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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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# Part I

## Thermodynamics

# Chapter 1

## Laws of Thermodynamics

### 1.1 The laws of Thermodynamics

In this chapter, we will recall some notions of Thermodynamics.

In Thermodynamics, a state is defined by a set of macroscopic 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 the volume and the number of particles change: extensive variables do scale with it whereas intensive ones do not. See Table 1.1. An equation of state is a functional relation among them.

Extensive	Intensive
Energy $E$	-
Entropy $S$	Temperature $T$
Volume $V$	Pressure $p$
Number of particles $N$	Chemical potential $\mu$
Polarization $\mathbf{P}$	Electric field $\mathbf{E}$
Magnetization $\mathbf{M}$	Magnetic field $\mathbf{B}$

Table 1.1: Extensive and intensive thermodynamical variables.

Thermodynamics is described by four laws.

#### Law 1.1 (0th)

*Two systems in thermal contact have the same empirical temperature  $T$  at equilibrium*

$$T_1 = T_2$$

**Law 1.2 (1st)**

The (generalised) principle of conservation of energy states that

$$dE = \delta Q - \delta L + \mu dN \quad (1.1)$$

where  $E$  is the internal energy,  $Q$  is the heat,  $L$  is the work,  $\mu$  is the chemical potential (the necessary energy to add or remove a particle) and  $N$  is the number of particles.

Recall that  $E$  is a exact differential, i.e  $\oint dE = 0$ , whereas heat and work are not, i.e  $\oint \delta Q \neq 0$  and  $\oint \delta L \neq 0$ .

**Law 1.3 (2nd)**

A system naturally evolves in order to maximize its entropy  $S$ . For reversible processes

$$dS = \frac{\delta Q}{T} \quad (1.2)$$

whereas for irreversible processes

$$dS \geq \frac{\delta Q}{T}$$

**Law 1.4 (3rd)**

For any reversible isothermal process

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

## 1.2 Thermodynamical potentials

For reversible processes, using (1.2) and  $\delta L = pdV$ , (1.1) can be expressed as

$$dE = TdS - pdV + \mu dN \quad (1.3)$$

Notice that the left variables are intensive and the right variables (those with the differential) are extensive.

$E$  is a function of  $S$ ,  $V$ ,  $N$ , hence it must be extensive and a homogeneous function of degree one, satisfying the property

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

where  $\lambda > 0$  is the scale factor. It can be proved that the only function is

$$E(S, V, N) = TS - pV + \mu N$$

Similar expression can be found for other thermodynamical quantities, simply exchanging the role of conjugate functions. See Table 1.2.

Potentials	Differential
Internal energy $E(S, V, N) = TS - pV + \mu N$	$dE = TdS - pdV + \mu dN$
Helmoltz free energy $F(T, V, N) = E - TS = -pV + \mu N$	$dF = -SdT - pdV + \mu dN$
Entalpy $H(S, p, N) = E + pV = St + \mu N$	$dH = TdS + Vdp + \mu dN$
Gibbs free energy $G(T, p, N) = E - TS + pV = \mu N$	$dG = -SdT + Vdp + \mu dN$
Granpotential $\Omega(T, V, \mu) = E - TS - \mu N = -pV$	$d\Omega = -SdT - pdV - Nd\mu$

Table 1.2: Thermodynamical potentials.

*Proof.* Maybe in the future.

q.e.d.

Fixing three of the thermodynamical variables to be constant, a system evolves in order to minimise the corresponding thermodynamical potential until it reaches its minimum, i.e the equilibrium state. Mathematically, it means that the first derivative must be vanishing and the hessian must be positive defined. See Table 1.3.

Inequality	Constant quantities
$dE \leq 0$	$S, V, N$
$dF \leq 0$	$T, V, N$
$dH \leq 0$	$S, p, N$
$dG \leq 0$	$T, p, N$
$d\Omega \leq 0$	$T, V, \mu$

Table 1.3: Thermodynamical variation principles.

*Proof.* Maybe in the future.

q.e.d.



## Part II

### Classical statistical mechanics

# Chapter 2

## 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 the 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 2.1 (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

## 2.1 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

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

## 2.2 Liouville's theorem

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

### Theorem 2.2 (Liouville)

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

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

*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

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

and the standard deviation is

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

## 2.3 Time-independent Hamiltonian

Consider a time-independent hamiltonian. Since the energy is constant for the theorem 2.1.

$$\int_{\mathcal{M}^N} \rho = 1 \tag{2.1}$$

$$\frac{\partial S}{\partial E} = \frac{1}{T} \tag{2.2}$$

$$\frac{\partial F}{\partial T} = -S \tag{2.3}$$

$$\Omega = -pV = E - TS - \mu N \tag{2.4}$$

# Chapter 3

## 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 (2.1)

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

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

$$\langle f(q^i, p_i) \rangle_{mc} = \int_{\mathcal{M}} d\Omega \rho_{mc} f = \int_{\mathcal{M}} d\Omega \frac{1}{\omega(E)}\delta(\mathcal{H} - E)f = \frac{1}{\omega(E)} \int_{S_E} dS_E f = \langle f \rangle_E \quad (3.1)$$

### 3.1 Thermodynamics potentials

The microcanonical entropy  $S_{mc}$  is defined by

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

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 sistems 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  $E_1 \in [0, E]$

$$\begin{aligned} \int_0^E dE_1 \omega_1(E_1) \omega_2(E_2 = E - E_1) &\leq \omega_1(E_1^*) \omega_2(E_2^* = E - E_1^*) \int_0^E dE_1 \\ &= \omega_1(E_1^*) \omega_2(E_2^* = E - E_1^*) E \end{aligned} \quad (3.2)$$

On the other hand, it is always possible to find a value for  $\Delta E$  in order to have

$$\Delta E \omega_1(E_1^*) \omega_2(E_2^*) \leq \omega(E) \quad (3.3)$$

Putting together (3.2) and (3.3)

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

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

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

Since the logarithm is a monotomic function

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

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

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

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

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  $E = const$ ,  $0 = \delta E = \delta E_1 + \delta E_2$ ,  $\delta E_2 = -\delta E_1$  and

$$\begin{aligned} 0 &= \frac{\partial\Gamma_1}{\partial E_1}\Big|_{E_1^*} \delta E_1 \Gamma_2(E_2^*) - \Gamma_1(E_1^*) \frac{\partial\Gamma_2}{\partial E_2}\Big|_{E_2^*} \delta E_1 \\ 0 &= \frac{\partial\Gamma_1}{\partial E_1}\Big|_{E_1^*} \Gamma_2(E_2^*) - \Gamma_1(E_1^*) \frac{\partial\Gamma_2}{\partial E_2}\Big|_{E_2^*} \\ \frac{\partial\Gamma_1}{\partial E_1}\Big|_{E_1^*} \Gamma_2(E_2^*) &= \Gamma_1(E_1^*) \frac{\partial\Gamma_2}{\partial E_2}\Big|_{E_2^*} \\ \frac{1}{\Gamma_1(E_1^*)} \frac{\partial\Gamma_1}{\partial E_1}\Big|_{E_1^*} &= \frac{1}{\Gamma_2(E_2^*)} \frac{\partial\Gamma_2}{\partial E_2}\Big|_{E_2^*} \end{aligned}$$

$$\left. \frac{\partial \log \Gamma_1}{\partial E_1} \right|_{E_1^*} = \left. \frac{\partial \log \Gamma_2}{\partial E_2} \right|_{E_2^*}$$

Using the thermodynamical relation (2.2)

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

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

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

*Proof.* Using (3.1),

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

q.e.d.



# Chapter 4

## 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

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

where  $\beta$  is

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

and  $Z_N$  is the partition function

$$Z_N[V, T] = \int_{\mathcal{M}^N} d\Omega \exp(-\beta \mathcal{H}(q^i, p_i)) \quad (4.1)$$

which depends on the temperature through  $\beta$  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  $\delta E_1$  to  $E_1$  to preserve equilibrium, the entropy transforms, using (2.2)

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

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

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

Putting together, dropping the indices

$$\rho_c = \frac{\omega_2(E_2)}{\omega(E)} = \frac{1}{\omega(E)} \exp\left(\frac{S_{mc}(E)}{k_B}\right) \exp\left(-\frac{E_1}{k_B T}\right) = C \exp\left(-\frac{E_1}{k_B T}\right) \quad (4.2)$$

where  $C$  is a normalisation constant, which can be evaluated by (2.1)

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

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  $\zeta_N$  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

$$Z_N = Z_{N_1} Z_{N_2} \quad (4.3)$$

*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

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

## 4.1 Thermodynamics variable

The canonical Helmotz free energy  $F$  is defined by

$$Z[V, T] = \exp(-\beta F[N, V, T]) \quad (4.4)$$

or, equivalently,

$$F[V, N, T] = -\frac{1}{\beta} \ln Z_N \quad (4.5)$$

Furthermore, the canonical internal energy is

$$E = \langle \mathcal{H} \rangle_c = \int d\Omega \frac{\exp(-\beta(\mathcal{H}))}{Z_N} \mathcal{H} \quad (4.6)$$

*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  $\beta$

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

Hence, using (2.3)

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

showing that is indeed the Helmotz free energy.

q.e.d.

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

$$S_c = \frac{E - F}{T} \quad (4.7)$$

The internal energy can also be written as

$$E = -\frac{\partial}{\partial \beta} \ln Z_N$$

*Proof.* Using (4.6),

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

q.e.d.

The universal Boltzmann's formula is still valid

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

*Proof.* Using (4.6) and (4.5)

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

q.e.d.

## 4.2 Equipartition theorem

### Theorem 4.1 (Generalised equipartition theorem)

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

$$\int \prod_{j \neq 1} d\xi_j [\xi_1 \exp(-\beta \mathcal{H})]_a^b = 0 \quad (4.8)$$

Then

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

*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 (4.8) are hamiltonians which depend on the square of momentum or confining potentials which go to infinity on the extremes  $a$  and  $b$ .

**Corollary 4.1 (Equipartition theorem)**

If  $\xi_1$  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 5

## 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.



## 5.1 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  $\mu$  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  $\Omega$  is the granpotential.

*Proof.* Using (4.4) 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 (2.4), 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

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

*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

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

*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

$$\Omega = -\frac{1}{\beta} \ln \mathcal{Z} \quad (5.1)$$

The universal Boltzmann's formula is still valid

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

*Proof.* Using (5.1),

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

q.e.d.

# Chapter 6

## 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 grancanonical and continuous energy levels. Therefore, the probability density distribution is (??)

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

where the canonical partition function (??) 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

$$H = - \sum_{i=1}^N p_i \ln p_i + \alpha \left( \sum_{i=1}^N p_i - 1 \right) + \beta \left( \sum_{i=1}^N p_i f(x_i) - \langle f \rangle \right) \quad (6.1)$$

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 (6.1), 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 7

## Microcanonical ensemble

### 7.1 Non-relativistic ideal gas

The hamiltonian of a non-relativistic ideal (non-interacting) gas is

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

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

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

The density state  $\omega(E)$  is

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

*Proof.* Maybe in the future.

q.e.d.

Now, consider a 3-dimensional ideal gas. The entropy is

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

or, in the thermodynamic limit,

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



*Proof.* Maybe in the future.

q.e.d.

The internal energy is

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

*Proof.* Maybe in the future.

q.e.d.

The pressure is

$$P = \frac{N k_B T}{V} .$$

*Proof.* Maybe in the future.

q.e.d.

## 7.2 Gas of harmonic oscillators

The hamiltonian of a non-relativistic ideal (non-interacting) gas 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

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

The density state  $\omega(E)$  is

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

*Proof.* Maybe in the future.

q.e.d.

# Chapter 8

## Canonical ensemble

### 8.1 Non-relativistic ideal gas

The canonical partition function is

$$Z = \frac{V^N}{N! \lambda_T^{3N}} = \frac{V^N}{N!} \sqrt{\frac{2\pi m k_B T}{h^2}} .$$

*Proof.* Maybe in the future.

q.e.d.

The internal energy is

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

*Proof.* Maybe in the future.

q.e.d.

The Helmholtz free energy is

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

*Proof.* Maybe in the future.

q.e.d.

The entropy is

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

*Proof.* Maybe in the future.

q.e.d.

The pression is

$$p = \frac{N}{V k_B} .$$

*Proof.* Maybe in the future.

q.e.d.

The chemical potential is

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

*Proof.* Maybe in the future.

q.e.d.

The specific heats per particle are

$$c_V = \frac{3}{2} k_B , \quad c_p = \frac{5}{2} k_B .$$

*Proof.* Maybe in the future.

q.e.d.

## 8.2 Gas of harmonic oscillators

The canonical partition function is

$$Z = \left( \frac{k_B T}{\hbar \omega} \right)^N .$$

*Proof.* Maybe in the future.

q.e.d.

The internal energy is

$$E = N k_B T .$$

*Proof.* Maybe in the future.

q.e.d.

The Helmolzt free energy is

$$F = N k_B T \ln \frac{\hbar \omega}{k_B T} .$$

*Proof.* Maybe in the future.

q.e.d.

The entropy is

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

*Proof.* Maybe in the future.

q.e.d.

The pression is

$$p = 0 .$$

*Proof.* Maybe in the future.

q.e.d.

The chemical potential is

$$\mu = k_B T \ln \frac{\hbar\omega}{k_B T} .$$

*Proof.* Maybe in the future.

q.e.d.

The specific heats per particle are

$$c_V = k_B , \quad c_p = k_B .$$

*Proof.* Maybe in the future.

q.e.d.

### 8.3 Ultra-relativistic ideal gas

The hamiltonian of an ultra-relativistic ideal (non-interacting) gas is

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

The canonical partition function is

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

*Proof.* Maybe in the future.

q.e.d.

The internal energy is

$$E = 3Nk_B T .$$

*Proof.* Maybe in the future.

q.e.d.

The Helmholtz free energy is

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

*Proof.* Maybe in the future.

q.e.d.

The entropy is

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

*Proof.* Maybe in the future.

q.e.d.

The pression is

$$p = \frac{N}{V\beta} .$$

*Proof.* Maybe in the future.

q.e.d.

The chemical potential is

$$\mu = -k_B T \ln \frac{8\pi(k_B T)^3}{n(hc)^3} .$$

*Proof.* Maybe in the future.

q.e.d.

The specific heats per particle are

$$c_V = 3k_B , \quad c_p = 4k_B .$$

*Proof.* Maybe in the future.

q.e.d.

## 8.4 Maxwell-Boltzmann velocity distribution

## 8.5 Magnetic solid

The hamiltonian of a magnetic solid is

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

The canonical partition function is

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

*Proof.* Maybe in the future.

q.e.d.

The internal energy is

$$E = -N\mu H \left( \coth(\beta\mu H) - \frac{1}{\beta\mu H} \right) .$$

*Proof.* Maybe in the future.

q.e.d.

The Helmholtz free energy is

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

*Proof.* Maybe in the future.

q.e.d.

The entropy is

$$S = .$$

*Proof.* Maybe in the future.

q.e.d.

# Chapter 9

## Grancanonical ensemble

### 9.1 Non-relativistic ideal gas

### 9.2 Gas of harmonic oscillators

### 9.3 Solid-vapor equilibrium phase

# Chapter 10

## Entropy

### 10.1 Maxwell-Boltzmann distribution

### 10.2 Fermi-Dirac distribution

### 10.3 Bose-Einstein distribution



## Part IV

Quantum statistical mechanics

# Chapter 11

## Quantum mechanics

## Part V

### Application of quantum statistical mechanics

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