathcal{H} - k_B \underbrace{\ln Z_N}_{\beta F} \underbrace{\int d\Omega \rho_c}_1 \\ &= \frac{E - F}{T} = S_c \end{aligned}$$

q.e.d.

6.2 Equipartition theorem

Theorem 6.1 (Generalised equipartition theorem)

Let $\xi \in [a, b]$ and ξ_j with $j \neq 1$ all the other coordinates or momenta. Suppose also

$$\int \prod_{j \neq 1} d\xi_j [\xi_1 \exp(-\beta \mathcal{H})]_a^b = 0 \quad (6.10)$$

Then

$$\langle \xi_1 \frac{\partial \mathcal{H}}{\partial \xi_1} \rangle_c = k_B T \quad (6.11)$$

Proof. By normalisation condition

$$1 = \int d\Omega \frac{\exp(-\beta \mathcal{H})}{Z_N} = \frac{1}{Z_N} \int \prod_{j \neq 1} d\xi_j \exp(-\beta \mathcal{H})$$

Using

$$d\xi_1(\xi_1 \exp(-\beta\mathcal{H})) = d\xi_1 \exp(-\beta\mathcal{H}) + \xi_1 \exp(-\beta\mathcal{H})(-\beta) \frac{\partial\mathcal{H}}{\partial\xi_1} d\xi_1$$

and integrating per parts

$$\begin{aligned} 1 &= \frac{1}{Z_N} \int \underbrace{\prod_{j \neq 1} d\xi_j [\xi_1 \exp(-\beta\mathcal{H})]_a^b}_0 + \frac{\beta}{Z_N} \int \prod_{j \neq 1} d\xi_j d\xi_1 \xi_1 \frac{\partial\mathcal{H}}{\partial\xi_1} \exp(-\beta\mathcal{H}) \\ &= \beta \int d\Omega_{\xi_1} \frac{\partial\mathcal{H}}{\partial\xi_1} \frac{\exp(-\beta\mathcal{H})}{Z_N} \\ &= \beta \langle \xi_1 \frac{\partial\mathcal{H}}{\partial\xi_1} \rangle_c \end{aligned}$$

Hence

$$\langle \xi_1 \frac{\partial\mathcal{H}}{\partial\xi_1} \rangle_c = \frac{1}{\beta} = k_B T$$

q.e.d.

Examples of system that satisfies the condition (6.10) are hamiltonians which depend on the square of momentum or confining potentials which go to infinity on the extremes a and b .

Corollary 6.1 (Equipartition theorem)

If ξ_1 appears quadratically in \mathcal{H} , then its contribution to E is $\frac{1}{2}k_B T$

Proof. Consider $\mathcal{H} = A\xi_1^2 + B\xi_j^2$ with $j \neq 1$, then by the previous theorem

$$\langle \xi_1 \frac{\partial\mathcal{H}}{\partial\xi_1} \rangle_c = \langle \xi_1 2A\xi_1 \rangle_c = k_B T$$

and

$$\langle A\xi_1^2 \rangle_c = \frac{1}{2}k_B T$$

q.e.d.

Chapter 7

Grancanonical ensemble

A grancanonical ensemble is a system which is immersed in a bigger environment or reservoir, which can exchange both energy and matter, so T , and V are fixed. Globally, both energy and number of particles are conserved, since the universe composed by the union of the system and the environment can be considered as a microcanonical ensemble. First, with the same method used in the previous chapter, microcanonical can be transformed into canonical. Now, the universe is canonical and, globally, the number of particles is conserved.

The grancanonical probability density distribution is

$$\rho_{gc}(q^i, p_i) = \frac{\exp(-\beta\mathcal{H}_1)}{N_1!h^{dN_1}} \frac{Z_{N_2}[T, V_2]}{Z_N[T, V]}$$

Proof. Consider the universe as a canonical ensemble. Its probability density distribution is

$$\rho_c(q_i^{(1)}, p_i^{(1)}, q_i^{(2)}, p_i^{(2)}) = \frac{\exp(-\beta\mathcal{H}(q_i^{(1)}, p_i^{(1)}, q_i^{(2)}, p_i^{(2)}))}{Z_N[T, V]}$$

where the total hamiltonian is

$$\mathcal{H}(q_i^{(1)}, p_i^{(1)}, q_i^{(2)}, p_i^{(2)}) = \mathcal{H}_1(q_i^{(1)}, p_i^{(1)}) + \mathcal{H}_2(q_i^{(2)}, p_i^{(2)})$$

Integrating it to all the possible state in the environment

$$\begin{aligned} \rho^{(1)} &= \int d\Omega_2 \rho_c \\ &= \int \prod_{i=1}^N \frac{d^d q_i^{(2)} d^d p_i^{(2)}}{N!h^{dN}} \frac{\exp(-\beta(\mathcal{H}_1 + \mathcal{H}_2))}{Z_N} \\ &= \frac{\exp(-\beta\mathcal{H}_1)}{N_1!h^{dN_1}} \frac{1}{Z_N} \int \prod_{i=1}^N \frac{d^d q_i^{(2)} d^d p_i^{(2)}}{N_2!h^{dN_2}} \exp(-\beta\mathcal{H}_2) \\ &= \frac{\exp(-\beta\mathcal{H}_1)}{N_1!h^{dN_1}} \frac{Z_{N_2}[T, V_2]}{Z_N[T, V]} \end{aligned}$$

q.e.d.

The normalisation condition becomes

$$\sum_{N_1=0}^N \int_{\mathcal{M}^{N_1}} d\Omega_1 \rho_{gc} = 1$$

Proof. Using the expression to evaluate the power of a sum

$$(a+b)^n = \sum_{i=1}^n \binom{n}{i} a^i b^{n-i}$$

and

$$\begin{aligned} \int_{\mathcal{M}^{N_1}} d\Omega_1 \rho_{gc} &= \int_{\mathcal{M}^{N_1}} d\Omega_1 \frac{\exp(-\beta\mathcal{H}_1)}{N_1! h^{dN_1}} \frac{Z_{N_2}[T, V_2]}{Z_N[T, V]} \\ &= \frac{N!}{N_1! N_2} \frac{\int_{\mathcal{M}^{N_1}} d\Omega_1 \exp(-\beta\mathcal{H}_1) \int_{\mathcal{M}^{N_2}} d\Omega_2 \exp(-\beta\mathcal{H}_2)}{\int_{\mathcal{M}^N} d\Omega \exp(-\beta\mathcal{H})} \\ &= \frac{N!}{N_1! N_2} \frac{\frac{\int_{\mathcal{M}^{N_1}} d\Omega_1 \exp(-\beta\mathcal{H}_1)}{(V_1)^{N_1}} \frac{\int_{\mathcal{M}^{N_2}} d\Omega_2 \exp(-\beta\mathcal{H}_2)}{(V_2)^{N_2}}}{\frac{\int_{\mathcal{M}^N} d\Omega \exp(-\beta\mathcal{H})}{V^N}} \frac{(V_1)^{N_1} (V_2)^{N_2}}{V^N} \end{aligned}$$

which in the thermodynamical limit

$$\lim_{td} \frac{\frac{\int_{\mathcal{M}^{N_1}} d\Omega_1 \exp(-\beta\mathcal{H}_1)}{(V_1)^{N_1}} \frac{\int_{\mathcal{M}^{N_2}} d\Omega_2 \exp(-\beta\mathcal{H}_2)}{(V_2)^{N_2}}}{\frac{\int_{\mathcal{M}^N} d\Omega \exp(-\beta\mathcal{H})}{V^N}} = 1$$

Hence

$$\int_{\mathcal{M}^{N_1}} d\Omega_1 \rho_{gc} = \frac{N!}{N_1! N_2} \frac{(V_1)^{N_1} (V_2)^{N_2}}{V^N}$$

and the normalisation condition becomes, using $N = N_1 + N_2$,

$$\sum_{N_1=0}^N \int_{\mathcal{M}^{N_1}} d\Omega_1 \rho_{gc} = \sum_{N_1=0}^N \frac{N!}{N_1! N_2!} \frac{(V_1)^{N_1} (V_2)^{N_2}}{V^N} = \sum_{N_1=0}^N \binom{N}{N_1} \left(\frac{V_1}{V}\right)^{N_1} \left(\frac{V_2}{V}\right)^{N-N_1} = \left(\frac{V_1 + V_2}{V}\right)^N$$

which in the thermodynamical limit is

$$\lim_{td} \left(\frac{V_1 + V_2}{V}\right)^N = 1$$

q.e.d.

7.1 Thermodynamical potentials

The grancanonical probability density distribution can be also written as

$$\rho_{gc}(q_i, p_i) = \frac{\exp(-\beta(\mathcal{H}(q_i, p_i) - \mu N))}{\mathcal{Z}}$$

where μ is the chemical potential and \mathcal{Z} is the grancanonical partition function

$$\mathcal{Z} = \sum_{N=0}^{\infty} z^N Z_N = \exp(-\beta\Omega)$$

where $z = \exp(\beta\mu)$ is the fugacity and Ω is the granpotential.

Proof. Using (6.5) and Taylor expanding to first order in $N_1 \ll N$ and $V_1 \ll V$,

$$\begin{aligned} \frac{Z_{N_2}[T, V]}{Z_N[T, V]} &= \frac{\exp(-\beta F(T, N_2, V_2))}{\exp(-\beta F(T, N, V))} \\ &= \exp(-\beta(F(T, N - N_1, V - V_1) - F(T, N, V))) \\ &\simeq \exp(-\beta(\underbrace{\left.\frac{\partial F}{\partial N}\right|_{T, V}}_{\mu}(-N_1) + \underbrace{\left.\frac{\partial F}{\partial V}\right|_{T, N}}_{-p}(-V_1))) \\ &= \exp(-\beta(-\mu N_1 + pV_1)) \end{aligned}$$

Hence, now all the degrees of freedom of the environment has been removed

$$\begin{aligned} \rho_{gc} &= \frac{\exp(\beta\mathcal{H})}{N!h^{dN}} \exp(-\beta(-\mu N + pV)) \\ &= \frac{\exp(\beta\mathcal{H})}{N!h^{dN}} \underbrace{\exp(\beta\mu)^N}_{z^N} \exp(-\beta pV) \\ &= \frac{z^N \exp(\beta\mathcal{H})}{N!h^{dN}} \exp(-\beta pV) \end{aligned}$$

where we introduced the fugacity.

Recall (4.5), the normalisation condition becomes

$$\begin{aligned}
1 &= \sum_{N=0}^{\infty} \int_{\mathcal{M}^N} d\Omega \rho_{gc} \\
&= \sum_{N=0}^{\infty} \int_{\mathcal{M}^N} d\Omega \frac{z^N \exp(\beta \mathcal{H})}{N! h^{dN}} \exp(-\beta pV) \\
&= \exp(-\beta pV) \sum_{N=0}^{\infty} z^N \frac{\int_{\mathcal{M}^N} d\Omega}{h^{dN} N!} \\
&= \exp(-\beta pV) \underbrace{\sum_{N=0}^{\infty} z^N Z_N}_{\mathcal{Z}} \\
&= \exp(-\beta pV) \mathcal{Z}
\end{aligned}$$

Hence

$$\mathcal{Z} = \sum_{N=0}^{\infty} z^N Z_N = \exp(\beta pV)$$

and

$$\rho_{gc}(q_i, p_i) = \frac{\exp(-\beta(\mathcal{H}(q_i, p_i) - \mu N))}{\mathcal{Z}} = \frac{\exp(-\beta \mathfrak{H}(q_i, p_i))}{\mathcal{Z}}$$

where $\mathfrak{H} = \mathcal{H} - \mu N$ is the grancanonical hamiltonian.

q.e.d.

Let $f(q^i, p_i)$ be an observable, then its grancanonical average is

$$\begin{aligned}
\langle f(q^i, p_i) \rangle_{gc} &= \sum_{N=0}^{\infty} \int_{\mathcal{M}} d\Omega \rho_{gc} f_N \\
&= \sum_{N=0}^{\infty} \int_{\mathcal{M}} d\Omega \frac{\exp(-\beta(\mathcal{H} - \mu N))}{\mathcal{Z}} f_N \\
&= \frac{1}{\mathcal{Z}} \sum_{N=0}^{\infty} z^N Z_N \int_{\mathcal{M}} d\Omega \frac{\exp(-\beta \mathcal{H})}{Z_N} f_N \\
&= \frac{1}{\mathcal{Z}} \sum_{N=0}^{\infty} z^N Z_N \langle f_N \rangle_c
\end{aligned}$$

The grancanonical internal energy is

$$E = - \frac{\partial}{\partial \beta} \ln \mathcal{Z} \Big|_z \quad (7.1)$$

Proof.

$$\begin{aligned}
-\frac{\partial}{\partial \beta} \ln \mathcal{Z} \Big|_z &= -\frac{1}{\mathcal{Z}} \frac{\partial}{\partial \beta} \mathcal{Z} \Big|_z \\
&= -\frac{1}{\mathcal{Z}} \frac{\partial}{\partial \beta} \sum_{N=0}^{\infty} z^N Z_N \Big|_z \\
&= -\sum_{N=0}^{\infty} \frac{z^N}{\mathcal{Z}} \frac{\partial}{\partial \beta} \int d\Omega \exp(-\beta \mathcal{H}) \\
&= \sum_{N=0}^{\infty} \int d\Omega \frac{\exp(-\beta(\mathcal{H} + \mu N))}{\mathcal{Z}} \mathcal{H} \\
&= \langle \mathcal{H} \rangle_{gc} = E
\end{aligned}$$

q.e.d.

The grancanonical number of particles is

$$\langle N \rangle_{gc} = z \frac{\partial}{\partial z} \ln \mathcal{Z} \Big|_T \quad (7.2)$$

Proof.

$$\begin{aligned}
z \frac{\partial}{\partial z} \ln \mathcal{Z} \Big|_T &= \frac{z}{\mathcal{Z}} \frac{\partial}{\partial z} \mathcal{Z} \Big|_T \\
&= \frac{z}{\mathcal{Z}} \frac{\partial}{\partial z} \sum_{N=0}^{\infty} z^N Z_N \\
&= \frac{z}{\mathcal{Z}} \sum_{N=0}^{\infty} N z^{N-1} Z_N \\
&= \sum_{N=0}^{\infty} z^N Z_N N = \langle N \rangle_{gc}
\end{aligned}$$

q.e.d.

The grancanonical potential is

$$\Omega = -\frac{1}{\beta} \ln \mathcal{Z} \quad (7.3)$$

The universal Boltzmann's formula is still valid

$$S_{gc} = -k_B \langle \ln \rho_{gc} \rangle_{gc}$$

Proof. Using (7.3),

$$\begin{aligned}
-k_B \langle \ln \rho_{gc} \rangle_{gc} &= -k_B \int d\Omega \rho_{gc} \ln \rho_{gc} \\
&= -k_B \sum_{N=0}^{\infty} \frac{z^N}{\mathcal{Z}} \int d\Omega \exp(-\beta \mathcal{H}) \ln \rho_{gc} \\
&== -k_B \sum_{N=0}^{\infty} \frac{z^N}{\mathcal{Z}} \int d\Omega \exp(-\beta \mathcal{H}) (-\beta \mathcal{H} + \beta \mu N + \ln \mathcal{Z}) \\
&= k_B \beta \underbrace{\sum_{N=0}^{\infty} \frac{z^N}{\mathcal{Z}} \int d\Omega \exp(-\beta \mathcal{H}) \mathcal{H}}_E - k_B \beta \mu \underbrace{\sum_{N=0}^{\infty} \frac{z^N}{\mathcal{Z}} \int d\Omega \exp(-\beta \mathcal{H}) N}_N \\
&\quad + k_B \ln \mathcal{Z} \underbrace{\sum_{N=0}^{\infty} \frac{z^N}{\mathcal{Z}} \int d\Omega \exp(-\beta \mathcal{H})}_1 \\
&= \frac{E - \mu N - \Omega}{T} = S
\end{aligned}$$

q.e.d.

Chapter 8

Entropy

The Boltzmann's universal law allows us to define entropy in terms of number of states

$$S = -k_B \langle \ln \rho \rangle = k_B \ln \Sigma = \lim_{TD} S_{TD}$$

Thermodynamics tells us that equilibrium corresponds to maximum entropy.

We consider a canonical ensemble with a discrete set of energy values, but it can be generalised for grandcanonical and continuous energy levels. Therefore, the probability density distribution is (6.1)

$$\rho_c(E_r) = \frac{\exp(-\beta E_r)}{Z_N}$$

where the canonical partition function (6.2) becomes

$$Z_N = \int_{\mathcal{M}^N} d\Omega \exp(-\beta \mathcal{H}(q^i, p_i)) = \int_0^\infty dE \int_{S_E} dS_E \exp(-\beta E) \simeq \sum_{r=1}^p g_r \exp(-\beta E_r)$$

where we foliated \mathcal{M}^N in energy surfaces S_E and g_r is the multiplicity or degeneracy, i.e. how many levels have the same energy.

So far, we have started from an a-priori probability density distribution and from it derive the entropy. From now on, we will change the picture and do the converse: the probability distribution is the one corresponding to maximum entropy, given the macroscopic constraints. To do so, we introduce the Shannon's information entropy

$$H = - \sum_{i=1}^N p_i \ln p_i$$

which is the only function with the following properties for a random variable x such that it has N possible outcomes x_i with probability p_i

1. it is continuous with p_i ,

2. is monotonically increasing with N ,
3. it is invariant under compositions of subsystems, i.e. change how we collect in group.

Inference problem

Given a certain constraint for a function $\langle f \rangle$, what is the expectation value for another function g ? The answer can be found with the principle of maximum entropy, subjected to Lagrange multipliers given by the constraints

$$\sum_{i=1}^N p_i = 1 \quad \sum_{i=1}^N p_i f(x_i) = \langle f(x) \rangle$$

Hence, the problem reduces to maximise the function

$$H = - \sum_{i=1}^N p_i \ln p_i + \alpha \left(\sum_{i=1}^N p_i - 1 \right) + \beta \left(\sum_{i=1}^N p_i f(x_i) - \langle f \rangle \right) \quad (8.1)$$

In particular, we need to count the number of ways $W_{\{n_r\}}$ we can find n_r systems with energy E_r , given a set of discrete energy levels E_r , each of degeneracy g_r on which we distribute n_r particles. Hence, the probability density distribution n_r^* is the one which maximises (8.1), with entropy

$$S = \ln W_{\{n_r\}}$$

and the constrains

$$N = \sum_r n_r \quad E = \sum_r n_r E_r$$

In order to count $W_{\{n_r\}}$, we need to take into account distinguishability or not of particles. Therefore, we decomposed it into

$$W_{\{n_r\}} = W_{\{n_r\}}^{(1)} W_{\{n_r\}}^{(2)}$$

where $W_{\{n_r\}}^{(1)}$ counts in how many we can put n_r particles in the energy level E_r and $W_{\{n_r\}}^{(2)}$ consider the degeneracy of these levels.

Boltzmann distribution

Bose-Einstein distribution

Fermi-Dirac distribution

Part III

Applications of classical statistical
mechanics

Chapter 9

Microcanonical ensemble

9.1 Non-relativistic ideal gas in d-dimensions

Consider a non-relativistic ideal (non-interacting) gas of N particles in an d -dimensional manifold with a finite volume V^N : $\mathcal{M}_N = V^N \times \mathbb{R}^{dN}$. Its hamiltonian is

$$H = \sum_i \frac{p_i^2}{2m} .$$

The number of states $\Sigma(E)$ is

$$\Sigma(E) = \frac{2V^N}{\xi_N dN \Gamma(dN/2)} \left(\frac{2\pi m E}{h^2} \right)^{dN/2} .$$

Proof. By definition,

$$\Sigma(E) = \int_{H(q_i, p_i) \leq E} d\Omega = \int_{H(q_i, p_i) \leq E} \frac{\prod_i d^d q_i d^d p_i}{h^{dN} \xi_N} = \frac{1}{h^{dN} \xi_N} \int_{H(q_i, p_i) \leq E} \prod_i d^d q_i d^d p_i .$$

From the energy,

$$H = \sum_i \frac{p_i^2}{2m} \leq E ,$$

$$\sum_i p_i^2 \leq 2mE .$$

Hence, by the volume of a dN -sphere of radius $\sqrt{2mE}$ (A.1),

$$\begin{aligned}
\Sigma(E) &= \frac{1}{h^{dN} \xi_N} \int_{\sum_i p_i^2 \leq 2mE} \prod_i d^d q_i d^d p_i \\
&= \frac{1}{h^{dN} \xi_N} \underbrace{\int_{V^N} \prod_i d^d q_i}_{V^N} \underbrace{\int_{\sum_i p_i^2 \leq 2mE} \prod_i d^d p_i}_{\frac{\pi^{dN/2} (2mE)^{dN/2}}{\Gamma(dN/2+1)}} \\
&= \frac{V^N}{h^{dN} \xi_N} \frac{\pi^{dN/2} (2mE)^{dN/2}}{\Gamma(dN/2+1)} \\
&= \left(\frac{2\pi m E}{h^2} \right)^{dN/2} \frac{2V^N}{\Gamma(dN/2) \xi_N dN} .
\end{aligned}$$

q.e.d.

The density state $\omega(E)$ is

$$\omega(E) = \frac{V^N}{\xi_N \Gamma(dN/2)} \left(\frac{2\pi m}{h^2} \right)^{\frac{dN}{2}} E^{dN/2-1} .$$

Proof. By definition,

$$\begin{aligned}
\omega(E) &= \frac{\partial \Sigma(E)}{\partial E} \\
&= \frac{\partial}{\partial E} \left(\frac{2\pi m E}{h^2} \right)^{dN/2} \frac{2V^N}{\Gamma(dN/2) \xi_N dN} \\
&= \left(\frac{2\pi m}{h^2} \right)^{dN/2} \frac{2V^N}{\Gamma(dN/2) \xi_N dN} \frac{\partial E^{dN/2}}{\partial E} \\
&= \left(\frac{2\pi m}{h^2} \right)^{dN/2} \frac{2V^N}{\Gamma(dN/2) \xi_N dN} \frac{dN}{2} E^{dN/2-1} \\
&= \frac{V^N}{\xi_N \Gamma(dN/2)} \left(\frac{2\pi m}{h^2} \right)^{\frac{dN}{2}} E^{dN/2-1} .
\end{aligned}$$

q.e.d.

Notice that

$$\omega(E) = \frac{dN}{2E} \Sigma(E) , \quad \Gamma(E) = \omega(E) \Delta E = \frac{dN}{2E} \Sigma(E) \Delta E . \quad (9.1)$$

As a consequence, in the thermodynamic limit, we have the following equivalent relations

$$\lim_{TD} \frac{\ln \Gamma(E)}{N} = \lim_{TD} \frac{\ln \omega(E)}{N} = \lim_{TD} \frac{\ln \Sigma(E)}{N} .$$

Proof. Observing (9.1), we find that the logarithmic expression differs only for factors $\ln \Delta E$ and $\ln \frac{dN}{2E}$, which are neglectible in the thermodynamic limit since they do not scale as N . q.e.d.

The entropy is

$$\frac{S}{k_B} = \ln \Gamma(E) = \ln \omega(E) + \ln \Delta E = \ln \Sigma(E) = \ln \Sigma(E) + \ln \frac{dN}{2E} + \ln \Delta E ,$$

In the thermodynamic limit, it becomes

$$S = k_B \begin{cases} \frac{d}{2}N + N \ln \left(V \left(\frac{4\pi m E}{dN h^2} \right)^{d/2} \right) & \text{for distinguishable particles} \\ \frac{d+2}{2}N + N \ln \left(\frac{V}{N} \left(\frac{4\pi m E}{dN h^2} \right)^{d/2} \right) & \text{for indistinguishable particles} \end{cases} .$$

Proof. By definition, using the Stirling approximation (B.1),

$$\begin{aligned} \frac{S}{k_B} &= \ln \Sigma(E) \\ &= \ln \left(\frac{2V^N}{\xi_N dN \Gamma(dN/2)} \left(\frac{2\pi m E}{h^2} \right)^{dN/2} \right) \\ &= \ln 2 + N \ln V - \ln \xi_N - \ln d - \ln N - \ln \Gamma(dN/2) + N \ln \left(\frac{2\pi m E}{h^2} \right)^{d/2} \\ &= N \ln V - \ln \xi_N - \underbrace{\ln \Gamma(dN/2)}_{\frac{dN}{2} \ln \frac{dN}{2} - \frac{dN}{2}} + N \ln \left(\frac{2\pi m E}{h^2} \right)^{d/2} \\ &= N \ln V - \ln \xi_N - \frac{dN}{2} \ln \frac{dN}{2} + \frac{dN}{2} + N \ln \left(\frac{2\pi m E}{h^2} \right)^{d/2} \\ &= N \ln V - \ln \xi_N - N \ln \left(\frac{dN}{2} \right)^{d/2} + \frac{dN}{2} + N \ln \left(\frac{2\pi m E}{h^2} \right)^{d/2} \\ &= -\ln \xi_N + \frac{dN}{2} + N \ln \left(V \left(\frac{4\pi m E}{dN h^2} \right)^{d/2} \right) . \end{aligned}$$

Now, we treat the distinguishable and indistinguishable case separately. For distinguishable particles $\xi_N = 1$, we find

$$\frac{S}{k_B} = -\ln 1 + \frac{dN}{2} + N \ln \left(V \left(\frac{4\pi m E}{dN h^2} \right)^{d/2} \right) = \frac{d}{2}N + N \ln \left(V \left(\frac{4\pi m E}{dN h^2} \right)^{d/2} \right) .$$

For indistinguishable particles $\xi_N = N!$, we find

$$\begin{aligned} \frac{S}{k_B} &= - \underbrace{\ln N!}_{N \ln N - N} + \frac{dN}{2} + N \ln \left(V \left(\frac{4\pi m E}{dN h^2} \right)^{d/2} \right) \\ &= -N \ln N + N + \frac{dN}{2} + N \ln \left(V \left(\frac{4\pi m E}{dN h^2} \right)^{d/2} \right) \\ &= \frac{d+2}{2}N + N \ln \left(\frac{V}{N} \left(\frac{4\pi m E}{dN h^2} \right)^{d/2} \right) . \end{aligned}$$

q.e.d.

The internal energy is

$$E = \frac{dNk_B T}{2} .$$

Proof. By (??)

$$\frac{1}{T} = \frac{\partial S}{\partial E} = k_B \frac{dN}{2} \frac{\partial}{\partial E} \ln E = k_B \frac{dN}{2E} ,$$

hence

$$E = \frac{dNk_B T}{2} .$$

q.e.d.

The equation of state is

$$pV = Nk_B T .$$

Proof. By (??)

$$\frac{p}{T} = \frac{\partial S}{\partial V} = k_B N \frac{\partial}{\partial V} \ln V = k_B \frac{N}{V} ,$$

hence

$$pV = Nk_B T .$$

q.e.d.

9.2 Non-relativistic ideal gas in 3-dimensions

Now, consider the case in which $d = 3$. The number of states $\Sigma(E)$ is

$$\Sigma(E) = \frac{2V^N}{\xi_N dN \Gamma(3N/2)} \left(\frac{2\pi m E}{h^2} \right)^{3N/2} .$$

The density state $\omega(E)$ is

$$\omega(E) = \frac{V^N}{\xi_N \Gamma(3N/2)} \left(\frac{2\pi m}{h^2} \right)^{\frac{3N}{2}} E^{3N/2-1} .$$

Notice that

$$\omega(E) = \frac{3N}{2E} \Sigma(E) , \quad \Gamma(E) = \omega(E) \Delta E = \frac{3N}{2E} \Sigma(E) \Delta E .$$

The entropy is

$$\frac{S}{k_B} = \ln \Gamma(E) = \ln \omega(E) + \ln \Delta E = \ln \Sigma(E) = \ln \Sigma(E) + \ln \frac{3N}{2E} + \ln \Delta E ,$$

In the thermodynamic limit, it becomes

$$S = k_B \begin{cases} \frac{3}{2}N + N \ln \left(V \left(\frac{4\pi m E}{3N h^2} \right)^{3/2} \right) & \text{for distinguishable particles} \\ \frac{5}{2}N + N \ln \left(\frac{V}{N} \left(\frac{4\pi m E}{3N h^2} \right)^{3/2} \right) & \text{for indistinguishable particles} \end{cases} .$$

The internal energy is

$$E = \frac{3}{2} N k_B T .$$

The equation of state is

$$pV = N k_B T .$$

9.3 Gas of harmonic oscillators in 3-dimensions

Consider a non-relativistic (non-interacting) gas of N particles in an d -dimensional manifold confined by an harmonic potential of frequency ω . Its hamiltonian is

$$H = \sum_i \left(\frac{p_i^2}{2m} + \frac{m\omega^2}{2} q_i^2 \right) .$$

The number of states $\Sigma(E)$ is

$$\Sigma(E) = \frac{1}{\xi_N \Gamma(dN/2) dN} \left(\frac{2\pi E}{h\omega} \right)^{dN} .$$

Proof. By definition,

$$\Sigma(E) = \int_{H(q_i, p_i) \leq E} d\Omega = \int_{H(q_i, p_i) \leq E} \frac{\prod_i d^d q_i d^d p_i}{h^{dN} \xi_N} = \frac{1}{h^{dN} \xi_N} \int_{H(q_i, p_i) \leq E} \prod_i d^d q_i d^d p_i .$$

We make a change of variable into x_j , with $j = 1, \dots, 2dN$,

$$p_i = \sqrt{2mE} x_j , \quad q_i = \sqrt{\frac{2E}{m\omega^2}} x_{dN+j} .$$

The differentials become

$$dp_i = \sqrt{2mE} dx_j , \quad dq_i = \sqrt{\frac{2E}{m\omega^2}} dx_{dN+j} .$$

From the energy,

$$H = \sum_i \left(\frac{p_i^2}{2m} + \frac{m\omega^2}{2} q_i^2 \right) \leq E ,$$

$$\sum_j x_j^2 \leq 1 .$$

Hence, by the volume of a $2dN$ -sphere of radius 1 (A.1),

$$\begin{aligned} \Sigma(E) &= \frac{1}{h^{dN} \xi_N} \int_{\sum_i \left(\frac{p_i^2}{2m} + \frac{m\omega^2}{2} q_i^2 \right) \leq E} \prod_i d^d q_i d^d p_i \\ &= \frac{1}{h^{dN} \xi_N} (2mE)^{dN/2} \left(\frac{2E}{m\omega^2} \right)^{dN/2} \underbrace{\int_{\sum_j x_j^2 \leq 1} \prod_j dx_j}_{\frac{\pi^{dN}}{\Gamma(dN+1)}} \\ &= \frac{1}{\xi_N \Gamma(dN+1)} \left(\frac{2\pi E}{h\omega} \right)^{dN} \\ &= \frac{1}{\xi_N \Gamma(dN) dN} \left(\frac{2\pi E}{h\omega} \right)^{dN} . \end{aligned}$$

q.e.d.

The density state $\omega(E)$ is

$$\omega(E) = \frac{1}{\xi_N \Gamma(dN)} \left(\frac{2\pi}{h\omega} \right)^{dN} E^{dN-1} .$$

Proof. By definition,

$$\begin{aligned} \omega(E) &= \frac{\partial \Sigma(E)}{\partial E} \\ &= \frac{\partial}{\partial E} \frac{1}{\xi_N \Gamma(dN) dN} \left(\frac{2\pi E}{h\omega} \right)^{dN} \\ &= \frac{1}{\xi_N \Gamma(dN) dN} \left(\frac{2\pi}{h\omega} \right)^{dN} \frac{\partial}{\partial E} E^{dN} \\ &= \frac{1}{\xi_N \Gamma(dN) dN} \left(\frac{2\pi}{h\omega} \right)^{dN} dN E^{dN-1} \\ &= \frac{1}{\xi_N \Gamma(dN)} \left(\frac{2\pi}{h\omega} \right)^{dN} E^{dN-1} . \end{aligned}$$

q.e.d.

In the thermodynamic limit, the entropy becomes

$$S = k_B \begin{cases} dN + N \ln \left(\frac{2\pi E}{h\omega dN} \right)^d & \text{for distinguishable particles} \\ (d+1)N + N \ln \left(\frac{1}{N} \left(\frac{2\pi E}{h\omega dN} \right)^d \right) & \text{for indistinguishable particles} \end{cases} .$$

Proof. By definition, using the Stirling approximation (B.1),

$$\begin{aligned} \frac{S}{k_B} &= \ln \Sigma(E) \\ &= \ln \frac{1}{\xi_N \Gamma(dN/2) dN} \left(\frac{2\pi E}{h\omega} \right)^{dN} \\ &= -\ln \xi_N - \ln \Gamma(dN) - \ln d - \ln N + N \ln \left(\frac{2\pi E}{h\omega} \right)^d \\ &= -\ln \xi_N - \underbrace{\ln \Gamma(dN)}_{dN \ln(dN) - dN} + N \ln \left(\frac{2\pi E}{h\omega} \right)^d \\ &= -\ln \xi_N - dN \ln(dN) + dN + N \ln \left(\frac{2\pi E}{h\omega} \right)^d \\ &= -\ln \xi_N + dN + N \ln \left(\frac{2\pi E}{h\omega dN} \right)^d . \end{aligned}$$

Now, we treat the distinguishable and indistinguishable case separately. For distinguishable particles $\xi_N = 1$, we find

$$\frac{S}{k_B} = -\ln 1 + dN + N \ln \left(\frac{2\pi E}{h\omega dN} \right)^d = dN + N \ln \left(\frac{2\pi E}{h\omega dN} \right)^d .$$

For indistinguishable particles $\xi_N = N!$, we find

$$\begin{aligned} \frac{S}{k_B} &= -\underbrace{\ln N!}_{N \ln N - N} + dN + N \ln \left(\frac{2\pi E}{h\omega dN} \right)^d \\ &= -N \ln N + N + dN + N \ln \left(\frac{2\pi E}{h\omega dN} \right)^d \\ &= (d+1)N + N \ln \left(\frac{1}{N} \frac{2\pi E}{h\omega dN} \right)^d . \end{aligned}$$

q.e.d.

The internal energy is

$$E = dN k_B T .$$

Proof. By (??)

$$\frac{1}{T} = \frac{\partial S}{\partial E} = k_B dN \frac{\partial}{\partial E} \ln E = k_B \frac{dN}{E} ,$$

hence

$$E = dN k_B T .$$

q.e.d.

Chapter 10

Canonical ensemble

10.1 Non-relativistic ideal gas in d-dimensions

Consider an indistinguishable non-relativistic ideal (non-interacting) gas of N particles in an d -dimensional manifold with a finite volume V^N : $\mathcal{M}_N = V^N \times \mathbb{R}^{dN}$. If we did not confine the particles to a finite volume, we would have found undesired divergences. Its hamiltonian is

$$H = \sum_i \frac{p_i^2}{2m} .$$

The canonical partition function Z is

$$Z = \frac{V^N}{\xi_N \lambda_T^{dN}} = \frac{V^N}{\xi_N} \left(\frac{2m\pi}{\beta h^2} \right)^{dN/2} .$$

Proof. By definition, using the gaussian integral (C.1),

$$\begin{aligned}
Z &= \int_{\mathcal{M}^N} d\Omega \exp(-\beta H(q_i, p_i)) \\
&= \int_{\mathcal{M}^N} \frac{\prod_i d^d q_i d^d p_i}{h^{dN} \xi_N} \exp(-\beta H(q_i, p_i)) \\
&= \frac{1}{h^{dN} \xi_N} \int_{\mathcal{M}^N} \prod_i d^d q_i d^d p_i \exp(-\beta H(q_i, p_i)) \\
&= \frac{1}{h^{dN} \xi_N} \underbrace{\int_{V^N} \prod_i d^d q_i}_{V^N} \underbrace{\prod_i \int_{\mathcal{M}^N} d^d p_i \exp(-\beta \frac{p_i^2}{2m})}_{(\frac{2m\pi}{\beta})^{dN/2}} \\
&= \frac{V^N}{h^{dN} \xi_N} \left(\frac{2m\pi}{\beta} \right)^{dN/2} \\
&= \frac{V^N}{\xi_N} \left(\frac{2m\pi}{\beta h^2} \right)^{dN/2} \\
&= \frac{V^N}{\xi_N \lambda_T^{dN}}
\end{aligned}$$

where we have defined the thermal wavelength

$$\lambda_T = \sqrt{\frac{\beta h^2}{2m\pi}} .$$

q.e.d.

For indistinguishable particles, the canonical partition function Z is

$$Z = \frac{V^N}{N! \lambda_T^{dN}} = \frac{V^N}{N!} \left(\frac{2m\pi}{\beta h^2} \right)^{dN/2} .$$

An useful intermediary formula is

$$\ln Z = N(1 - \ln(\frac{N}{V} \lambda_T^d)) = N(1 - \ln(n \lambda_T^d)) .$$

Proof. In fact, using the Stirling approximation (B.1),

$$\begin{aligned}
 \ln Z &= \ln \frac{V^N}{N! \lambda_T^{dN}} \\
 &= N \ln(V \lambda_T^d) - \underbrace{\ln N!}_{N \ln N - N} \\
 &= N - N \frac{V \lambda_T^d}{N} \\
 &= N(1 - \ln(\frac{N}{V} \lambda_T^d)) \\
 &= N(1 - \ln(n \lambda_T^d)) ,
 \end{aligned}$$

where we have defined the density

$$n = \frac{N}{V} .$$

q.e.d.

The internal energy E is

$$E = \frac{d}{2} N k_B T .$$

Proof. By (6.9)

$$\begin{aligned}
 E &= - \frac{\partial \ln Z}{\partial \beta} \\
 &= - \frac{\partial}{\partial \beta} N(1 - \ln(n \lambda_T^d)) \\
 &= - N d \frac{\partial}{\partial \beta} \ln(\lambda_T) \\
 &= - N d \frac{\partial}{\partial \beta} \ln(\beta^{1/2}) \\
 &= \frac{N d}{2} \frac{1}{\beta} \\
 &= \frac{d}{2} N k_B T .
 \end{aligned}$$

q.e.d.

The Helmholtz free energy F is

$$F = \frac{N}{\beta} (\ln(n \lambda_T^d) - 1) .$$

Proof. By (6.6)

$$F = -\frac{\ln Z}{\beta} = \frac{N}{\beta}(\ln(n\lambda_T^d) - 1) .$$

q.e.d.

The entropy S is

$$S = Nk_B \left(\frac{d+2}{2} - \ln(n\lambda_T^d) \right) .$$

Proof. By (6.8)

$$\begin{aligned} S &= \frac{E - F}{T} \\ &= \frac{1}{T} \left(\frac{d}{2} Nk_B T - \frac{N}{\beta} (\ln(n\lambda_T^d) - 1) \right) \\ &= \frac{N}{\beta T} \left(\frac{d+2}{2} - \ln(n\lambda_T^d) \right) \\ &= Nk_B \left(\frac{d+2}{2} - \ln(n\lambda_T^d) \right) \end{aligned}$$

q.e.d.

Entropy becomes negative at a certain critical temperature

$$T_c = \frac{2m\pi k_B}{h^2} e^{(d+2)/2} n^{-2/d} .$$

Proof. In fact, $S < 0$ for

$$Nk_B \left(\frac{d+2}{2} - \ln(n\lambda_T^d) \right) < 0 ,$$

$$\frac{d+2}{2} - \ln(n\lambda_T^d) < 0 ,$$

$$\frac{d+2}{2} < \ln(n\lambda_T^d) ,$$

$$e^{(d+2)/2} < n\lambda_T^d ,$$

$$e^{(d+2)/2} < n \left(\frac{h^2 \beta}{2m\pi} \right)^{d/2} ,$$

$$e^{(d+2)/d} n^{2/d} < \frac{h^2 \beta}{2m\pi} ,$$

$$\frac{2m\pi}{h^2} e^{(d+2)/2} n^{-2/d} < \beta ,$$

hence

$$T < \frac{2m\pi k_B}{h^2} e^{(d+2)/2} n^{-2/d} = T_c .$$

q.e.d.

The equation of state is

$$pV = Nk_B T . \quad (10.1)$$

Proof. By (??)

$$p = -\frac{\partial F}{\partial V} = -\frac{\partial}{\partial V} \frac{N}{\beta} (\ln(n\lambda_T^d) - 1) = \frac{N}{\beta} \frac{\partial}{\partial V} \ln V = \frac{N}{V\beta} ,$$

hence

$$pV = Nk_B T .$$

q.e.d.

The chemical potential μ is

$$\mu = \frac{1}{\beta} \ln(n\lambda_T^d) .$$

Proof. By (??)

$$\mu = \frac{\partial F}{\partial N} = \frac{\partial}{\partial N} \frac{N}{\beta} (\ln(n\lambda_T^d) - 1) = \frac{1}{\beta} (\ln(n\lambda_T^d) - 1) + \frac{1}{\beta} = \frac{1}{\beta} \ln(n\lambda_T^d) .$$

q.e.d.

The specific heats C_V and C_p are

$$C_V = N \frac{d}{2} k_B , \quad C_p = N \frac{d+2}{2} k_B .$$

Proof. At V constant

$$C_V = \frac{\partial E}{\partial T} = \frac{\partial}{\partial T} \frac{d}{2} N k_B T = N \frac{d}{2} k_B .$$

At p constant, using (10.1)

$$C_p = C_V + p \frac{\partial V}{\partial T} = C_V + p \frac{\partial}{\partial T} \frac{N k_B T}{p} = N \frac{d}{2} k_B + N k_B = \frac{d+2}{2} k_B .$$

q.e.d.

10.2 Non-relativistic ideal gas in 3-dimensions

Now, consider the case in which $d = 3$. For indistinguishable particles, the canonical partition function Z is

$$Z = \frac{V^N}{N! \lambda_T^{3N}} = \frac{V^N}{N!} \left(\frac{2m\pi}{\beta h^2} \right)^{3N/2} .$$

The internal energy E is

$$E = \frac{3}{2} N k_B T .$$

The Helmholtz free energy F is

$$F = \frac{N}{\beta} (\ln(n \lambda_T^3) - 1) .$$

The entropy S is

$$S = N k_B \left(\frac{5}{2} - \ln(n \lambda_T^3) \right) .$$

Entropy becomes negative at a certain critical temperature

$$T_c = \frac{2m\pi k_B}{h^2} e^{3/2} n^{-2/3} .$$

A plot of this is in Figure 10.1.

The equation of state is

$$pV = N k_B T .$$

The chemical potential μ is

$$\mu = \frac{1}{\beta} \ln(n \lambda_T^3) .$$

A plot of this is in Figure 10.2.

The specific heats C_V and C_p are

$$C_V = N \frac{3}{2} k_B , \quad C_p = N \frac{5}{2} k_B .$$

Notice that there are two problems: entropy cannot be negative and the specific heat $C_V \rightarrow 0$ for $T \rightarrow 0$, by thermodynamics. This means that this model is not correct and we must go quantum.

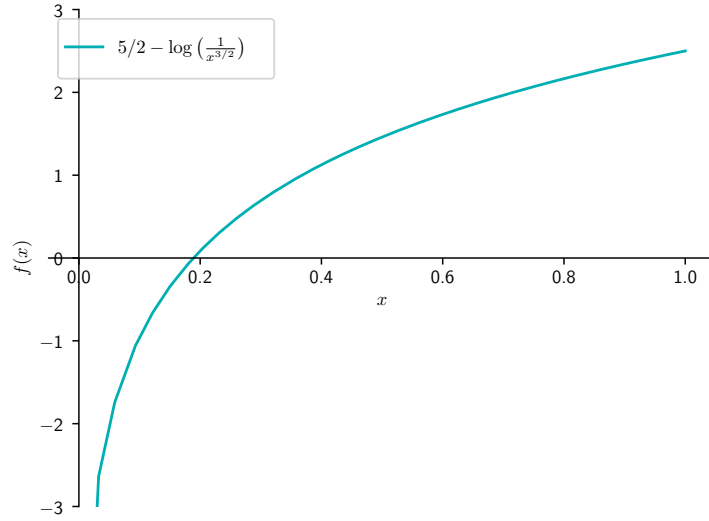


Figure 10.1: A plot of the entropy S as a function of T . We have used $x = \frac{2\pi mk_B T n^{2/3}}{h^2}$ and $f(x) = \frac{S}{Nk_B}$.

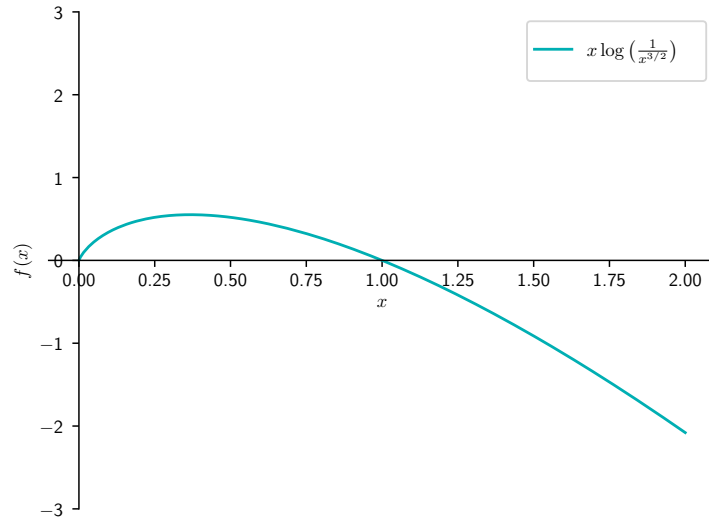


Figure 10.2: A plot of the chemical potential μ as a function of T . We have used $x = \frac{2\pi mk_B T n^{2/3}}{h^2}$ and $f(x) = \frac{2\pi m \mu}{h^2 n^{3/2}}$.

10.3 Gas of harmonic oscillators in d-dimensions

Consider a distinguishable non-relativistic (non-interacting) gas of N particles in an d -dimensional manifold confined by an harmonic potential of frequency ω . Its hamiltonian is

$$H = \sum_i \left(\frac{p_i^2}{2m} + \frac{m\omega^2}{2} q_i^2 \right) .$$

The canonical partition function Z is

$$Z = \frac{1}{\xi_N} \left(\frac{2\pi k_B T}{h\omega} \right)^{dN} .$$

Proof. By definition, using the gaussian integral (C.1),

$$\begin{aligned} Z &= \int_{\mathcal{M}^N} d\Omega \exp(-\beta H(q_i, p_i)) \\ &= \int_{\mathcal{M}^N} \frac{\prod_i d^d q_i d^d p_i}{h^{dN} \xi_N} \exp(-\beta H(q_i, p_i)) \\ &= \frac{1}{h^{dN} \xi_N} \int_{\mathcal{M}^N} \prod_i d^d q_i d^d p_i \exp(-\beta H(q_i, p_i)) \\ &= \frac{1}{h^{dN} \xi_N} \underbrace{\int_{V^N} \prod_i d^d q_i \exp(-\beta \frac{m\omega^2}{2} q_i^2)}_{(\frac{2\pi}{m\omega\beta})^{dN/2}} \underbrace{\prod_i \int_{\mathcal{M}^N} d^d p_i \exp(-\beta \frac{p_i^2}{2m})}_{(\frac{2m\pi}{\beta})^{dN/2}} \\ &= \frac{1}{h^{dN} \xi_N} \left(\frac{2\pi}{m\omega\beta} \right)^{dN/2} \left(\frac{2m\pi}{\beta} \right)^{dN/2} \\ &= \frac{1}{\xi_N} \left(\frac{2\pi}{h\omega\beta} \right)^{dN} \\ &= \frac{1}{\xi_N} \left(\frac{2\pi k_B T}{h\omega} \right)^{dN} . \end{aligned}$$

q.e.d.

For distinguishable particles, the canonical partition function Z is

$$Z = \left(\frac{2\pi k_B T}{h\omega} \right)^{dN} = (Z_1)^N .$$

An useful intermediary formula is

$$\ln Z = dN \ln \frac{2\pi k_B T}{h\omega} .$$

Proof. In fact, using the Stirling approximation (B.1),

$$\ln Z = \ln \left(\frac{2\pi k_B T}{h\omega} \right)^{dN} = dN \ln \frac{2\pi k_B T}{h\omega} .$$

q.e.d.

The internal energy E is

$$E = dN k_B T .$$

Proof. By (6.9)

$$E = -\frac{\partial \ln Z}{\partial \beta} = -\frac{\partial}{\partial \beta} dN \ln \frac{2\pi}{h\omega\beta} = dN \frac{1}{\beta} = dN k_B T .$$

q.e.d.

The Helmholtz free energy F is

$$F = \frac{dN}{\beta} \ln \frac{h\omega}{2\pi k_B T} .$$

Proof. By (6.6)

$$F = -\frac{\ln Z}{\beta} = -\frac{dN}{\beta} \ln \frac{2\pi k_B T}{h\omega} = \frac{dN}{\beta} \ln \frac{h\omega}{2\pi k_B T} .$$

q.e.d.

The entropy S is

$$S = dN k_B (1 - \ln \frac{h\omega}{2\pi k_B T}) .$$

Proof. By (6.8)

$$S = \frac{E - F}{T} = \frac{1}{T} \left(dN k_B T - \frac{dN}{\beta} \ln \frac{h\omega}{2\pi k_B T} \right) = dN k_B (1 - \ln \frac{h\omega}{2\pi k_B T}) .$$

q.e.d.

Entropy becomes negative at a certain critical temperature

$$T_c = \frac{h\omega}{2\pi k_B e} .$$

Proof. In fact, $S < 0$ for

$$dNk_B(1 - \ln \frac{h\omega}{2\pi k_B T}) < 0 ,$$

$$1 - \ln \frac{h\omega}{2\pi k_B T} < 0 ,$$

$$1 < \ln \frac{h\omega}{2\pi k_B T} ,$$

$$e < \frac{h\omega}{2\pi k_B T} ,$$

hence

$$T < \frac{h\omega}{2\pi k_B e} = T_c .$$

q.e.d.

The equation of state is

$$p = 0 . \quad (10.2)$$

Proof. By (??)

$$p = -\frac{\partial F}{\partial V} = 0 .$$

q.e.d.

The chemical potential μ is

$$\mu = \frac{d}{\beta} \ln \frac{h\omega}{2\pi k_B T} .$$

Proof. By (??)

$$\mu = \frac{\partial F}{\partial N} = \frac{\partial}{\partial N} \frac{dN}{\beta} \ln \frac{h\omega}{2\pi k_B T} = \frac{d}{\beta} \ln \frac{h\omega}{2\pi k_B T} .$$

q.e.d.

The specific heats C_V and C_p are

$$C_V = dNk_B , \quad C_p = dNk_B .$$

Proof. At V constant

$$C_V = \frac{\partial E}{\partial T} = \frac{\partial}{\partial T} dNk_B T = dNk_B .$$

At p constant, using (10.2)

$$C_p = C_V + p \frac{\partial V}{\partial T} = C_V + p \frac{\partial}{\partial T} \frac{Nk_B T}{p} = C_V = dNk_B .$$

q.e.d.

10.4 Gas of harmonic oscillators in 1-dimension

Now, consider the case in which $d = 1$. For distinguishable particles, the canonical partition function Z is

$$Z = \left(\frac{2\pi k_B T}{h\omega} \right)^N = (Z_1)^N .$$

The internal energy E is

$$E = Nk_B T .$$

The Helmholtz free energy F is

$$F = \frac{N}{\beta} \ln \frac{h\omega}{2\pi k_B T} .$$

The entropy S is

$$S = Nk_B \left(1 - \ln \frac{h\omega}{2\pi k_B T} \right) .$$

Entropy becomes negative at a certain critical temperature

$$T_c = \frac{h\omega}{2\pi k_B e} .$$

A plot of this is in Figure 10.3.

The equation of state is

$$p = 0 .$$

The chemical potential μ is

$$\mu = \frac{1}{\beta} \ln \frac{h\omega}{2\pi k_B T} .$$

A plot of this is in Figure 10.4.

The specific heats C_V and C_p are

$$C_V = Nk_B , \quad C_p = Nk_B .$$

Notice that also here there are two problems: entropy cannot be negative and the specific heat $C_V \rightarrow 0$ for $T \rightarrow 0$, by thermodynamics. This means that this model is not correct and we must go quantum.

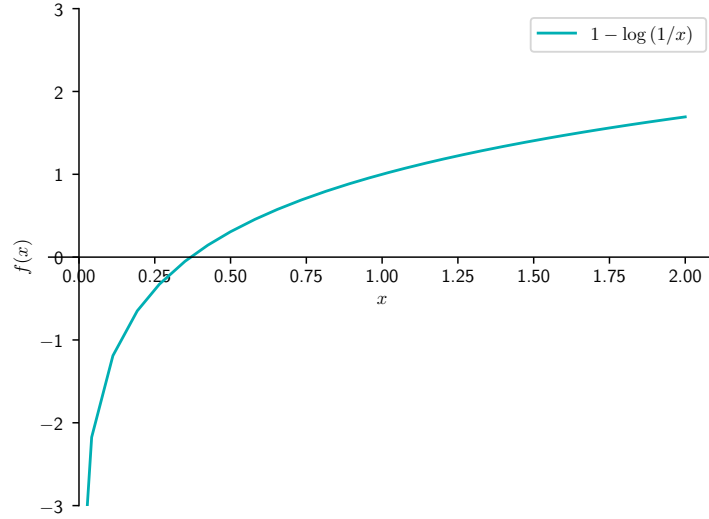


Figure 10.3: A plot of the entropy S as a function of T . We have used $x = \frac{2\pi k_B T}{h\omega}$ and $f(x) = \frac{S}{Nk_B}$.

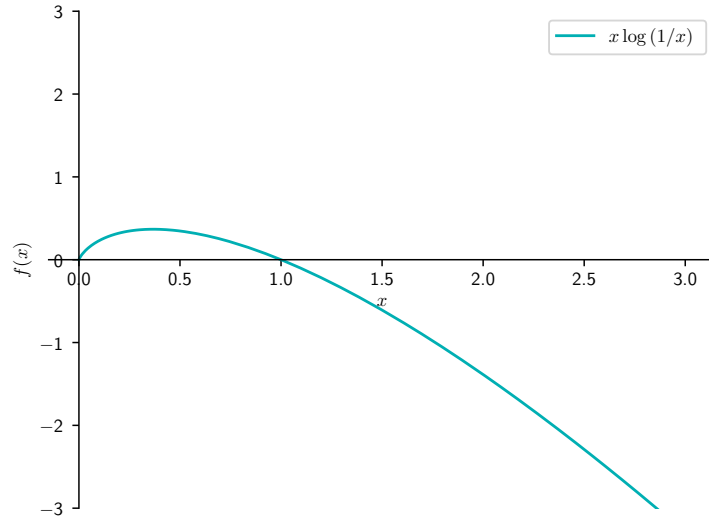


Figure 10.4: A plot of the chemical potential μ as a function of T . We have used $x = \frac{2\pi k_B T}{h\omega}$ and $f(x) = \frac{2\pi\mu}{h\omega}$.

10.5 Ultra-relativistic ideal gas

Consider an indistinguishable ultra-relativistic ideal (non-interacting) gas of N particles in an d -dimensional manifold with a finite volume V^N : $\mathcal{M}_N = V^N \times \mathbb{R}^{dN}$. Its hamiltonian is

$$H = \sum_i c|p_i| .$$

The canonical partition function Z is

$$Z = \frac{1}{\xi_N} \left(\frac{8\pi V}{(\beta hc)^3} \right)^N .$$

Proof. By definition, using the gaussian integral (C.1),

$$\begin{aligned} Z &= \int_{\mathcal{M}^N} d\Omega \exp(-\beta H(q_i, p_i)) \\ &= \int_{\mathcal{M}^N} \frac{\prod_i d^3 q_i d^3 p_i}{h^{3N} \xi_N} \exp(-\beta H(q_i, p_i)) \\ &= \frac{1}{h^{3N} \xi_N} \int_{\mathcal{M}^N} \prod_i d^3 q_i d^3 p_i \exp(-\beta H(q_i, p_i)) \\ &= \frac{1}{h^{3N} \xi_N} \int_{V^N} \underbrace{\prod_i d^d q_i}_{V^N} \prod_i \int_{\mathcal{M}^N} d^d p_i \exp(-\beta c p_i) \\ &= \frac{V^N}{h^{3N} \xi_N} \prod_i \int_{\mathcal{M}^N} d^d p_i \exp(-\beta c p_i) . \end{aligned}$$

Now, in order to evaluate the integral, we use the polar coordinates in the momentum space (p, θ, ϕ)

$$\prod_i \int_{\mathcal{M}^N} d^3 p_i \exp(-\beta c p_i) = \prod_i 4\pi \int_0^\infty dp p^2 \exp(-\beta c p_i)$$

We change variable

$$z = \beta c p , \quad dz = -\beta c dp ,$$

and we find

$$\prod_i \frac{4\pi}{(\beta c)^3} \underbrace{\int_0^\infty dz z^2 \exp(-z)}_{\Gamma(3)} = \prod_i \frac{4\pi}{(\beta c)^3} \underbrace{\Gamma(3)}_2 = \prod_i \frac{8\pi}{(\beta c)^3} = \left(\frac{8\pi}{(\beta c)^3} \right)^N .$$

Therefore

$$Z = \frac{V^N}{h^{3N} \xi_N} \left(\frac{8\pi}{(\beta c)^3} \right)^N = \frac{1}{\xi_N} \left(\frac{8\pi V}{(\beta hc)^3} \right)^N .$$

q.e.d.

For indistinguishable particles, the canonical partition function Z is

$$Z = \frac{1}{N!} \left(\frac{8\pi V}{(\beta hc)^3} \right)^N .$$

An useful intermediary formula is

$$\ln Z = N \left(1 - \ln \frac{n(\beta hc)^3}{8\pi} \right) .$$

Proof. In fact, using the Stirling approximation (B.1),

$$\begin{aligned} \ln Z &= \ln \frac{1}{N!} \left(\frac{8\pi V}{(\beta hc)^3} \right)^N \\ &= - \underbrace{\ln N!}_{N \ln N - N} + N \ln \frac{8\pi V}{(\beta hc)^3} \\ &= N \left(1 - \ln \frac{N(\beta hc)^3}{8\pi V} \right) \\ &= N \left(1 - \ln \frac{n(\beta hc)^3}{8\pi} \right) , \end{aligned}$$

where we have defined the density

$$n = \frac{N}{V} .$$

q.e.d.

The internal energy E is

$$E = 3Nk_B T .$$

Proof. By (6.9)

$$E = -\frac{\partial \ln Z}{\partial \beta} = -\frac{\partial}{\partial \beta} N \left(1 - \ln \frac{n(\beta hc)^3}{8\pi} \right) = N \frac{\partial}{\partial \beta} \ln(\beta^3) = 3N \frac{\beta^2}{\beta^3} = 3N \frac{1}{\beta} = 3Nk_B T .$$

As an aside, it can be also derived from the generalised equipartition theorem (6.11). In fact

$$k_B T = \left\langle p_i \frac{\partial H}{\partial p_i} \right\rangle = \left\langle p_i \frac{\partial}{\partial p_i} c \sqrt{p_1^2 + p_2^2 + p_3^2} \right\rangle = \left\langle c \frac{p_i^2}{\sqrt{p_1^2 + p_2^2 + p_3^2}} \right\rangle ,$$

hence

$$\langle H \rangle = \left\langle c \frac{p_1^2 + p_2^2 + p_3^2}{\sqrt{p_1^2 + p_2^2 + p_3^2}} \right\rangle = \sum_{i=1}^3 \underbrace{\left\langle c \frac{p_i^2}{\sqrt{p_1^2 + p_2^2 + p_3^2}} \right\rangle}_{k_B T} = 3k_B T .$$

q.e.d.

The Helmholtz free energy F is

$$F = \frac{N}{\beta} \left(\ln \frac{n(\beta hc)^3}{8\pi} - 1 \right) .$$

Proof. By (6.6)

$$F = -\frac{\ln Z}{\beta} = \frac{N}{\beta} \left(\ln \frac{n(\beta hc)^3}{8\pi} - 1 \right) .$$

q.e.d.

The entropy S is

$$S = Nk_B \left(4 - \ln \frac{n(\beta hc)^3}{8\pi} \right) .$$

Proof. By (6.8)

$$\begin{aligned} S &= \frac{E - F}{T} \\ &= \frac{1}{T} \left(3Nk_B T - \frac{N}{\beta} \left(\ln \frac{n(\beta hc)^3}{8\pi} - 1 \right) \right) \\ &= 3Nk_B - Nk_B \left(\ln \frac{n(\beta hc)^3}{8\pi} - 1 \right) \\ &= Nk_B \left(4 - \ln \frac{n(\beta hc)^3}{8\pi} \right) . \end{aligned}$$

q.e.d.

Entropy becomes negative at a certain critical temperature

$$T_c = \frac{hc}{k_B} \left(\frac{n}{8\pi e^4} \right)^{1/3} .$$

Proof. In fact, $S < 0$ for

$$Nk_B \left(4 - \ln \frac{n(\beta hc)^3}{8\pi} \right) < 0 ,$$

$$4 - \ln \frac{n(\beta hc)^3}{8\pi} < 0 ,$$

$$4 < \ln \frac{n(\beta hc)^3}{8\pi} ,$$

$$e^4 < \frac{n(\beta hc)^3}{8\pi} ,$$

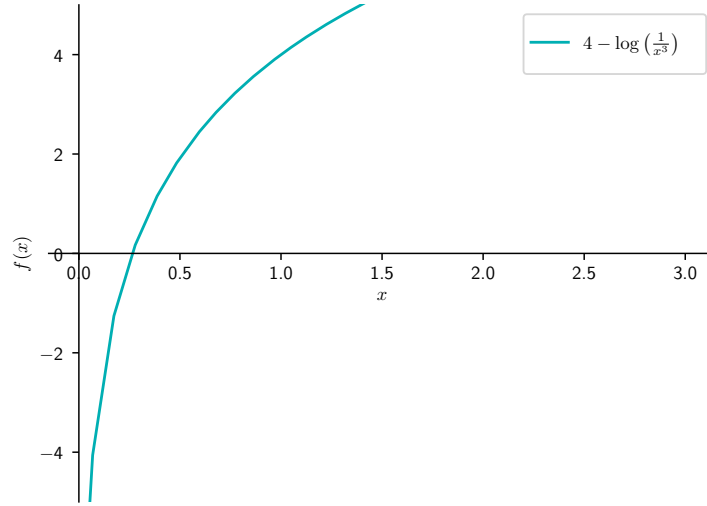


Figure 10.5: A plot of the entropy S as a function of T . We have used $x = \frac{(8\pi)^{1/3} k_B T}{hc n^{1/3}}$ and $f(x) = \frac{S}{N k_B}$.

$$e^4 < \frac{n(hc)^3}{8\pi k_B^3 T^3} ,$$

$$T^3 < \frac{n(hc)^3}{8\pi k_B^3 e^4} ,$$

hence

$$T < \frac{hc}{k_B} \left(\frac{n}{8\pi e^4} \right)^{1/3} = T_c .$$

q.e.d.

A plot of this is in Figure 10.5.

The equation of state is

$$pV = N k_B T . \quad (10.3)$$

Proof. By(??)

$$p = -\frac{\partial F}{\partial V} = -\frac{\partial}{\partial V} \frac{N}{\beta} \left(\ln \frac{n(\beta hc)^3}{8\pi} - 1 \right) = \frac{N}{\beta} \frac{\partial}{\partial V} \ln V = \frac{N}{V\beta} ,$$

hence

$$pV = N k_B T .$$

q.e.d.

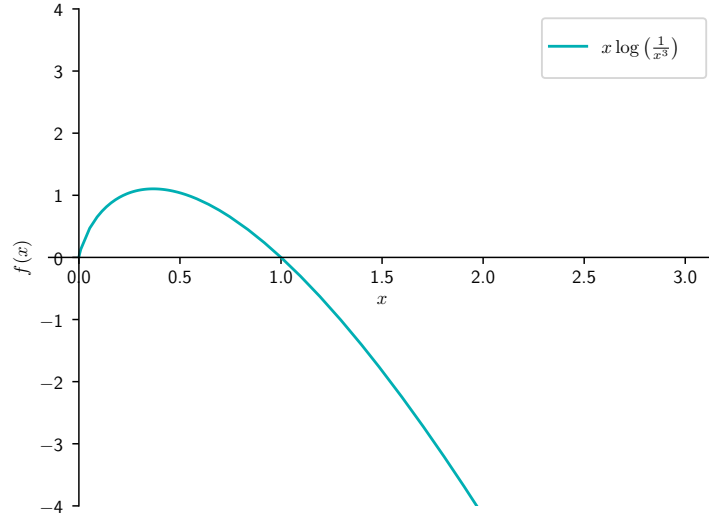


Figure 10.6: A plot of the chemical potential μ as a function of T . We have used $x = \frac{(8\pi)^{1/3}k_B T}{hcn^{1/3}}$ and $f(x) = \frac{(8\pi)^{1/3}\mu}{hcn^{1/3}}$.

The chemical potential μ is

$$\mu = \frac{1}{\beta} \ln \frac{n(\beta hc)^3}{8\pi} .$$

Proof. By (??)

$$\mu = \frac{\partial F}{\partial N} = \frac{\partial}{\partial N} \frac{N}{\beta} \left(\ln \frac{n(\beta hc)^3}{8\pi} - 1 \right) = \frac{1}{\beta} \left(\ln \frac{n(\beta hc)^3}{8\pi} - 1 + 1 \right) = \frac{1}{\beta} \ln \frac{n(\beta hc)^3}{8\pi} .$$

q.e.d.

A plot of this is in Figure 10.6.

The specific heats C_V and C_p are

$$C_V = 3Nk_B , \quad C_p = 4Nk_B T .$$

Proof. At V constant

$$C_V = \frac{\partial E}{\partial T} = \frac{\partial}{\partial T} 3Nk_B T = 3Nk_B .$$

At p constant, using (10.3)

$$C_p = C_V + p \frac{\partial V}{\partial T} = C_V + p \frac{\partial}{\partial T} \frac{Nk_B T}{p} = 3Nk_B + Nk_B = 4Nk_B T .$$

q.e.d.

10.6 Maxwell-Boltzmann velocity distribution

Consider a non-relativistic ideal (non-interacting) gas of N particles in an 3-dimensional manifold $\mathcal{M}^N = \mathbb{R}^6$, confined into a potential $V(q_i)$. In this discussion we put $\hbar = 1$. Its hamiltonian is

$$H = \sum_i \left(\frac{p_i^2}{2m} + V(q_i) \right) .$$

The probability distribution density ρ_c for each particle is

$$\rho_c(q_i, p_i) = \frac{\exp(-\beta(\frac{p_i^2}{2m} + V(q_i)))}{(\frac{2\pi m}{\beta})^{3/2} \int_{\mathbb{R}^3} d^3 q \exp(-\beta V(q))} .$$

Proof. By definition,

$$\rho_c(q_i, p_i) = \mathcal{N} \exp(-\beta(\frac{p_i^2}{2m} + V(q_i))) ,$$

where the normalisation constant is, using the gaussian integral (C.1)

$$\begin{aligned} 1 &= \int_{\mathbb{R}^6} \prod_i d^3 q \, d^3 p \mathcal{N} \exp(-\beta(\frac{p^2}{2m} + V(q))) \\ &= \mathcal{N} \int_{\mathbb{R}^3} d^3 q \exp(-\beta V(q)) \underbrace{\int_{\mathbb{R}^3} d^3 p \exp(-\beta \frac{p^2}{2m})}_{\left(\frac{2\pi m}{\beta}\right)^{3/2}} \\ &= \mathcal{N} \left(\frac{2\pi m}{\beta}\right)^{3/2} \int_{\mathbb{R}^3} d^3 q \exp(-\beta V(q)) , \end{aligned}$$

hence

$$\mathcal{N} = \left(\left(\frac{2\pi m}{\beta}\right)^{3/2} \int_{\mathbb{R}^3} d^3 q \exp(-\beta V(q)) \right)^{-1} .$$

q.e.d.

The marginal probability density distribution is

$$\rho(q_i) = \frac{\exp(-\beta V(q_i))}{\int_{\mathbb{R}^3} d^3q \exp(-\beta V(q))} .$$

Proof. By definition,

$$\begin{aligned} \rho(q_i) &= \int_{\mathbb{R}^3} d^3p \rho_c(q_i, p) \\ &= \int_{\mathbb{R}^3} d^3p \frac{\exp(-\beta(\frac{p^2}{2m} + V(q_i)))}{(\frac{2\pi m}{\beta})^{3/2} \int_{\mathbb{R}^3} d^3q \exp(-\beta V(q))} \\ &= \frac{\exp(-\beta V(q_i))}{\int_{\mathbb{R}^3} d^3q \exp(-\beta V(q))} \frac{\int_{\mathbb{R}^3} d^3p \exp(-\beta \frac{p^2}{2m})}{(\frac{2\pi m}{\beta})^{3/2}} \\ &= \frac{\exp(-\beta V(q_i))}{\int_{\mathbb{R}^3} d^3q \exp(-\beta V(q))} . \end{aligned}$$

q.e.d.

If we have a potential defined as

$$V(q) = \begin{cases} 0 & \text{inside a region } \mathcal{A} \\ \infty & \text{outside a region } \mathcal{A} \end{cases} ,$$

the probability is null outside this region and uniform inside it.

Proof. In fact

$$\rho(q_i) = \frac{1}{\int_{\mathcal{A}} d^3q} = \frac{1}{\mathcal{A}} .$$

q.e.d.

The momentum probability density distribution is

$$\rho(p) = (2\pi m k_B T)^{-3/2} \exp(-\beta \frac{p^2}{2m}) = \prod_i (2\pi m k_B T)^{-1/2} \exp(-\beta \frac{p_i^2}{2m}) .$$

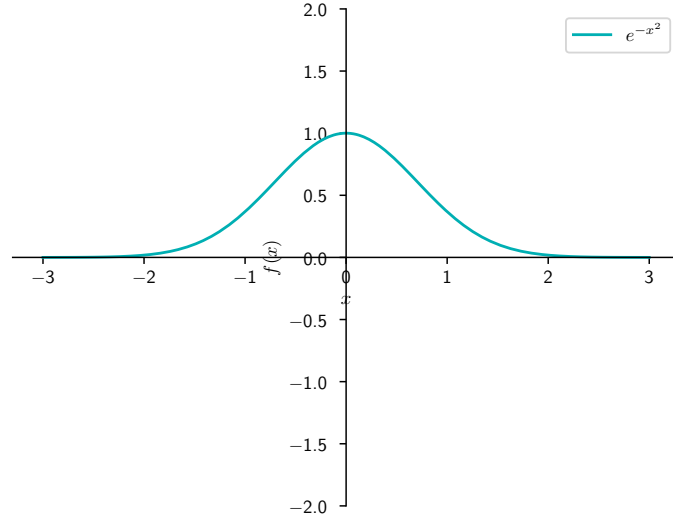


Figure 10.7: A plot of the momentum probability density distribution. We have used $x = \sqrt{\frac{\beta}{2m}}p$ and $f(x) = (2\pi mk_B T)^{3/2} \rho$.

Proof. By definition,

$$\begin{aligned}
 \rho(p) &= \int_{\mathbb{R}^3} d^3q \, \rho_c(q, p) \\
 &= \int_{\mathbb{R}^3} d^3q \, \frac{\exp(-\beta(\frac{p^2}{2m} + V(q)))}{(\frac{2\pi m}{\beta})^{3/2} \int_{\mathbb{R}^3} d^3q' \exp(-\beta V(q'))} \\
 &= \frac{\exp(-\beta \frac{p^2}{2m})}{(\frac{2\pi m}{\beta})^{3/2}} \frac{\int_{\mathbb{R}^3} d^3q \exp(-\beta V(q))}{\int_{\mathbb{R}^3} d^3q' \exp(-\beta V(q'))} \\
 &= \frac{\exp(-\beta \frac{p^2}{2m})}{(\frac{2\pi m}{\beta})^{3/2}} \\
 &= (2\pi mk_B T)^{-3/2} \exp(-\beta \frac{p^2}{2m}) \\
 &= \prod_i (2\pi mk_B T)^{-1/2} \exp(-\beta \frac{p_i^2}{2m}) .
 \end{aligned}$$

q.e.d.

A plot of this is in Figure 10.7.

The velocity probability density distribution is

$$\rho(p) = \left(\frac{m}{2\pi k_B T}\right)^{1/2} \exp\left(-\beta \frac{mv_i^2}{2}\right) .$$

Proof. With a change of variable

$$p_i = mv_i , \quad \rho(v_i)dv_i = \rho(p_i)dp_i = \rho(p_i)m dv_i ,$$

we find

$$\rho(v_i) = m\rho(p_i) = \left(\frac{2\pi k_B T}{m}\right)^{-1/2} \exp\left(-\beta \frac{m^2 v_i^2}{2m}\right) = \left(\frac{m}{2\pi k_B T}\right)^{1/2} \exp\left(-\beta \frac{mv_i^2}{2}\right) .$$

q.e.d.

The velocity modulus probability density distribution is

$$\rho(v) = \left(\frac{m}{2\pi k_B T}\right)^{3/2} 4\pi v^2 \exp\left(-\beta \frac{mv^2}{2}\right) .$$

Proof. With a change of variable into the polar coordinates (v, θ, ϕ)

$$\rho(v_1, v_2, v_3)dv_1 dv_2 dv_3 = \rho(v_1, v_2, v_3)v^2 \sin \theta d\theta d\phi dv = \rho(\theta, \phi, v)d\theta d\phi dv ,$$

we find

$$\rho(v) = 4\pi v^2 \prod_i \rho(v_i) = \left(\frac{m}{2\pi k_B T}\right)^{3/2} 4\pi v^2 \exp\left(-\beta \frac{mv^2}{2}\right) .$$

q.e.d.

A plot of this is in Figure 10.8.

The most probable velocity value is

$$v_p = \sqrt{\frac{2k_B T}{m}} .$$

Proof. By definition,

$$0 = \frac{d\rho(v)}{dv} = 2v \exp\left(-\beta \frac{mv^2}{2}\right) - \beta m v^3 \exp\left(-\beta \frac{mv^2}{2}\right) ,$$

hence

$$v_p = \sqrt{\frac{2k_B T}{m}} .$$

q.e.d.

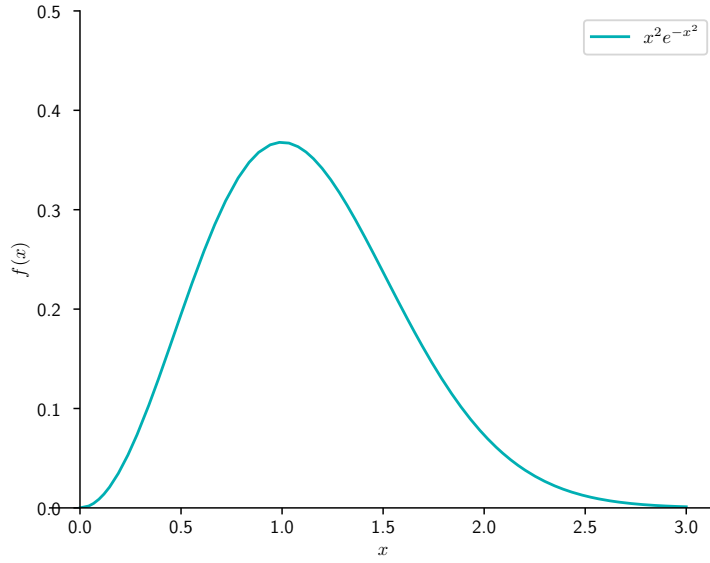


Figure 10.8: A plot of the velocity modulus probability density distribution. We have used $x = \sqrt{\frac{\beta m}{2}}v$ and $f(x) = \rho$.

The mean velocity value is

$$\langle v \rangle = \sqrt{\frac{8k_B T}{\pi m}} .$$

Proof. By definition,

$$\langle v \rangle = \int_{\mathbb{R}^3} dv \, \rho(v) v = \left(\frac{m}{2\pi k_B T} \right)^{3/2} 4\pi \int_{\mathbb{R}^3} dv \, v^3 \exp\left(-\beta \frac{mv^2}{2}\right) .$$

We make a change of variables

$$t = \frac{m\beta v^2}{2} , \quad dt = m\beta v dv ,$$

hence

$$\begin{aligned} \langle v \rangle &= \left(\frac{m}{2\pi k_B T} \right)^{3/2} 4\pi \left(\frac{2}{m\beta} \right) \frac{1}{m\beta} \underbrace{\int_0^\infty dt \, t \exp(-t)}_{\Gamma(2)} \\ &= \sqrt{\frac{8}{m\pi\beta}} \underbrace{\Gamma(2)}_1 = \sqrt{\frac{8}{m\pi\beta}} \\ &= \sqrt{\frac{8k_B T}{\pi m}} . \end{aligned}$$

q.e.d.

The mean square velocity value is

$$\langle v^2 \rangle = \frac{3k_B T}{m} .$$

Proof. By definition,

$$\langle v^2 \rangle = \int_{\mathbb{R}^3} dv \, \rho(v) v^2 = \left(\frac{m}{2\pi k_B T} \right)^{3/2} 4\pi \int_{\mathbb{R}^3} dv \, v^4 \exp\left(-\beta \frac{mv^2}{2}\right) .$$

We make a change of variables

$$t = \frac{m\beta v^2}{2} , \quad dt = m\beta v dv ,$$

hence

$$\begin{aligned} \langle v^2 \rangle &= \left(\frac{m}{2\pi k_B T} \right)^{3/2} 4\pi \left(\frac{2}{m\beta} \right) \frac{1}{m\beta} \left(\frac{2}{m\beta} \right)^{1/2} \underbrace{\int_0^\infty dt \, t^{3/2} \exp(-t)}_{\Gamma(5/2)} \\ &= \frac{4}{\sqrt{\pi} m \beta} \underbrace{\Gamma(5/2)}_{\frac{3\sqrt{\pi}}{4}} \\ &= \frac{3}{\beta m} \\ &= \frac{3k_B T}{m} . \end{aligned}$$

q.e.d.

10.7 Magnetic solid

A solid, composed by N atoms/molecules with an intrinsic magnetic moment $\boldsymbol{\mu}$ in an external magnetic field \mathbf{B} , can be modelled by an hamiltonian

$$H = - \sum_i \boldsymbol{\mu} \cdot \mathbf{B} = -\mu B \sum_i \cos \theta_i .$$

where the phase space coordinates are ϕ_i and θ_i .

The canonical partition function Z is

$$Z = \left(\frac{4\pi \sinh(\beta\mu B)}{\beta\mu B} \right)^N .$$

Proof. By definition

$$\begin{aligned}
 Z &= \int_{\mathcal{M}} d\Omega \exp(-\beta H(\theta_i)) \\
 &= \underbrace{\prod_i \int_0^{2\pi} d\phi_i}_{(2\pi)^N} \prod_i \int_0^\pi d\theta_i \sin \theta_i \exp(-\beta \mu B \cos \theta_i) \\
 &= (2\pi)^N \prod_i \int_0^\pi d\theta_i \sin \theta_i \exp(-\beta \mu B \cos \theta_i) .
 \end{aligned}$$

We make a change of variable

$$x_i = \cos \theta_i , \quad dx_i = -\sin \theta_i d\theta_i ,$$

with extremis

$$\theta_i = 0 \rightarrow x_i = 1 , \quad \theta_i = \pi \rightarrow x_i = -1 ,$$

hence

$$\begin{aligned}
 Z &= (2\pi)^N \prod_i \int_{-1}^1 dx_i \exp(-\beta \mu B x_i) \\
 &= (2\pi)^N \left(\frac{\exp(-\beta \mu B x)}{-\beta \mu B} \Big|_{-1}^1 \right)^N \\
 &= (2\pi)^N \left(\frac{1}{-\beta \mu B} \underbrace{(\exp(-\beta \mu B) - \exp(\beta \mu B))}_{-2 \sinh \beta \mu B} \right)^N \\
 &= (2\pi)^N \left(\frac{1}{\beta \mu B} (2 \sinh(\beta \mu B)) \right)^N \\
 &= \left(\frac{4\pi \sinh(\beta \mu B)}{\beta \mu B} \right)^N .
 \end{aligned}$$

q.e.d.

An useful intermediary formula is

$$\ln Z = N \ln \frac{4\pi \sinh(\beta \mu B)}{\beta \mu B} .$$

Proof. In fact,

$$\ln Z = \ln \left(\frac{4\pi \sinh(\beta \mu B)}{\beta \mu B} \right)^N = N \ln \frac{4\pi \sinh(\beta \mu B)}{\beta \mu B} .$$

q.e.d.

The internal energy E is

$$E = -N\mu B(\coth(\beta\mu B) + \frac{1}{\beta\mu B}) .$$

Proof. By (6.9)

$$\begin{aligned} E &= -\frac{\partial \ln Z}{\partial \beta} \\ &= -\frac{\partial}{\partial \beta} N \ln \frac{4\pi \sinh(\beta\mu B)}{\beta\mu B} \\ &= -N \frac{\partial}{\partial \beta} \ln \sinh(\beta\mu B) + N \frac{\partial}{\partial \beta} \ln \beta \\ &= -N\mu B \coth(\beta\mu B) + \frac{N}{\beta} \\ &= -N\mu B(\coth(\beta\mu B) + \frac{1}{\beta\mu B}) . \end{aligned}$$

To study the limit for $\beta \rightarrow 0$ or $T \rightarrow \infty$, we Taylor expand for the variable $x = \beta\mu B$

$$\lim_{x \rightarrow 0} \frac{E}{N\mu B}(x) \simeq 0 ,$$

hence

$$E \xrightarrow{T \rightarrow \infty} 0 .$$

To study the limit for $\beta \rightarrow \infty$ or $T \rightarrow 0$, we Taylor expand for the variable $x = \beta\mu B$

$$\lim_{x \rightarrow \infty} \frac{E}{N\mu B}(x) \simeq -1 ,$$

hence

$$E \xrightarrow{T \rightarrow 0} -N\mu B .$$

q.e.d.

A plot of this is in Figure 10.9.

The Helmholtz free energy F is

$$F = -\frac{N}{\beta} \ln \frac{4\pi \sinh(\beta\mu B)}{\beta\mu B} .$$

Proof. By (6.6)

$$F = -\frac{\ln Z}{\beta} = -\frac{N}{\beta} \ln \frac{4\pi \sinh(\beta\mu B)}{\beta\mu B} .$$

q.e.d.

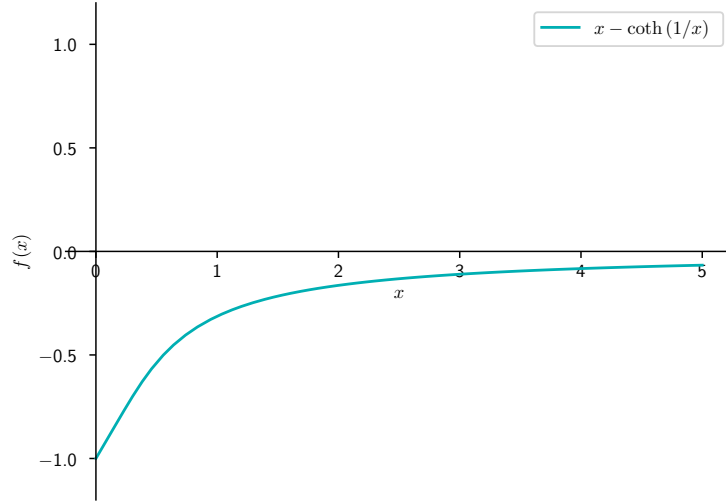


Figure 10.9: A plot of the internal energy E as a function of T . We have used $x = \frac{1}{\beta\mu B}$ and $f(x) = \frac{E}{N\mu B}$.

The entropy S is

$$S = Nk_B \left(\ln \frac{4\pi \sinh(\beta\mu B)}{\beta\mu B} - \beta\mu B (\coth(\beta\mu B) - \frac{1}{\beta\mu B}) \right).$$

Proof. By (6.8)

$$\begin{aligned} S &= \frac{E - F}{T} = \frac{1}{T} \left(-N\mu B (\coth(\beta\mu B) + \frac{1}{\beta\mu B}) + \frac{N}{\beta} \ln \frac{4\pi \sinh(\beta\mu B)}{\beta\mu B} \right) \\ &= Nk_B \left(\ln \frac{4\pi \sinh(\beta\mu B)}{\beta\mu B} - \beta\mu B (\coth(\beta\mu B) - \frac{1}{\beta\mu B}) \right). \end{aligned}$$

q.e.d.

The intrinsic magnetic moment \mathbf{M} is

$$\mathbf{M} = (0, 0, N\mu (\coth(\beta\mu B) - \frac{1}{\beta\mu B})).$$

Proof. By definition, since we have oriented $\mathbf{B} = (0, 0, B)$,

$$M_x = -\frac{\partial F}{\partial B_x} = M_y = -\frac{\partial F}{\partial B_y} = 0,$$

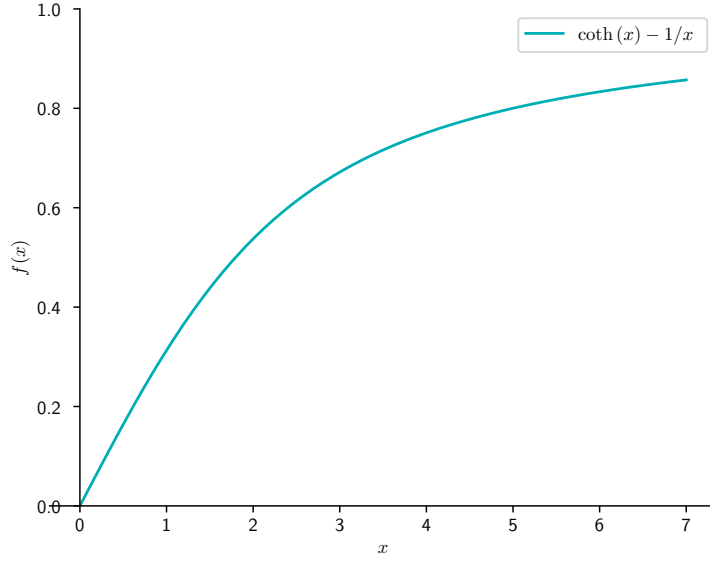


Figure 10.10: A plot of the intrinsic magnetic moment \mathbf{M} as a function of β . We have used $x = \beta\mu B$ and $f(x) = \frac{M_z}{N\mu}$.

but

$$\begin{aligned}
 M_z &= -\frac{\partial F}{\partial B} \\
 &= \frac{\partial}{\partial B} \frac{N}{\beta} \ln \frac{4\pi \sinh(\beta\mu B)}{\beta\mu B} \\
 &= \frac{N}{\beta} \frac{\partial}{\partial \beta} \ln \sinh(\beta\mu B) - \frac{N}{\beta} \frac{\partial}{\partial B} \ln B \\
 &= N\mu \coth(\beta\mu B) - \frac{N}{\beta B} \\
 &= N\mu \left(\coth(\beta\mu B) - \frac{1}{\beta\mu B} \right) .
 \end{aligned}$$

q.e.d.

A plot of this is in Figure 10.10.

The isothermal susceptibility χ_β is

$$\chi_\beta = N\mu^2 \beta \left(\frac{1}{(\beta\mu B)^2} - \frac{1}{\sinh^2(\beta\mu H)} \right) .$$

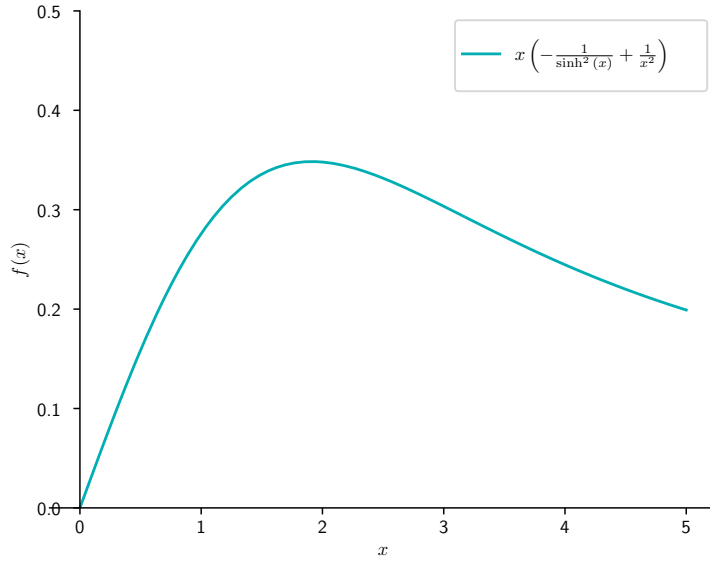


Figure 10.11: A plot of the intrinsic magnetic moment \mathbf{M} as a function of β . We have used $x = \beta\mu B$ and $f(x) = \frac{B\chi_\beta}{N\mu}$.

Proof. By definition

$$\begin{aligned}
 \chi_\beta &= \frac{\partial M}{\partial B} \\
 &= \frac{\partial}{\partial B} N\mu(\coth(\beta\mu B) - \frac{1}{\beta\mu B}) \\
 &= N\mu\left(-\frac{\beta\mu}{\sinh^2(\beta\mu B)} + \frac{\beta\mu}{(\beta\mu B)^2}\right) \\
 &= N\mu^2\beta\left(\frac{1}{(\beta\mu B)^2} - \frac{1}{\sinh^2(\beta\mu H)}\right).
 \end{aligned}$$

q.e.d.

A plot of this is in Figure 10.11.

For $T \rightarrow \infty$, the Curie law is

$$\chi_\beta = \frac{C}{T},$$

where the Curie constant is

$$C = \frac{N\mu^2}{3k_B}.$$

Proof. To study the limit for $\beta \rightarrow 0$ or $T \rightarrow \infty$, we Taylor expand for the variable $x = \beta\mu B$

$$\frac{B\chi_\beta}{N\mu}(x) \simeq \frac{x}{3} + O(x^2) \ ,$$

hence

$$\frac{B\chi_\beta}{N\mu} = \frac{\beta\mu B}{3} \ ,$$

which means

$$\chi_\beta = \frac{N\mu^2}{3k_B} \frac{1}{T} = \frac{C}{T} \ .$$

q.e.d.

Chapter 11

Grancanonical ensemble

11.1 Non-relativistic ideal gas in d-dimensions

Consider an indistinguishable non-relativistic ideal (non-interacting) gas of N particles in an d -dimensional manifold with a finite volume V^N : $\mathcal{M}_N = V^N \times \mathbb{R}^{dN}$.

Recall that the canonical partition function is

$$Z = \frac{V^N}{N! \lambda_T^{dN}} = \frac{V^N}{N!} \left(\frac{2m\pi}{\beta h^2} \right)^{dN/2} .$$

The grancanonical partition function is

$$\mathcal{Z} = \exp\left(\frac{zV}{\lambda_T^d}\right) .$$

Proof. By definition, using the Taylor expansion of the exponential,

$$\mathcal{Z} = \sum_{N=0}^{\infty} z^N Z_N = \sum_{N=0}^{\infty} \frac{1}{N!} \left(\frac{zV}{\lambda_T^d} \right)^N = \exp\left(\frac{zV}{\lambda_T^d}\right) .$$

q.e.d.

The internal energy E is

$$E = \frac{zV}{\lambda_T^d} \frac{d}{2\beta} .$$

Proof. By (7.1)

$$\begin{aligned}
 E &= -\frac{\partial \ln \mathcal{Z}}{\partial \beta} \Big|_z \\
 &= -\frac{\partial}{\partial \beta} \ln \exp\left(\frac{zV}{\lambda_T^d}\right) \\
 &= -\frac{\partial}{\partial \beta} \frac{zV}{\lambda_T^d} \\
 &= -\frac{1}{zV} \frac{\partial}{\partial \beta} \left(\frac{2m\pi}{\beta h^2}\right)^{d/2} \\
 &= -\frac{1}{zV} \left(\frac{2m\pi}{h^2}\right)^{d/2} \frac{\partial}{\partial \beta} \beta^{-d/2} \\
 &= \frac{1}{zV} \left(\frac{h^2}{2m\pi}\right)^{d/2} \frac{d}{2} \beta^{-d/2-1} \\
 &= \frac{zV}{\lambda_T^d} \frac{d}{2\beta} .
 \end{aligned}$$

q.e.d.

The number of particle N is

$$N = \frac{V}{\lambda_T^d} .$$

Proof. By (7.2)

$$N = z \frac{\partial}{\partial z} \ln \mathcal{Z} = z \frac{\partial}{\partial z} \frac{zV}{\lambda_T^d} = \frac{V}{\lambda_T^d} .$$

q.e.d.

The equation of state is

$$p = \frac{z}{\beta \lambda_T^d} .$$

Proof. By definition

$$p = \frac{1}{\beta V} \ln \mathcal{Z} = \frac{1}{\beta V} \frac{zV}{\lambda_T^d} = \frac{z}{\beta \lambda_T^d} .$$

q.e.d.

11.2 Van der Waals potential

Chapter 12

Entropy

12.1 Maxwell-Boltzmann distribution

We can distribute N particle in p boxes in ways

$$W_{n_r}^{(1)} = \frac{N!}{n_1! \dots n_p!} ,$$

whereas there is no restriction for the states

$$W_{n_r}^{(2)} = \prod_r g_r^{n_r} ,$$

hence

$$W_{n_r} = \frac{N!}{n_1! \dots n_p!} \prod_{r=1}^p g_r^{n_r} = N! \prod_{r=1}^p \frac{g_r^{n_r}}{n_r!} .$$

Maximising the constrained entropy, we find the Boltzmann canonical distribution

$$p_r^* = \frac{n_r^*}{N} = \frac{g_r \exp(-\beta E_r)}{\sum_r g_r \exp(-\beta E_r)} .$$

Proof. The entropy is, using the Stirling approximation (B.1),

$$\begin{aligned}
S &= \ln W_{n_r} \\
&= \ln \left(N! \prod_{r=1}^p \frac{g_r^{n_r}}{n_r!} \right) \\
&= \ln N! + \sum_{r=1}^p \ln \frac{g_r^{n_r}}{n_r!} \\
&= \underbrace{\ln N!}_{N \ln N - N} + \sum_{r=1}^p (\ln g_r^{n_r} - \underbrace{\ln n_r!}_{n_r \ln n_r - n_r}) \\
&= N \ln N - N + \sum_{r=1}^p n_r \ln g_r - \sum_{r=1}^p n_r \ln n_r - \cancel{\sum_{r=1}^p n_r} \\
&= N \ln N + \sum_{r=1}^p n_r \ln g_r + \sum_{r=1}^p n_r \ln n_r .
\end{aligned}$$

The constrained entropy is

$$S = N \ln N + \sum_{r=1}^p n_r \ln g_r - \sum_{r=1}^p n_r \ln n_r + \alpha \left(N - \sum_{r=1}^p n_r \right) + \beta \left(E - \sum_{r=1}^p n_r E_r \right) .$$

The maximum is

$$0 = \frac{\partial S}{\partial n_r} = \ln g_r - \ln n_r - 1 - \alpha - \beta E_r ,$$

hence

$$n_r^* = \frac{g_r \exp(-\beta E_r)}{\exp(1 + \alpha)} .$$

We find α by the normalisation condition

$$N = \sum_r n_r^* = \sum_r \frac{g_r \exp(-\beta E_r)}{\exp(1 + \alpha)} ,$$

hence

$$\exp(1 + \alpha) = \frac{\sum_r g_r \exp(-\beta E_r)}{N} .$$

Finally, if we identify $\beta = 1/k_B T$ the probability distribution density is

$$p_r^* = \frac{n_r^*}{N} = \frac{g_r \exp(-\beta E_r)}{\sum_r g_r \exp(-\beta E_r)} .$$

q.e.d.

A plot of this is in Figure 12.1.

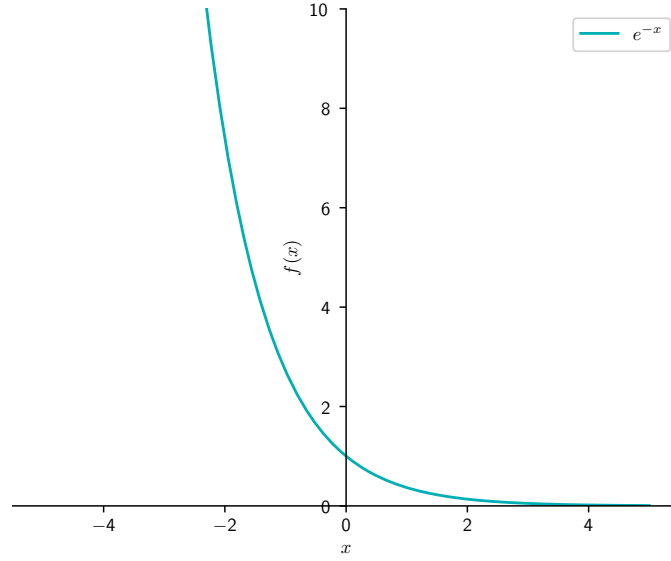


Figure 12.1: A plot of the probability density distribution p_r^* as a function of βE_r . We have used $x = \beta E_r$ and $f(x) = p_r^* \frac{\sum_r g_r \exp(-\beta E_r)}{g_r}$.

12.2 Fermi-Dirac distribution

For fermions, there is no restriction on how we can distribute N particle in p boxes in ways, since they are indistinguishable

$$W_{n_r}^{(1)} = 1 ,$$

whereas we can distribute n_r objects in g_r boxes

$$W_{n_r}^{(2)} = \prod_r \binom{g_r}{n_r} = \prod_r \frac{g_r!}{n_r!(g_r - n_r)!} ,$$

hence

$$W_{n_r} = \prod_r \binom{g_r}{n_r} = \prod_r \frac{g_r!}{n_r!(g_r - n_r)!} .$$

Maximising the constrained entropy, we find the Bose-Einstein distribution

$$n_r^* = \frac{g_r}{\exp(\alpha + \beta E_r) + 1} .$$

Proof. The entropy is, using the Stirling approximation (B.1),

$$\begin{aligned}
S &= \ln W_{n_r} \\
&= \ln \left(\prod_r \frac{g_r!}{n_r!(g_r - n_r)!} \right) \\
&= \sum_r \left(\underbrace{\ln g_r!}_{g_r \ln g_r - g_r} - \underbrace{\ln n_r!}_{n_r \ln n_r - n_r} - \underbrace{\ln(g_r - n_r)!}_{(g_r - n_r) \ln(g_r - n_r) - g_r + n_r} \right) \\
&= \sum_r \left(g_r \ln g_r - g_r - n_r \ln n_r + n_r - (g_r - n_r) \ln(g_r - n_r) + g_r - n_r \right) \\
&= \sum_r \left(g_r \ln g_r - n_r \ln n_r - (g_r - n_r) \ln(g_r - n_r) \right) .
\end{aligned}$$

The constrained entropy is

$$S = \sum_r \left(g_r \ln g_r - n_r \ln n_r - (g_r - n_r) \ln(g_r - n_r) \right) + \alpha \left(N - \sum_{r=1}^p n_r \right) + \beta \left(E - \sum_{r=1}^p n_r E_r \right) .$$

The maximum is

$$\begin{aligned}
0 &= \frac{\partial S}{\partial n_r} \\
&= -\ln n_r - 1 + \ln(g_r - n_r) + 1 - \alpha - \beta E_r \\
&= -\ln n_r + \ln(g_r - n_r) - \alpha - \beta E_r \\
&= \ln\left(\frac{g_r}{n_r} - 1\right) - \alpha - \beta E_r ,
\end{aligned}$$

hence

$$\begin{aligned}
\frac{g_r}{n_r} - 1 &= \exp(\alpha + \beta E_r) , \\
n_r^* &= \frac{g_r}{\exp(\alpha + \beta E_r) + 1} .
\end{aligned}$$

q.e.d.

A plot of this is in Figure 12.2.

12.3 Bose-Einstein distribution

For bosons, there is no restriction on how we can distribute N particle in p boxes in ways, since they are indistinguishable

$$W_{n_r}^{(1)} = 1 ,$$

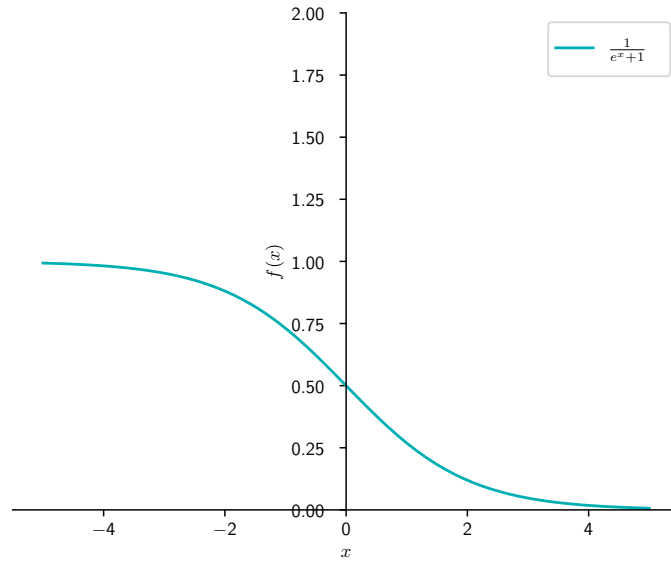


Figure 12.2: A plot of the Fermi-Dirac distribution n_r^* as a function of $\alpha + \beta E_r$. We have used $x = \alpha + \beta E_r$ and $f(x) = \frac{n_r^*}{g_r}$.

whereas we can distribute n_r objects in g_r boxes

$$W_{n_r}^{(2)} = \prod_r \binom{n_r + g_r - 1}{n_r} = \prod_r \frac{(n_r + g_r - 1)!}{n_r! (g_r - 1)!} ,$$

hence

$$W_{n_r} = \prod_r \frac{(n_r + g_r - 1)!}{n_r! (g_r - 1)!} .$$

Maximising the constrained entropy, we find the Bose-Einstein distribution

$$n_r^* = \frac{g_r}{\exp(\alpha + \beta E_r) - 1} .$$

Proof. The entropy is, using the Stirling approximation (B.1),

$$\begin{aligned}
S &= \ln W_{n_r} \\
&= \ln \prod_r \frac{(n_r + g_r - 1)!}{n_r! (g_r - 1)!} \\
&= \sum_r \left(\underbrace{\ln(n_r + g_r - 1)!}_{(n_r + g_r - 1) \ln(n_r + g_r - 1) - n_r - g_r + 1} - \underbrace{\ln n_r!}_{n_r \ln n_r - n_r} - \underbrace{\ln(g_r - 1)!}_{(g_r - 1) \ln(g_r - 1) - g_r + 1} \right) \\
&= \sum_r \left((n_r + g_r - 1) \ln(n_r + g_r - 1) - \cancel{n_r} - \cancel{g_r} + \cancel{1} \right. \\
&\quad \left. - n_r \ln n_r + \cancel{n_r} - (g_r - 1) \ln(g_r - 1) + \cancel{g_r} - \cancel{1} \right) \\
&= \sum_r \left((n_r + g_r - 1) \ln(n_r + g_r - 1) \ln n_r - n_r \ln n_r - (g_r - 1) \ln(g_r - 1) \right)
\end{aligned}$$

The constrained entropy is

$$\begin{aligned}
S &= \sum_r \left((n_r + g_r - 1) \ln(n_r + g_r - 1) \ln n_r - n_r \ln n_r - (g_r - 1) \ln(g_r - 1) \right) \\
&\quad + \alpha \left(N - \sum_{r=1}^p n_r \right) + \beta \left(E - \sum_{r=1}^p n_r E_r \right) .
\end{aligned}$$

The maximum is

$$\begin{aligned}
0 &= \frac{\partial S}{\partial n_r} \\
&= \ln(n_r + g_r - 1) + \cancel{1} - \ln n_r - \cancel{1} - \alpha - \beta E_r \\
&= \ln(n_r + g_r - 1) - \ln n_r - \alpha - \beta E_r \\
&= \ln\left(\frac{g_r - 1}{n_r} + 1\right) - \alpha - \beta E_r ,
\end{aligned}$$

hence, for $g_r \gg 1$,

$$\frac{g_r - 1}{n_r} + 1 = \exp(\alpha + \beta E_r) ,$$

$$n_r^* = \frac{g_r - 1}{\exp(\alpha + \beta E_r) - 1} \simeq \frac{g_r}{\exp(\alpha + \beta E_r) - 1} .$$

q.e.d.

A plot of this is in Figure 12.3.

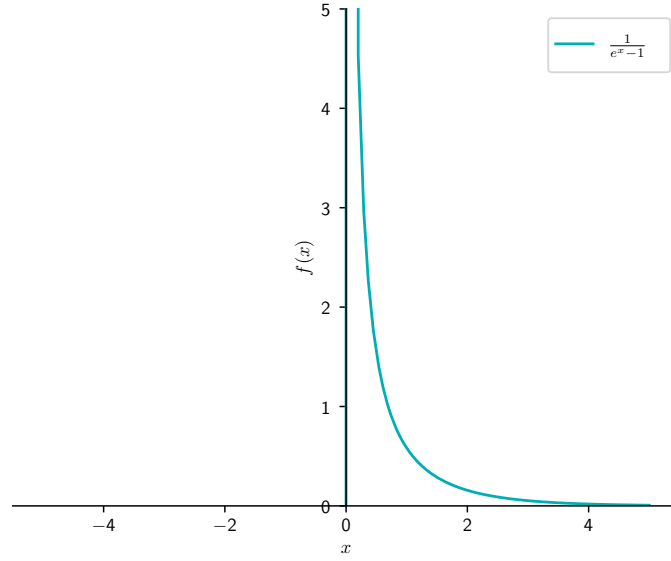


Figure 12.3: A plot of the Bose-Einstein distribution n_r^* as a function of $\alpha + \beta E_r$. We have used $x = \alpha + \beta E_r$ and $f(x) = \frac{n_r^*}{g_r}$.

12.4 Two-levels system

Consider a system composed by 2 levels of energies $\epsilon_+ = +\epsilon$ and $\epsilon_- = -\epsilon$. The constraints are

$$E = \epsilon(n_+ - n_-) , \quad N = n_+ + n_- .$$

They can be inverted as

$$n_+ = \frac{N}{2} + \frac{E}{2\epsilon} , \quad n_- = \frac{N}{2} - \frac{E}{2\epsilon} . \quad (12.1)$$

We can distribute N objects in n_+ boxes

$$\Omega(E) = \binom{N}{n_+} = \frac{N!}{n_+!(N - n_+)!} = \frac{N!}{n_+!n_-!} .$$

The entropy is

$$S = -Nk_B \left(\left(\frac{1}{2} + \frac{E}{2\epsilon N} \right) \ln \left(\frac{1}{2} + \frac{E}{2\epsilon N} \right) + \left(\frac{1}{2} - \frac{E}{2\epsilon N} \right) \ln \left(\frac{1}{2} - \frac{E}{2\epsilon N} \right) \right) .$$

Proof. By definition, using the Stirling approximation (B.1),

$$\begin{aligned}
\frac{S}{k_B} &= \ln \Omega(E) \\
&= \ln \frac{N!}{n_+! n_-!} \\
&= \underbrace{\ln N!}_{N \ln N - N} - \underbrace{\ln n_+!}_{n_+ \ln n_+ - n_+} - \underbrace{\ln n_-!}_{n_- \ln n_- - n_-} \\
&= N \ln N - N - n_+ \ln n_+ + n_+ - n_- \ln n_- + n_- \\
&= N \ln N - n_+ \ln n_+ - n_- \ln n_- \\
&= (n_+ + n_-) \ln N - n_+ \ln n_+ - n_- \ln n_- \\
&= n_+ \ln \frac{N}{n_+} + n_- \ln \frac{N}{n_-} \\
&= \left(\frac{N}{2} + \frac{E}{2\epsilon}\right) \ln \frac{N}{\frac{N}{2} + \frac{E}{2\epsilon}} + \left(\frac{N}{2} - \frac{E}{2\epsilon}\right) \ln \frac{N}{\frac{N}{2} - \frac{E}{2\epsilon}} \\
&= N \left(\left(\frac{1}{2} + \frac{E}{2\epsilon N}\right) \ln \frac{1}{\frac{1}{2} + \frac{E}{2\epsilon N}} + \left(\frac{1}{2} - \frac{E}{2\epsilon N}\right) \ln \frac{1}{\frac{1}{2} - \frac{E}{2\epsilon N}} \right) \\
&= -N \left(\left(\frac{1}{2} + \frac{E}{2\epsilon N}\right) \ln \left(\frac{1}{2} + \frac{E}{2\epsilon N}\right) + \left(\frac{1}{2} - \frac{E}{2\epsilon N}\right) \ln \left(\frac{1}{2} - \frac{E}{2\epsilon N}\right) \right).
\end{aligned}$$

q.e.d.

A plot of this is in Figure 12.4.

The temperature is

$$T = \frac{2\epsilon}{k_B} \frac{1}{\ln \frac{\frac{1}{2} - \frac{E}{2\epsilon N}}{\frac{1}{2} + \frac{E}{2\epsilon N}}}.$$

Proof. Using (4.3)

$$\begin{aligned}
T &= \left(\frac{\partial S}{\partial E}\right)^{-1} \\
&= -\left(\frac{k_B}{2\epsilon} \ln \left(\frac{1}{2} + \frac{E}{2\epsilon N}\right) + \frac{k_B}{2\epsilon} \ln \left(\frac{1}{2} - \frac{E}{2\epsilon N}\right) - \frac{k_B}{2\epsilon}\right)^{-1} \\
&= -\left(\frac{k_B}{2\epsilon} \ln \frac{\frac{1}{2} + \frac{E}{2\epsilon N}}{\frac{1}{2} - \frac{E}{2\epsilon N}}\right)^{-1} \\
&= -\frac{2\epsilon}{k_B} \frac{1}{\ln \frac{\frac{1}{2} + \frac{E}{2\epsilon N}}{\frac{1}{2} - \frac{E}{2\epsilon N}}} \\
&= \frac{2\epsilon}{k_B} \frac{1}{\ln \frac{\frac{1}{2} - \frac{E}{2\epsilon N}}{\frac{1}{2} + \frac{E}{2\epsilon N}}}.
\end{aligned}$$

q.e.d.

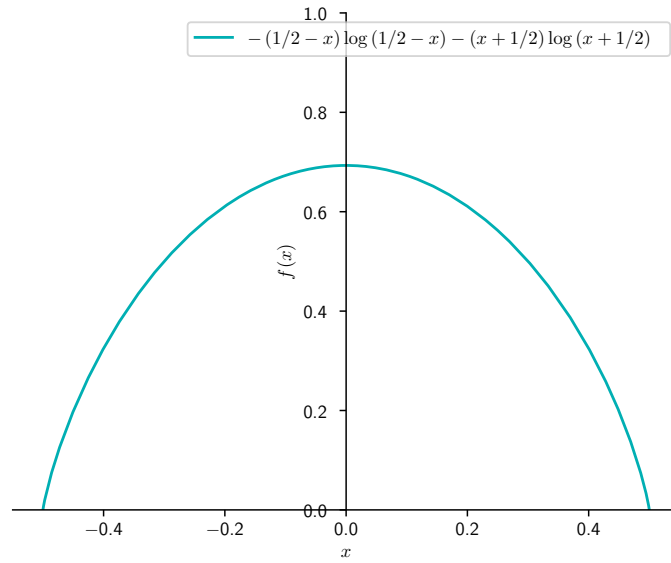


Figure 12.4: A plot of the entropy S as a function of E . We have used $x = \frac{E}{2\epsilon N}$ and $f(x) = \frac{S}{Nk_B}$.

A plot of this is in Figure 12.5.

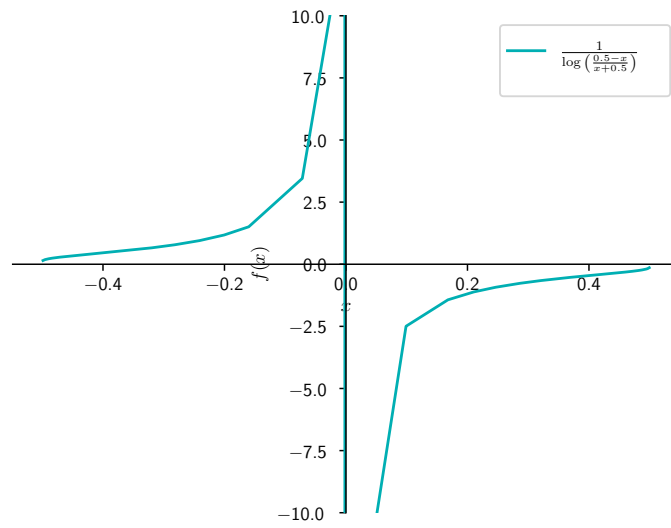


Figure 12.5: A plot of the temperature T as a function of E . We have used $x = \frac{E}{2\epsilon N}$ and $f(x) = \frac{k_B T}{2\epsilon}$.

Part IV

Quantum mechanics

Chapter 13

Quantum Mechanics

In this chapter, we will study the mathematical framework necessary to study quantum statistical mechanics.

13.1 States

In quantum mechanics, a pure state of a quantum particle is represented by a normalised vector in a Hilbert space $|\psi\rangle \in \mathcal{H}$. An Hilbert space \mathcal{H} is a vector space on \mathbb{C} , i.e. in which a linear superposition of is still in the space

$$\lambda|\psi\rangle + \mu|\phi\rangle \in \mathcal{H} , \quad \forall |\psi\rangle, |\phi\rangle \in \mathcal{H} , \forall \lambda, \mu \in \mathbb{C} ,$$

endowed with a scalar product $\langle\psi|\phi\rangle$. In particular, via the scalar product, it is possible to associate a norm to the state, which is set to zero by the probability interpretation $||\psi||^2 = \langle\psi|\psi\rangle = 1$. In the Schrodinger representation, this means that the wave function is a square-integrable function $\psi(t, \mathbf{x}) \in L^2(\mathbb{R}^d)$, the probability interpretation tells us that $|\psi(t, \mathbf{x})|^2$ is the probability density to find the particle in a volume element $d^d x$ at time t and the normalisation condition that the total probability to find the particle in the whole \mathbb{R}^d is 1

$$\int_{\mathbb{R}^d} d^d x |\psi(t, x)|^2 = 1 .$$

However, by the normalisation condition, a state is not associated to a single vector, but a class of equivalence of them, called a ray in the Hilbert space, since two states are physically equivalent if $|\psi'\rangle = \exp(i\varphi)|\psi\rangle$, because their norms are the same.

Projectors

To remove this ambiguity, we introduce the notion of projection operators or projectors, which uniquely determine a state

$$P_\psi = \frac{|\psi\rangle\langle\psi|}{\langle\psi|\psi\rangle} ,$$

which for normalisation states becomes

$$P_\psi = |\psi\rangle\langle\psi| . \quad (13.1)$$

Proof. If $|\psi'\rangle = \exp(i\varphi)|\psi\rangle$ and $\langle\psi'| = \exp(-i\varphi)\langle\psi|$, we have

$$P_{\psi'} = |\psi'\rangle\langle\psi'| = \exp(i\varphi)|\psi\rangle\exp(-i\varphi)\langle\psi| = |\psi\rangle\langle\psi| = P_\psi .$$

q.e.d.

It projects onto the 1-dimensional subspace $\mathcal{H}_\psi = \{\lambda|\psi\rangle : \lambda \in \mathbb{C}\}$ generated by the state $|\psi\rangle$

$$P_\psi : \mathcal{H} \rightarrow \mathcal{H}_\psi .$$

Proof. In fact, $\forall |\phi\rangle \in \mathcal{H}$, we decomposed the Hilbert space into the direct orthogonal sum of the subspace spanned by \mathcal{H}_ψ and its orthogonal complement \mathcal{H}^\perp :

$$|\phi\rangle = \alpha|\psi\rangle + \beta|\psi^\perp\rangle ,$$

where $|\psi\rangle \in \mathcal{H}_\psi$, $|\psi^\perp\rangle \in \mathcal{H}^\perp$ and $\langle\psi|\psi^\perp\rangle = 0$. Therefore, the action of the projector is

$$P_\psi|\phi\rangle = \alpha P_\psi|\psi\rangle + \beta P_\psi|\psi^\perp\rangle = \alpha|\psi\rangle \underbrace{\langle\psi|\psi\rangle}_1 + \beta|\psi\rangle \underbrace{\langle\psi|\psi^\perp\rangle}_0 = \alpha|\psi\rangle \in \mathcal{H}_\psi .$$

q.e.d.

Moreover, since the projectors is orthogonal, we can define the projector onto the orthogonal subspace as $P_\psi^\perp = \mathbb{I} - P_\psi$ such that it satisfies $P_\psi P_\psi^\perp = P_\psi^\perp P_\psi = 0$. This can be generalised for a generic set of orthogonal subspaces. In fact, given an orthonormal basis $\{|e_n\rangle\}$, a projector onto an element of this basis is $P_n = |e_n\rangle\langle e_n|$ and the orthonormality condition reads as $P_n P_m = P_m P_n = 0$ for $n \neq m$.

It satisfies the following properties

1. boundness, i.e.

$$\|P_\psi\| < \infty ,$$

2. hermiticity, i.e.

$$P_\psi^\dagger = P_\psi ,$$

3. idempotence, i.e.

$$P_\psi^2 = P_\psi , \quad (13.2)$$

4. positive defined, i.e. $\forall |\phi\rangle \in \mathcal{H}$

$$\langle\phi|P_\psi|\phi\rangle \geq 0 ,$$

5. trace equals to 1, i.e.

$$\text{tr } P_\psi = 1 .$$

Actually, there is a theorem that ensures that an operators such that it satisfies these 5 conditions is indeed a projector.

Proof. For the boundness, $\forall |\phi\rangle \in \mathcal{H}$

$$||P_\psi|\phi\rangle||^2 = \langle\phi|P_\psi^\dagger P_\psi|\phi\rangle = \langle\phi|\psi\rangle \underbrace{\langle\psi|\psi\rangle}_1 \langle\psi|\phi\rangle = |\langle\psi|\phi\rangle|^2 \leq ||\phi||^1 ,$$

hence

$$||P_\psi|| = \frac{||P_\psi|\phi\rangle||}{||\phi||} \leq 1 .$$

For the hermiticity

$$P_\psi^\dagger = (|\psi\rangle\langle\psi|)^\dagger = \langle\psi|^\dagger|\psi\rangle^\dagger = |\psi\rangle\langle\psi| = P_\psi .$$

For the idempotence

$$P_\psi^2 = (|\psi\rangle\langle\psi|)^2 = |\psi\rangle \underbrace{\langle\psi|\psi\rangle}_1 \langle\psi| = |\psi\rangle\langle\psi| = P_\psi .$$

For the positive definedness

$$\langle\phi|P_\psi|\phi\rangle = \langle\phi|\psi\rangle\langle\psi|\phi\rangle = |\langle\psi|\phi\rangle|^2 \geq 0 .$$

For the trace, since it is independent from the choice of the basis, we choose $|\psi\rangle = |\psi_1\rangle$ such that $\langle\psi|\psi_n\rangle = \delta_{n,1}$ and

$$\text{tr } P_\psi = \sum_{n=0}^{\infty} \langle\psi_n|P_\psi|\psi_n\rangle = \sum_{n=0}^{\infty} \underbrace{\langle\psi_n|\psi\rangle}_{\delta_{n,1}} \langle\psi|\psi_n\rangle = \sum_{n=0}^{\infty} \underbrace{\delta_{n,1}}_{n=1} \langle\psi|\psi_n\rangle = \langle\psi|\psi_1\rangle = \langle\psi_1|\psi_1\rangle = 1 .$$

q.e.d.

Trace

Given an orthonormal basis $\{|e_n\rangle\}_{n=1}^{\infty}$ of a separable Hilbert space, the trace is defined as

$$\text{tr } A = \sum_{n=1}^{\infty} A_{nn} = \sum_{n=1}^{\infty} \langle e_n|A|e_n\rangle .$$

It may happen that this series is not convergent. If it is convergent, the operator A is called a trace-class operator. Furthermore, if it is absolute convergent, the trace is independent on the choice of the basis. Recall that in the finite-dimensional case, the trace of a matrix is always convergent and independent on the choice of the basis.

13.2 Observables

An observable is a linear hermitian operator \hat{A} acting on the Hilbert space. We require the self-adjointness because, by the spectral theorem, they are always diagonalisable with a positive spectrum. This means that its eigenvalues are real and it always admit an orthonormal eigenbasis $\{|\psi_n\rangle\}$

$$A|\psi_n\rangle = \lambda_n|\psi_n\rangle , \quad (13.3)$$

where $\lambda_n \in \mathbb{R}$. In this way, $\forall |\phi\rangle \in \mathcal{H}$, we can expand it into the eigenbasis

$$|\phi\rangle = \sum_{n=1}^{\infty} c_n |\psi_n\rangle , \quad (13.4)$$

where $c_n \in \mathbb{C}$.

The eigenprojectors, defined as

$$P_n = |\psi_n\rangle\langle\psi_n| ,$$

satisfy the following properties

1. self-adjointness, i.e.

$$P_n^\dagger = P_n ,$$

2. orthonormality, i.e.

$$P_n P_m = \delta_{nm} P_n ,$$

3. completeness relation, i.e.

$$\sum_{n=0}^{\infty} P_n = \mathbb{I} , \quad (13.5)$$

4. spectral decomposition, i.e.

$$\hat{A} = \sum_{n=0}^{\infty} \lambda_n P_n . \quad (13.6)$$

Measure

Prepare a quantum system in a state $|\psi\rangle$. A measurement of an observable \hat{A} has outcomes corresponding to its eigenvalues λ_n with probability $p_n = |c_n|^2$. Recall that λ_n are the coefficients in (13.3) and c_n in (13.4). Its average value is

$$\langle A \rangle = \langle \psi | \hat{A} | \psi \rangle = \sum_n \lambda_n |c_n|^2 = \sum_n \lambda_n p_n , \quad (13.7)$$

whereas its standard deviation is

$$(\Delta A)^2 = \langle A^2 \rangle - \langle A \rangle^2 .$$

Proof. In fact, using (13.3) and (13.4)

$$\begin{aligned} \langle A \rangle &= \langle \psi | \hat{A} | \psi \rangle \\ &= \sum_{m=0}^{\infty} c_m^* \langle \psi_m | \sum_{n=0}^{\infty} c_n \underbrace{\hat{A} | \psi_n \rangle}_{\lambda_n | \psi_n \rangle} \\ &= \sum_{n=0}^{\infty} \sum_{m=0}^{\infty} \lambda_n c_m^* c_n \underbrace{\langle \psi_m | \psi_n \rangle}_{\delta_{nm}} \\ &= \sum_{n=0}^{\infty} \sum_{m=0}^{\infty} \lambda_n c_m^* c_n \underbrace{\delta_{nm}}_{n=m} \\ &= \sum_{n=0}^{\infty} \lambda_n \underbrace{c_n^* c_n}_{|c_n|^2} \\ &= \sum_n \lambda_n |c_n|^2 . \end{aligned}$$

q.e.d.

Notice that measurement in quantum mechanics is a destructive process, since the wave function collapses into one of the eigenstates.

Time evolution

Time evolution of a quantum system is governed by a special observable, the hamiltonian \hat{H} , through the Schroedinger equation

$$i\hbar \frac{\partial}{\partial t} |\psi(t)\rangle = \hat{H} |\psi(t)\rangle .$$

Notice that this equation is linear, consistent with the superposition principle. It is also at first-order in time, meaning that once the initial condition is fixed, $|\psi(t)\rangle$ is completely determined.

Moreover, for a time-independent hamiltonian, time evolution can be equivalently expressed by a unitary operator $\hat{U}(t)$

$$|\psi(t)\rangle = \hat{U}(t) |\psi(0)\rangle , \tag{13.8}$$

where $\hat{U}(t) = \exp(\frac{i}{\hbar}\hat{H}t)$. Since it is unitary

$$\hat{U}^\dagger(t) = \exp(-\frac{i}{\hbar}\hat{H}t) = \hat{U}(-t) = \hat{U}^{-1}(t) ,$$

it preserves the probability.

13.3 Density matrices and mixed states

The projector (13.1) is also called a density matrix ρ_ψ . In terms of the density matrix, the average value (13.7) of an operator \hat{A} is

$$\langle A \rangle = \langle \psi | \hat{A} | \psi \rangle = \text{tr}(\hat{A} \rho_\psi) .$$

Proof. In fact, using (13.5)

$$\begin{aligned} \langle A \rangle &= \langle \psi | \hat{A} | \psi \rangle \\ &= \langle \psi | \mathbb{I} \hat{A} | \psi \rangle \\ &= \sum_{n=0}^{\infty} \langle \psi | P_n \hat{A} | \psi \rangle \\ &= \sum_{n=0}^{\infty} \langle \psi | \psi_n \rangle \langle \psi_n | \hat{A} | \psi \rangle \\ &= \sum_{n=0}^{\infty} \langle \psi_n | \hat{A} | \psi \rangle \underbrace{\langle \psi | \psi_n \rangle}_{\rho_\psi} \\ &= \sum_{n=0}^{\infty} \langle \psi_n | \hat{A} \rho_\psi | \psi_n \rangle \\ &= \text{tr}(\hat{A} \rho_\psi) , \end{aligned}$$

where we have exchanged brackets because they are only numbers.

q.e.d.

The time evolution of the density matrix is

$$\rho_\psi(t) = \exp(-\frac{i}{\hbar}\hat{H}t) \rho_\psi(0) \exp(\frac{i}{\hbar}\hat{H}t) .$$

Proof. In fact, using (13.8)

$$\rho_\psi(t) = |\psi(t)\rangle \langle \psi(t)| = \exp(-\frac{i}{\hbar}\hat{H}t) \underbrace{|\psi(0)\rangle \langle \psi(0)|}_{\rho_\psi(0)} \exp(\frac{i}{\hbar}\hat{H}t) = \exp(-\frac{i}{\hbar}\hat{H}t) \rho_\psi(0) \exp(\frac{i}{\hbar}\hat{H}t) .$$

q.e.d.

Mixed states

A mixed state belonging to a classical mixture is a system which can be found in a state $|\psi_n\rangle$ with a probability p_n

$$\{|\psi_n\rangle, p_n\} ,$$

where $p_n \geq 0$ and $\sum_{n=0}^{\infty} p_n = 1$. The difference from a pure state is that, in a mixed state, the system is in a classical fixed state before the measurement whereas in a pure state, the state is in a quantum superposition. The density matrix of a mixed state is

$$\rho = \sum_n p_n |\psi_n\rangle \langle \psi_n| = \sum_n p_n \rho_n , \quad (13.9)$$

It defines a statistical ensemble.

Similarly to the pure state case, it satisfies the following properties

1. boundness, i.e.

$$||\rho|| < \infty ,$$

2. hermiticity, i.e.

$$\rho^\dagger = \rho ,$$

3. positive defined, i.e. $\forall |\phi\rangle \in \mathcal{H}$

$$\langle \phi | \rho | \phi \rangle \geq 0 ,$$

4. trace equals to 1, i.e.

$$\text{tr } \rho = 1 .$$

However, the idempotence property (13.2) is a particular property of only pure states. There is a theorem that states that a state is pure if and only if $\rho^2 = \rho$.

Proof. In the simple case of orthogonal states $|\psi_n\rangle$, i.e. $\langle \psi_n | \psi_m \rangle = \delta_{nm}$, we have

$$\begin{aligned} \rho^2 &= \sum_n p_n |\psi_n\rangle \langle \psi_n| \sum_m p_m |\psi_m\rangle \langle \psi_m| \\ &= \sum_n \sum_m p_n p_m |\psi_n\rangle \underbrace{\langle \psi_n | \psi_m \rangle}_{\delta_{nm}} \langle \psi_m| \\ &= \sum_n \sum_m p_n p_m |\psi_n\rangle \underbrace{\delta_{nm}}_{n=m} \langle \psi_m| \\ &= \sum_n p_n^2 |\psi_n\rangle \langle \psi_n| \\ &= \sum_n p_n \rho_n . \end{aligned}$$

This means that if $\rho^2 = \rho$, we obtain

$$p_n^2 = p_n ,$$

which means that $p_{\bar{n}} = 1$ for a single \bar{n} and for all the others $p_n = 0$ for $n \neq \bar{n}$, but this is indeed a pure state $\rho = |\psi_{\bar{n}}\rangle\langle\psi_{\bar{n}}|$. q.e.d.

However, the average value of an observable is the same as the pure states

$$\langle\hat{A}\rangle = \langle\psi|\hat{A}|\psi\rangle = \text{tr}(\rho\hat{A}) . \quad (13.10)$$

Proof. In fact,

$$\langle\hat{A}\rangle = \sum_n p_n \langle\hat{A}\rangle_n = \sum_n p_n \text{tr}(\hat{A}\rho_n) = \text{tr}(\hat{A} \underbrace{\sum_n p_n \rho_n}_{\rho}) = \text{tr}(\hat{A}\rho) ,$$

where we have used the linearity of the trace. q.e.d.

Notice that in the classical case, the average value of an observable is (4.1)

$$\langle f \rangle = \int_{\mathcal{M}} d^d x f(x) \rho(x) ,$$

which shows that, in the quantum case, we have substituted the integral with the trace, the function with the observable operator and the density distribution with the density matrix.

13.4 Composite systems

Consider a quantum system composed by 2 particles. The total Hilbert space is the tensor product between the 2 single particle Hilbert spaces

$$\mathcal{H}_{tot} = \mathcal{H}_1 \otimes \mathcal{H}_2 .$$

Given an orthonormal basis for each Hilbert space $\{|\psi_n\rangle\} \in \mathcal{H}_1$ and $\{|\phi_m\rangle\} \in \mathcal{H}_2$, the orthonormal basis for the total Hilbert space is

$$\{|\psi_n\rangle_1 |\phi_m\rangle_2 = |\psi_n \phi_m\rangle\} ,$$

such that a generic state can be expanded into this basis, $\forall |\phi\rangle \in \mathcal{H}_{tot}$

$$|\phi\rangle = \sum_n \sum_m \alpha_{nm} |\psi_n \phi_m\rangle ,$$

where $\alpha_{nm} \in \mathbb{C}$ and the normalisation condition reads $\sum_{nm} |\alpha_{nm}|^2 = 1$.

If the 2 particle are identical, we have $\mathcal{H}_1 = \mathcal{H}_2 = \mathcal{H}$. Therefore $\mathcal{H}_{tot} = \mathcal{H}^{\otimes 2}$.

The scalar product between two sparable is

$$\langle \psi_n \phi_m | \psi_{n'} \phi_{m'} \rangle = \langle \psi_n | \psi_{n'} \rangle_1 \langle \phi_m | \phi_{m'} \rangle_2 ,$$

such that if the two states are orthonormal we have

$$\langle \psi_n \phi_m | \psi_{n'} \phi_{m'} \rangle = \langle \psi_n | \psi_{n'} \rangle_1 \langle \phi_m | \phi_{m'} \rangle_2 = \delta_{nn'} \delta_{mm'} .$$

By lincerity, we can generalised this construction for N particles. However, for infinite dimensional Hilbert spaces, we need the convergence of $\sum_{nm} |\alpha_{nm}|^2$ in order to remain in a Hilbert space. The total Hilbert space is

$$\mathcal{H}_{tot} = \mathcal{H}_1 \otimes \dots \otimes \mathcal{H}_N ,$$

its orthonormal basis is

$$|e_{n_1}\rangle \dots |e_{n_N}\rangle$$

and its scalar product is

$$\langle \cdot | \cdot \rangle = \prod_k \langle \cdot | \cdot \rangle_k .$$

A generic state can be expanded into the orthonormal basis, $\forall |\phi\rangle \in \mathcal{H}_{tot}$

$$|\phi\rangle = \sum_{n_1, \dots, n_N} \alpha_{n_1, \dots, n_N} |e_{n_1}\rangle \dots |e_{n_N}\rangle .$$

If all the particles are identical, we have $\mathcal{H}_1 = \dots = \mathcal{H}_N = \mathcal{H}$. Therefore $\mathcal{H}_{tot} = \mathcal{H}^{\otimes N}$.

N particles

Explicitly, a single particle lives in \mathbb{R}^3 and its Hilbert space is $\mathcal{H} = L^2(\mathbb{R}^3) \ni \psi(x)$. The scalar product is

$$\langle \psi | \phi \rangle = \int d^3x \psi^*(x) \phi(x) ,$$

where the normalisation condition is

$$||\psi||^2 = \langle \psi | \psi \rangle = \int_{\mathbb{R}^3} d^3x |\psi(x)|^2 < \infty .$$

For N distinguishable particles, the total Hilbert space is $\mathcal{H}_{tot} = \mathcal{H} \otimes \dots \otimes \mathcal{H}$ and a generic state is $|\psi_{n_1} \dots \psi_{n_N}\rangle$ where $|\psi_{n_j}\rangle$ is a single particle state. Explicitly, N distinguishable particle live in \mathbb{R}^{3N} and their Hilbert space is $\mathcal{H}_N = L^2(\mathbb{R}^3) \otimes \dots \otimes L^2(\mathbb{R}^3) = L^2(\mathbb{R}^{3N}) \ni \psi(x_1, \dots, x_N)$. Therefore, an orthonormal basis is $\{u_{\alpha_1(x_1)} \dots u_{\alpha_N(x_N)} = u_{\alpha_1 \dots \alpha_N}(x_1, \dots, x_N)\}$ where $\{u_{\alpha}(x)\}$ is the single particle orthonormal basis. A generic state can be expanded in this basis as

$$\psi(x_1, \dots, x_N) = \sum_{\alpha_1 \dots \alpha_N} c_{\alpha_1 \dots \alpha_N} u_{\alpha_1 \dots \alpha_N}(x_1, \dots, x_N) .$$

Distinguishable and indistinguishable particles

Choosing $\alpha_1 = a$ and $\alpha_2 = b$ or viceversa, we obtain

$$u_{\alpha_1=a}(x_1)u_{\alpha_2=b}(x_2) \neq u_{\alpha_1=b}(x_1)u_{\alpha_2=a}(x_2) ,$$

but if the particles are indistinguishable, we have

$$u_{\alpha_1=a}(x_1)u_{\alpha_2=b}(x_2) \propto u_{\alpha_1=b}(x_1)u_{\alpha_2=a}(x_2) ,$$

where the proportionality factor is due to the fact that states are the same up to a global phase factor. This means that they are invariant under permutations

$$\psi(P(x_1, \dots, x_N)) = \exp(i\alpha_P)\psi(x_1, \dots, x_N) ,$$

since in this way

$$|\psi(P(x_1, \dots, x_N))|^2 = |\psi(x_1, \dots, x_N)|^2 ,$$

where P belongs to the permutation group. In the next chapter, we will evaluate the phase factor α_P .

Chapter 14

Permutation group

14.1 Permutation as consecutive transpositions

The permutation of N elements form a group P_N . The composition of 2 permutations PP' is defined as the permutation obtained by applying first P and then P' . The identity permutation \mathbb{I} does not change anything and the inverse is the permutation such that $PP^{-1} = \mathbb{I}$.

Example 14.1. Given 4 numbers $(1, 2, 3, 4)$,

1. the identity is

$$(1, 2, 3, 4) \xrightarrow{\mathbb{I}} (1, 2, 3, 4) ,$$

2. the inverse of

$$(1, 2, 3, 4) \xrightarrow{P} (4, 2, 1, 3) ,$$

is

$$(1, 2, 3, 4) \xrightarrow{P^{-1}} (3, 2, 4, 1) ,$$

since

$$(1, 2, 3, 4) \xrightarrow{P} (4, 3, 1, 3) \xrightarrow{P^{-1}} (1, 2, 3, 4) ,$$

This group is generated by transposition, since any permutation $P \in P_N$ can be decomposed and written as a consecutive swap of two near elements

$$\sigma_i: (1, 2, \dots, i, i+1, \dots, N) \mapsto (1, 2, \dots, i+1, i, \dots, N) ,$$

in the following way

$$P = \sigma_{\alpha_1} \dots \sigma_{\alpha_k} . \tag{14.1}$$

However, this decomposition is not unique but the number of transposition in its decomposition is always even or odd. Therefore, we can define the sign of a permutation $\forall P \in P_N$

$$\text{sgn}(P) = \begin{cases} +1 & \text{even number of transposition in its decomposition} \\ -1 & \text{odd number of transposition in its decomposition} \end{cases} .$$

Transpositions follow the properties

1. if $|i - j| > 2$, which means that they are not next to each other,

$$\sigma_i \sigma_j = \sigma_j \sigma_i , \quad (14.2)$$

- 2.

$$\sigma_i \sigma_{i+1} \sigma_i = \sigma_{i+1} \sigma_i \sigma_{i+1} , \quad (14.3)$$

- 3.

$$(\sigma_i)^2 = \mathbb{I} . \quad (14.4)$$

Proof. A transposition can be pictorially seen in Figure 14.1. The proofs can be seen in Figure 14.2, Figure 14.4 and Figure 14.3. q.e.d.

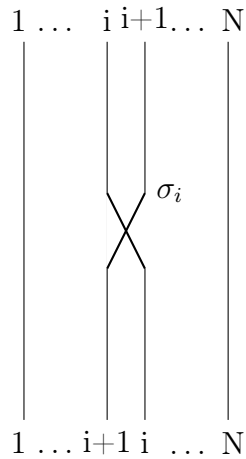


Figure 14.1: A pictorial diagram of a transposition σ_i

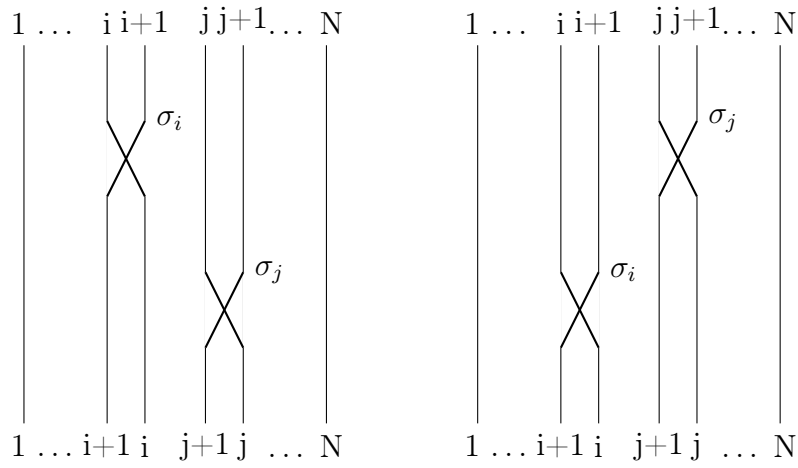


Figure 14.2: A pictorial diagram of a transposition $\sigma_i \sigma_j$ on the left and $\sigma_j \sigma_i$ on the right, where $|i - j| > 2$

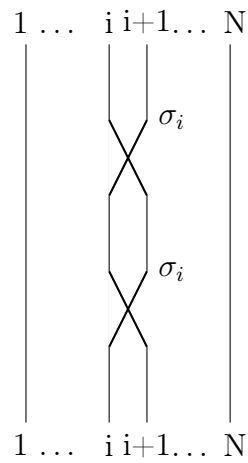


Figure 14.3: A pictorial diagram of a transposition $(\sigma_i)^2$

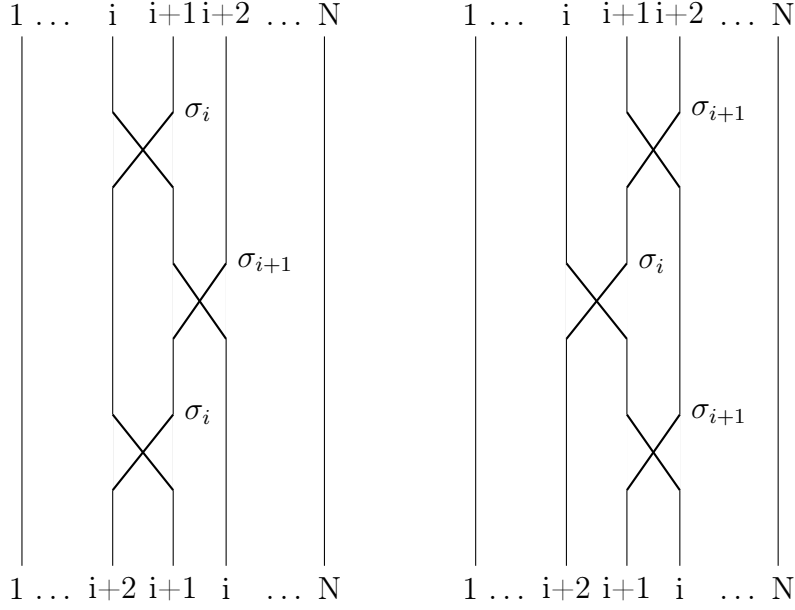


Figure 14.4: A pictorial diagram of a transposition $\sigma_i \sigma_{i+1} \sigma_i$ on the left and $\sigma_{i+1} \sigma_i \sigma_{i+1}$ on the right

14.2 Bosons and fermions

Hence, we can calculate explicitly (13.4), which is

$$\alpha_P = \alpha_1 + \dots \alpha_N , \quad (14.5)$$

where α_i is the phase factor of a transposition σ_{α_i}

$$\psi(\sigma_{\alpha_i}(x_1, \dots x_N)) = \exp(i\alpha_i) \psi(x_1, \dots x_N) .$$

Proof. In fact, using (14.1)

$$\begin{aligned} \psi(P(x_1, \dots x_N)) &= \psi((\sigma_{\alpha_1} \dots \sigma_{\alpha_N})(x_1, \dots x_N)) \\ &= \exp(i\alpha_1) \psi((\sigma_{\alpha_2} \dots \sigma_{\alpha_N})) \\ &\quad \vdots \\ &= \exp(i\alpha_1) \dots \exp(i\alpha_N) \psi(x_1, \dots x_N) \\ &= \exp(i(\alpha_1 + \dots \alpha_N)) \psi(x_1, \dots x_N) \\ &= \exp(i\alpha_P) \psi(x_1, \dots x_N) . \end{aligned}$$

q.e.d.

Furthermore, there are only two possibilities for α_P

1. $\alpha_P = 0$ and $\exp(i\alpha_P) = 1$, which correspond respectively to a bosonic totally symmetric wavefunction, i.e. under P

$$\psi(x_1, \dots, x_N) \xrightarrow{P} (+1)\psi(x_1, \dots, x_N) ,$$

2. $\alpha_P = \pi$ and $\exp(i\alpha_P) = \text{sgn}(P)$, which correspond respectively to a fermionic totally antisymmetric wavefunction, i.e. under P

$$\psi(x_1, \dots, x_N) \xrightarrow{P} \text{sgn}(P)\psi(x_1, \dots, x_N) = \begin{cases} +\psi(x_1, \dots, x_N) & \text{sign}(P) = +1 \\ -\psi(x_1, \dots, x_N) & \text{sign}(P) = -1 \end{cases} .$$

By the spin-statistic theorem, bosons, which have symmetric wavefunctions, are associated to integer spin particles and fermions, which have antisymmetric wavefunctions, are associated to half-integer spin particles.

Proof. Using (14.2)

$$\psi(x_1, \dots, x_N) \xrightarrow{\sigma_i} \exp(i\alpha_i)\psi(x_1, \dots, x_N) \xrightarrow{\sigma_i\sigma_j} \exp(i\alpha_i) \exp(i\alpha_j)\psi(x_1, \dots, x_N) ,$$

$$\psi(x_1, \dots, x_N) \xrightarrow{\sigma_j} \exp(i\alpha_j)\psi(x_1, \dots, x_N) \xrightarrow{\sigma_j\sigma_i} \exp(i\alpha_j) \exp(i\alpha_i)\psi(x_1, \dots, x_N) ,$$

hence

$$\exp(i\alpha_i) \exp(i\alpha_j) = \exp(i\alpha_j) \exp(i\alpha_i) ,$$

which means that they commute

$$\alpha_i + \alpha_j = \alpha_j + \alpha_i . \quad (14.6)$$

Using (14.3)

$$\begin{aligned} \psi(x_1, \dots, x_N) &\xrightarrow{\sigma_i} \exp(i\alpha_i)\psi(x_1, \dots, x_N) \\ &\xrightarrow{\sigma_i\sigma_{i+1}} \exp(i\alpha_i) \exp(i\alpha_{i+1})\psi(x_1, \dots, x_N) \\ &\xrightarrow{\sigma_i\sigma_{i+1}\sigma_i} \exp(i\alpha_i) \exp(i\alpha_{i+1}) \exp(i\alpha_i)\psi(x_1, \dots, x_N) , \end{aligned}$$

$$\begin{aligned} \psi(x_1, \dots, x_N) &\xrightarrow{\sigma_{i+1}} \exp(i\alpha_{i+1})\psi(x_1, \dots, x_N) \\ &\xrightarrow{\sigma_{i+1}\sigma_i} \exp(i\alpha_{i+1}) \exp(i\alpha_i)\psi(x_1, \dots, x_N) \\ &\xrightarrow{\sigma_{i+1}\sigma_i\sigma_{i+1}} \exp(i\alpha_{i+1}) \exp(i\alpha_i) \exp(i\alpha_{i+1})\psi(x_1, \dots, x_N) , \end{aligned}$$

hence

$$\exp(i\alpha_i) \exp(i\alpha_{i+1}) \exp(i\alpha_i) = \exp(i\alpha_{i+1}) \exp(i\alpha_i) \exp(i\alpha_{i+1}) ,$$

which means that

$$\alpha_i + \alpha_{i+1} + \alpha_i = \alpha_{i+1} + \alpha_i + \alpha_{i+1} . \quad (14.7)$$

Putting together this two properties (14.6) and (14.7), we have

$$\alpha_i + \alpha_{i+1} + \alpha_i = \alpha_{i+1} + \alpha_i + \alpha_{i+1} ,$$

$$\cancel{\alpha_{i+1}} + \cancel{\alpha_i} + \alpha_i = \cancel{\alpha_{i+1}} + \cancel{\alpha_i} + \alpha_{i+1} ,$$

$$\alpha_i = \alpha_{i+1} .$$

Therefore, $\forall i = 1, \dots, N-1$ and $\alpha_i \in [0, 2\pi[$ we have $\alpha_i = \alpha_{i+1} = \alpha$.

Using (14.4)

$$\exp(i\alpha)^2 = \exp(2i\alpha) = \mathbb{I} = \exp(0) ,$$

which means that

$$\alpha = 0, \pi .$$

Finally, there are only two possibilities

$$\psi(x_1, \dots, x_N) \xrightarrow{\sigma_i} \underbrace{\exp(i0)}_{+1} \psi(x_1, \dots, x_N) = \psi(x_1, \dots, x_N)$$

and

$$\psi(x_1, \dots, x_N) \xrightarrow{\sigma_i} \underbrace{\exp(i\pi)}_{-1} \psi(x_1, \dots, x_N) = -\psi(x_1, \dots, x_N) .$$

q.e.d.

The Hilbert space of indistinguishable particle is smaller than the distinguishable one, because we have seen that the phase factor can only have two possible values. In the next chapters, we will see how we can describe such spaces, in terms of the symmetrised or antisymmetrised Hilbert space $\mathcal{H}_{S/A}$ in the language of first quantisation and in terms of the Fock space $\mathcal{F}_{B/F}$ in the language of second quantisation.

Chapter 15

First quantisation

15.1 Symmetric/antisymmetric Hilbert space

2 particles

Consider 2 particles. If they are distinguishable, the total Hilbert space is

$$\mathcal{H}_{tot} = \mathcal{H} \otimes \mathcal{H} ,$$

whereas if the particles are indistinguishable, we can decomposed the Hilbert space into

$$\mathcal{H}_{tot} = \mathcal{H}_S \oplus_{\perp} \mathcal{H}_A .$$

Proof. In fact, given two states $|a\rangle_1 \in \mathcal{H}_1$ and $|b\rangle_2 \in \mathcal{H}_2$, we have

$$\begin{aligned} |a\rangle_1 |b\rangle_2 &= \frac{2}{2} |a\rangle_1 |b\rangle_2 + \frac{1}{2} |b\rangle_1 |a\rangle_2 - \frac{1}{2} |b\rangle_1 |a\rangle_2 \\ &= \underbrace{\frac{|a\rangle_1 |b\rangle_2 + |b\rangle_1 |a\rangle_2}{2}}_{|\psi_S\rangle} + \underbrace{\frac{|a\rangle_1 |b\rangle_2 - |b\rangle_1 |a\rangle_2}{2}}_{|\psi_A\rangle} \\ &= |\psi_S\rangle + |\psi_A\rangle . \end{aligned}$$

Furthermore, the permutation group for 2 particles is $P_2 = \{\mathbb{I}, \sigma\}$. The symmetric part $|\psi_S\rangle \in \mathcal{H}_S$, since

$$\sigma |\psi_S\rangle = \sigma \frac{|a\rangle_1 |b\rangle_2 + |b\rangle_1 |a\rangle_2}{2} = \frac{|b\rangle_1 |a\rangle_2 + |a\rangle_1 |b\rangle_2}{2} = \frac{|a\rangle_1 |b\rangle_2 + |b\rangle_1 |a\rangle_2}{2} = |\psi_S\rangle ,$$

where we used the commutativity property. The antisymmetric part is $|\psi_A\rangle \in \mathcal{H}_A$, since

$$\sigma |\psi_A\rangle = \sigma \frac{|a\rangle_1 |b\rangle_2 - |b\rangle_1 |a\rangle_2}{2} = \frac{|b\rangle_1 |a\rangle_2 - |a\rangle_1 |b\rangle_2}{2} = -\frac{|a\rangle_1 |b\rangle_2 - |b\rangle_1 |a\rangle_2}{2} = -|\psi_A\rangle ,$$

where we used the commutativity property. Finally, the decomposition is orthogonal, since

$$\begin{aligned}\langle \psi_S | \psi_A \rangle &= \frac{\langle a|_1 \langle b|_2 + \langle b|_1 \langle a|_2}{2} \frac{|a\rangle_1 |b\rangle_2 - |b\rangle_1 |a\rangle_2}{2} \\ &= \frac{1}{4} (\underbrace{\langle a|a\rangle_1}_1 \underbrace{\langle b|b\rangle_2}_1 - \langle a|b\rangle_1 \langle b|a\rangle_2 + \langle b|a\rangle_1 \langle a|b\rangle_2 - \underbrace{\langle b|b\rangle_1}_1 \underbrace{\langle a|a\rangle_2}_1) \\ &= \frac{1}{4} (-\langle a|b\rangle_1 \langle b|a\rangle_2 + \langle b|a\rangle_1 \langle a|b\rangle_2)\end{aligned}$$

and

$$\begin{aligned}-\langle \psi_S | \psi_A \rangle &= -\frac{\langle a|_1 \langle b|_2 + \langle b|_1 \langle a|_2}{2} \frac{|a\rangle_1 |b\rangle_2 - |b\rangle_1 |a\rangle_2}{2} \\ &= -\frac{1}{4} (\underbrace{\langle a|a\rangle_1}_1 \underbrace{\langle b|b\rangle_2}_1 - \langle a|b\rangle_1 \langle b|a\rangle_2 + \langle b|a\rangle_1 \langle a|b\rangle_2 - \underbrace{\langle b|b\rangle_1}_1 \underbrace{\langle a|a\rangle_2}_1) \\ &= -\frac{1}{4} (-\langle a|b\rangle_1 \langle b|a\rangle_2 + \langle b|a\rangle_1 \langle a|b\rangle_2) \\ &= \frac{1}{4} (-\langle a|b\rangle_2 \langle b|a\rangle_1 + \langle b|a\rangle_2 \langle a|b\rangle_1),\end{aligned}$$

which means that $\langle \psi_S | \psi_A \rangle = -\langle \psi_S | \psi_A \rangle$. Therefore, the only solution is $\langle \psi_S | \psi_A \rangle = 0$. q.e.d.

Notice that Pauli's exclusion principle is encoded into the antisymmetric part, because if $a = b$ we have $|\psi_A\rangle = 0$.

The decomposition is equivalent to define two orthogonal projectors: the symmetriser

$$\hat{S}: \mathcal{H} \rightarrow \mathcal{H}_S$$

and the antisymmetriser

$$\hat{A}: \mathcal{H} \rightarrow \mathcal{H}_A,$$

such that they satisfy the properties

$$\hat{S}^\dagger = \hat{S}, \quad \hat{A}^\dagger = \hat{A}, \quad \hat{S}^2 = \hat{S}, \quad \hat{A}^2 = \hat{A}, \quad \hat{S}\hat{A} = \hat{A}\hat{S} = 0. \quad (15.1)$$

N particles

Generalising for N particles, if they are distinguishable, the total Hilbert space is

$$\mathcal{H}_{tot} = \mathcal{H} \otimes \dots \mathcal{H}$$

and a state is $|\psi\rangle = |a_1\rangle_1 \dots |a_N\rangle_N = |1, \dots, N\rangle$ where $|a_j\rangle \in \mathcal{H}$.

However, if the particles are indistinguishable, similarly to the 2 particles case, we can define the symmetriser

$$\hat{S}: |\psi\rangle \mapsto \frac{1}{N!} \sum_{P \in P_N} |P(1), \dots, P(N)\rangle$$

and the antisymmetriser

$$\hat{A}: |\psi\rangle \mapsto \frac{1}{N!} \sum_{P \in P_N} \text{sgn}(P) |P(1), \dots, P(N)\rangle ,$$

where $P(1, \dots, N) \mapsto (P(1), \dots, P(N))$. They satisfy the orthogonal projector properties (15.1). Notice that for $N > 2$ particles, the total Hilbert space is $\mathcal{H}_{tot} = \mathcal{H}_S \otimes \mathcal{H}_A \otimes \mathcal{H}'$, where bosons work only in \mathcal{H}_S , fermions work only in \mathcal{H}_A and \mathcal{H}' is not physical.

Example 15.1. For $N = 3$, we can have $\psi_S \in \text{mathcal{H}}_S$

$$\psi_S = \hat{S}\psi = \psi(1, 2, 3) + \psi(1, 3, 2) + \psi(2, 1, 3) + \psi(2, 3, 1) + \psi(3, 1, 2) + \psi(3, 2, 1)$$

and $\psi_A \in \text{mathcal{H}}_A$

$$\psi_A = \hat{A}\psi = \psi(1, 2, 3) - \psi(1, 3, 2) - \psi(2, 1, 3) + \psi(2, 3, 1) + \psi(3, 1, 2) - \psi(3, 2, 1) .$$

However, we can also have $\psi' \in \mathcal{H}'$ such that

$$\psi' = \psi(1, 2, 3) + \psi(1, 3, 2) - \psi(2, 1, 3) - \psi(2, 3, 1) + \psi(3, 1, 2) + \psi(3, 2, 1) .$$

For N distinguishable particles, consider an orthonormal basis for the total Hilbert space

$$\{u_{\alpha_1}(x_1) \dots u_{\alpha_N}(x_N)\}_{\alpha_1, \dots, \alpha_N=0}^{\infty} ,$$

where $\{u_{\alpha_K}(x_K)\}_{\alpha_K=1}^{\infty}$ is an orthonormal basis for a single Hilbert space \mathcal{H}_1 . Notice that they are labelled by the ordered set $(\alpha_1, \dots, \alpha_N)$ and we are specifying which particle is in which states.

In order to construct an orthonormal basis for \mathcal{H}_A and \mathcal{H}_S for N indistinguishable particles, we project the distinguishable orthonormal basis respectively with the antisymmetriser and the symmetriser

$$|n_1, \dots, n_j, \dots\rangle = C \begin{cases} \hat{S} \\ \hat{A} \end{cases} u_{\alpha_1}(x_1) \dots u_{\alpha_N}(x_N) .$$

By the properties of the projectors, they are orthonormal but they are not normalised. Therefore, we need to choose a normalisation constant

$$C = \begin{cases} \sqrt{\frac{N!}{n_1! \dots n_k! \dots}} & \mathcal{H}_S \\ \sqrt{N!} & \mathcal{H}_A \end{cases} .$$

On the contrary for the distinguishable case, now we lose information, because we know only how many particle are in each state and not anymore which is in which state. We label the states with n_1, \dots, n_k, \dots with $j = 1, \dots, \infty$, which are the occupation number. For bosons, we have $n_k = 0, 1, \dots, \infty$, whereas for fermions, we have $n_k = 0, 1$. For both cases, there is the constrain $N = \sum_k n_k$, which is an infinite sum but mostly are zero occupied. Moreover, given the set α_k , we uniquely determine the occupation number n_k , but given the occupation number n_k , we use the symmetric or antisymmetric property to uniquely determine the state, because it is in 1 – 1 correspondence to the set n_k .

Nonetheless, there is another way to describe fermionic or bosonic quantum space in a intrinsic way, called the second quantisation because we make a further quantisation, promoting fields to operators.

Chapter 16

Second quantisation

16.1 Bosonic case

We define bosonic creation and annihilation operators such that they satisfies the properties

$$[\hat{a}, \hat{a}^\dagger]_- = \hat{a}\hat{a}^\dagger - \hat{a}^\dagger\hat{a} = \mathbb{I} . \quad (16.1)$$

Furthermore, the number operator $\hat{N} = \hat{a}^\dagger\hat{a}$ such that

$$[\hat{N}, \hat{a}] = -\hat{a} , \quad [\hat{N}, \hat{a}^\dagger] = \hat{a}^\dagger .$$

By analogy with the harmonic oscillator, the ground state is the vacuum

$$\hat{a}|0\rangle = 0 ,$$

and a generic state is defined by the ladder operators

$$|\psi\rangle = \frac{1}{\sqrt{n!}}(\hat{a}^\dagger)^N|0\rangle .$$

16.2 Fermionic case

We define fermionic creation and annihilation operators such that they satisfies the properties

$$[\hat{a}, \hat{a}^\dagger]_+ = \hat{a}\hat{a}^\dagger + \hat{a}^\dagger\hat{a} = \mathbb{I} . \quad (16.2)$$

Furthermore, the number operator $\hat{N} = \hat{a}^\dagger\hat{a}$ such that

$$[\hat{N}, \hat{a}] = -\hat{a} , \quad [\hat{N}, \hat{a}^\dagger] = \hat{a}^\dagger .$$

The properties can be obtained from the Pauli matrices

$$\sigma_\pm = \sigma_1 \pm i\sigma_2 ,$$

such that

$$(\sigma_+)^{\dagger} = \sigma_- , \quad (\sigma_-)^{\dagger} = \sigma_+ , \quad (\sigma_+)^2 = (\sigma_-)^2 = 0 , \quad [\sigma_-, \sigma_+]_+ = \mathbb{I} .$$

By analogy with the harmonic oscillator, the ground state is the vacuum

$$\hat{a}|0\rangle = 0 ,$$

and a generic state is defined by the ladder operators

$$|\psi\rangle = \frac{1}{\sqrt{n!}} (\hat{a}^{\dagger})^N |0\rangle .$$

However, the anticommutator relation ensures the validity of the Pauli's exclusion principle. In fact, we have

$$a^2 = (\hat{a}^{\dagger})^2 = 0 .$$

16.3 Fock space

Consider a single particle Hilbert space \mathcal{H} with an orthonormal basis $\{|e_n\rangle\}_{n=1}^{\infty}$. To each $|e_n\rangle$, we associate an annihilation and a creation operators

$$|e_n\rangle \mapsto \{\hat{a}_n, \hat{a}_n^{\dagger}\}_{n=1}^{\infty} ,$$

such that they satisfy

$$[\hat{a}_n, \hat{a}_m]_{\pm} = [\hat{a}_n^{\dagger}, \hat{a}_m^{\dagger}]_{\pm} = 0 , \quad [\hat{a}_n, \hat{a}_m^{\dagger}]_{\pm} = \delta_{nm} ,$$

where the minus sign corresponds to the commutator (bosons) (16.1) and the plus sign to the anticommutator (fermions) (16.2).

The normalised vacuum state, which describes a no particle state, is defined by the annihilation of every annihilation operator

$$\hat{a}_n|0\rangle = 0 \quad \forall n .$$

It generates a subspace of dimension 1

$$\mathcal{H}_{S/A}^{(0)} = \{\lambda|0\rangle : \lambda \in \mathbb{C}\} . \quad (16.3)$$

Now, we define the one-particle state. For each $|e_n\rangle$, we associate a number operator $\hat{n}_k = \hat{a}_k^{\dagger} \hat{a}_k$ such that

$$\hat{n}_k \hat{a}_k^{\dagger} |0\rangle = 1 \hat{a}_k^{\dagger} |0\rangle , \quad \hat{n}_{k'} \hat{a}_k^{\dagger} |0\rangle = 0 \quad k' \neq k .$$

For a n particle state, we have

$$\hat{a}_k^\dagger |0\rangle = |n_1 = 0, \dots, n_k = 1, \dots, n_N = 0\rangle = |e_k\rangle .$$

However, for

$$\hat{a}_{k_1}^\dagger \hat{a}_{k_2}^\dagger |0\rangle = |e_{k_1}\rangle |e_{k_2}\rangle$$

we have for fermions, if $k_1 = k_2 = k$

$$(\hat{a}_k^\dagger)^2 |0\rangle = 0 ,$$

whereas for bosons

$$(\hat{a}_k^\dagger)^2 |0\rangle \neq 0 .$$

Furthermore, if $k_1 \neq k_2$, we have for fermions

$$\hat{a}_{k_1}^\dagger \hat{a}_{k_2}^\dagger |0\rangle = -\hat{a}_{k_2}^\dagger \hat{a}_{k_1}^\dagger |0\rangle ,$$

whereas for bosons

$$\hat{a}_{k_1}^\dagger \hat{a}_{k_2}^\dagger |0\rangle = \hat{a}_{k_2}^\dagger \hat{a}_{k_1}^\dagger |0\rangle .$$

There is a 1 – 1 correspondence between the orthonormal basis $\{|e_n\rangle\}_{n=1}^\infty$ of \mathcal{H} and the orthonormal basis $\{\hat{a}_k |0\rangle\}_{k=1}^\infty$ of $\mathcal{H}_{S/A}$. Hence for N particles, we have

$$\mathcal{H}_{S/A}^{(N)} = \{|n_1, \dots, n_k, \dots\rangle = \frac{1}{\sqrt{\prod_j n_j}} (\hat{a}_1^\dagger)^{n_1} \dots (\hat{a}_k^\dagger)^{n_k} \dots |0\rangle\} .$$

If N is not fixed, like the passage from canonical to grancanonical ensemble, the total Fock space is

$$\mathcal{F} = \bigoplus_{N=0}^{\infty} \mathcal{H}_{S/A}^{(N)} .$$

It satisfies the following properties

1. orthonormality, i.e.

$$\langle n'_1, \dots, n'_k, \dots | n_1, \dots, n_k, \dots \rangle = \delta_{n'_1, n_1} \dots \delta_{n'_k, n_k} \dots ,$$

2. annihilation $\hat{a}_k: \mathcal{H}_{S/A}^{(N)} \rightarrow \mathcal{H}_{S/A}^{(N-1)}$, i.e.

$$\hat{a}_k |n_1, \dots, n_k, \dots\rangle = \eta_k \sqrt{n_k} |n_1, \dots, (n_k - 1), \dots\rangle ,$$

where for bosons $\eta_k = 1$ and for fermions $\eta_k = (-1)^{\sum_{j < k} n_j}$,

3. creation $\hat{a}_k^\dagger: \mathcal{H}_{S/A}^{(N)} \rightarrow \mathcal{H}_{S/A}^{(N+1)}$, i.e. for bosons

$$\hat{a}_k^\dagger |n_1, \dots, n_k, \dots\rangle = \sqrt{n_k + 1} |n_1, \dots, (n_k + 1), \dots\rangle ,$$

and for fermions

$$\hat{a}_k^\dagger |n_1, \dots, n_k, \dots\rangle = \eta_k \sqrt{1 - n_k} |n_1, \dots, (n_k + 1), \dots\rangle ,$$

4. number operator $\hat{n}_k = \hat{a}_k^\dagger \hat{a}_k$ such that

$$\hat{n}_k |n_1, \dots, n_k, \dots\rangle = n_k |n_1, \dots, n_k, \dots\rangle$$

and the total number operator $\hat{N} = \sum_k \hat{n}_k = \sum_k \hat{a}_k^\dagger \hat{a}_k$ such that

$$\hat{N} |n_1, \dots, n_k, \dots\rangle = \left(\sum_k n_k \right) |n_1, \dots, n_k, \dots\rangle .$$

16.4 Field operators

In the first quantisation, we quantise observables to operators, while, in the second quantisation, we quantise fields to operators. Now, a generic particle state is represented by $|f\rangle = \sum_k f_k |e_k\rangle \in \mathcal{H}$, which is equivalent to $\sum_k f_k \hat{a}_k^\dagger |0\rangle$. Hence, we define the field operators

$$\hat{\psi}^\dagger(f) = \sum_k f_k \hat{a}_k^\dagger , \quad \hat{\psi}(f) = \sum_k f_k^* \hat{a}_k ,$$

in order to get a state $\hat{\psi}(f)|0\rangle$. The related commutator relations become

$$[\hat{\psi}(f), \hat{\psi}^\dagger(g)]_\pm = \langle f|g\rangle \mathbb{I} .$$

Proof. In fact,

$$\begin{aligned} [\hat{\psi}(f), \hat{\psi}^\dagger(g)]_\pm &= \left[\sum_k f_k^* \hat{a}_k, \sum_m g_m \hat{a}_m^\dagger \right]_\pm \\ &= \sum_k \sum_m f_k^* g_m \underbrace{[\hat{a}_k, \hat{a}_m^\dagger]}_{\delta_{km} \mathbb{I}} \\ &= \sum_k \sum_m f_k^* g_m \underbrace{\delta_{km}}_{k=m} \mathbb{I} \\ &= \sum_k f_k^* g_k \mathbb{I} \\ &= \langle f|g\rangle \mathbb{I} . \end{aligned}$$

where we have used

$$|f\rangle = \sum_k f_k |e_k\rangle , \quad |g\rangle = \sum_m g_m |e_m\rangle , \quad \langle f|g\rangle = \sum_k \sum_m f_k^* g_m \underbrace{\langle e_k|e_m\rangle}_{\delta_{km}} = \sum_k f_k^* g_k .$$

q.e.d.

Consider a single particle state in $\mathcal{H} = L^2(\mathbb{R}^d) \ni \psi(x)$ with an orthonormal basis $u_k(x)$ such that to each ket there are ladder operators \hat{a}_k and \hat{a}_k^\dagger . Hence $L^2(\mathbb{R}^d) \ni f(x) = \sum_k f_k u_k(x)$ and we define field operators

$$\hat{\psi}(x) = \sum_k u_k^*(x) \hat{a}_k, \quad \hat{\psi}^\dagger(x) = \sum_k u_k(x) \hat{a}_k^\dagger,$$

which is a linear superposition of annihilation and creation operators. Actually, it is called an operator-valued function because its output is an operator. In fact

$$\int_{\mathbb{R}^d} d^d x \hat{\psi}^\dagger(x) \sum_k u_k^*(x) \hat{a}_k^\dagger = \sum_k \hat{a}_k^\dagger \int_{\mathbb{R}^d} d^d x u_k^*(x) f(x) = \sum_k \hat{a}_k^\dagger f_k,$$

where we have exchanged sum and integral because they are convergent.

The commutation relations are

$$[\psi(x), \psi^\dagger(y)] = \mathbb{I} \delta(x - y).$$

Proof. In fact,

$$\begin{aligned} [\hat{\psi}(f), \hat{\psi}^\dagger(g)]_\pm &= \left[\int d^d x f^*(x) \hat{\psi}(x), \int d^d y g(y) \hat{\psi}^\dagger(y) \right]_\pm \\ &= \int d^d x \int d^d y f^*(x) g(y) [\psi(x), \psi^\dagger(y)], \end{aligned}$$

which must be equal to

$$\langle f | g \rangle = \int d^d x f^*(x) g(x).$$

Hence

$$[\psi(x), \psi^\dagger(y)] = \mathbb{I} \delta(x - y).$$

q.e.d.

For instance, a plane wave $u(x) = \exp(i\mathbf{k} \cdot \mathbf{x})$ and $\hat{\psi}(x) = \sum_k \hat{a}_k^\dagger \exp(i\mathbf{k} \cdot \mathbf{x})$.

Notice that field operators are basis independent

16.5 Operators

Consider a Fock space $\mathcal{F} = \bigoplus_{N=0}^{\infty} \mathcal{H}_{B/F}^{(N)}$ with orthonormal basis $|n_1, \dots, n_k, \dots\rangle = \frac{1}{\sqrt{\prod_j n_j!}} (\hat{a}_1^\dagger)^{n_1} \dots (\hat{a}_k^\dagger)^{n_k} \dots |0\rangle$, which is in 1 – 1 correspondence to the orthonormal basis $\psi_{n_1 \dots n_k \dots}(x_1, \dots, x_k, \dots) = c_N \left[\begin{smallmatrix} \hat{S} \\ \hat{A} \end{smallmatrix} \right] u_{\alpha_1}(x_1) \dots u_{\alpha_k}(x_k) \dots$, where *hatS* is the symmetriser and \hat{A} is the antisymmetriser.

We define a one-body operator, associated to a system in which all the particles are the same, as

$$\hat{O}^{(1)} = \sum_{j=1}^N \hat{O}(\hat{p}_j, \hat{x}_j) .$$

Since it is self-adjoint, it exists an orthonormal basis of eigenvalues $\{u_\alpha(x)\}$, such that

$$\hat{O}(\hat{p}, \hat{x})u_\alpha(x) = \epsilon_\alpha u_\alpha(x) .$$

Since

$$\begin{aligned} \hat{O}^{(1)}\psi_{n_1\dots n_k\dots}(x_1, \dots x_k, \dots) &= \left(\sum_{j=1}^{\infty} \hat{O}(\hat{p}_j, \hat{x}_j) \right) \psi_{n_1\dots n_k\dots}(x_1, \dots x_k, \dots) \\ &= \left(\sum_{j=1}^{\infty} \hat{O}(\hat{p}_j, \hat{x}_j) \right) c_N \left[\begin{smallmatrix} \hat{S} \\ \hat{A} \end{smallmatrix} \right] u_{\alpha_1}(x_1) \dots u_{\alpha_k}(x_k) \dots \\ &= c_N \left[\begin{smallmatrix} \hat{S} \\ \hat{A} \end{smallmatrix} \right] \left(\sum_{j=1}^{\infty} \hat{O}(\hat{p}_j, \hat{x}_j) u_{\alpha_1}(x_1) \dots u_{\alpha_k}(x_k) \dots \right) \\ &= c_N \left[\begin{smallmatrix} \hat{S} \\ \hat{A} \end{smallmatrix} \right] \left(\sum_{j=1}^{\infty} u_{\alpha_1}(x_1) \dots \underbrace{\hat{O}(\hat{p}_j, \hat{x}_j) u_{\alpha_j}(x_j)}_{\epsilon_{\alpha_j} u_{\alpha_j}(x_j)} \dots \right) \\ &= \left(\sum_{j=1}^{\infty} \epsilon_j n_j \right) \psi_{n_1\dots n_k\dots}(x_1, \dots x_k, \dots) . \end{aligned}$$

For the Fock space, we have

$$\hat{O}_F^{(1)} = \sum_{j=1}^{\infty} \epsilon_j \hat{n}_j = \sum_{j=1}^{\infty} \epsilon_j \hat{a}_j^\dagger \hat{a}_j ,$$

where

$$\epsilon_j = \langle u_j(x) | \hat{O}(\hat{p}_j, \hat{x}_j) | u_j(x) \rangle .$$

Hence

$$\hat{O}_F^{(1)} = \sum_{j=1}^{\infty} \langle u_j(x) | \hat{O}(\hat{p}_j, \hat{x}_j) | u_j(x) \rangle \hat{a}_j^\dagger \hat{a}_j .$$

Since it is dependent of the basis, because we choose the eigenbasis, we choose a different arbitrary basis

$$\psi^\dagger(x) = \sum_k u_k(x) \hat{a}_k^\dagger = \sum_m v_m(x) b_m^\dagger ,$$

and we define the one-body operator

$$\hat{O}_F^{(1)} = \int d^d x \, \hat{\varphi}^\dagger(x) \hat{O}(\hat{p}, \hat{x}) \hat{\varphi}(x) , \quad (16.4)$$

which this time is basis independent.

Proof. In fact

$$\begin{aligned}
\int d^d x \, \hat{\varphi}^\dagger(x) \hat{O}(\hat{p}, \hat{x}) \hat{\varphi}(x) &= \int d^d x \left(\sum_k u_k(x) \hat{a}_k^\dagger(x) \right) \hat{O}(\hat{p}, \hat{x}) \left(\sum_m u_m^*(x) \hat{a}_m(x) \right) \\
&= \sum_k \sum_m \hat{a}_k^\dagger \hat{a}_m \int d^d x \, u_k(x) \underbrace{\hat{O}(\hat{p}, \hat{x}) u_m^*(x)}_{\epsilon_m u_m^*(x)} \\
&= \sum_k \sum_m \hat{a}_k^\dagger \hat{a}_m \epsilon_m \underbrace{\int d^d x \, u_k(x) u_m^*(x)}_{\delta_{km}} \\
&= \sum_k \sum_m \hat{a}_k^\dagger \hat{a}_m \epsilon_m \underbrace{\delta_{km}}_{k=m} \\
&= \sum_k \hat{a}_k^\dagger \hat{a}_k \epsilon_k = \hat{O}_F^{(1)} .
\end{aligned}$$

q.e.d.

It can be written as

$$\hat{O}_F^{(1)} = \sum_k \sum_m t_{km} \hat{b}_k^\dagger \hat{h}_m ,$$

where the transition amplitude is

$$t_{km} = \langle v_k | \hat{O}(\hat{p}, \hat{x}) | v_m \rangle .$$

To summarise, the onebody operator is

$$\hat{O}_F^{(1)} = \begin{cases} \sum_{mm'} t_{mm'} \hat{b}_m^\dagger \hat{b}_m & \text{arbitrary basis} \\ \sum_k \epsilon_k \hat{a}_k^\dagger \hat{a}_k & \text{eigenbasis} \end{cases} .$$

16.6 Examples

The density operator of a single particle j is

$$\hat{\rho}_j = \delta(x - x_j)$$

and the corresponding field operator is

$$\hat{\varphi}(x_j) = \int d^d x \, \psi(x) \delta(x - x_j) .$$

The onebody operator is

$$\hat{\rho}^{(1)} = \sum_{j=1}^N \delta(x - x_j) ,$$

which in the basis independent definition (16.4) on the Fock space

$$\hat{\rho}_F = \int d^d y \hat{\psi}^\dagger(y) \delta(x-y) \hat{\psi}(y) = \hat{\psi}^\dagger \hat{\psi} = \sum_{kk'} u_k^*(x) u_{k'}(x) \hat{a}_k^\dagger \hat{a}_{k'} .$$

The number of particle operator is

$$\begin{aligned} \hat{N} &= \int d^d x \hat{\rho}_F^{(1)}(x) \\ &= \int d^d x \sum_{kk'} u_k^*(x) u_{k'}(x) \hat{a}_k^\dagger \hat{a}_{k'} \\ &= \sum_{kk'} \hat{a}_k^\dagger \hat{a}_{k'} \underbrace{\int d^d x u_k^*(x) u_{k'}(x)}_{\delta_{kk'}} \\ &= \sum_{kk'} \hat{a}_k^\dagger \hat{a}_{k'} \underbrace{\delta_{kk'}}_{k=k'} \\ &= \sum_k \hat{a}_k^\dagger \hat{a}_k = \sum_k \hat{n}_k , \end{aligned}$$

which is consistent with the definition of ρ since it can be seen as a density of particle whose intergral is indeed the number of particles.

Free non-relativistic 3-dimensional particles

Consider the hamiltonian of a single particle described by the wave function $\psi(x) \in L^2(\mathbb{R}^d)$

$$\hat{H}_1 = \frac{\hbar^2 \hat{p}^2}{2m} = -\frac{\hbar^2}{2m} \nabla_x^2 .$$

The onebody operator for N particles is

$$\hat{H} = \sum_{j=1}^N \frac{\hbar^2 \hat{p}_j^2}{2m} = -\sum_{j=1}^N \frac{\hbar^2}{2m} \nabla_{x_j}^2 .$$

On the Fock space, it becomes

$$H = \sum_k \epsilon_k \hat{a}_k^\dagger \hat{a}_k ,$$

where

$$\hat{H}_1 u_k(x) = -\frac{\hbar^2}{2m} u_k(x) = \epsilon_k u_k(x) .$$

However, wave plane solutions do not belong in $u_k(x) \sim \exp(i\mathbf{k} \cdot \mathbf{x}) \notin L^2(\mathbb{R}^d)$, because they are not normalisable. The trick is to go into a finite volume V and

consider the space $L^2(V)$. The simple example is the particle in a cube of length L describer by the coordinates $(x, y, z) \in [0, L]$. The Schrodinger's equation becomes

$$-\frac{\hbar^2}{2m} \nabla_x^2 u_k(x, y, z) = \epsilon_k u_k(x, y, z) .$$

Now, we do not choose the Dirichlet or the Neumann boundary condition, but we choose the periodic boundary conditions

$$\begin{cases} u(x=0, y, z) = u(x=L, y, z) \\ u(x, y=0, z) = u(x, y=L, z) \\ u(x, y, z=0) = u(x, y, z=L) \end{cases} ,$$

which transforms the cube into a 3-torus.

The ansatz solution is

$$u_\alpha(\mathbf{x}) = c \exp(i\mathbf{k} \cdot \mathbf{x}) ,$$

where c is a normalisation constant and

$$\nabla^2 u_\alpha(\mathbf{x}) = -(k_x^2 + k_y^2 + k_z^2) u_\alpha(\mathbf{x}) = \epsilon_{\mathbf{k}} = -k^2 .$$

Imposing the periodic boundary conditions, we obtain

$$u_{\mathbf{k}}(0, y, z) = c \exp(i(\cancel{k_y y} + \cancel{k_z z})) = u_{\mathbf{k}}(L, y, z) = c \exp(i(k_x L + \cancel{k_y y} + \cancel{k_z z})) ,$$

hence

$$\exp(ik_x L) = 1$$

and

$$k_x = \frac{2\pi}{L} n_x ,$$

where $n \in \mathbb{Z}$ is an integer number. Simiarly for y and z , we have

$$\mathbf{k} = (k_x, k_y, k_z) = \frac{2\pi}{L} \mathbf{n} = \frac{2\pi}{L} (n_x, n_y, n_z)$$

where $n_x, n_y, n_z \in \mathbb{Z}$. Finally, the energy eigenvalues are

$$\epsilon_{n_x, n_y, n_z} = -\frac{4\pi^2}{L^2} (n_x^2 + n_y^2 + n_z^2)$$

and the eigenstates are

$$u_{n_x, n_y, n_z} = c \exp(i \frac{2\pi}{L} (n_x x + n_y y + n_z z)) \in L^2(V) .$$

The normalisation constant is

$$C = \frac{1}{\sqrt{V}} .$$

In fact

$$1 = \|u_{n_x, n_y, n_z}\|^2 = \int_V dx dy dz |c|^2 |\exp(i\mathbf{k} \cdot \mathbf{x})|^2 = |c|^2 V .$$

Hence, the onebody operator is

$$\hat{O}^{(1)} = \sum_{j=1}^N \frac{\hat{p}_j^2}{2m} = - \sum_{j=1}^N \frac{\hbar^2}{2m} \nabla_{\mathbf{x}_j}^2 ,$$

and choosing the orthonormal basis of wavefunctions, we have in the Fock space

$$\hat{O}_F = \sum_{\mathbf{k}} \epsilon_{\mathbf{k}} \hat{a}_{\mathbf{k}}^\dagger \hat{a}_{\mathbf{k}} ,$$

where $\mathbf{k} = \frac{2\pi}{L}(n_x, n_y, n_z)$ and $\epsilon_{\mathbf{k}} = \frac{\hbar^2}{2m} k^2$.

Interaction potential

Now, consider a different operator than the onebody one: the two particles operator, used to describe interaction potential

$$\hat{O}^{(2)} = \sum_{i < j} V(x_i, x_j) = \frac{1}{2} \int dx \int dy V(x, y) \hat{\psi}^\dagger(x) \hat{\psi}^\dagger(y) \psi(y) \psi(x) ,$$

or

$$\hat{O}^{(2)} = \frac{1}{2} \sum_{ijkl} V_{ijkl} \hat{a}_i^\dagger \hat{a}_j^\dagger \hat{a}_k \hat{a}_l ,$$

where the first expression is basis-independent, while in the last one it is in the eigenbasis.

Proof. Maybe in the future.

q.e.d.

Part V

Quantum statistical mechanics

Chapter 17

Microcanonical ensemble

The microcanonical ensemble is characterised by constant volume, energy and number of particle. Since N is fixed, we can work in the Hilbert space \mathcal{H}_{tot} . Given a time-independent hamiltonian operator \hat{H} , we find the energy eigenbasis $|\psi_j\rangle \in \mathcal{H}_{tot}$

$$\hat{H}|\psi_j\rangle = E_j|\psi_j\rangle .$$

However, there could be some degeneracy we want to consider, i.e. $E_{j,\alpha} = E_{j,\beta}$ for $|\psi_{j,\alpha}\rangle \neq |\psi_{j,\beta}\rangle$. Therefore, we have

$$\hat{H}|\psi_{j,\alpha}\rangle = E_j|\psi_{j,\alpha}\rangle , \quad (17.1)$$

where $\alpha = 1, \dots, n_j$.

The density operator for mixed states is (13.9)

$$\rho_{mc} = \sum_{\alpha=1}^{n_j} p_{\alpha} |\psi_{j,\beta}\rangle \langle \psi_{j,\beta}| ,$$

where p_{α} is the probability for the eigenstate $|\psi_{j,\beta}\rangle$. Since $E = E_j$ is fixed, all the eigenstates have the same probability to occur. Therefore $p_{\alpha} = \frac{1}{n_j}$ and

$$\rho_{mc} = \frac{1}{n_j} \sum_{\alpha=1}^{n_j} |\psi_{j,\alpha}\rangle \langle \psi_{j,\alpha}| = \frac{1}{n_j} \hat{P}_j ,$$

where

$$\hat{P}_j = \sum_{\alpha=1}^{n_j} |\psi_{j,\alpha}\rangle \langle \psi_{j,\alpha}|$$

is the projector onto the energy eigenspace. Notice that we can expand the hamiltonian using (13.6)

$$\hat{H} = \sum_j E_j \hat{P}_j \quad (17.2)$$

or the total number operator

$$\hat{N} = \sum_j n_j \hat{P}_j . \quad (17.3)$$

The average of an observable \hat{A} in the microcanonical ensemble is

$$\langle A \rangle_{mc} = \frac{1}{n_j} \sum_{\alpha=1}^{n_j} \langle \psi_{n,\alpha} | \hat{A} | \psi_{n,\alpha} \rangle .$$

Proof. In fact, choosing an orthonormal basis $|e_j\rangle$, the trace is

$$\text{tr}_{\mathcal{H}_{tot}} \hat{A} = \sum_j \langle e_j | \hat{A} | e_j \rangle .$$

Therefore, using (13.10)

$$\begin{aligned} \langle A \rangle_{mc} &= \text{tr}_{\mathcal{H}_{tot}} (\hat{A} \rho_{mc}) \\ &= \text{tr}_{\mathcal{H}_{tot}} \left(\hat{A} \frac{1}{n_j} \sum_{\alpha=1}^{n_j} |\psi_{j,\alpha}\rangle \langle \psi_{j,\alpha}| \right) \\ &= \frac{1}{n_j} \sum_{\alpha=1}^{n_j} \text{tr}_{\mathcal{H}_{tot}} \left(\hat{A} |\psi_{j,\alpha}\rangle \langle \psi_{j,\alpha}| \right) \\ &= \frac{1}{n_j} \sum_{\alpha=1}^{n_j} \langle \psi_{j,\alpha} | \hat{A} | \psi_{j,\alpha} \rangle . \end{aligned}$$

q.e.d.

The entropy in the microcanonical ensemble is

$$S_{mc} = k_B \log n_j ,$$

where n_j is the number of states with $E = E_j$. Notice that it is similar to the classical case (5.2).

Proof. In fact, using (5.5)

$$S_{mc} = -k_B \langle \log \rho_{mc} \rangle_{mc} = -k_B \text{tr}_{\mathcal{H}_{tot}} (\rho_{mc} \log \rho_{mc}) .$$

In matrix notation, the density operator is

$$\rho_{mc} = \begin{bmatrix} \begin{bmatrix} \frac{1}{n_1} & 0 & \dots & 0 \\ 0 & \frac{1}{n_1} & \dots & 0 \\ \dots & \dots & \dots & \dots \\ 0 & 0 & \dots & \frac{1}{n_1} \end{bmatrix} & 0 & \dots & 0 & \dots & \dots \\ 0 & \begin{bmatrix} \frac{1}{n_2} & 0 & \dots & 0 \\ 0 & \frac{1}{n_2} & \dots & 0 \\ \dots & \dots & \dots & \dots \\ 0 & 0 & \dots & \frac{1}{n_2} \end{bmatrix} & \dots & 0 & \dots & \dots \\ \dots & \dots & \dots & \dots & \dots & \dots & \dots \\ 0 & 0 & \dots & \begin{bmatrix} \frac{1}{n_j} & 0 & \dots & 0 \\ 0 & \frac{1}{n_j} & \dots & 0 \\ \dots & \dots & \dots & \dots \\ 0 & 0 & \dots & \frac{1}{n_j} \end{bmatrix} & \dots & \dots & \dots \\ \dots & \dots & \dots & \dots & \dots & \dots & \dots \\ \dots & \dots & \dots & \dots & \dots & \dots & \dots \end{bmatrix} .$$

$$= \sum_j \begin{bmatrix} 0 & 0 & \dots & 0 & \dots & \dots \\ 0 & 0 & \dots & 0 & \dots & \dots \\ \dots & \dots & \dots & \dots & \dots & \dots \\ 0 & 0 & \dots & \begin{bmatrix} \frac{1}{n_j} & 0 & \dots & 0 \\ 0 & \frac{1}{n_j} & \dots & 0 \\ \dots & \dots & \dots & \dots \\ 0 & 0 & \dots & \frac{1}{n_j} \end{bmatrix} & \dots & \dots \\ \dots & \dots & \dots & \dots & \dots & \dots \end{bmatrix}$$

In order to compute the logarithm of 0, we use a trick: we define a small parameter ϵ and we make it go to zero. In this way, the limit becomes $\epsilon \log \epsilon \xrightarrow{\epsilon \rightarrow 0} 0$. Finally, we compute the trace

$$\text{tr}_{\mathcal{H}_{tot}}(\rho_{mc} \log \rho_{mc}) = \text{tr} \begin{bmatrix} 0 & 0 & \dots & 0 & \dots & \dots \\ 0 & 0 & \dots & 0 & \dots & \dots \\ \dots & \dots & \dots & \dots & \dots & \dots \\ 0 & 0 & \dots & \begin{bmatrix} \frac{1}{n_j} \log \frac{1}{n_j} & 0 & \dots & 0 \\ 0 & \frac{1}{n_j} \log \frac{1}{n_j} & \dots & 0 \\ \dots & \dots & \dots & \dots \\ 0 & 0 & \dots & \frac{1}{n_j} \log \frac{1}{n_j} \end{bmatrix} & \dots & \dots \\ \dots & \dots & \dots & \dots & \dots & \dots \end{bmatrix}$$

$$= \sum_j \frac{1}{n_j} \log \frac{1}{n_j} = n_j \frac{1}{n_j} \log \frac{1}{n_j} = -\log n_j .$$

Hence,

$$S_{mc} = -k_B \operatorname{tr}_{\mathcal{H}_{tot}}(\rho_{mc} \log \rho_{mc}) = k_B \log n_j .$$

q.e.d.

Notice that entropy is always a positive function, since there is at least one state occupied $n_j \geq 1$, which implies $S \geq 0$.

Chapter 18

Canonical ensemble

The canonical ensemble is characterised by constant volume, temperature and number of particle. Energy, which can be exchange in an external reservoir, can be in one of the eigenstates (17.1) with probability

$$p_j \propto \exp(-\beta E_j) . \quad (18.1)$$

Consider a family of projectors $\{\hat{P}_j\}$, the density matrix of a mixed states is

$$\rho_c = \frac{1}{Z_N} \sum_j \exp(-\beta E_j) \hat{P}_j = \frac{\exp(-\beta \hat{H})}{Z_N} ,$$

where the quantum canonical partition function is

$$Z_N(T, V) = \text{tr}_{\mathcal{H}_{tot}} \left(\frac{\exp(-\beta \hat{H})}{Z_N} \right) .$$

Proof. For a mixed state, the density matrix is (13.9)

$$\rho_c = \sum_j p_j \hat{P}_j = C \sum_j \exp(-\beta E_j) \hat{P}_j ,$$

where the probability is given by (18.1) and C is a normalisation function.

Moreover, using (17.2)

$$\begin{aligned}
\rho_c &= C \sum_j \exp(-\beta E_j) \hat{P}_j \\
&= C \sum_j \sum_k \frac{1}{k!} (-\beta E_j)^k \underbrace{\hat{P}_j}_{(P_j)^k} \\
&= C \sum_j \sum_k \frac{1}{k!} (-\beta E_j \hat{P}_j)^k \\
&= C \sum_k \frac{1}{k!} (-\beta \sum_j E_j \hat{P}_j)^k \\
&= C \exp(-\beta \underbrace{\sum_j E_j \hat{P}_j}_{\hat{H}}) \\
&= C \exp(-\beta \hat{H}) ,
\end{aligned}$$

where we have used the Taylor expansion of the exponential, one of the properties of the projectors (13.2) and we have exchanged the two series.

Finally, We set $C = \frac{1}{Z_N}$, where Z_N is the quantum canonical partition function, and by the normalisation condition

$$1 = \text{tr}_{\mathcal{H}_{tot}} \rho_c = \frac{1}{Z_N} \text{tr}_{\mathcal{H}_{tot}} \exp(-\beta \hat{H}) ,$$

hence

$$Z_N = \text{tr}_{\mathcal{H}_{tot}} \exp(-\beta \hat{H}) .$$

q.e.d.

We define the Helmholtz free energy

$$Z_N = \exp(-\beta F) ,$$

or equivalently

$$F = -\frac{1}{\beta} \log Z_N .$$

The average energy is

$$E = \langle \hat{H} \rangle_c = -\frac{\partial}{\partial \beta} \log Z_N .$$

Proof. In fact,

$$\begin{aligned}
 E &= \langle \hat{H} \rangle_c \\
 &= \text{tr}_{\mathcal{H}_{tot}}(\hat{H} \rho_c) \\
 &= \text{tr}_{\mathcal{H}_{tot}} \left(\hat{H} \frac{\exp(-\beta \hat{H})}{Z_N} \right) \\
 &= \frac{1}{Z_N} \text{tr}_{\mathcal{H}_{tot}} \left(- \frac{\partial}{\partial \beta} \exp(-\beta \hat{H}) \right) \\
 &= - \frac{1}{Z_N} \frac{\partial}{\partial \beta} \underbrace{\text{tr}_{\mathcal{H}_{tot}} \exp(-\beta \hat{H})}_{Z_N} \\
 &= - \frac{1}{Z_N} \frac{\partial}{\partial \beta} Z_N \\
 &= - \frac{\partial}{\partial \beta} \log Z_N .
 \end{aligned}$$

q.e.d.

The entropy is

$$S = \frac{E - F}{T} = \frac{\partial F}{\partial T} .$$

Proof. In fact, using (5.5)

$$\begin{aligned}
 S_c &= -k_B \langle \log \rho_c \rangle_c \\
 &= -k_B \text{tr}_{\mathcal{H}_{tot}}(\rho_c \log \rho_c) \\
 &= -k_B \text{tr}_{\mathcal{H}_{tot}} \left(\frac{\exp(-\beta \hat{H})}{Z_N} \log \frac{\exp(-\beta \hat{H})}{Z_N} \right) \\
 &= -k_B \text{tr}_{\mathcal{H}_{tot}} \left(\frac{\exp(-\beta \hat{H})}{Z_N} (\log \exp(-\beta \hat{H}) - \log Z_N) \right) \\
 &= -k_B \text{tr}_{\mathcal{H}_{tot}} \left(\frac{\exp(-\beta \hat{H})}{Z_N} (-\beta \hat{H} - \log Z_N) \right) \\
 &= k_B \beta \underbrace{\text{tr}_{\mathcal{H}_{tot}} \left(\frac{\exp(-\beta \hat{H})}{Z_N} \hat{H} \right)}_E + k_B \text{tr}_{\mathcal{H}_{tot}} \left(\frac{\exp(-\beta \hat{H})}{Z_N} \underbrace{\log Z_N}_{-\beta F} \right) \\
 &= \frac{E}{T} - k_B \beta F \frac{1}{Z_N} \underbrace{\text{tr}_{\mathcal{H}_{tot}}(\exp(-\beta \hat{H}))}_{Z_N} \\
 &= \frac{E - F}{T} .
 \end{aligned}$$

q.e.d.

Notice that the entropy is well defined because the trace of the exponential of the energy eigenvalues diverges only if they are negative. Thus, we assume that $E_j \geq \min E_j = 0$.

Chapter 19

Grancanonical ensemble

The grancanonical ensemble is characterised by constant volume, temperature and chemical potential. Since N is not fixed, we work in the full Fock space \mathcal{F}_N . However, we restrict the hamiltonian operator in the Fock space to the condition that it conserves the number of particles, i.e. $[\hat{H}, \hat{N}] = 0$

$$\hat{H}\Big|_{\mathcal{F}_N} = \hat{H}_N .$$

An example of physical system which does not satisfy this condition is a photons absorbed by an electron. Energy can be in one of the eigenstates, each for a fixed N

$$\hat{H}^{(N)}|\psi_{j,\alpha}^{(N)}\rangle = E_j^{(N)}|\psi_{j,\alpha}^{(N)}\rangle ,$$

with probability

$$p_j^{(N)} \propto \exp(-\beta(E_j - \mu N)) . \quad (19.1)$$

Consider a family of projectors $\{\hat{P}_j^{(N)}\}$

$$\hat{P}_j^N = \sum_{\alpha} |\psi_{j,\alpha}^{(N)}\rangle \langle \psi_{j,\alpha}^{(N)}| ,$$

the density matrix of a mixed states is

$$\rho_{gc} = \frac{1}{\mathcal{Z}} \sum_N \sum_j \exp(-\beta(E_j - \mu N)) \hat{P}_j^{(N)} = \frac{\exp(-\beta(\hat{H} - \mu \hat{N}))}{\mathcal{Z}} ,$$

where $z = \exp(\beta\mu)$ is the fugacity and the quantum grancanonical partition function is

$$\mathcal{Z} = \sum_{N=0}^{\infty} \text{tr}_{\mathcal{H}_{tot}} \left(\exp(-\beta(\hat{H} - \mu \hat{N})) \right) = \sum_{N=0}^{\infty} z^N Z_N .$$

Proof. For a mixed state, the density matrix is (13.9)

$$\rho_{gc} = \sum_N \sum_j p_j \hat{P}_j^{(N)} = C \sum_N \sum_j \exp(-\beta(E_j^{(N)} - \mu N)) \hat{P}_j^{(N)},$$

where the probability is given by (19.1) and C is a normalisation function.

Moreover, using (17.2) and (17.3)

$$\begin{aligned} \rho_{gc} &= C \sum_N \sum_j \exp(-\beta(E_j - \mu N)) \hat{P}_j^{(N)} \\ &= C \sum_N \sum_j \sum_k \frac{1}{k!} (-\beta(E_j^{(N)} - \mu N))^k \underbrace{\hat{P}_j^{(N)}}_{(P_j^{(N)})^k} \\ &= C \sum_j \sum_k \frac{1}{k!} (-\beta(E_j^{(N)} \hat{P}_j^{(N)} - \mu N P_j^{(N)}))^k \\ &= C \sum_k \frac{1}{k!} (-\beta \sum_N \sum_j (E_j^{(N)} \hat{P}_j^{(N)} - \mu N P_j^{(N)}))^k \\ &= C \exp(-\beta(\underbrace{\sum_j \sum_N E_j^{(N)} \hat{P}_j^{(N)}}_{\hat{H}}) - \mu \underbrace{\sum_j \sum_N N P_j^{(N)}}_{\hat{N}}) \\ &= C \exp(-\beta(\hat{H} - \mu \hat{N})), \end{aligned}$$

where we have used the Taylor expansion of the exponential, one of the properties of the projectors (13.2) and we have exchanged the two series.

Finally, We set $C = \frac{1}{\mathcal{Z}}$, where \mathcal{Z} is the quantum canonical partition function, and by the normalisation condition

$$1 = \text{tr}_{\mathcal{F}} \rho_{gc} = \sum_N \frac{1}{\mathcal{H}_{tot}} \text{tr}_{\mathcal{F}} \exp(-\beta(\hat{H} - \mu \hat{N})),$$

hence

$$\begin{aligned} \mathcal{Z} &= \text{tr}_{\mathcal{F}} \exp(-\beta(\hat{H} - \mu \hat{N})) \\ &= \sum_{N=0}^{\infty} \text{tr}_{\mathcal{H}_{tot}} \exp(-\beta(\hat{H} - \mu \hat{N})) \\ &= \sum_{N=0}^{\infty} z^N \underbrace{\text{tr}_{\mathcal{H}_{tot}} \exp(-\beta \hat{H})}_{Z_N} \\ &= \sum_N z^N Z_N. \end{aligned}$$

q.e.d.

Consider an observable \hat{A} such that it conserves the number of particles, i.e. $[\hat{A}, \hat{N}] = 0$, the average value is

$$\langle \hat{A} \rangle_{gc} = \text{tr}_{\mathcal{F}}(\hat{A} \rho_{gc}) = \frac{1}{\mathcal{Z}} \sum_{N=0}^{\infty} z^N Z_N \langle \hat{A} \rangle_c .$$

Proof. In fact,

$$\begin{aligned} \langle \hat{A} \rangle_{gc} &= \text{tr}_{\mathcal{F}}(\hat{A} \rho_{gc}) \\ &= \sum_{N=0}^{\infty} \text{tr}_{\mathcal{H}_{tot}} \left(\hat{A} \frac{z^N \exp(-\beta \hat{H})}{\mathcal{Z}} \right) = \frac{1}{\mathcal{Z}} \sum_{N=0}^{\infty} z^N \text{tr}_{\mathcal{H}_{tot}}(\hat{A} \exp(-\beta \hat{H})) \\ &= \frac{1}{\mathcal{Z}} \sum_{N=0}^{\infty} z^N Z_N \underbrace{\frac{\text{tr}_{\mathcal{H}_{tot}}(\hat{A} \exp(-\beta \hat{H}))}{Z_N}}_{\langle \hat{A} \rangle_c} \\ &= \frac{1}{\mathcal{Z}} \sum_{N=0}^{\infty} z^N Z_N \langle \hat{A} \rangle_c . \end{aligned}$$

q.e.d.

We define the granpotential

$$\Omega = -\frac{1}{\beta} \log \mathcal{Z} ,$$

the energy in the grancanonical is

$$E - \mu N = \langle \hat{H} - \mu \hat{N} \rangle = -\frac{\partial}{\partial \beta} \log \mathcal{Z} .$$

Proof. In fact

$$\begin{aligned} E - \mu N &= \langle \hat{H} - \mu \hat{N} \rangle \\ &= \text{tr}_{\mathcal{F}} \left((\hat{H} - \mu \hat{N}) \frac{\exp(-\beta(\hat{H} - \mu \hat{N}))}{\mathcal{Z}} \right) \\ &= -\frac{1}{\mathcal{Z}} \frac{\partial}{\partial \beta} \underbrace{\text{tr}_{\mathcal{F}}(\exp(-\beta(\hat{H} - \mu \hat{N})))}_{\mathcal{Z}} \\ &= -\frac{1}{\mathcal{Z}} \frac{\partial}{\partial \beta} \mathcal{Z} \\ &= -\frac{\partial}{\partial \beta} \log \mathcal{Z} . \end{aligned}$$

q.e.d.

The entropy in the grancanonical ensemble is

$$S = \frac{E - \mu N - \Omega}{T} .$$

Proof. In fact

$$\begin{aligned}
 S &= -k_B \langle \log \rho_{gc} \rangle_{gc} \\
 &= -k_B \operatorname{tr}_{\mathcal{F}} (\rho_{gc} \log \rho_{gc}) \\
 &= -k_B \operatorname{tr}_{\mathcal{F}} \left(\frac{\exp(-\beta(\hat{H} - \mu\hat{N}))}{\mathcal{Z}} \log \frac{\exp(-\beta(\hat{H} - \mu\hat{N}))}{\mathcal{Z}} \right) \\
 &= -k_B \operatorname{tr}_{\mathcal{F}} \left(\frac{\exp(-\beta(\hat{H} - \mu\hat{N}))}{\mathcal{Z}} (\log \exp(-\beta(\hat{H} - \mu\hat{N})) - \log \mathcal{Z}) \right) \\
 &= k_B \beta \underbrace{\operatorname{tr}_{\mathcal{F}} \frac{\exp(-\beta(\hat{H} - \mu\hat{N}))}{\mathcal{Z}} (\hat{H} - \mu\hat{N})}_{E - \mu N} + k_B \underbrace{\operatorname{tr}_{\mathcal{F}} \log \mathcal{Z}}_{-\beta\Omega} \\
 &= \frac{E - \mu N - \Omega}{T} .
 \end{aligned}$$

q.e.d.

Chapter 20

Quantum gas

20.1 Generic quantum gas

Consider a quantum gas. The hamiltonian operator of one particle, labelled by k is

$$\hat{H}_k = \epsilon_k \hat{n}_k = \epsilon_k \hat{a}_k^\dagger \hat{a}_k ,$$

where $\hat{n}_k = \hat{a}_k^\dagger \hat{a}_k$ is the number operator and ϵ_k is the energy eigenvalue associated to the eigenbasis $|u_k(x)\rangle$ by the eigenvalue relation

$$\hat{H}_k |u_k(x)\rangle = \epsilon_k |u_k(x)\rangle .$$

Therefore, the hamiltonian one-body operator in the Fock space \mathcal{F} , created by the ladder operators \hat{a}_k^\dagger each associated to the element of the eigenbasis $|u_k(x)\rangle$, is

$$\hat{H} = \sum_k \hat{H}_k = \sum_k \epsilon_k \hat{n}_k = \sum_k \epsilon_k \hat{a}_k^\dagger \hat{a}_k .$$

In \mathcal{F} , the total number onebody operator is

$$\hat{N} = \sum_k \hat{n}_k ,$$

where their eigenvalues are given with respect to an orthonormal basis $|n_1, \dots, n_k, \dots\rangle$ by the eigenvalue relation

$$\hat{n}_k |n_1, \dots, n_k, \dots\rangle = n_k |n_1, \dots, n_k, \dots\rangle .$$

In particular, we distinguish the bosonic and the fermionic case

$$n_k = \begin{cases} 0, 1, 2, \dots & \text{bosons} \\ 0, 1 & \text{fermions} \end{cases} .$$

We exploit the grancanonical ensemble. The grancanonical partition function is

$$\mathcal{Z} = \text{tr}_{\mathcal{F}} \exp(-\beta(\hat{H} - \mu\hat{N})) = \prod_k \sum_{n_1, \dots, n_k, \dots} \exp(-\beta(\epsilon_k - \mu)n_k) .$$

Proof. In fact,

$$\begin{aligned} \mathcal{Z} &= \text{tr}_{\mathcal{F}} \exp(-\beta(\hat{H} - \mu\hat{N})) \\ &= \sum_{n_1, \dots, n_k, \dots} \langle n_1, \dots, n_k, \dots | \exp(-\beta \sum_k (\epsilon_k - \mu) \hat{n}_k) | n_1, \dots, n_k, \dots \rangle \\ &= \sum_{n_1, \dots, n_k, \dots} \langle n_1, \dots, n_k, \dots | \underbrace{\exp(-\beta \sum_k (\epsilon_k - \mu) n_k)}_{\prod_k \exp} | n_1, \dots, n_k, \dots \rangle \\ &= \sum_{n_1, \dots, n_k, \dots} \prod_k \exp(\beta(\epsilon_k - \mu)n_k) \langle n_1, \dots, n_k, \dots | n_1, \dots, n_k, \dots \rangle \\ &= \prod_k \sum_{n_1, \dots, n_k, \dots} \exp(-\beta(\epsilon_k - \mu)n_k) , \end{aligned}$$

where in the last passage, we have switched the product with the sum because n_1, \dots, n_k, \dots are independent. q.e.d.

Furthermore, for bosons and fermions, it becomes

$$\mathcal{Z} = \begin{cases} \prod_k \frac{1}{1 - \exp(-\beta(\epsilon_k - \mu))} & \text{bosons} \\ \prod_k (1 + \exp(-\beta(\epsilon_k - \mu))) & \text{fermions} \end{cases} ,$$

or, in compact notation,

$$\mathcal{Z}_{\mp} = \prod_k \left(1 \mp \exp(-\beta(\epsilon_k - \mu)) \right)^{\mp} ,$$

where the minus is associated to bosons and the plus to fermions.

Proof. For fermions, $n_k = 0, 1$

$$\mathcal{Z}_+ = \prod_k \sum_{n_1, \dots, n_k, \dots=0}^1 \exp(-\beta(\epsilon_k - \mu)n_k) = \prod_k (1 + \exp(-\beta(\epsilon_k - \mu))) .$$

For bosons, $n_k = 0, 1, 2, \dots$

$$\begin{aligned}
 \mathcal{Z}_- &= \prod_k \sum_{n_1, \dots, n_k, \dots=0}^{\infty} \exp(-\beta(\epsilon_k - \mu)n_k) \\
 &= \prod_k \underbrace{\sum_{n_1, \dots, n_k, \dots=0}^{\infty} \exp(-\beta(\epsilon_k - \mu))^{n_k}}_{\text{geometrical series}} \\
 &= \prod_k \frac{1}{1 - \exp(-\beta(\epsilon_k - \mu))} .
 \end{aligned}$$

Notice that the condition of convergence of the geometrical series is $\mu < \min \epsilon_k = 0$, which we have set to zero for convenience. q.e.d.

The grancanonical potential is

$$\Omega_{\mp} = -\frac{1}{\beta} \log \mathcal{Z}_{\mp} = \pm \frac{1}{\beta} \sum_k \log \left(1 \mp \exp(-\beta(\epsilon_k - \mu)) \right) .$$

Proof. In fact

$$\begin{aligned}
 \Omega_{\mp} &= -\frac{1}{\beta} \log \mathcal{Z}_{\mp} \\
 &= -\frac{1}{\beta} \log \left(\underbrace{\prod_k (1 \mp \exp(-\beta(\epsilon_k - \mu)))^{\mp}}_{\sum_k \log} \right) \\
 &= -(\mp) \sum_k \log \left(1 \mp \exp(-\beta(\epsilon_k - \mu)) \right) \\
 &= \pm \frac{1}{\beta} \sum_k \log \left(1 \mp \exp(-\beta(\epsilon_k - \mu)) \right) .
 \end{aligned}$$

q.e.d.

The grancanonical average number of particle in an energy level state \bar{k} is

$$\langle \hat{n}_{\bar{k}} \rangle_{gc} = \text{tr}_{\mathcal{F}} \left(\hat{n}_{\bar{k}} \frac{\exp(-\beta \sum_k (\epsilon_k - \mu) \hat{n}_k)}{\mathcal{Z}} \right) = \frac{\partial \Omega}{\partial \epsilon_{\bar{k}}} = \frac{1}{\exp(\beta(\epsilon_{\bar{k}} \mp 1))} .$$

Proof. In fact

$$\begin{aligned}
\langle \hat{n}_{\bar{k}} \rangle_{gc} &= \text{tr}_{\mathcal{F}} \left(\hat{n}_{\bar{k}} \frac{\exp(-\beta \sum_k (\epsilon_k - \mu) \hat{n}_k)}{\mathcal{Z}} \right) \\
&= \frac{1}{\mathcal{Z}} \text{tr}_{\mathcal{F}} \left(-\frac{1}{\beta} \frac{\partial}{\partial \epsilon_{\bar{k}}} \exp(-\beta \sum_k (\epsilon_k - \mu) \hat{n}_k) \right) \\
&= -\frac{1}{\beta \mathcal{Z}} \frac{\partial}{\partial \epsilon_{\bar{k}}} \underbrace{\text{tr}_{\mathcal{F}} \left(\exp(-\beta \sum_k (\epsilon_k - \mu) \hat{n}_k) \right)}_{\mathcal{Z}} \\
&= -\frac{1}{\beta \mathcal{Z}} \frac{\partial}{\partial \epsilon_{\bar{k}}} \mathcal{Z} = \\
&= \frac{\partial}{\partial \epsilon_{\bar{k}}} \underbrace{\left(-\frac{\log \mathcal{Z}}{\beta} \right)}_{\Omega} \\
&= \frac{\partial}{\partial \epsilon_{\bar{k}}} \Omega .
\end{aligned}$$

Therefore,

$$\begin{aligned}
\frac{\partial}{\partial \epsilon_{\bar{k}}} \Omega &= \frac{\partial}{\partial \epsilon_{\bar{k}}} \left(\pm \frac{1}{\beta} \sum_k \log(1 \mp \exp(-\beta(\epsilon_k - \mu))) \right) \\
&= \pm \frac{1}{\beta} (-\beta) \frac{\exp(-\beta(\epsilon_k - \mu))}{1 \mp \exp(-\beta(\epsilon_k - \mu))} \\
&= \mp \frac{1}{1 \mp \exp(\beta(\epsilon_k - \mu))} \\
&= \frac{1}{\exp(\beta(\epsilon_k - \mu)) \mp 1} .
\end{aligned}$$

q.e.d.

The average total number of particle is

$$N = \langle \hat{N} \rangle_{gc} = \left\langle \sum_k \hat{n}_k \right\rangle_{gc} = \sum_k \frac{1}{\exp(\beta(\epsilon_k - \mu)) \mp 1} .$$

The average energy is

$$E = \langle \hat{H} \rangle_{gc} = \text{tr}_{\mathcal{F}} \left(\hat{H} \frac{\exp(-\beta(\hat{H} - \mu \hat{N}))}{\mathcal{Z}} \right) = \sum_k \epsilon_k \langle \hat{n}_k \rangle$$

Proof. In fact

$$\begin{aligned}
E &= \langle \hat{H} \rangle_{gc} \\
&= \text{tr}_{\mathcal{F}} \left(\hat{H} \frac{\exp(-\beta(\hat{H} - \mu\hat{N}))}{\mathcal{Z}} \right) \\
&= \frac{1}{\mathcal{Z}} \text{tr}_{\mathcal{F}} \left(-\frac{\partial}{\partial\beta} \exp(-\beta(\hat{H} - \mu\hat{N})) \right) \\
&= -\frac{1}{\mathcal{Z}} \frac{\partial}{\partial\beta} \underbrace{\text{tr}_{\mathcal{F}} \left(\exp(-\beta(\hat{H} - \mu\hat{N})) \right)}_{\mathcal{Z}} \\
&= -\frac{\partial}{\partial\beta} \Big|_z \log \mathcal{Z} \\
&= -\frac{\partial}{\partial\beta} \Big|_z \left(\mp \sum_k \log \left(1 \mp \exp(-\beta(\epsilon_k - \mu)) \right) \right) \\
&= \mp \sum_k \frac{\epsilon_k \exp(-\beta(\epsilon_k - \mu))}{1 \mp \exp(-\beta(\epsilon_k - \mu))} \\
&= \sum_k \frac{\epsilon_k}{\exp(\beta(\epsilon_k - \mu)) \mp 1} \\
&= \sum_k \epsilon_k \langle \hat{n}_k \rangle
\end{aligned}$$

where we have kept the fugacity z constant.

q.e.d.

20.2 Non-relativistic non-interacting quantum gas

So far, we have made computations for a generic quantum gas. From now on, we will deal with non-relativistic non-interacting quantum gas. The finite-volume energy eigenvalues are

$$\epsilon_k = \frac{\hbar^2 k^2}{2m} \quad \mathbf{k} = \frac{2\pi}{L} \mathbf{n} ,$$

where $\mathbf{n} = (n_1, n_2, n_3) \in \mathbb{Z}^3$. In the thermodynamic limit, the spectrum \mathbf{k} becomes continuous, but \mathbf{n} not, because

$$\Delta K_i = \frac{2\pi}{L} (n_i + 1 - n_i) = \frac{2\pi}{L} .$$

Therefore, sums in k becomes integrals in dk

$$\sum_k = \sum_{n_1, n_2, n_3=-\infty}^{\infty} \rightarrow \frac{V}{2\pi^2} \int dk \, k^2 .$$

Proof. In fact, in 1-dimensional

$$\sum_{n_1} \underbrace{\Delta n_1}_1 = \sum_{k_1} \frac{L}{2\pi} \Delta k_1 \rightarrow \frac{L}{2\pi} \int dk_1 .$$

Similarly, in the 3-dimensional case

$$\begin{aligned} \sum_{n_1, n_2, n_3 = -\infty}^{\infty} \underbrace{\Delta n_1 \Delta n_2 \Delta n_3}_1 &\rightarrow \left(\frac{L}{2\pi}\right)^3 \int dk_1 dk_2 dk_3 \\ &= \left(\frac{L}{2\pi}\right)^3 \int dk_1 dk_2 dk_3 \\ &= \left(\frac{L}{2\pi}\right)^3 \int dk^3 \\ &= \left(\frac{L}{2\pi}\right)^3 4\pi \int dk k^2 \\ &= \frac{V}{2\pi^2} \int dk k^2 . \end{aligned}$$

q.e.d.

The grandcanonical potential is

$$\Omega_{\mp} = \mp \frac{2}{3} AV \int_0^{\infty} d\epsilon^{\frac{3}{2}} \frac{1}{\exp(\beta(\epsilon - \mu)) \mp 1} .$$

Proof. In fact

$$\begin{aligned} \Omega_{\mp} &= \pm \frac{1}{\beta} \sum_k \log \left(1 \mp \exp(-\beta(\epsilon_k - \mu)) \right) \\ &\rightarrow \pm \frac{1}{\beta} \frac{V}{2\pi^2} \int_{-\infty}^{\infty} dk k^2 \log \left(1 \mp \exp(-\beta(\epsilon_k - \mu)) \right) . \end{aligned}$$

Under a change of variable

$$\epsilon = \frac{\hbar^2 k^2}{2m} , \quad k^2 dk = \frac{1}{2} \left(\frac{2m}{\hbar^2} \right)^{\frac{3}{2}} \sqrt{\epsilon} d\epsilon ,$$

we obtain

$$\begin{aligned}
\Omega_{\mp} &= \pm \frac{AV}{\beta} \int_0^{\infty} \underbrace{d\epsilon \sqrt{\epsilon}}_{\frac{2}{3} d\epsilon^{\frac{3}{2}}} \log(1 \mp \exp(-\beta(\epsilon_k - \mu))) \\
&= \pm \frac{2}{3} \frac{AV}{\beta} \int_0^{\infty} d\epsilon^{\frac{3}{2}} \log(1 \mp \exp(-\beta(\epsilon_k - \mu))) \\
&= \pm \frac{2}{3} \frac{AV}{\beta} \underbrace{\epsilon^{\frac{3}{2}}}_{0 \text{ for } \epsilon=0} \underbrace{\log(1 \mp \exp(-\beta(\epsilon_k - \mu)))}_{0 \text{ for } \epsilon=\infty} \Big|_0^{\infty} \\
&\quad \mp \frac{2}{3} \frac{AV}{\beta} \int_0^{\infty} d\epsilon^{\frac{3}{2}} \frac{1}{\exp(\beta(\epsilon - \mu)) \mp 1} \\
&= \mp \frac{2}{3} AV \int_0^{\infty} d\epsilon^{\frac{3}{2}} \frac{1}{\exp(\beta(\epsilon - \mu)) \mp 1} \\
&= \mp \frac{2}{3} AV \int_0^{\infty} d\epsilon^{\frac{3}{2}} \frac{1}{\exp(\beta(\epsilon - \mu)) \mp 1} .
\end{aligned}$$

where we have integrated by parts and, introducing the degeneracy ($g = 2s + 1$ for spin particles), we have called

$$A = \frac{g}{4\pi^2} \left(\frac{2m}{\hbar^2} \right)^{\frac{3}{2}} .$$

q.e.d.

The equation of state reads as

$$\Omega = -pV = -\frac{2}{3}E .$$

Furthermore, we have the formulas

$$\begin{aligned}
N &= AV \int_0^{\infty} d\epsilon \epsilon^{\frac{1}{2}} n(\epsilon) , \\
P &= \frac{2}{3} \frac{E}{V} = \frac{2}{3} A \int_0^{\infty} d\epsilon \epsilon^{\frac{3}{2}} n(\epsilon) .
\end{aligned}$$

20.3 Expanding with respect to fugacity z

We can expand the density with respect to the fugacity $z = \exp(\beta\mu) \geq 0$

$$n = \frac{g}{\lambda_T^3} f_{\frac{3}{2}}^{\mp} ,$$

where

$$f_l^{\mp} = \begin{cases} \sum_{n=0}^{\infty} \frac{2^{n+1}}{(n+1)^l} & f^- \text{ for bosons} \\ \sum_{n=0}^{\infty} \frac{(-1)^n 2^{n+1}}{(n+1)^l} & f^+ \text{ for fermions} \end{cases} .$$

Proof. Under a change of variable

$$x^2 = \beta\epsilon, \quad \beta d\epsilon = 2x dx,$$

we obtain

$$\begin{aligned} n &= A \int_0^\infty dx \frac{2x}{\beta} \frac{x}{\sqrt{(\beta)(\exp(x^2)z^{-1}) \mp 1}} \\ &= \frac{4g}{\sqrt{\pi}\lambda_T^3} \int_0^\infty dx \frac{x^2 z}{\exp(x^2) \mp 2} \\ &= \frac{4g}{\sqrt{\pi}\lambda_T^3} \int_0^\infty dx x^2 z \exp(-x^2) \sum_{n=0}^\infty (\pm 1)^n z^n \exp(-nx^2) \\ &= \frac{4g}{\sqrt{\pi}\lambda_T^3} \sum_{n=0}^\infty (\pm 1)^n z^{n+1} \underbrace{\int_0^\infty dx x^2 \exp(-x^2(n+1))}_{\frac{\sqrt{\pi}}{4(n+1)^{\frac{3}{2}}}} \\ &= \frac{g}{\lambda_T^3} \sum_{n=0}^\infty (\pm 1)^n \frac{z^{n+1}}{(n+1)^{\frac{3}{2}}} \\ &= \frac{g}{\lambda_T^3} f_{\frac{3}{2}}^\mp. \end{aligned}$$

q.e.d.

Notice that for bosons, the convergence of the series implies $z < 1$, which means $\mu > 0$.

20.4 Classical limit

20.5 Semiclassical limit

Chapter 21

Fermions

In this chapter, we restrict ourselves with the fermionic case. The equations of state are

$$n = \frac{g}{\lambda_T^3} f_{\frac{3}{2}}^+(z) , \quad \beta p = \frac{g}{\lambda_T^3} f_{\frac{5}{2}}^+(z) ,$$

where

$$f_l^+(z) = \sum_{n=0}^{\infty} \frac{(-1)^n z^{n+1}}{(n+1)^l}$$

which is an alternate-sign power series in $z = \exp(\beta\mu) > 0$, always positive. It absolutely converges for $z < 1$ and pointwisely converges for $z > 1$. Moreover, it is a monotonic function in z .

It is interesting to study its behaviour for $z \ll 1$. In fact, in the classical limit

$$f_{\frac{3}{2}}(z) \sim f_{\frac{5}{2}}(z) \sim z ,$$

and in the semiclassical limit

$$f_{\frac{3}{2}}(z) \sim z - \frac{z^2}{2^{\frac{3}{2}}} , \quad f_{\frac{5}{2}}(z) \sim z - \frac{z^2}{2^{\frac{5}{2}}} .$$

21.1 Low temperature limit

For the zero temperature limit $T = 0$, the Fermi-Dirac distribution becomes

$$n(\epsilon) = \frac{1}{\exp(\beta(\epsilon - \mu)) + 1} \xrightarrow{T \rightarrow 0} \begin{cases} 0 & \epsilon > \mu \\ \frac{1}{2} & \epsilon = \mu \\ 1 & \epsilon < \mu \end{cases} .$$

It is a step function in $\epsilon = \mu$. This energy value is called Fermi energy ϵ_F . Physically, it means that all the states below this energy level are occupied. Hence, for $\epsilon < \epsilon_F$,

we have as many states as particles. If we add a particle, we increase ϵ_F , whereas if we remove a particle, we decrease ϵ_F . This is the procedure to dope a material.

For small T , it is no longer a step function, but it can be accurately approximate to it for a certain range $\Delta\epsilon$. Physically, more energetic particle are transfered over ϵ_F . We define Fermi temperature T_F

$$\epsilon_F = \lim_{T \rightarrow 0} \mu(T) = k_B T_F .$$

In fact, if $\Delta\epsilon \ll \epsilon_F$, which means $T \ll T_F$, we can approximate $n(\epsilon)$ with a step function without making a big error.

21.2 Fermi Energy for a non-relativistic non-interacting quantum gas

In the 3-dimensional case, the density is

$$n = A \frac{2}{3} \epsilon_F^{\frac{3}{2}} .$$

Proof. In fact, using $n(\epsilon) = \theta(-\epsilon_F)$

$$\begin{aligned} n &= A \int_0^\infty d\epsilon \epsilon^{\frac{1}{2}} n(\epsilon) \\ &= A \int_0^{\epsilon_F} d\epsilon \epsilon^{\frac{1}{2}} \\ &= A \frac{2}{3} \epsilon_F^{\frac{3}{2}} . \end{aligned}$$

q.e.d.

The energy is

$$E = AV \frac{2}{5} \epsilon_F^{\frac{5}{2}} .$$

Proof. In fact, using $n(\epsilon) = \theta(-\epsilon_F)$

$$\begin{aligned} n &= A \int_0^\infty d\epsilon \epsilon^{\frac{3}{2}} n(\epsilon) \\ &= A \int_0^{\epsilon_F} d\epsilon \epsilon^{\frac{3}{2}} \\ &= A \frac{2}{5} \epsilon_F^{\frac{5}{2}} . \end{aligned}$$

q.e.d.

Notice that at $T = 0$, there is a positive pressure

$$p = \frac{2}{5}n\epsilon_F > 0 .$$

This can be seen visually, because at $T = 0$, there are particles with energy $\epsilon \neq 0$, unlike the classical case, in which $p = 0$.

Proof. In fact

$$p = \frac{2}{3} \frac{E}{V} = \frac{2}{3} \frac{E}{N} \frac{N}{V} = \frac{2}{5}n\epsilon_F .$$

q.e.d.

Chapter 22

Bosons

In this chapter, we restrict ourselves with the bosonic case. The equations of state are

$$n = \frac{g}{\lambda_T^3} f_{\frac{3}{2}}^-(z) , \quad \beta p = \frac{g}{\lambda_T^3} f_{\frac{5}{2}}^-(z)$$

where

$$f_l^-(z) = \sum_{n=0}^{\infty} \frac{z^{n+1}}{(n+1)^l}$$

which is a positive-terms power series in $z = \exp(\beta\mu) > 0$, always positive. It absolutely converges for $z < 1$ and converges for $z > 1$ only if $l < 2$. Moreover, it is a monotonic function in z . At $z = 1$, it becomes the Riemann zeta

$$g_{\frac{3}{2}}(z = 1) = \sum_{n=0}^{\infty} \frac{1}{n+1}^{\frac{3}{2}} = \zeta\left(\frac{3}{2}\right) .$$

Notice that in $z = 1$, it has a vertical derivative, and for $z > 1$, it is not defined according to the physical $\mu > 0$ in the grandcanonical ensemble.

We can study the behaviour of the chemical potential μ . It goes to $-\infty$ for $T \rightarrow \infty$ but the equilibrium condition implies that $\frac{\partial \mu}{\partial T} < 0$, therefore, it cannot increase.

22.1 Low temperature limit

22.2 Bose-Einstein condensate

Part VI

Application of quantum statistical
mechanics

Chapter 23

Quantum ensemble

23.1 Magnetic 1/2-spin

Consider a system composed by N distinguishable magnetic dipoles in an external magnetic field along the z -axis, with spin $S = 1/2$. Its hamiltonian is

$$\hat{H} = \sum_i S_i^{(z)} B ,$$

where

$$S_i^{(z)} = \frac{1}{2} \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix} .$$

The canonical partition function is

$$Z = \left(2 \cosh \frac{\beta B}{2} \right)^N .$$

Proof. By definition, for distinguishable particles,

$$\begin{aligned} Z &= (Z_1)^N \\ &= \left(\text{tr}_{\mathcal{H}} \exp(-\beta \hat{H}_1) \right)^N \\ &= \left(\text{tr}_{\mathcal{H}} \exp(-\beta B \begin{bmatrix} \frac{1}{2} & 0 \\ 0 & -\frac{1}{2} \end{bmatrix}) \right)^N \\ &= \left(\text{tr}_{\mathcal{H}} \begin{bmatrix} \exp(-\frac{\beta B}{2}) & 0 \\ 0 & \exp(\frac{\beta B}{2}) \end{bmatrix} \right)^N \\ &= \left(\exp(-\frac{\beta B}{2}) + \exp(\frac{\beta B}{2}) \right)^N \\ &= \left(2 \cosh \frac{\beta B}{2} \right)^N . \end{aligned}$$

q.e.d.

The Helmholtz free energy F is

$$F = -Nk_B T \ln \left(2 \cosh \frac{\beta B}{2} \right) .$$

Proof. By definition,

$$F = -\frac{\ln Z}{\beta} = -Nk_B T \ln \left(2 \cosh \frac{\beta B}{2} \right) .$$

q.e.d.

The internal energy E is

$$E = -N \frac{B}{2} \tanh \frac{\beta B}{2} .$$

Proof. By definition,

$$E = -\frac{\partial \ln Z}{\partial \beta} = -N \frac{\partial \beta}{\partial \ln} \left(2 \cosh \frac{\beta B}{2} \right) = -N \frac{B}{2} \tanh \frac{\beta B}{2} .$$

q.e.d.

A plot of this is in Figure 23.1.

The magnetisation M is

$$M = -\frac{N}{2} \tanh \frac{\beta B}{2} .$$

Proof. By definition,

$$M = \frac{\partial F}{\partial B} = -Nk_B T \frac{\partial}{\partial B} \ln \left(2 \cosh \frac{\beta B}{2} \right) = -\frac{N}{2} \tanh \frac{\beta B}{2} .$$

q.e.d.

A plot of this is in Figure 23.2.

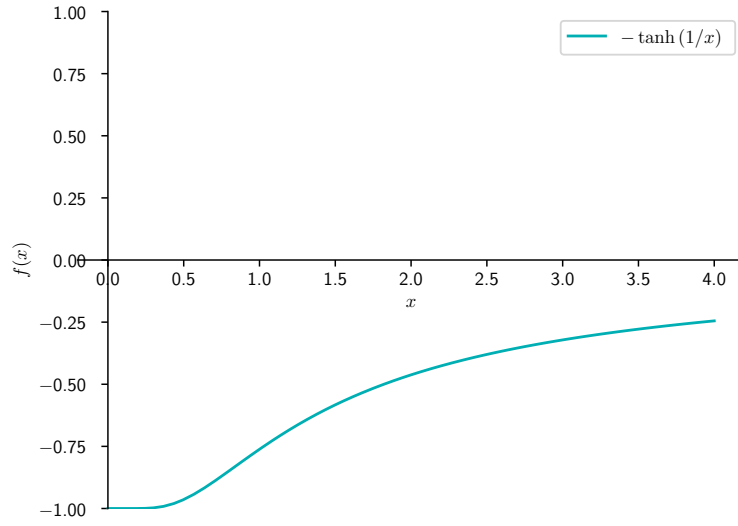


Figure 23.1: A plot of the energy E as a function of T . We have used $x = \frac{2k_B T}{B}$ and $f(x) = \frac{2E}{BN}$.

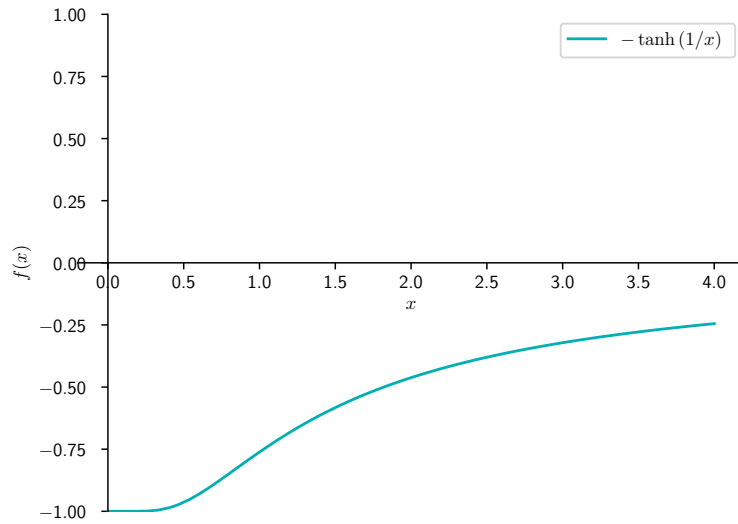


Figure 23.2: A plot of the magnetisation M as a function of T . We have used $x = \frac{2k_B T}{B}$ and $f(x) = \frac{2M}{N}$.

23.2 Magnetic 1-spin

Consider a system composed by N distinguishable magnetic dipoles in an external magnetic field along the z -axis, with spin $S = 1$. Its hamiltonian is

$$\hat{H} = \sum_i S_i^{(z)} B ,$$

where

$$S_i^{(z)} = \begin{bmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -1 \end{bmatrix} .$$

The canonical partition function is

$$Z = \left(2 \cosh(\beta B) + 1 \right)^N .$$

Proof. By definition, for distinguishable particles,

$$\begin{aligned} Z &= (Z_1)^N \\ &= \left(\text{tr}_{\mathcal{H}} \exp(-\beta \hat{H}_1) \right)^N \\ &= \left(\text{tr}_{\mathcal{H}} \exp(-\beta B \begin{bmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -1 \end{bmatrix}) \right)^N \\ &= \left(\text{tr}_{\mathcal{H}} \begin{bmatrix} \exp(-\beta B) & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & \exp(\beta B) \end{bmatrix} \right)^N \\ &= \left(\exp(-\beta B) + 1 + \exp(\beta B) \right)^N \\ &= \left(2 \cosh(\beta B) + 1 \right)^N . \end{aligned}$$

q.e.d.

The Helmholtz free energy F is

$$F = -Nk_B T \ln \left(2 \cosh(\beta B) + 1 \right) .$$

Proof. By definition,

$$F = -\frac{\ln Z}{\beta} = -Nk_B T \ln \left(2 \cosh(\beta B) + 1 \right) .$$

q.e.d.

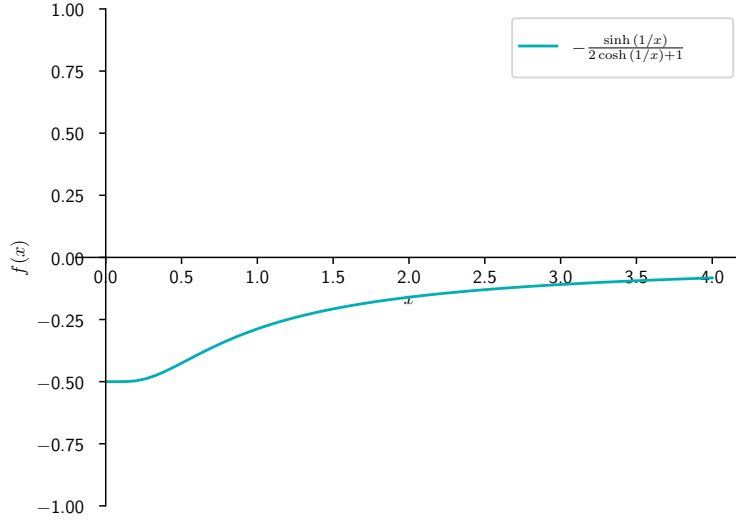


Figure 23.3: A plot of the energy E as a function of T . We have used $x = \frac{k_B T}{B}$ and $f(x) = \frac{2E}{BN}$.

The internal energy E is

$$E = -2NB \frac{\sinh(\beta B)}{2 \cosh(\beta B) + 1} .$$

Proof. By definition,

$$E = -\frac{\partial \ln Z}{\partial \beta} = -N \frac{\partial}{\partial \beta} \ln \left(2 \cosh(\beta B) + 1 \right) = -2NB \frac{\sinh(\beta B)}{2 \cosh(\beta B) + 1} .$$

q.e.d.

A plot of this is in Figure 23.3.

The magnetisation M is

$$M = -\frac{N}{2} \frac{\sinh(\beta B)}{2 \cosh(\beta B) + 1} .$$

Proof. By definition,

$$M = \frac{\partial F}{\partial B} = -Nk_B T \frac{\partial}{\partial B} \ln \left(2 \cosh(\beta B) + 1 \right) = -2N \frac{\sinh(\beta B)}{2 \cosh(\beta B) + 1} .$$

q.e.d.

A plot of this is in Figure 23.4.

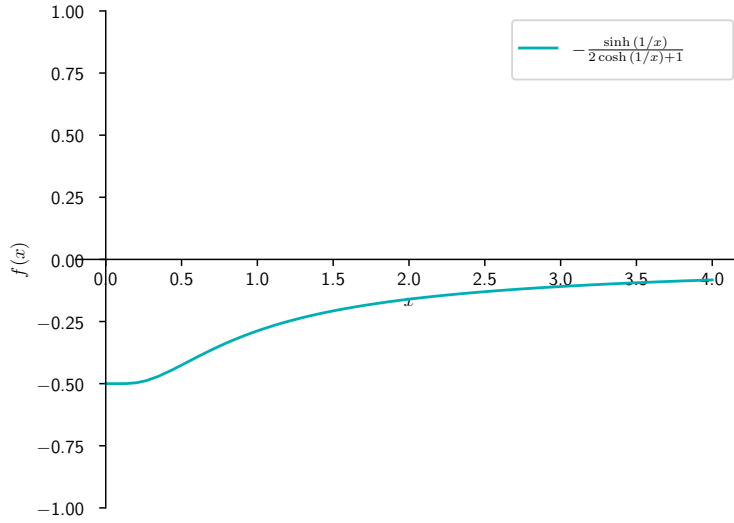


Figure 23.4: A plot of the magnetisation M as a function of T . We have used $x = \frac{2k_B T}{B}$ and $f(x) = \frac{M}{2N}$.

23.3 Quantum harmonic oscillators

Consider a system composed by N distinguishable quantum harmonic oscillators. Its hamiltonian is

$$\hat{H} = \sum_i \hbar\omega(\hat{a}_i\hat{a}_i^\dagger + \frac{1}{2}) .$$

The canonical partition function is

$$Z = \left(\frac{\exp(-\frac{\beta\hbar\omega}{2})}{1 - \exp(-\beta\hbar\omega)} \right)^N .$$

Proof. By definition, for distinguishable particles,

$$\begin{aligned}
 Z &= (Z_1)^N \\
 &= \left(\text{tr}_{\mathcal{H}} \exp(-\beta \hat{H}_1) \right)^N \\
 &= \left(\text{tr}_{\mathcal{H}_i} \exp(-\beta \hbar \omega (\hat{a}_i \hat{a}_i^\dagger + \frac{1}{2})) \right)^N \\
 &= \left(\sum_i \langle n_i | \exp(-\beta \hbar \omega (\hat{a}_i \hat{a}_i^\dagger + \frac{1}{2})) | n_i \rangle \right)^N \\
 &= \left(\sum_i \exp(-\beta \hbar \omega (n_i + \frac{1}{2})) \right)^N \\
 &= \left(\exp(-\frac{\beta \hbar \omega}{2}) \sum_i \exp(-\beta \hbar \omega n_i) \right)^N \\
 &= \left(\exp(-\frac{\beta \hbar \omega}{2}) \frac{1}{1 - \exp(-\beta \hbar \omega)} \right)^N \\
 &= \left(\frac{\exp(-\frac{\beta \hbar \omega}{2})}{1 - \exp(-\beta \hbar \omega)} \right)^N .
 \end{aligned}$$

q.e.d.

The Helmholtz free energy F is

$$F = N k_B T \left(\frac{\beta \hbar \omega}{2} + \ln(1 - \exp(-\beta \hbar \omega)) \right) .$$

Proof. By definition,

$$\begin{aligned}
 F &= -\frac{\ln Z}{\beta} \\
 &= -N k_B T \ln \left(\frac{\exp(-\frac{\beta \hbar \omega}{2})}{1 - \exp(-\beta \hbar \omega)} \right) \\
 &= -N k_B T \left(\ln \exp(-\frac{\beta \hbar \omega}{2}) - \ln(1 - \exp(-\beta \hbar \omega)) \right) \\
 &= -N k_B T \left(-\frac{\beta \hbar \omega}{2} - \ln(1 - \exp(-\beta \hbar \omega)) \right) \\
 &= N k_B T \left(\frac{\beta \hbar \omega}{2} + \ln(1 - \exp(-\beta \hbar \omega)) \right) .
 \end{aligned}$$

q.e.d.

The internal energy E is

$$E = N \left(\frac{\hbar \omega}{2} + \frac{\hbar \omega}{\exp(-\beta \hbar \omega) - 1} \right) .$$

Proof. By definition,

$$\begin{aligned}
 F &= -\frac{\partial \ln Z}{\partial \beta} \\
 &= -N \frac{\partial}{\partial \beta} \ln \left(\frac{\exp(-\frac{\beta \hbar \omega}{2})}{1 - \exp(-\beta \hbar \omega)} \right) \\
 &= -N \frac{\partial}{\partial \beta} (\ln \exp(-\frac{\beta \hbar \omega}{2}) - \ln(1 - \exp(-\beta \hbar \omega))) \\
 &= -N \frac{\partial}{\partial \beta} (-\frac{\beta \hbar \omega}{2} - \ln(1 - \exp(-\beta \hbar \omega))) \\
 &= N \frac{\partial}{\partial \beta} (\frac{\beta \hbar \omega}{2} + \ln(1 - \exp(-\beta \hbar \omega))) \\
 &= N (\frac{\hbar \omega}{2} - \frac{\hbar \omega}{1 - \exp(-\beta \hbar \omega)}) \\
 &= N (\frac{\hbar \omega}{2} + \frac{\hbar \omega}{\exp(-\beta \hbar \omega) - 1}) .
 \end{aligned}$$

q.e.d.

The specific heat is

$$C_V = N \frac{\hbar^2 \omega^2}{k_B T^2} \frac{\exp(\beta \hbar \omega)}{(\exp(\beta \hbar \omega) - 1)^2} .$$

Proof. In fact

$$C_V = \frac{\partial E}{\partial T} = N \frac{\partial}{\partial T} (\frac{\hbar \omega}{2} + \frac{\hbar \omega}{\exp(-\beta \hbar \omega) - 1}) = N \frac{\hbar^2 \omega^2}{k_B T^2} \frac{\exp(\beta \hbar \omega)}{(\exp(\beta \hbar \omega) - 1)^2} .$$

q.e.d.

Chapter 24

Fermions

24.1 White dwarf

A white dwarf is an helium star with mass $M \sim 10^{30}kg$ and a density of $\rho = 10^{10}kg/m^3$ at a temperature of 10^7K . Our approximated model is composed by N electrons and $N/2$ helium nuclei.

Assuming $M = N(m_e + 2m_p) \sim 2Nm_p$, the electronic density is

$$n = 3 \times 10^{36}m^{-3} .$$

Proof. In fact

$$n = \frac{N}{V} = \frac{N\rho}{M} = \frac{N\rho}{2Nm_p} = \frac{\rho}{2m_p} = \frac{10^{10}}{2 \times 1.6 \times 10^{-27}} = 3 \times 10^{36}m^{-3} .$$

q.e.d.

The Fermi momentum p_F is

$$p_F = h \left(\frac{3n}{4\pi g} \right)^{1/3} = 6.63 \times 10^{-34} \times \left(\frac{3 \times 10^{36}}{4 \times 3.14 \times 2} \right)^{1/3} = 0.88MeV/c .$$

Proof. In fact, using $p = \hbar k$

$$N = g \sum_n \rightarrow g \frac{V}{(2\pi)^3} \int d^3k = g \frac{V}{(2\pi\hbar)^3} \int d^3p = g \frac{4\pi V}{(2\pi\hbar)^3} \int_0^{p_F} dp p^2 = g \frac{4\pi V}{(2\pi\hbar)^3} \frac{p_F^3}{3} ,$$

hence

$$p_F = h \left(\frac{3n}{4\pi g} \right)^{1/3} = 6.63 \times 10^{-34} \times \left(\frac{3 \times 10^{36}}{4 \times 3.14 \times 2} \right)^{1/3} = 0.88MeV/c .$$

q.e.d.

The Fermi energy ϵ_F is

$$\epsilon_F = \sqrt{(p_F c)^2 + (mc^2)^2} - mc^2 = 0.5 \text{ Mev} .$$

The Fermi temperature T_F is

$$T_F = \frac{\epsilon_F}{k_B} = 10^{10} K ,$$

which means that we are in the regime $T \ll T_F$ and we can use $T = 0$.

The internal energy E is

$$E = \frac{\pi V m^4 c^5}{\pi^2 \hbar^3} f(x_F) .$$

Proof. In fact,

$$\begin{aligned} E &= g \sum_n \epsilon \rightarrow g \frac{V}{(2\pi)^3} \int d^3 k \epsilon \\ &= g \frac{V}{(2\pi \hbar)^3} \int d^3 p \epsilon \\ &= g \frac{4\pi V}{(2\pi \hbar)^3} \int_0^{p_F} dp p^2 \epsilon \\ &= g \frac{4\pi V}{(2\pi \hbar)^3} \int_0^{p_F} dp p^2 c \sqrt{p^2 + (mc)^2} . \end{aligned}$$

Now we make a change of variable

$$x = \frac{p}{mc} , \quad dp = mc dx ,$$

hence

$$\begin{aligned} E &= g \frac{4\pi V}{(2\pi \hbar)^3} c (mc)^3 \int_0^{x_F} dx x^2 (mc) \sqrt{x^2 + 1} \\ &= \frac{4g\pi V m^4 c^5}{h^3} \int_0^{x_F} dx x^2 \sqrt{x^2 + 1} \\ &= \frac{4g\pi V m^4 c^5}{h^3} f(x_F) \\ &= \frac{V m^4 c^5}{\pi^2 \hbar^3} f(x_F) , \end{aligned}$$

where

$$f(x_F) = \int_0^{x_F} dx x^2 \sqrt{x^2 + 1} .$$

q.e.d.

The pressure P is

$$P = \frac{m^4 c^5}{\pi^2 \hbar^3} \left(\frac{x_F^3}{3} \sqrt{1 + x_F^2} - f(x_F) \right) .$$

Proof. In fact,

$$\begin{aligned} P &= - \frac{\partial E}{\partial V} \\ &= - \frac{\partial}{\partial V} \frac{V m^4 c^5}{\pi^2 \hbar^3} f(x_F) \\ &= - \frac{\pi m^4 c^5}{\pi^2 \hbar^3} f(x_F) - \frac{V m^4 c^5}{\pi^2 \hbar^3} \frac{\partial x_F}{\partial V} \frac{\partial f(x_F)}{\partial x_F} \\ &= - \frac{m^4 c^5}{\pi^2 \hbar^3} f(x_F) - \frac{V m^4 c^5}{\pi^2 \hbar^3} \frac{\partial}{\partial V} \left(\frac{h}{mc} \left(\frac{3N}{4\pi g V} \right)^{1/3} \right) \frac{\partial f(x_F)}{\partial x_F} \\ &= - \frac{m^4 c^5}{\pi^2 \hbar^3} f(x_F) - \frac{V m^4 c^5}{\pi^2 \hbar^3} \left(\frac{h}{mc} \left(\frac{3N}{4\pi g V} \right)^{1/3} \right) \frac{\partial}{\partial V} V^{-1/3} \frac{\partial f(x_F)}{\partial x_F} \\ &= - \frac{m^4 c^5}{\pi^2 \hbar^3} f(x_F) + \frac{1}{3} \frac{V m^4 c^5}{\pi^2 \hbar^3} \left(\frac{h}{mc} \left(\frac{3N}{4\pi g V} \right)^{1/3} \right) V^{-4/3} \frac{\partial f(x_F)}{\partial x_F} \\ &= \frac{m^4 c^5}{\pi^2 \hbar^3} \left(\frac{x_F^3}{3} \sqrt{1 + x_F^2} - f(x_F) \right) . \end{aligned}$$

q.e.d.

Now, we solve the integral

$$f(x) = \frac{x^5}{4\sqrt{x^2+1}} + \frac{3x^3}{8\sqrt{x^2+1}} + \frac{x}{8\sqrt{x^2+1}} - \frac{\operatorname{asinh}(x)}{8} .$$

In the non-relativistic limit, $x_F \ll 1$, we can make the approximations

$$g(x) = \frac{x^3}{3} \sqrt{1+x^2} = \frac{x^3}{3} + \frac{x^5}{6} + O(x^6)$$

and

$$f(x_F) = \frac{x_F^3}{3} + \frac{x_F^5}{10} + O(x_F^6) .$$

In the ultra-relativistic limit, $x_F \gg 1$ or equivalently $y_F = 1/x_F \ll 1$, we can make the approximations

$$g(1/x) = \frac{1}{3x^4} + \frac{1}{6x^2} + O\left(\frac{1}{x}\right)$$

and

$$f(1/x) = \frac{1}{4x^4} + \frac{1}{4x^2} + O\left(\frac{1}{x}\right) .$$

Imposing the equilibrium condition $dE = 0$, between the gravitational and the pressure forces, and the structure of a sphere, the pressure must be

$$P = \frac{\alpha GM^2}{4\pi R^4}$$

and the Fermi momentum is

$$p_F = \frac{\hbar}{R} \left(\frac{9\pi M}{8m_p} \right)^{1/3} .$$

Proof. For the gravitational force

$$E_g = -\alpha \frac{GM^2}{R} , \quad dE_g = \alpha \frac{GM^2}{R^2} dR .$$

For the pressure force

$$E_p = -pV = -p \frac{4}{3} \pi R^3 , \quad dE_p = -4\pi p R^2 dR .$$

Imposing the equilibrium condition,

$$0 = dE = dE_g + dE_p = \alpha \frac{GM^2}{R^2} dR - 4\pi p R^2 dR ,$$

hence

$$p = \frac{\alpha GM^2}{4\pi R^4} .$$

The Fermi momentum is

$$p_F = h \left(\frac{3n}{4\pi g} \right)^{1/3} = h \left(\frac{3}{8\pi} \frac{M}{2m_p \frac{4}{3} \pi R^3} \right)^{1/3} = \frac{\hbar}{R} \left(\frac{9\pi M}{8m_p} \right)^{1/3} .$$

q.e.d.

In the ultra-relativistic limit

$$P = \frac{m^4 c^5}{12\pi \hbar^3} (x_F^4 - x_F^2) = \frac{\alpha GM^2}{4\pi R^4} .$$

Chapter 25

Bosons

25.1 Blackbody radiation

Part VII

Phase transition

Chapter 26

Classical phase transitions

Consider the phase diagram of the water. Microscopically, they all have the same hamiltonian, however, the macroscopical variables changes. There are three phases

1. solid, i.e. it has its own shape and volume,
2. liquid, i.e. it has its own volume but it has the shape of the container,
3. gas, i.e. it has the shape and volume of the container.

There are lines, called coexistence lines, along which 2 phases are in equilibrium. They are lines because, other than $T_1 = T_2$ and $p_1 = p_2$, we have a constrain

$$\mu_1(p, T) = \mu_2(p, T) .$$

This reduce to a line.

Furthermore, there are points, called coexistence points or triple point, in which 3 phases are in equilibrium. They are points because, other than $T_1 = T_2 = T_3$ and $p_1 = p_2 = p_3$, we have the constrains

$$\mu_1(p, T) = \mu_2(p, T) = \mu_3(p, T) .$$

This reduce to a point.

26.1 Symmetries

We can use symmetries of the system to study it. We can distinguish solid from fluid by the translation or rotations invariance. In fact, solid has only discrete invariance, whereas fluid has continuous invariance. However, we cannot distinguish with symmetries between gas and liquid.

26.2 Clausius-Clapeyron equation

On the coexistence line, the Clausius-Clapeyron equation is

$$\frac{dp}{dT} = \frac{s_2 - s_1}{v_1 - v_2} = \frac{\Delta q}{T \Delta v} .$$

Proof. In order to remain on the coexistence line, the constrain relation holds

$$\mu_1(p, T) = \mu_2(p, T) .$$

We differentiate it, keeping in mind that $p = p(T)$

$$\left. \frac{\partial \mu_1}{\partial T} \right|_p + \left. \frac{d\mu_1}{dp} \right|_T \frac{dp}{dT} = \left. \frac{\partial \mu_2}{\partial T} \right|_p + \left. \frac{d\mu_2}{dp} \right|_T \frac{dp}{dT} ,$$

hence

$$\frac{dp}{dT} = \frac{\left. \frac{\partial \mu_1}{\partial T} \right|_p - \left. \frac{\partial \mu_2}{\partial T} \right|_p}{\left. \frac{\partial \mu_2}{\partial p} \right|_T - \left. \frac{\partial \mu_1}{\partial p} \right|_T} .$$

At fixed number of particle, we can use the Gibbs free energy $G(p, T, N) = \mu(p, T)N$ or the Gibbs free energy per particle

$$g = \frac{G}{N} = \mu(p, T) .$$

Using the relations (??)

$$\left. \frac{\partial \mu}{\partial p} \right|_T = \left. \frac{\partial g}{\partial p} \right|_T = \frac{1}{N} \left. \frac{\partial G}{\partial p} \right|_T = \frac{V}{N} = v ,$$

$$\left. \frac{\partial \mu}{\partial T} \right|_p = \left. \frac{\partial g}{\partial T} \right|_p = \frac{1}{N} \left. \frac{\partial G}{\partial T} \right|_p = -\frac{S}{N} = -s ,$$

we obtain

$$\frac{dp}{dT} = \frac{\left. \frac{\partial \mu_1}{\partial T} \right|_p - \left. \frac{\partial \mu_2}{\partial T} \right|_p}{\left. \frac{\partial \mu_2}{\partial p} \right|_T - \left. \frac{\partial \mu_1}{\partial p} \right|_T} = \frac{s_2 - s_1}{v_2 - v_1} .$$

Furthermore, when there is a phase change, temperature remains constant whereas the thermal energy put in the system is transformed into latent heat

$$\Delta s = \frac{\Delta q}{T} .$$

Therefore

$$\frac{dp}{dT} = \frac{\Delta q}{T \Delta v} .$$

q.e.d.

There could be 2 different kind of phase transitions

1. 1st order phase transitions, i.e. those in which the 1st derivatives of thermodynamic potentials are discontinuous;
2. continuous phase transitions, i.e. those in which the higher derivatives of thermodynamic potentials are discontinuous.

In our case, the former are those in which there is a jump $v_2 \neq v_1$ and $s_2 \neq s_1$ and the latter are those in which $v_2 = v_1$ and $s_2 = s_1$.

Chapter 27

Theorems of Lee and Young

Consider a classical fluid in a volume $V \subset \mathbb{R}^3$. We treat it in the grancanonical ensemble. The grancanonical partition function is

$$\mathcal{Z}[V, T, z] = \sum_{N=0}^{\infty} z^N Z_N[T, V] .$$

The canonical partition function is

$$\begin{aligned} Z_N &= \frac{1}{N! h^{3N}} \int_{V^N} \prod_i d^3 q^i \underbrace{\int_{\mathbb{R}^{3N}} \prod_i d^3 p^i \exp(-\beta \sum_j \frac{p_j^2}{2m} + U_N(q^i))}_{\frac{1}{\lambda_T^{3N}}} \\ &= \frac{1}{N! \lambda_T^{3N}} \underbrace{\int_{V^N} \prod_i d^3 q^i \exp(-\beta U_N(q^i))}_{Q_N(T, V)} = \frac{Q_N(T, V)}{N! \lambda_T^{3N}} . \end{aligned}$$

Notice that $Q_N(T, V) > 0$.

Therefore

$$\mathcal{Z} = \sum_{N=0}^{\infty} z^N \frac{Q_N}{N! \lambda_T^{3N}} .$$

which is a power series in z . Now we promote z into a complex variables, keeping in mind that the physical states are only the ones for $z \in \mathbb{R}^+$

We make the assumption that $U_N \geq -BN$ with $B > 0$, which means that it grows no more than N order. This implies that

$$\exp(-\beta U_N) \leq \exp(\beta BN) ,$$

hence

$$Q_N = \int_{V^N} \prod_i d^3 q^i \exp(-\beta U_N(q^i)) \leq \exp(\beta BN) \int_{V^N} \prod_i d^3 q^i = \exp(\beta BN) V^N$$

and

$$Z_N = \frac{Q_N}{N! \lambda_T^{3N}} \leq \frac{V^N}{N! \lambda_T^{3N}} \exp(\beta B N) .$$

Therefore

$$|\mathcal{Z}| \leq \sum_{N=0}^{\infty} \frac{|z|^N}{N! \lambda_T^{3N}} V^N \exp(\beta B N) = \exp\left(\frac{V \exp(\beta B) |z|}{\lambda_T^3}\right) ,$$

which, given the fact that it is an exponential, has an infinite convergence radius. We have proved that it is analytical $\forall z \in \mathbb{C}$, in particular for $z \in \mathbb{R}^+$. Furthermore, \mathcal{Z} cannot become vanishing since it is convergent and it is a sum of positive terms. We introduce the grandpotential

$$\Omega = -\frac{1}{\beta} \ln \mathcal{Z} ,$$

which is well defined, since $\mathcal{Z} \neq 0$ and analytical $\forall z \in \mathbb{R}^+$.

Finally, all thermodynamic functions are analytical for $z \in \mathbb{R}^+$ and there are no phase transitions. How is it possible? We have not yet computed the thermodynamic limit.

Consider a system composed of hard sphere occupying a finite volume v . The maximum number of particles is $M = \frac{V}{v}$. \mathcal{Z} is a polynomial function in z of degree M and, by the fundamental theorem of algebra, it has M zeroes but none in \mathbb{R}^+ . If we go into the thermodynamic limit, $V \rightarrow \infty$, $M \rightarrow \infty$ and the number of zeroes increases. However, it holds that

1. $\forall V, M$, there exists an open subset of \mathbb{R}^+ which does not contain zeroes, i.e. $\mathcal{Z}(V \rightarrow \infty, T, \mu)$ has no zeroes on \mathbf{R}^+ ,
2. if zeroes accumulate towards a certain $z = z_c$, then $\mathcal{Z}(V \rightarrow \infty, T, \mu)$ has a zero in $z = z_c$.

This means that Ω is no longer analytical at $z = z_c$. Now, there is no more equilibrium and phase transitions come up from the dark.

This statements can be written down in terms of

$$\psi = \lim_{td} \frac{\ln \mathcal{Z}}{V} .$$

Hence

$$p\beta = \psi , \quad n = z \frac{\partial}{\partial z} \psi .$$

Proof. For the first

$$\Omega = -pV = -\frac{1}{\beta} \ln \mathcal{Z} ,$$

hence

$$p\beta = \frac{\ln \mathcal{Z}}{V} = \psi .$$

For the second,

$$N = z \frac{\partial}{\partial z} \ln \Omega = -\frac{z}{\beta} \frac{\partial}{\partial z} \ln \Omega ,$$

hence

$$n = \frac{N}{V} = -\frac{z}{\beta} \frac{\partial}{\partial z} \frac{\ln \Omega}{V} = -\frac{z}{\beta} \frac{\partial}{\partial z} \psi .$$

q.e.d.

Theorem 27.1

If $U_N \geq -BN$ with $B > 0$, if the boundary of the volume does not increases faster than $V^{2/3}$, in order to neglect surface terms, then ψ exists, it is continuous and monotonically increasing.

Theorem 27.2

Given an open subset of an interval of \mathbb{R}^+ such that it doesn't contain zeroes, then ψ exists and it is analytic.

Corollary 27.1

A phase transition may appear at $z = z_c$ if it is an accumulation point of zeroes. This point divides \mathbb{R}^+ into 2 regions corresponding to 2 different phases. Furthermore, ψ is continuous but it is not analytic: 1st order phase transitions or continuous phase transitions.

Chapter 28

Ising model

Consider a system composed by a discrete lattice, for example an hypercubic lattice of dimension d . For each vertex, there is a degree of freedom, characterised by the approximation of a spin that could have only two values $\sigma = \pm 1$. Two adjacent vertexes are called nearest neighborhood. Each site has therefore z nearest neighborhood, called the coordination number. For a dimension d hypercube, $z = 2d$. A possible configuration state is defined as $\{\sigma_i\}_{i \in \mathcal{L}}$. The phase space is a discrete space composed by 2^N states $\{\{\sigma_i\}_{i \in \mathcal{L}}, \sigma_i = \pm 1\}$.

The hamiltonian of the system is

$$H(\sigma_i) = H_{int} + H_{field} ,$$

where

$$H_{int} = -J \sum_{i \text{ near } j} \sigma_i \sigma_j$$

and

$$H_{field} = -B \sum_{i=1}^N \sigma_i .$$

B is an external magnetic field and J is the interaction constant. Notice that in order to have a phase transitions, we have to allow interactions. For $J > 0$, the minimum energy configuration is the one in which all the spins are aligned $\sigma_i = \sigma_j$, $\forall i, j$. This model is called ferromagnetic model. For $J < 0$, the minimum energy configuration is the one in which all the spins are antialigned $\sigma_i = -\sigma_j$, $\forall i, j$. This model is called antiferromagnetic model. For $B > 0$, the minimum energy configuration is the one in which all the spins are aligned upwards $\sigma_i = +1$. For $B < 0$, the minimum energy configuration is the one in which all the spins are aligned downwards $\sigma_i = -1$.

In the canonical ensemble, the partition function is

$$Z_N = \sum_{\sigma_i = \pm 1} \exp(-\beta H(\sigma_i)) ,$$

where the sum is made over all the 2^N states. The Helmholtz free energy is

$$F = E - TS = -\frac{1}{\beta} \ln Z_N .$$

where $E = \langle H \rangle_c$. The thermodynamic equilibrium correspond to the configuration of minimum free energy.

Suppose the external magnetic field is shut down. The ground state is the one with minimum energy and the entropy is small, because there are only 2 states possible with all aligned spins. The excited state is the one with minimum energy and the entropy is big, because all spins point in all direction. Recall that entropy is $S = k_B \ln \Gamma(E)$. The minimal configuration of free energy is therefore at low T with minimum E , i.e. all aligned, and at high T with large S , i.e. random alignment.

To estimate the alignment, we introduce the magnetisation

$$M = \langle \sum_{i=1}^N \sigma_i \rangle_c = \sum_{i=1}^N \langle \sigma_i \rangle_c ,$$

where we have used the translation invariance.

Computing the phase diagram, we find that along the T -axis at $B = 0$, $m \neq 0$ for $T < T_c$ and $m = 0$ for $T > T_c$, where T_c is the critical temperature. The former is called the ferromagnetic phase and the latter is called the paramagnetic phase. We can use the magnetisation as an order parameter, since when it is zero there is disorder and when it is different from zero, there is order. In the neighborhood of T_c , we have the behaviour for $T < T_c$

$$M \sim (T - T_c)^\beta ,$$

where β is a parameter. It characterise the phase transition, since it tells which speed $M \rightarrow 0$ when approaching $T \rightarrow T_c$.

28.1 Correlation

28.2 Symmetry breaking

Consider the Ising model hamiltonian. The first term is invariant under

$$\sigma_i \rightarrow -\sigma_i , \quad \sigma_i \rightarrow \sigma_i .$$

Therefore, it is invariant under the global symmetry group \mathbb{Z}_2 . If it were the only term, i.e. with $B = 0$, the whole hamiltonian would be invariant under this group. However, the second term breaks explicitly the symmetry. Moreover, noticing that under this symmetry $m \rightarrow -m$, the only possible value of m would be zero. Hence,

for $T > T_c$ there is indeed $m = 0$. But for $T < T_c$, the equilibrium state is no longer invariant under this symmetry. The hamiltonian remains the same but states are not invariant. There is a spontaneous symmetry breaking, spontaneous because H is stil invariant under \mathbb{Z}_2 .

Chapter 29

Mean-field treatment

Consider the Ising model. In general, it is difficult to compute the canonical partition function,

$$Z_N = \sum_{\{\sigma_i = \pm 1\}} \exp(-\beta H) \neq (Z_1)^N ,$$

since it is interacting. However, we can make an useful approximation

$$\sigma_i \sigma_j = ((\sigma_i - m) + m)((\sigma_j - m) + m) = m^2 + m(\sigma_i - m) + m(\sigma_j - m) + (\sigma_i - m)(\sigma_j - m) ,$$

in which we keep only the first constant term and the second linear, but we neglect the last quadratic fluctuation term. Therefore

$$\sigma_i \sigma_j = m^2 + m(\sigma_i - m) + m(\sigma_j - m) = m^2 + m\sigma_i - m^2 + m\sigma_j - m^2 = -m^2 + m(\sigma_i + \sigma_j) .$$

The hamiltonian in the mean-field approximation becomes

$$H_{mf} = -J \sum_{\text{inn}j} (-m^2 + m(\sigma_i + \sigma_j)) - B \sum_i \sigma_i = m^2 J \sum_{\text{inn}j} 1 - Jm \sum_{\text{inn}j} (\sigma_i + \sigma_j) - B \sum_i \sigma_i .$$

The number of links, given the coordination number z which tells how many neighboring sites, is $NZ/2$. Hence

$$H_{mf} = \frac{m^2 z N J}{2} - Jmz \sum_i \sigma_i - B \sum_i \sigma_i = \frac{m^2 z N J}{2} - (Jmz + B) \sum_i \sigma_i .$$

The physical interpretation of the mean-field treatment is that we do not have to compute every links with respect to each others but only with respect to the mean field m . It is valid only if fluctuations are smaller than the mean-field. The partition

function is

$$\begin{aligned}
Z_N^{mf} &= \sum_{\{\sigma_i=\pm 1\}} \exp(-\beta H_{mf}) \\
&= \exp(-\beta \frac{Jznm^2}{2}) \sum_{\{\sigma_i=\pm 1\}} \exp(\beta(B + Jmz) \sum_i \sigma_i) \\
&= \exp(-\beta \frac{Jznm^2}{2}) \left(\sum_{\{\sigma_i=\pm 1\}} \exp(\beta(B + Jmz) \sigma_i) \right)^N \\
&= \exp(-\beta \frac{Jznm^2}{2}) (\exp(\beta(B + Jmz)) + \exp(-\beta(B + Jmz)))^N \\
&= \exp(-\beta \frac{Jznm^2}{2}) (2 \cosh(\beta(B + Jmz))) .
\end{aligned}$$

The Helmholtz free energy is

$$\begin{aligned}
F &= -\frac{1}{\beta} \ln Z_N^{mf} \\
&= -\frac{1}{\beta} \left(-\beta \frac{JzNm^2}{2} \right) N \ln(2 \cosh(\beta(B + Jmz))) \\
&= \frac{JzNm^2}{2} N \ln(2 \cosh(\beta(B + Jmz))) .
\end{aligned}$$

The magnetisation is

$$\begin{aligned}
m &= \frac{1}{N} \langle \sum_i \sigma_i \rangle_c \\
&= \frac{1}{N} \sum_{\{\sigma_i=\pm 1\}} \sum_i \sigma_i \exp(-\beta H) \\
&= -\frac{1}{\beta N} \sum_{\{\sigma_i=\pm 1\}} \frac{1}{Z_N} \frac{\partial}{\partial \beta} \exp(-\beta H) \\
&= -\frac{1}{\beta N} \frac{\partial \ln Z_N}{\partial \beta} .
\end{aligned}$$

Hence

$$m = \tanh(\beta(B + Jmz)) .$$

We have a self-consistent equation for m to solve. The condition for solution is $B > 0$ then $m > 0$ and $B < 0$ then $m < 0$. Particular attention we can study for $B = 0$ then

$$m = \tanh \frac{Jmz}{k_B T} = \tanh \frac{T_c m}{T} ,$$

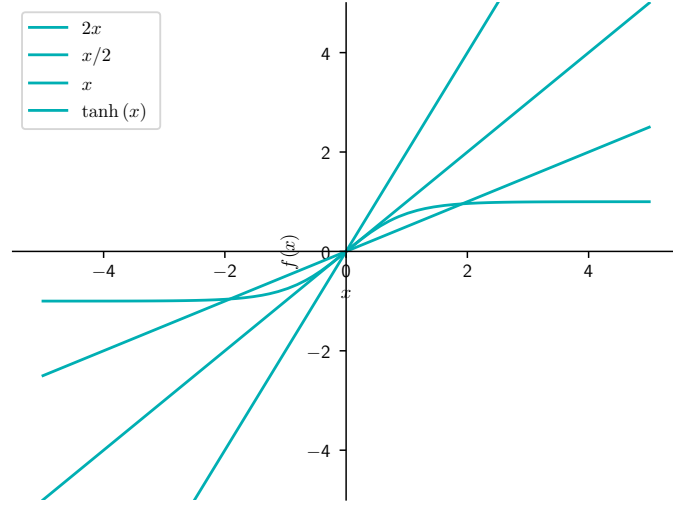


Figure 29.1: A plot of the graphical solution of $T\tilde{m}/T_c = \tanh \tilde{m}$.

where $T_c = Jz/k_B$ is the critical temperature. Notice that it depends on z . Calling $\tilde{m} = T_c m/T$, we have

$$\frac{T\tilde{m}}{T_c} = \tanh \tilde{m} .$$

The solutions are points that intersect a straight line and an hyperbolic tangent. If $T > T_c$, there is only one solution $m = 0$. If $T < T_c$, there are two solutions $\pm m_0$, one positive and one negative. See Figure 29.1. We have proven that m is indeed an order parameter.

Part VIII

Appendix

Chapter A

Volume of an N-dimensional sphere

In this appendix chapter, we will prove that the volume of an N -dimensional sphere of radius R is

$$V_n(R) = \frac{\pi^{n/2} R^n}{\Gamma(n/2 + 1)} . \quad (\text{A.1})$$

Proof. Consider the rotationally invariant function f

$$f(x_1, \dots, x_n) = \exp\left(-\frac{1}{2} \sum_{i=1}^n x_i^2\right) = \prod_{i=1}^n \exp\left(-\frac{1}{2} x_i^2\right) = .$$

Using the Gaussian integral, this function can be integrated over all \mathbb{R}^n , with volume element $dV = dx_1 \dots dx_n$, and it gives

$$\begin{aligned} \int_{\mathbb{R}^n} dV f &= \int_{\mathbb{R}^n} \prod_{i=1}^n dx_i f = \int_{\mathbb{R}^n} \prod_{i=1}^n dx_i \exp\left(-\frac{1}{2} \sum_{i=1}^n x_i^2\right) \\ &= \prod_{i=1}^n \underbrace{\left(\int_{\mathbb{R}} dx_i \exp\left(-\frac{1}{2} x_i^2\right) \right)}_{(2\pi)^{1/2}} = \prod_{i=1}^n (2\pi)^{1/2} = (2\pi)^{n/2} . \end{aligned}$$

Exploiting the rotational invariant property, we can decomposed the volume element into a surface element dA , which integrated gives an $(n-1)$ -dimensional sphere $S^{n-1}(r)$ of radius r , multiplied by a length element dr , i.e.

$$\int_{\mathbb{R}^n} dV f = \int_0^\infty dr \int_{S^{n-1}(r)} dA f .$$

Since the area is proportional to the radius, e.g. for $n=3$ the area is $A \propto r^2$, the radius-dependence of the area is given by $A_{n-1}(r) = r^{n-1} A_{n-1}(1)$. Therefore, putting it inside the integral, we obtain

$$A_{n-1}(1) \int_0^\infty dr r^{n-1} \exp\left(-\frac{1}{2} r^2\right) .$$

Now, we make a change of variables into

$$t = \frac{r^2}{2} , \quad r = (2t)^{1/2} , \quad dr = 2^{-1/2} t^{-1/2} dt$$

to have the integral of the gamma function

$$\begin{aligned} \int_0^\infty dr \, r^{n-1} \exp\left(-\frac{1}{2}r^2\right) &= 2^{(n-1)/2} 2^{-1/2} \int_0^\infty dt \, t^{(n-1)/2} t^{-1/2} \exp(-t) \\ &= 2^{n/2-1} \underbrace{\int_0^\infty dt \, t^{n/2-1} \exp(-t)}_{\Gamma(n/2)} = 2^{n/2-1} \Gamma(n/2) . \end{aligned}$$

Now, we combine the two results together to obtain the surface

$$(2\pi)^{n/2} = A_{n-1}(1) 2^{n/2-1} \Gamma(n/2) ,$$

hence

$$A_{n-1}(1) = \frac{2\pi^{n/2}}{\Gamma(n/2)} .$$

Finally, in order to find the volume we need to integrate from 0 to R

$$\begin{aligned} V_n(R) &= \int_0^R dr \, A_{n-1}(r) = \int_0^R dr \, A_{n-1}(1) r^{n-1} = \frac{2\pi^{n/2}}{\Gamma(n/2)} \int_0^R dr \, r^{n-1} \\ &= \frac{2\pi^{n/2}}{\Gamma(n/2)} \frac{r^n}{n} \Big|_0^R = \frac{2\pi^{n/2}}{n\Gamma(n/2)} R^n = \frac{\pi^{n/2} R^n}{\Gamma(n/2 + 1)} . \end{aligned}$$

q.e.d.

Chapter B

Stirling approximation

In this appendix chapter, we will prove the Stirling approximation

$$\ln n! \simeq n \ln n - n . \quad (\text{B.1})$$

Proof. The factorial can be expressed in integral form via the gamma function

$$\Gamma(n+1) = n! = \int_0^\infty dt \, t^n \exp(-t) .$$

Now, we make a change of variables into

$$t = nx , \quad x = \frac{t}{n} , \quad dx = \frac{dt}{n} ,$$

to have

$$\begin{aligned} \int_0^\infty dt \, t^n \exp(-t) &= n \int_0^\infty dx \, \exp(\ln t^n) \exp(-t) \\ &= \int_0^\infty dt \, \exp(n \ln t - t) \\ &= n \int_0^\infty dx \, \exp(n \ln(nx) - nx) \\ &= n \int_0^\infty dx \, \exp(n \ln x + n \ln n - nx) \\ &= n \exp(n \ln n) \int_0^\infty dx \, \exp(n(\ln x - x)) . \end{aligned}$$

In the limit for which n is large, we can use the Laplace approximation method

$$\int_a^b dx \, \exp(nf(x)) \simeq \exp(nf(x_0)) \sqrt{\frac{2\pi}{n|f''(x_0)|}} .$$

where $x_0 \in [a, b]$ is a stationary point of $f(x)$. A simple sketch of the proof is given by means of the Taylor expansion around x_0

$$f(x) \simeq f(x_0) - \frac{1}{2}|f''(x_0)|(x - x_0)^2 ,$$

hence, integrating the Gaussian integral,

$$\begin{aligned} \int_a^b dx \exp(nf(x)) &\simeq \exp(nf(x_0)) \int_a^b dx \exp(-\frac{n}{2}|f''(x_0)|(x - x_0)^2) \\ &= \sqrt{\frac{2\pi}{n|f''(x_0)|}} . \end{aligned}$$

In our case, $a = 0$, $b = \infty$ and $f(x) = \ln x - x$, which has a maximum in $x_0 = 1$ and second derivatives equals to $|f''(x)| = 1/x^2$. Therefore

$$\int_0^\infty dx \exp(n(\ln x - x)) \simeq \exp(n(\ln x_0 - x_0)) \sqrt{\frac{2\pi x_0^2}{n}} \Big|_{x_0=1} = \exp(-n) \sqrt{\frac{2\pi}{n}} .$$

Now, we combine the two results together

$$n! \simeq n \exp(n \ln n) \exp(-n) \sqrt{\frac{2\pi}{n}} = \exp(n \ln n - n) \sqrt{2\pi n} = n^n \exp(-n) \sqrt{\frac{2\pi}{n}} ,$$

which can be rewritten in terms of logarithms rather than exponentials

$$\ln n! \simeq \ln(n^n \exp(-n) \sqrt{\frac{2\pi}{n}}) = n \ln n - n + O(\ln n) .$$

q.e.d.

Chapter C

Gaussian integral

In this appendix chapter, we will prove that the Gaussian integral is

$$\int_{-\infty}^{\infty} dx \exp(-x^2) = \sqrt{\pi} . \quad (\text{C.1})$$

Proof. We start from the square Gaussian integral, which it is the square same integral for the mute properties of the integration variables

$$\begin{aligned} \left(\int_{-\infty}^{\infty} dx \exp(-x^2) \right)^2 &= \int_{-\infty}^{\infty} dx \exp(-x^2) \int_{-\infty}^{\infty} dy \exp(-y^2) \\ &= \int_{-\infty}^{\infty} dx \int_{-\infty}^{\infty} dy \exp(-(x^2 + y^2)) . \end{aligned}$$

Now, we make a change of variables and we use polar coordinates (r, θ)

$$r^2 = x^2 + y^2 , \quad \theta = \arctan \frac{y}{x} , \quad dx \, dy = r \, dr \, d\theta , \quad (r, \theta) \in [0, \infty) \times [0, 2\pi] ,$$

to obtain

$$\begin{aligned} \int_{-\infty}^{\infty} dx \int_{-\infty}^{\infty} dy \exp(-(x^2 + y^2)) &= \underbrace{\int_0^{2\pi} d\theta}_{2\pi} \int_0^{\infty} dr \, r \exp(-r^2) \\ &= 2\pi \int_0^{\infty} dr \, r \exp(-r^2) \\ &= \pi \int_0^{\infty} dr \, 2r \exp(-r^2) \\ &= \pi \exp(-r^2) \Big|_0^{\infty} = \pi . \end{aligned}$$

Now, we combine the two results together

$$\left(\int_{-\infty}^{\infty} dx \exp(-x^2) \right)^2 = \pi ,$$

hence

$$\int_{-\infty}^{\infty} dx \exp(-x^2) = \sqrt{\pi} .$$

q.e.d.

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