

Theory & Notes

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1. QUANTUM PHASE TRANSITIONS

Quantum Phase Transitions

Consider a Hamiltonian $\mathbf{H}(g)$, whose degrees of freedom reside on the sites of a lattice $\Lambda \subseteq \mathbb{Z}^d$, and which varies a function of a dimensionless coupling g . Then,

- for the case of a finite lattice, the ground state energy will be, in general, a smooth, analytic function of g , $\varepsilon_\Omega(g) \in C^\omega(\Lambda)$. The main exception comes from the case when g couples only to a conserved quantity, ie.

$$\mathbf{H}(g) = \mathbf{H}_0 + g\mathbf{H}_1 \text{ where } [\mathbf{H}_0, \mathbf{H}_1] = 0.$$

This entails that these Hamiltonians can be simultaneously diagonalized and so, the eigenfunctions are independent of g even though the eigenvalues vary with g ¹. Then, there may be a level crossing where an excited level becomes the ground state at $g = g_c$, creating a point of non-analyticity of the ground state energy as a function of g .

- For an infinite lattice, the possibilities are richer. An avoided level-crossing between the ground and an excited state in a finite lattice could become progressively sharper as the lattice size increases, leading to a non-analyticity at $g = g_c$ in the infinite-lattice limit. Any point of non-analyticity in the ground state energy of the infinite lattice system as a quantum phase transition. The non-analyticity could be either the limiting case of an avoided level crossing, or an actual level crossing, the first situation being the most common. The phase transition is usually accompanied by a qualitative change in the nature of the correlations in the ground state.

In particular, second order quantum phase transitions are transitions at which the characteristic energy scale of fluctuations above the ground state vanishes as $g \rightarrow g_c$. Let Δ be the scale characterizing some significant spectral density of fluctuations at zero temperature for $g \neq g_c$. Thus, Δ could be the energy of the lowest excitation above the ground state. If this is non-zero (ie. there is an energy gap Δ), or if there are excitations at arbitrarily low energies in the infinite lattice limit (ie. the energy spectrum is gapless), Δ is the scale at which there is a qualitative change in the nature of the frequency spectrum from its lowest frequency to its higher frequency behaviour. In most cases, it is the case that

$$\Delta_\pm \underset{g \rightarrow g_c^\pm}{\sim} J|g - g_c|^{zv}, \quad \text{where } J \text{ is the energy scale of a characteristic microscopic coupling,} \\ \text{and where } zv \text{ is the critical phase exponent, which is universal.}$$

A critical phase exponent is independent of most of the microscopic details of the $\mathbf{H}(g)$ -Hamiltonian. The previous behaviour holds both for $g > g_c$ and $g < g_c$ with the same value of the exponent zv , but with different non-universal constant of proportionality. In addition to a vanishing energy scale, second order quantum phase

¹Two complex-valued matrices \mathbf{A}, \mathbf{B} are simultaneously diagonalizable if and only if there exists a single matrix $P \in \text{GL}(\mathbb{C})$ which diagonalizes them both at the same time. Then, if $[\mathbf{H}_0, g\mathbf{H}_1] = 0$, then the P -matrix must be independent of g . Since P is constructed from eigenvectors, this implies that the eigenvectors themselves do not depend on g , but the eigenvalues of the individual matrices might.

transitions invariably have a diverging characteristic length scale ζ , which could be the length scale determining the exponential decay of equal time correlations in the ground state or the length scale at which some characteristic crossover occurs to the correlation at the longest distances. This length diverges as

$$\zeta^{-1} \sim \Lambda |g - g_c|^v, \quad \text{where } v \text{ is a critical exponent,} \\ \text{and where } \Lambda \text{ is an inverse length scale ("a momentum cutoff") of order the inverse lattice spacing.}$$

The ratio of the exponents for energy and length is z , the dynamic critical exponent: the characteristic energy scale vanishes as the z -th power of the characteristic inverse length scale

$$\Delta \sim \zeta^{-z}.$$

It is important to notice that the previous discussion concerns only the ground state of the system. Thus, quantum phase transitions occur only at zero temperature. Since all experiments are at necessarily non-zero temperature, a central task of the theory of quantum phase transitions is to describe the consequences of this $T = 0$ -singularity on physical properties at $T > 0$. It turns out that working outward from the critical point $g = g_c$, and $T = 0$, is a powerful way of understanding and describing the thermodynamic properties of numerous systems over a broad range of values $|g - g_c|$ and T ².

1.0.0.1. Quantum versus classical phase transitions. There are two important possibilities for the $T > 0$ phase diagram of a system near a quantum critical point,

- either the thermodynamic singularity is present only at zero temperature and all $T > 0$ properties are analytic as a function of g near $g = g_c$,
- or there is a curve of $T > 0$ second order phase transitions (this is a line at which the thermodynamic free energy is not analytic) which terminates at the $T = 0$ -quantum critical point, $g = g_c$.

In particular, for the second kind of transitions, in the vicinity of such curve, the typical frequency at which the important long distance degrees of freedom fluctuate, ω_{typ} , satisfies

$$\hbar \omega_{\text{typ}} \ll k_B T.$$

Under these conditions, the classical description can be applied to the important degrees of freedom. Consequently, the ultimate critical singularity along the line of $T > 0$ is ultimately described by the theory of second order classical phase transitions. In said classical systems, the phase transitions are driven only by thermal fluctuations, as classical systems usually freeze into a fluctuationless ground state at $T = 0$. In contrast, quantum systems have fluctuations driven by the Heisenberg uncertainty principle even in the ground state, and these can drive interesting phase transitions at zero temperature. The $T > 0$ region, in the vicinity of a quantum critical point, therefore offers a fascinating interplay of both quantum and thermal fluctuations.

²Indeed, in many cases it is not even necessary that the system of interest ever have its microscopic coupling reach a value such that $g = g_c$, since it can be very useful to develop a theory based on a physically inaccessible coupling and to develop a description based in the deviation $|g - g_c|$.

2. SPIN CHAIN MODELS

Spin Chain Models Spin Chain Models

2.1. Classical Spin Systems. Classical spin systems are idealized versions of magnets. Although many magnetic phenomena in materials are inherently quantum mechanical, many properties are well described at least qualitatively by classical spin systems.

Each degree of freedom of a spin system models a magnetic moment at a fixed location. A classical spin is simply the angular momentum vector of a spinning particle, and so can be represented by an arrow, which is assumed to have fixed length. For a spin system to be well-defined, the following requirements are needed

- constraints on where the arrow is allowed to point. For example, in the XY-model, the arrow is constrained to point in a two-dimensional plane, and in the Ising model, the arrow is allowed to point in only directions. More complicated examples are the sigma-models, where instead of an arrow, the degrees of freedom take values on some manifold. In this language, the spins in the XY model take values on a circle.
- Where the spins physically are. For example, they may be at each point of some lattice in some dimension or they could be continuously distributed (ie. an arrow for every point in space eg. a field theory).
- How the spins themselves interact with one another ie. the energies associated with the possible spin configurations. Two major types of interaction are the ferromagnetic, where the energy is lowered when two spin are aligned, and antiferromagnetic, where the energy is lowered when they point in opposite directions. A given model can include interactions of both types.

Many many-body physical systems can effectively be treated as being on a lattice. For example, many of these systems are often well treated by utilizing a tight-binding model, where each electron is treated as being located at a fixed nucleus and so live at particular points in space. In some situations, the physics itself arises from the interplay between the degrees of freedom and the particular lattice they live on eg. geometrical frustration in the two-dimensional antiferromagnetic Ising model on a triangular lattice. Calling the two allowed directions in the Ising model as "up" and "down", antiferromagnetic interactions make adjacent spins different. On the square lattice, it is possible to make a low-energy state by alternating up and down spins, but on a triangular lattice this is not the case. Around a triangle, there must be at least two mutually up or down spins adjacent to each other. Such a bond is said to be unsatisfied and so the spins are frustrated, yielding completely new properties different from the unfrustrated model.

2.1.1. Formalization of Statistical Mechanics. A formal definition for a classical general spin chain, as a dynamical system, requires the use of measure theory terminology.

2.1.1.1. Some technical aspects. An introduction to Measure Theory

A **measure** is a mathematical device which reflects the notion of quantity for a given set. Let X be a set, then each subset $U \in X$ is assigned a positive real number $\mu[U]$. Thus, the measure is a function

$$\mu : \text{Dom}(\mu) \subset 2^X \rightarrow \mathbb{R},$$

where $2^X = \{S \in X\}$ is the power set of X . However, it's usually impossible to define a satisfactory notion of quantity for all subsets of X . Therefore, it is useful to consider $\text{Dom}(\mu) \in X$ ie. only some subsets of X will be measurable. Before defining a proper measure, it's domain must be specified first. This domain will be a collection of subsets of the space X , called a **σ -algebra**.

Let X be a nonempty set, which in the following sections will be a sample space. A **σ -algebra** over X is a collection $\mathcal{F} = \{U_i\}_{i \in \mathbb{N}}$ of subsets of X with the following properties:

- \mathcal{F} contains the set X : $X \in \mathcal{F}$.
- \mathcal{F} is closed under complementation: if $U \in \mathcal{F}$, then $U^c = (X/U) \in \mathcal{F}$.

- \mathcal{F} is closed under countable unions ie. $\{\mathbf{U}_i\}_{i \in \mathbb{N}} \mid \mathbf{U}_i \subset \mathcal{F} \Rightarrow \bigcup_{i \in \mathbb{N}} \mathbf{U}_i \subset \mathcal{F}$.
- \mathcal{F} is closed under countable intersections ie. $\{\mathbf{U}_i\}_{i \in \mathbb{N}} \mid \mathbf{U}_i \subset \mathcal{F} \Rightarrow \bigcap_{i \in \mathbb{N}} \mathbf{U}_i \subset \mathcal{F}$.

Technically, the fourth property, is a corollary and follows directly from the second and third properties and due to De Morgan's laws.

A **measurable space** is an ordered pair $(\mathbf{X}, \mathcal{F})$, where \mathbf{X} is a set and where \mathcal{F} is a σ -algebra on \mathbf{X} . Some of the most common examples of σ -algebras are

- Two examples of trivial σ -algebras are:
 - for any set \mathbf{X} the collection $\{\emptyset, \mathbf{X}\}$ is a σ -algebra. This first example is far too small to be of any use.
 - The power set $2^{\mathbf{X}}$ is also a σ -algebra. Note that for large sets, the power set becomes too large to be manageable.
- A more manageable σ -algebra is the one induced by the (co-)countable sets. Let \mathcal{M} be the most conservative collection of "manageable" sets, this is

$$\mathcal{M} = \left\{ \{x\} : x \in \mathbf{X} \right\},$$

ie. the set of all the singleton subsets of \mathbf{X} . Then $\mathcal{C} = \sigma(\mathcal{M})$ is the σ -algebra of **countable and co-countable sets**

$$\mathcal{C} = \{C \subset \mathbf{X} \mid \text{either } C \text{ is countable, or } \mathbf{X}/C \text{ is countable}\}.$$

If \mathbf{X} is itself finite or countable, then $\mathcal{C} = \mathcal{P}(\mathbf{X})$.

- Another example are the partition algebras. Let \mathbf{X} be a set. Then a **partition** of \mathbf{X} is a collection $\mathcal{P} = \{\mathbf{P}_i\}_{i=1}^N$ of disjoint subsets, such that $\mathbf{X} = \bigsqcup_{i=1}^N \mathbf{P}_i$. These subsets \mathbf{P}_i are called the **atoms** of the partition. Then, the σ -algebra generated by \mathcal{P} is the collection of all possible unions of \mathcal{P} -atoms:

$$\sigma(\mathcal{P}) = \left\{ \bigsqcup_{j=1}^k \mathbf{P}_{n_j} \mid \{n_j\}_{j=1}^k \in \mathbb{N}_{[1, \dots, N]} \right\}.$$

Therefore if $\text{card}[\mathbf{P}] = N$, then $\text{card}[\sigma(\mathbf{P})] = 2^N$.

If \mathcal{Q} is another partition, we say \mathcal{Q} redefines \mathcal{P} ($\mathcal{Q} \prec \mathcal{P}$) if, for every $\mathbf{P} \in \mathcal{Q}$, there are $\{\mathbf{Q}_i\}_{i=1}^N \in \mathcal{Q}$ so that $\mathbf{P} = \bigsqcup_{j=1}^N \mathbf{Q}_j$. In said case we have

$$\mathcal{P} \prec \mathcal{Q} \Leftrightarrow \sigma(\mathcal{P}) \subset \sigma(\mathcal{Q}).$$

- **Borel σ -algebra of \mathbb{R} :**

Let $\mathbf{X} = \mathbb{R}$ be the real numbers and let \mathcal{M} be the set of all open intervals in \mathbb{R} :

$$\mathcal{M} = \{(a, b) : -\infty \leq a < b \leq \infty\},$$

then the σ -algebra $\mathcal{B} = \sigma(\mathcal{M})$ contains all open subsets of \mathbb{R} , all closed subset, all countable intersections of open subsets, countable unions of closed subsets, etc. For example, \mathcal{B} contains, as elements, the set \mathbb{Z} of integers, the set \mathbb{Q} of rationals and the set \mathbb{I} of irrationals. Then, \mathcal{B} is called the **Borel σ -algebra** of \mathbb{R} .

In general, let \mathbf{X} be a topological space and let \mathcal{M} be the set of all open subsets of \mathbf{X} . The σ -algebra $\sigma(\mathcal{M})$ is the **Borel σ -algebra** of \mathbf{X} and is denoted by $\mathcal{B}(\mathbf{X})$. It contains all open sets and closed subsets of \mathbf{X} , all countable intersections of open sets (called $G\delta$ sets), all countable unions of closed sets (called $F\sigma$ sets). For example, if \mathbf{X} is Hausdorff, then $\mathcal{B}(\mathbf{X})$ contains all countable and co-countable sets.

Let $(\mathbf{X}, \mathcal{F})$ be a measurable space. A **measure** on \mathcal{F} is a map $\mu : \mathcal{F} \rightarrow \mathbb{R}_+$, which is **countably additive** ie.

$$\text{If } \{\mathbf{Y}_i\}_{i=1}^\infty \mid \mathbf{Y}_i \in \mathcal{F} \rightarrow \mu\left[\sqcup_{n=1}^\infty \mathbf{Y}_n\right] = \sum_{n=1}^\infty \mu[\mathbf{Y}_n].$$

Then a **measure space** is an ordered triple $(\mathbf{X}, \mathcal{F}, \mu)$, where \mathbf{X} is a set, \mathcal{F} is a σ -algebra and μ is a measure on \mathcal{F} . Thus, μ assigns a size to the \mathcal{F} -measurable subsets of \mathbf{X} . Some important measures and measureable spaces are

- **The counting measure** assigns, to any set, the cardinality of that set.

$$\mu[\mathbf{S}] = \text{card}[\mathbf{S}].$$

This measure provides no means of distinguishing between sets of the same cardinality, only being useful in finite measure spaces.

- A **finite measure space** is made up by a finite set \mathbf{X} , and a σ -algebra $\mathcal{F} = \mathbf{X}$. Then a measure μ on \mathbf{X} is entirely defined by some function $f : \mathbf{X} \rightarrow \mathbb{R}_+$. For any subset $\{x_i\}_{i=1}^N$ we then define

$$\mu\left(\{x_i\}_{i=1}^N\right) = \sum_{i=1}^N f(x_i).$$

Discrete probability theory is based on at most countable sample spaces Ω , where all subsets of Ω may be treated as events, and thus elements of the σ -algebra \mathcal{F} .

- **Discrete measures:** If, instead, the sample space is uncountable, (eg. it has a discrete part a non-discrete part), it can be decomposed into **atoms**. The atoms are an at most countable (maybe even empty) set, whose probability is the sum of probabilities of all atoms. If this sum is equal to 1, then all other points can safely be excluded from the sample space, returning to the discrete case. Otherwise, if the sum of probabilities of all atoms is between 0 and 1, then the measure space decomposes into a discrete atomic part and a non-atomic part. If $(\mathbf{X}, \mathcal{F}, \mu)$ is a measure space, then an **atom** of μ is a subset $\mathbf{A} \in \mathcal{F}$ such that
 - it has positive measure $\mu[\mathbf{A}] = A > 0$.
 - and it contains no set of smaller positive measure, ie. $\forall \mathbf{B} \subset \mathbf{A}$, either $\mu[\mathbf{B}] = A$ or $\mu[\mathbf{B}] = 0$.

For example, in the finite measure space above, the singleton set $\{x_n\}$ is an atom if $f(x_n) > 0$. The measure space $(\mathbf{X}, \mathcal{F}, \mu)$ is called discrete if it can be written as

$$\mathbf{X} = \mathbf{Z} \sqcup \left(\sqcup_{n=1}^\infty A_n\right),$$

where $\mu[\mathbf{Z}] = 0$ and where $\{A_n\}_{n=1}^\infty$ is a collection of atoms. Note that any finite measure space is discrete.

For example, consider the set $\mathbf{X} = \{1, 2, \dots, 9, 10\}$ and let the σ -algebra Σ be the power set of \mathbf{X} . Define the measure μ of a set to be its cardinality, that is, the number of elements in the set. Then, each of the singletons i , for $i = 1, 2, \dots, 9, 10$ is an atom. On the other hand, consider the Lebesgue measure on the real line. This measure has no atoms.

- **Density Functions:** Let $\rho : \mathbb{R}^n \rightarrow \mathbb{R}$ be a positive, integrable function on \mathbb{R}^n . For any $\mathbf{B} \in \mathcal{B}(\mathbb{R}^n)$, then a measure can be defined as

$$\mu_\rho(\mathbf{B}) = \int_{\mathbf{B}} \rho,$$

where ρ is the **density function** for μ .

Then, the ordered triple $(\mathbf{X}, \mathcal{F}, \mu)$ is a **probability space** if μ is a **probability measure**. This is, a measure μ on \mathbf{X} is a **probability measure** if $\mu[\mathbf{X}] = 1$ ³.

2.1.1.2. **Gibbs Random Fields**. A **Gibbs random field** is defined on a lattice, as follows

- A lattice, which is a countable set \mathbb{L} and a set \mathcal{L} , the set of all finite subsets of \mathbb{L} .
- A single-spin space, which is a probability space $(S, \mathcal{S}, \lambda)$, where S is the set of all possible spin-values and is assumed to be discrete, \mathcal{S} is the σ -algebra associated with this set. By definition, $\mathcal{S} \subseteq 2^S$, where 2^S is the powerset of S , and where λ is the probability measure on this single-spin space.
- A configuration space (Ω, \mathcal{F}) , where $\Omega = S^{\mathbb{L}}$ is the event space and where $\mathcal{F} = \mathcal{S}^{\mathbb{L}}$.
- Given a configuration $\omega \in \Omega$ and a subset $\Lambda \subset \mathbb{L}$, the restriction of ω to Λ is

$$\omega_{\Lambda} = (\omega(t))_{t \in \Lambda}.$$

In particular, if $A_1 \cap A_2 = \emptyset$ and $A_1 \cup A_2 = \mathbb{L}$, then the configuration $\omega_{\Lambda_1} \omega_{\Lambda_2}$ is the configuration whose restrictions to Λ_1 and Λ_2 are ω_{Λ_1} and ω_{Λ_2} , respectively.

- For each subset $\Lambda \subset \mathbb{L}$, \mathcal{F}_{Λ} is the σ -algebra generated by the family of functions

$$(\sigma(t))_{t \in \Lambda}, \text{ where } \sigma(t)(\omega) = \omega(t).$$

The union of these σ -algebras as Λ varies over \mathcal{L} is the algebra of cylinder sets on the lattice.

- The potential is a family of functions

$$\Phi = (\Phi_A)_{A \in \mathcal{L}} \text{ where } \Phi_A : \Omega \rightarrow \mathbb{R},$$

such that

- 1) $\forall A \in \mathcal{L}$, Φ_A is \mathcal{F}_{Λ} -measurable, meaning it depends only on the restriction ω_A , and does so measurably.

³**Stochastic processes** are a particular kind of probability measures, which represent a system randomly evolving in time. Let S be any randomly-evolving complex system, let \mathbf{X} be the set of all possible states of the system S and let \mathbb{T} be a set representing time. For example,

- If S is a rolling die, then $\mathbf{X} = \{1, 2, 3, 4, 5, 6\}$ and $\mathbb{T} = \mathbb{N}$ indexes the successive dice rolls.
- If S is a publically traded stock, then its state is its price. Thus $\mathbf{X} = \mathbb{R}$. If we assume trading occurs continuously when the market is open, and let each trading day have length $c < 1$, then one representation of market time is $\mathbb{T} = \sqcup_{n=1}^{\infty} [n, n + c]$.
- If S is a weather system, then its state can be represented by a large array of data $\mathbf{x} = [x_1, \dots, x_n]$. Thus $\mathbf{X} = \mathbb{R}^n$ and since the weather evolves continuously $\mathbb{T} = \mathbb{R}$.

The random evolution of S is represented by assigning a probability to every possible *history*. A history is an assignment of a state in \mathbf{X} to every moment in time (ie. in \mathbb{T}). In other words, it's a function $f : \mathbb{T} \rightarrow \mathbf{X}$. The set of all possibly histories is $\mathbf{H} = \mathbf{X}^{\mathbb{T}}$ ie. the set of all functions $f : \mathbb{T} \rightarrow \mathbf{X}$. The σ -algebra on \mathbf{H} is usually a **cylinder algebra**.

- Let $(\mathbf{X}_{\lambda}, \mathcal{X}_{\lambda})$ be measurable spaces for all $\lambda \in \Lambda$, where Λ is some (possibly uncountably infinite) indexing set. Consider the cartesian product $\prod_{\lambda \in \Lambda} \mathbf{X}_{\lambda}$. Let

$$\mathcal{M} = \left\{ \prod_{\lambda \in \Lambda} \mathbf{U}_{\lambda} \mid \forall \lambda \in \Lambda \text{ and } \mathbf{U}_{\lambda} \in \mathcal{F}_{\lambda} \text{ and } \mathbf{U}_{\lambda} = \mathbf{X}_{\lambda} \text{ for all but finitely many } \lambda \right\},$$

such subsets are called **cylinder sets** in \mathbf{X} and $\sigma(\mathcal{M})$ is the **cylinder σ -algebra** denoted by $\prod_{\lambda \in \Lambda} \mathbf{X}_{\lambda} \mathcal{F}_{\lambda}$. If the \mathbf{X}_{λ} are topological spaces with Borel σ -algebras \mathcal{F}_{λ} , and we endow \mathbf{X} with the Tychonoff product topology, then $\prod_{\lambda \in \Lambda} \mathcal{F}_{\lambda}$ is the Borel σ -algebra of \mathbf{X} .

Suppose that \mathbf{X} has a σ -algebra \mathcal{F} , then it follows \mathbf{H} 's σ -algebra is $\mathcal{H} = \prod_{t \in \mathbb{T}} \mathcal{F}_t$. An **event** is an element of \mathcal{F} and thus corresponds to a cylinder set, a countable union of cylinder sets etc. Suppose, for all $t \in \mathbb{T}$, that $\mathbf{U}_t \in \mathcal{F}$ with $\mathbf{U}_t = \mathbf{X}$ for all but finitely many t . The cylinder set $\mathbf{U} = \prod_{t \in \mathbb{T}} \mathbf{U}_t$ thus corresponds to the assertion: "for every $t \in \mathbb{T}$, at time t , the state of S was inside \mathbf{U}_t ". A probability measure on $(\mathbf{H}, \mathcal{H})$ is then a way of assigning probabilities to such assertions.

2) $\forall A \in \mathcal{L}, \forall \omega \in \Omega$, then

$$\exists H_{\Lambda}^{\Phi}(\omega) = \sum_{\substack{A \in \mathcal{L} \\ A \cap \Lambda \neq \emptyset}} \Phi_A(\omega),$$

where Φ_A is interpreted as the contribution to the total energy, ie. the Hamiltonian, associated to the interaction among all the points of finite set A . Then, $H_{\Lambda}^{\Phi}(\omega)$ is the contribution to the total energy of all the finite sets A that meet Λ . Typically, the total energy may be infinite.

3) The Hamiltonian in $\Lambda \in \mathcal{L}$ with boundary conditions $\bar{\omega}$, for the potential Φ , is defined by

$$H_{\Lambda}^{\Phi}(\omega|\bar{\omega}) = H_{\Lambda}^{\Phi}(\omega_{\Lambda}\bar{\omega}_{\Lambda^c}) \text{ where } \Lambda^c = \mathbb{L}/\Lambda.$$

- The partition function in $\Lambda \in \mathbb{L}$ with boundary conditions $\bar{\omega}$ and inverse temperature $\beta > 0$, for the Φ -potential and λ previously defined, is given by

$$\mathcal{Z}_{\Lambda}^{\Phi}(\bar{\omega}) = \int \lambda^{\Lambda}(d\omega) \exp(-\beta H_{\Lambda}^{\Phi}(\omega|\bar{\omega})), \text{ where } \lambda^{\Lambda}(d\omega) = \prod_{t \in \Lambda} \lambda(d\omega(t)) \text{ is the product measure.}$$

As dynamical systems, the Ising model or lattice gas can be obtained by first ignoring the momenta and assuming that the particles' positions are restricted to a discrete subset of \mathbb{R}^d . In general, this subset is taken to be the **d -dimensional cubic lattice**. Then,

$$(2.1.1.1) \quad \text{the lattice is simply } \mathbb{L} = \mathbb{Z}^d, \text{ where } \mathbb{Z}^d = \{\mathbf{i} = (i_1, \dots, i_d) \in \mathbb{R}^d \mid i_k \in \mathbb{Z} \forall k \in \mathbb{N}_{[1,d]}\}.$$

In other words, the system is imagined as living on \mathbb{R}^d and said space is composed of small cells, so that each cell can accommodate at most one particle. On each one of these cells, a single-spin space is then defined as $S_i = \{-1, 1\}^i$, $i \in \mathbb{Z}$. To describe the model's microstates, consider a finite region $\Lambda \subset \mathbb{Z}^d$, representing the vessel and an **occupation number** n_i defined as

$$(2.1.1.2) \quad n_i : \Lambda \rightarrow \{0, 1\} \Rightarrow \text{then the set of microstates is } \Omega_{\Lambda} = \{0, 1\}^{\times \Lambda}.$$

The model automatically includes a short-range repulsion between the particles since no two particles are allowed to share a same cell. The attractive part of the interaction can then be included into the Hamiltonian:

$$\forall \mathbf{n} = (n_i)_{i \in \Lambda}, n_i \in \Omega_{\Lambda}, \mathbf{H}(\mathbf{n}) = \sum_{\{i,j\} \subset \Lambda} \mathcal{J}(i-j) n_i n_j,$$

which is completely similar to a classical, continuous, Hamiltonian $\mathbf{H}(\mathbf{q})$ and where the function $\mathcal{J} : \mathbb{Z}^d \rightarrow \mathbb{R}$ depends, *a priori*, only on the distance between the i -th and j -th cells. Note that the contribution of a pair of cells $\{i, j\}$ is zero if they do not both contain a particle. The number of particles in Λ is given by

$$N_{\Lambda}(\mathbf{n}) = \sum_{i \in \Lambda} n_i.$$

Thus the basic object classical statistical mechanics is the Boltzmann weight $\mathbb{P}(n)$. In thermal equilibrium, this is a probability measure that gives the probability that a system will be in a certain configuration $\mathbf{n} \in \Omega_{\Lambda}$, in terms of said configuration's energy and system's temperature. In other words,

$$(2.1.1.3) \quad \mathbb{P}_{\Lambda, \beta, N} : \mathbf{n} \in \Omega_{\Lambda} \rightarrow \mathbb{R}_{[0,1]} \text{ where } \mathbb{P}_{\Lambda, \beta, N}(\mathbf{n}) = \frac{e^{-\mathbf{H}(\mathbf{n})}}{\mathcal{Z}_{\Lambda, \beta, N}},$$

where $\mathcal{Z}_{\Lambda, \beta, N}$ is the partition function defined on the finite region $\Lambda \subset \mathbb{Z}^d$, at inverse temperature β with N particles. The Boltzmann weight is defined by the requirement that the probabilities sum to one. Then, the two relevant ensemble partition functions can be defined, as follows

$$\mathcal{Q}_{\Lambda,\beta,N} = \sum_{\eta \in \Omega_{\Lambda,N}} e^{-\beta \mathbf{H}(\eta)} \quad \text{where } \Omega_{\Lambda,N} = \{n \in \Omega_{\Lambda} | N_{\Lambda}(n) = N\}$$

$$\Theta_{\Lambda,\beta,N} = \sum_N e^{\beta \mu N} \sum_{\eta \in \Omega_{\Lambda,N}} e^{-\beta \mathbf{H}(\eta)}$$

Note that if the degrees of freedom take on continuous values, or the model is defined in the continuum, then the sums are replaced by integrals. Note that, indeed, the probability of a given configuration increases as the energy gets lower, and conversely, that as the temperature gets higher and higher, the energies involved must get larger and larger to make a difference. Note as well that if all the energies are shifted by some constant E_0 , the probabilities are left unchanged, since both the numerator and denominator are multiplied by the same constant, namely $e^{-\beta E_0}$.

Consider then the Ising model, with nearest-neighbour interactions (coupling constant J) and a magnetic field (h). Its regular crystalline structure, corresponding to the positions of the magnet's atoms by a finite, non-oriented graph $G = (\Lambda, \mathcal{E})$, whose set of vertices is $\Lambda \subset \mathbb{Z}^d$. Then, the **box of radius** n is

$$B(n) = \{-n, \dots, n\}^d$$

The edges of the graph will most often be between nearest neighbours, that is, pairs of vertices i, j such that $\|j - i\|_1 = 1$ ⁴. Two nearest-neighbours spins are denoted by $i \sim j$. Then, the set of edges in the box $B(n)$ is $\{\{i, j\} \subset B(n) : i \sim j\}$.

The Ising model is defined by first assuming that a spin is located at each vertex of a graph $G = (\Lambda, \mathcal{E})$, such that each single-spin space has cardinality two, ie. each spin can only take two values. It follows, then, that to describe a microstate, a variable σ_i taking two possible values (± 1) is associated to each vertex $i \in \Lambda$, which is called the **spin** at i . A microstate of the system is called a **configuration**, and is thus an element

$$\sigma \in \Omega_{\Lambda} = \{0, 1\}^{\times \Lambda}.$$

An N -site chain, with finite N , has a configuration space, which is a discrete space, which is a 0-manifold. This 0-manifold has zero topological dimension and 2^N cardinality. The microscopic interactions among the spins are defined such that

- 1) There are only interactions between pairs of spins located at neighbouring vertices. That is, it is assumed that the spins at two distinct vertices $i, j \in \Lambda$ interact if and only if the pair $\{i, j\}$ is an edge of the graph.
- 2) The interaction favours agreement of spin values. In the simplest form of the Ising model, this is done in the simplest possible way: a pair of spins at the endpoints i and j of an edge, decreases the overall configuration's energy if they agree $\sigma_i = \sigma_j$ and increases if they differ. More precisely, the spins at the endpoints of the $\{i, j\}$ -edge contributes to the total energy by an amount $-\sigma_i \sigma_j$. Therefore, configurations in which most pairs of neighbours are aligned to have smaller energy.
- 3) Spins align with the external magnetic field. Assume that a constant external magnetic field of $h \in \mathbb{R}$ intensity, oriented along the same direction as the spins, acts on the system. Its interaction with the i -th spin contributes an $-h \sigma_i$ to the total energy. That is, when the magnetic field is positive, the configurations with most of their spins equal to $+1$ have smaller energy.

Then, the energy of a given configuration $\sigma \in \Omega_{\Lambda}$ is obtained by summing the interactions over all pairs and by adding the interaction of each spin with the external magnetic field, thus yielding the **Ising Hamiltonian**

$$(2.1.1.4) \quad \mathbf{H}_{\Lambda;h}(\sigma) = -J \sum_{\substack{i,j \in \Lambda \\ i \sim j}} \sigma_i \sigma_j - h \sum_{i \in \Lambda} \sigma_i.$$

⁴In this discrete setting, the norm is defined such that $\|i\|_1 = \sum_{k=1}^d |i_k|$

Since it favours local alignment of the spins, the Hamiltonian of the model is said to be ferromagnetic, which is not to say it behaves like a ferromagnet. The Gibbs measure/distribution is then denoted by

$$(2.1.1.5) \quad \mathbb{P}_{\Lambda, \beta, h} : \mathbf{n} \in \Omega_{\Lambda} \rightarrow \mathbb{R}_{[0,1]} \text{ where } \mathbb{P}_{\Lambda, \beta, h}(\sigma) = \frac{e^{-\beta \mathbf{H}_{\Lambda; h}(\sigma)}}{\mathcal{Z}_{\Lambda, \beta, h}}.$$

Note that, in the absence of a magnetic field, even though local spin alignment is favoured by the Hamiltonian, neither of the orientations (+1 or -1) is globally favoured. Namely, if $-\sigma$ denotes the spin-flipped configuration in which $(-\sigma)_i = -\sigma_i$, then $\mathbf{H}_{\Lambda; 0}(\sigma) = \mathbf{H}_{\Lambda; 0}(-\sigma)$, this implies

$$\mathbb{P}_{\Lambda, \beta, 0}(-\sigma) = \mathbb{P}_{\Lambda, \beta, 0}.$$

The model is then said to be invariant under global spin flip. When $h \neq 0$, this symmetry no longer holds.

The expectation value of an observable $f : \Omega_{\Lambda} \rightarrow \mathbb{R}$, under $\mu_{\Lambda, \beta, h}$, is denoted by $\langle f \rangle_{\Lambda, \beta, h}$. The parameter J is the coupling strength, if $J > 0$ this models describes a ferromagnetic coupling, while $J < 0$ describes an antiferromagnetic coupling. From the partition function, the expectation values of physical quantities can be calculated. For example

- Internal energy $\langle E \rangle_{\Lambda, \beta, h} = \frac{1}{\mathcal{Z}} \sum_n E_n e^{-\beta E_n} = -\frac{\partial}{\partial \beta} \log \mathcal{Z}$.
- Specific heat $C = \frac{\partial}{\partial T} \langle E \rangle_{\Lambda, \beta, h} = k \frac{\partial}{\partial T} T^2 \frac{\partial \log \mathcal{Z}}{\partial T}$.
- Two-point correlator in the Ising model,

$$\langle \sigma_a \sigma_b \rangle_{\Lambda, \beta, h} = \sum_{\sigma_i} \sigma_a \sigma_b \mathbb{P}_{\Lambda, \beta, h}(\sigma_i) = \frac{1}{\mathcal{Z}} \sum_{\sigma_i = \pm 1} \sigma_a \sigma_b e^{-\beta \mathbf{H}_{\Lambda; h}(\sigma)},$$

which describes how the degrees of freedom at one point are affected by the degrees of freedom at another point. If $\langle \sigma_a \sigma_b \rangle = 1$, the two spins are aligned, while if $\langle \sigma_a \sigma_b \rangle = -1$ the two spins are anti-aligned, and if $\langle \sigma_a \sigma_b \rangle = 0$, they are uncorrelated.

2.1.2. *Exact solution to the classical Ising Models.* Many quantum phase transition in d dimensions are intimately connected to certain well-studied finite temperature phase transition in classical statistical mechanics models in $d + 1$ dimensions.

Physical Context

The $d = 0$ single site quantum Ising and rotor models are then mapped to classical statistical mechanics models in $D = 1$. These $D = 1$ are very simple and do not have phase transitions. Nevertheless, it is quite useful to examine them thoroughly as they do have regions in which the correlation length ζ becomes very large. Furthermore, the properties of these regions are very similar to those in the vicinity of the phase transition points in higher dimension. In particular, the central ideas of the scaling limit and universality arise in this model. These classical models are equivalent to zero-dimensional quantum models, said mapping becoming exact in the scaling limit.

2.1.2.1. **Zero-dimensional Classical Ising model.** A system with a finite number of spins, no matter in which directions, is commonly referred to a zero-dimensional model. In the $(D = 0)$ -Ising model, there are just two degrees of freedom, called s_1 and s_2 , both of which can only take two values $s_i \in S = \{\pm 1\}$ and where the configuration space is given by $(\Omega_\Lambda, \mathcal{F})$, where $\Omega_\Lambda = \{-1, 1\}^{\times 2}$ and where the σ -algebra $\mathcal{F} = 2^{\Omega_\Lambda}$. To compute the partition function, $\mathcal{Z}_{\Lambda, \beta, B}$, the system's energy must be known for every configuration $(s_1, s_2) \in \Omega_\Lambda$, which is assumed to be

$$(2.1.2.1) \quad E(\mathbf{s}) = -J s_1 s_2 - B(s_1 + s_2).$$

This energy models a system made of magnetic moments. Although, in general, the moments can point in any direction, in the Ising model they point up or down one axis - say the z -axis - along which an external magnetic field B is applied. The constant J has units of energy and B has units of magnetic moment. If $J > 0$, the model is ferromagnetic and the energy favours aligned spins while if $J < 0$ it models an antiferromagnetic system, which favours anti-parallel spins. In both cases, $B > 0$ is chosen so that the spins tend to align with the external field.

The partition function may be evaluated as

$$(2.1.2.2) \quad \mathcal{Z}_{\Lambda, K, h} = \sum_{\mathbf{s} \in \Omega_\Lambda} e^{K s_1 s_2 + h(s_1 + s_2)} \quad \text{where} \quad \begin{matrix} K = \beta J \\ h = \beta B \end{matrix} \Rightarrow \mathcal{Z}_{\Lambda, K, h} = \sum_{m=-\frac{1}{2}}^{\frac{1}{2}} e^{\beta h m} 2 \cosh\left(\frac{h + K m}{2}\right) \\ = 2 \cosh(2h) e^K + 2e^{-K}.$$

At very low temperatures (very low T or very high β), the state of lowest energy will dominate the sum and thus, the spins are expected to be aligned with B and also with each other. At very high temperatures and vanishing β , all four states get equal weight and the spins will fluctuate independently of each other and the applied field. In effect, let the system's (spin-) average magnetization be $M = \frac{1}{2}(s_1 + s_2)$, in any given configuration⁵. The thermal average can then be found to be

$$(2.1.2.3) \quad \langle M \rangle_{\Lambda, K, h} = \frac{1}{\mathcal{Z}_{\Lambda, \beta, h}} \sum_{\mathbf{s} \in \Omega_\Lambda} \frac{1}{2}(s_1 + s_2) e^{K s_1 s_2 + h(s_1 + s_2)} = \frac{1}{2 \mathcal{Z}_{\Lambda, K, h}} \frac{\partial \mathcal{Z}_{\Lambda, K, h}}{\partial h} \\ = \frac{1}{2} \frac{\partial \log \mathcal{Z}_{\Lambda, K, h}}{\partial h}.$$

In terms of the free energy $F(K, h)$, defined as $\mathcal{Z}_{\Lambda, K, h} = e^{-\beta F}$, where $-\beta F = \log(2 \cosh(2h) e^K + 2e^{-K})$ it yields

⁵Note that this is the average over all spins in the system, and not the thermal average

$$(2.1.2.4) \quad \langle M \rangle_{\Lambda, K, h} = \frac{\partial[-\beta F(K, h)]}{\partial h} = \frac{\sinh 2h}{\cosh 2h + e^{-2K}} \text{ which, as expected } \begin{matrix} \langle M \rangle_{h, K \rightarrow \infty} \rightarrow 1 \\ \langle M \rangle_{h, K \rightarrow 0} \rightarrow h. \end{matrix}$$

The thermal average of a particular spin can also be calculated, provided a source term of the form $h_1 s_1 + h_2 s_2$ is added in the Boltzmann weight, which couples the i -th spin to its own independent field h_i . Then,

$$(2.1.2.5) \quad \mathcal{Z}_{\Lambda, K, h} = \sum_{\mathbf{s} \in \Omega_{\Lambda}} e^{K s_1 s_2 + \mathbf{h} \cdot \mathbf{s}} = e^{-\beta F(K, \mathbf{h})} \Rightarrow \langle s_i \rangle = \frac{\partial[-\beta F(K, \mathbf{h})]}{\partial h_i} = \frac{\partial \log \mathcal{Z}_{\Lambda, K, h}}{\partial h_i} = \frac{1}{\mathcal{Z}_{\Lambda, K, h}} \frac{\partial \mathcal{Z}_{\Lambda, K, h}}{\partial h_i}.$$

Taking the mixed derivative with respect to both h_1 and h_2 yields

$$(2.1.2.6) \quad \begin{aligned} \frac{\partial^2[-\beta F(K, \mathbf{h})]}{\partial h_1 \partial h_2} &= \frac{\partial}{\partial h_1} \left(\frac{1}{\mathcal{Z}_{\Lambda, K, h}} \frac{\partial \mathcal{Z}_{\Lambda, K, h}}{\partial h_2} \right) \\ &= \frac{1}{\mathcal{Z}_{\Lambda, K, h}} \frac{\partial^2 \mathcal{Z}_{\Lambda, K, h}}{\partial h_1 \partial h_2} - \frac{1}{\mathcal{Z}^2} \frac{\partial \mathcal{Z}_{\Lambda, K, h}}{\partial h_1} \frac{\partial \mathcal{Z}_{\Lambda, K, h}}{\partial h_2} \\ &= \langle s_1 s_2 \rangle - \langle s_1 \rangle \langle s_2 \rangle \\ &= \langle s_1 s_2 \rangle_c \text{ which is the connected correlation function.} \end{aligned}$$

Similarly, if there are four spins coupled to four independent magnetic field h_1, h_2, h_3, h_4 then

$$(2.1.2.7) \quad \begin{aligned} \frac{\partial^4[-\beta F(K, \mathbf{h})]}{\partial h_1 \partial h_2 \partial h_3 \partial h_4} \Big|_{h=0} &= \langle s_1 s_2 s_3 s_4 \rangle_c \\ &= \langle s_1 s_2 s_3 s_4 \rangle - \langle s_1 s_2 \rangle \langle s_3 s_4 \rangle - \langle s_1 s_3 \rangle \langle s_2 s_4 \rangle - \langle s_1 s_4 \rangle \langle s_2 s_3 \rangle. \end{aligned}$$

Since, at the end of the calculation, $h = 0$, there will be no correlators with an odd number of spins. Note that the previous derivations are valid even if s is replaced by some other variable, with eg. continuous values, coupled to the corresponding magnetic field in the Boltzmann weight, since the fact that $s = \pm 1$ was not used. The thermal averages may be computed in the case of non-zero zeros. In particular, if the field is uniform or zero. In that case, the h_i -derivatives are calculated and then set to $h_i = h$ or 0, for all i . Then, for example, if there were a uniform external field h , the connected correlation function is

$$(2.1.2.8) \quad \langle s_1 s_2 \rangle - \langle s_1 \rangle \langle s_2 \rangle = \langle s_1 s_2 \rangle_c = \frac{\partial^2[-\beta F(K, \mathbf{h})]}{\partial h_1 \partial h_2} \Big|_{h_i=h, \forall i}.$$

The correlation function $\langle s_1 s_2 \rangle$ as a K -derivative,

$$(2.1.2.9) \quad \begin{aligned} \frac{\partial[-\beta F(K, \mathbf{h})]}{\partial K} &= \frac{\partial \log \mathcal{Z}_{\Lambda, K, h}}{\partial K} = \frac{1}{\mathcal{Z}_{\Lambda, K, h}} \frac{\partial \mathcal{Z}_{\Lambda, K, h}}{\partial K} = \frac{1}{\mathcal{Z}_{\Lambda, K, h}} \frac{\partial}{\partial K} \sum_{\mathbf{s} \in \Omega_{\Lambda}} e^{K s_1 s_2 + h(s_1 + s_2)} \\ &= \frac{1}{\mathcal{Z}_{\Lambda, K, h}} \sum_{\mathcal{M} \times 2} s_1 s_2 e^{K s_1 s_2 + h(s_1 + s_2)} = \langle s_1 s_2 \rangle \\ &= \frac{e^K \cosh 2h - e^{-K}}{e^K \cosh 2h + e^{-K}}, \end{aligned}$$

which may interpreted as quantifying the average interaction energy, which happens to coincide with the correlation of neighbouring spins. In a model with more spins, correlations of not-neighbouring spins (eg. the correlation between s_5 and s_{92}) cannot be obtained via a K -derivative. The distinction between the average interaction energy and generic spin correlations is blurred in this toy-model.

Higher-order derivatives with respect to K or h (for uniform h) give additional information about fluctuations about the mean. Thus,

$$(2.1.2.10) \quad \chi = \frac{1}{2} \frac{\partial \langle M \rangle}{\partial h} = \frac{1}{4} \frac{\partial^2 [-\beta F]}{\partial h^2} = \langle M^2 \rangle_{\Lambda, K, h} - \langle M \rangle_{\Lambda, K, h}^2$$

measures the magnetization's fluctuation about its mean and is the magnetic susceptibility, and quantifies the rate of change of the average magnetization with the applied field.

2.1.2.2. One-dimensional Classical Ising model. The standard way for solving the one-dimensional Ising model is to use the transfer matrix. The physical motivation behind it is very elegant: it consists on singling out a given spatial direction, say the x -direction. Then the partition function is built up by "evolving" the system from one value of $x = x_0$ to the next value $x = x_1$. Repeating this process yields the full partition function after all values of x are summed over. This method is not only useful for classical statistical mechanics, but by replacing the spatial evolution with a real time-evolution, yields a connection between quantum and classical statistical mechanics.

Now, any real system is finite, one may ask why we should focus on infinite systems. The answer lies that systems with a very large number of degrees of freedom look more like infinite systems than finite ones in the following sense. Phenomenologically, system at a temperature lower than some Curie temperature T_C will magnetize. But it can be shown rigorously that a finite system of Ising spins can never magnetize: if it is polarized one way (say the z -axis) for some time, it can jump to the opposite polarization after some time, so that on average it is not magnetized. This rigorous result can be reconciled with reality by noting that the time to flip can be as large as the age of the universe, so that in human time, magnetization is possible. The nice thing about infinite system is that remote possibilities are rendered impossible, making them a better model of real life.

In the ($D = 1$)-dimensional Ising model, the energy of a given configuration is given by

$$(2.1.2.11) \quad E = -J \sum_{i=0}^{N-1} \sigma_i \sigma_{i+1},$$

In said model, at a given x -value, there is only a single site, so that the system "evolves" from one spin to the next. This model, as it stands, is called ferromagnetic, since there is an overall minus sign and thus favours local spin alignment. The reader should be aware that the model may not behave like a ferromagnet.

Given a one-dimensional Ising spin chain. Consider two nearest-neighbour sites in said chain, there are, then, four configurations on these two sites, $|++\rangle$, $|+-\rangle$, $| -+\rangle$, $|--\rangle$. The Boltzmann weights $P(\sigma_1 \sigma_2)$ for these four configurations are

$$(2.1.2.12) \quad P(++) = P(--) = e^{\beta J} \text{ and } P(+-) = P(-+) = e^{-\beta J}.$$

Now, consider for example an Ising model on three sites, with the simplest boundary conditions, where both spins σ_1 and σ_3 are both fixed at specific values. Summing over all values of the middle spin $\sigma_2 = \pm 1$, the partition function with both end spins fixed is

$$(2.1.2.13) \quad \sum_{\sigma_2=\pm 1} P(+\sigma_2)P(\sigma_2+) = e^{2J} + e^{-2J}.$$

The general expression for the partition function of an N -site Ising model with fixed boundary conditions can be found to be

$$\text{If } N = 3 \text{ then } \mathcal{Z}_3(\text{fixed}) = \sum_{\sigma_2=\pm 1} P(\sigma_1 \sigma_2) P(\sigma_2 \sigma_3).$$

$$\text{If } N = 4 \text{ then } \mathcal{Z}_4(\text{fixed}) = \sum_{\substack{\sigma_2=\pm 1 \\ \sigma_3=\pm 1}} P(\sigma_1 \sigma_2) P(\sigma_2 \sigma_3) P(\sigma_3 \sigma_4).$$

These expressions look exactly like matrix multiplications, explicitly written out in terms of the matrix elements. The matrix being multiplied is called the transfer matrix and includes all the nearest-neighbour interactions. For the one-dimensional classical Ising model it is given by

$$(2.1.2.14) \quad \mathcal{T} = \begin{pmatrix} \mathbb{P}(++) & \mathbb{P}(+-) \\ \mathbb{P}(-+) & \mathbb{P}(--) \end{pmatrix} = \begin{pmatrix} e^J & e^{-J} \\ e^{-J} & e^J \end{pmatrix}.$$

The partition function for other systems, with different boundary conditions, can be readily calculated from the transfer matrix. To compute the partition function with fixed boundaries conditions, from the transfer matrix, the boundary conditions are treated as vectors, on which the transfer \mathcal{T} -matrix acts on. Here, the basis elements are the two spin values $+$ and $-$. These two vectors span a \mathbb{C}^2 -vector space. The partition function of an N -site system with fixed boundary conditions can be found out to be

- if $N = 2$, $\mathcal{Z}_2(\text{fixed}) = \langle \sigma_1 | \mathcal{T} | \sigma_2 \rangle = \mathbb{P}(\sigma_1 \sigma_2)$,
- if $N = 3$, $\mathcal{Z}_3(\text{fixed}) = \langle \sigma_1 | \mathcal{T}^2 | \sigma_3 \rangle$,
- in general, $\mathcal{Z}_N(\text{fixed}) = \langle \sigma_1 | \mathcal{T}^{N-1} | \sigma_N \rangle$.

Each time the transfer matrix is applied, it corresponds to adding one bond to the system. The transfer matrix can now be used to find analogous expressions for other boundary conditions. For example,

- for free boundary conditions, allowing the extreme-most spins to take on all allowed values, the partition function is given by

$$(2.1.2.15) \quad \begin{aligned} \mathcal{Z}_N(\text{free}) &= \sum_{\substack{\sigma_1=\pm 1 \\ \sigma_N=\pm 1}} \langle \sigma_1 | \mathcal{T}^{N-1} | \sigma_N \rangle \\ &= \begin{pmatrix} 1 & 0 \end{pmatrix} \mathcal{T}^{N-1} \begin{pmatrix} 1 \\ 0 \end{pmatrix} + \begin{pmatrix} 1 & 0 \end{pmatrix} \mathcal{T}^{N-1} \begin{pmatrix} 0 \\ 1 \end{pmatrix} \\ &\quad + \begin{pmatrix} 0 & 1 \end{pmatrix} \mathcal{T}^{N-1} \begin{pmatrix} 1 \\ 0 \end{pmatrix} + \begin{pmatrix} 0 & 1 \end{pmatrix} \mathcal{T}^{N-1} \begin{pmatrix} 0 \\ 1 \end{pmatrix} \\ &= \begin{pmatrix} 1 & 1 \end{pmatrix} \mathcal{T}^{N-1} \begin{pmatrix} 1 \\ 1 \end{pmatrix}. \end{aligned}$$

- For periodic boundary conditions, where an extra interaction $-J\sigma_N\sigma_1$ is added to the energy, including a bond between the two most-extreme spins. In this case, the model is translationally-invariant, shifting all the spins by one site mod N doesn't change the energy. The transfer \mathcal{T} must therefore evolve the first spin back to the N -th one. This is equivalent to having fixed boundary conditions on $(N+1)$ -sites where $\sigma_{N+1} = \sigma_1$, and then summing over both possible values, thus yielding

$$(2.1.2.16) \quad \begin{aligned} \mathcal{Z}_N(\text{periodic}) &= \sum_{\sigma_1=\pm 1} \langle \sigma_1 | \mathcal{T} | \sigma_1 \rangle = \begin{pmatrix} 1 & 0 \end{pmatrix} \mathcal{T}^N \begin{pmatrix} 1 \\ 0 \end{pmatrix} + \begin{pmatrix} 0 & 1 \end{pmatrix} \mathcal{T}^N \begin{pmatrix} 0 \\ 1 \end{pmatrix} \\ &= \text{Tr}_{\mathbb{C}^2} \mathcal{T}^N. \end{aligned}$$

The usefulness of the transfer matrix in a $(D = 1)$ -dimensional system lies in that the partition function can be computed simply by diagonalizing the transfer matrix. This diagonalization yields

$$(2.1.2.17) \quad \mathcal{T} = \begin{pmatrix} 1 & 1 \\ -1 & 1 \end{pmatrix} \begin{pmatrix} 2 \cosh(\beta J) & 0 \\ 0 & 2 \sinh(\beta J) \end{pmatrix} \begin{pmatrix} 1 & 1 \\ -1 & 1 \end{pmatrix}^{-1},$$

and given that the transfer matrix is hermitian, the eigenvalues are real. Then, the partition functions for the different boundary conditions can then be found

- For the free case, the $|\sigma_j\rangle$ -vectors can be decomposed in terms of the \mathcal{T} -matrix's eigenvectors, which can be written as

$$v_{\pm} = \frac{1}{2} \left(\begin{pmatrix} 1 & 1 \end{pmatrix}^T + \begin{pmatrix} 1 & -1 \end{pmatrix}^T \right)$$

Therefore, plugging in these results in [equation \(2.1.2.15\)](#) yields,

$$(2.1.2.18) \quad \mathcal{Z}_N(\text{free}) = \begin{pmatrix} 1 & 1 \end{pmatrix} \mathcal{T}^{N-1} \begin{pmatrix} 1 \\ 1 \end{pmatrix} = 2(2 \cosh \beta J)^{N-1}.$$

- For the periodic-boundary conditions, the partition function may be calculated as the sum of $\mathcal{Z}_N(++)$ and $\mathcal{Z}_N(--)$, which can be easily calculated from [equation \(2.1.2.17\)](#), as follows

$$(2.1.2.19) \quad \begin{aligned} \mathcal{Z}_N(++) &= \begin{pmatrix} 1 & 0 \end{pmatrix} \mathcal{T}^N \begin{pmatrix} 1 \\ 0 \end{pmatrix} = \frac{1}{2} \left[(2 \cosh \beta J)^N + (2 \sinh \beta J)^N \right] \\ \mathcal{Z}_N(+-) &= \begin{pmatrix} 1 & 0 \end{pmatrix} \mathcal{T}^N \begin{pmatrix} 0 \\ 1 \end{pmatrix} = \frac{1}{2} \left[(2 \cosh \beta J)^N - (2 \sinh \beta J)^N \right] \\ \mathcal{Z}_N(-+) &= \begin{pmatrix} 0 & 1 \end{pmatrix} \mathcal{T}^N \begin{pmatrix} 1 \\ 0 \end{pmatrix} = \frac{1}{2} \left[(2 \cosh \beta J)^N - (2 \sinh \beta J)^N \right] \\ \mathcal{Z}_N(--) &= \begin{pmatrix} 0 & 1 \end{pmatrix} \mathcal{T}^N \begin{pmatrix} 0 \\ 1 \end{pmatrix} = \frac{1}{2} \left[(2 \cosh \beta J)^N + (2 \sinh \beta J)^N \right] \end{aligned}$$

From these, it is easy to see that partition functions with different boundary conditions are simply related, eg. $\mathcal{Z}_N(\text{periodic}) = \mathcal{Z}_N(++) + \mathcal{Z}_N(--)$. Then, partition function for a system with the periodic boundary condition is given by

$$(2.1.2.20) \quad \begin{aligned} \mathcal{Z}_N(\text{periodic}) &= (2 \cosh \beta J)^N + (2 \sinh \beta J)^N \\ &= (2 \cosh \beta J)^N \left(1 + (2 \tanh \beta J)^N \right) \\ &\approx (2 \cosh \beta J)^N \text{ for } N \gg 1. \end{aligned}$$

Note as well that the free energy per site is proportional to $\frac{\log \mathcal{Z}}{N}$, and is independent of boundary conditions when $N \gg 1$. In effect, the free energy per site in the thermodynamic limit may be written as

$$(2.1.2.21) \quad F(\beta, K) = - \lim_{N \rightarrow \infty} \frac{\log \mathcal{Z}}{N}.$$

It is straight forward to compute the two-point function in terms of these partition functions. Consider the following special cases.

- First, consider an Ising spin chain with free boundary conditions. Then, the two-point function between the two end-spins can be computed by summing over the four possible configurations. But these spins are opposite, then

$$(2.1.2.22) \quad \langle \sigma_1 \sigma_N \rangle = \frac{\mathcal{Z}_N(++) + \mathcal{Z}_N(--) - \mathcal{Z}_N(+-) - \mathcal{Z}_N(-+)}{\mathcal{Z}_{\text{free}}} = (\tanh \beta J)^{N-1},$$

which falls off exponentially with distance, so even for small distances, the end spins are effectively uncorrelated.

- Now, consider an Ising chain with periodic boundary conditions. Then, the the two-point function between the j -th and the i -th spins, with $j > i$, may be written as

$$(2.1.2.23) \quad \langle \sigma_j \sigma_i \rangle_{\text{periodic}} = \frac{2\mathcal{Z}_{|j-i|}(++)\mathcal{Z}_{N-|j-i|}(--) + 2\mathcal{Z}_{|j-i|}(+-)\mathcal{Z}_{N-|j-i|}(-+)}{\mathcal{Z}_{N\text{periodic}}},$$

The general two-point correlator is similar and is given by

$$(2.1.2.24) \quad \langle \sigma_j \sigma_i \rangle_{j>i} = \frac{1}{\mathcal{Z}} \sum_k \sigma_j \sigma_i e^{\sum_k \beta J \sigma_k \sigma_{k+1}},$$

which measures how likely σ_i and σ_j are, on average, to point in the same direction. In any ferromagnetic system, the average will be positive for a pair of neighbouring spins since the Boltzmann weight is biased towards parallel, ie. aligned, values. Surely, if a spin can influence its neighbours to be parallel to it, they in turn will act similarly on their own neighbours, therefore long-range correlations may be present. In addition, if there is an external magnetic field present, it will enhance this correlation further. Naturally, if $j > i$ it holds that

$$\sigma_j \sigma_i = \sigma_i \sigma_j = \sigma_i \sigma_{i+1} \sigma_{i+2} \cdots \sigma_{j-1} \sigma_j = \mathbf{t}_i \mathbf{t}_{i+1} \cdots \mathbf{t}_{j-1} \text{ where } \mathbf{t}_i = \sigma_{i+1},$$

then

$$(2.1.2.25) \quad \langle \sigma_j \sigma_i \rangle = \langle \mathbf{t}_i \rangle \langle \mathbf{t}_{i+1} \rangle \cdots \langle \mathbf{t}_{j-1} \rangle.$$

Note that the answer factorizes over i since the Boltzmann weight factorizes over i as well, when written in terms of the \mathbf{t}_i ⁶. Then, the thermal average for any one \mathbf{t} can be easily calculated as

$$(2.1.2.27) \quad \langle \mathbf{t} \rangle = \frac{1}{\mathcal{Z}} \sum_{\mathbf{t}=\pm 1} \mathbf{t} e^{E(\{\mathbf{t}\})} \Rightarrow \langle \mathbf{t} \rangle = \frac{1e^{\beta J} - 1e^{-\beta J}}{e^{\beta J} + e^{-\beta J}} = \tanh \beta J$$

$$\Rightarrow \langle \sigma_j \sigma_i \rangle = (\tanh \beta J)^{|j-i|} = \exp \left(|j-i| \log \tanh \beta J \right).$$

Consider the form of the correlations given by **equation (2.1.2.27)** in the thermodynamic limit $N \rightarrow \infty$, in which case at any finite βJ , since $\tanh \beta J < 1$, $\langle \sigma_j \sigma_i \rangle \xrightarrow{|j-i| \rightarrow \infty} 0$. Only at $T = 0$ or $\beta J \rightarrow \infty$,⁷, does the correlation not exponentially decay but remain flat. In other words, note that

$$(2.1.2.28) \quad \langle \sigma_j \sigma_i \rangle \sim e^{\frac{-|j-i|}{\zeta}} \text{ where } \zeta = -\frac{1}{\log \tanh \beta J} \text{ is the correlation length.}$$

Here ζ always represents the actual correlation length at zero external magnetic field. In the large βJ -limit, the correlation length becomes much larger than the lattice spacing, $|i-j|$,

$$(2.1.2.29) \quad \frac{\zeta}{|i-j|} \approx \frac{1}{2} e^{\beta J} \gg 1, \text{ for } \beta J \gg 1.$$

Thus, the one-dimensional Ising model is in a **disordered phase for any non-zero temperature**. The zero-temperature case is pathological, since there is no meaning of equilibrium in a zero-temperature classical system.

⁶In effect, note that the partition function may be written as

$$(2.1.2.26) \quad \mathcal{Z} = \sum_{\sigma_i=\pm 1} \exp \left(\sum_{i=0}^{N-1} \beta J \sigma_i \sigma_{i+1} \right) = \sum_{\mathbf{t}_i=\pm 1} \exp \left(\sum_{i=0}^{N-1} \beta J \mathbf{t}_i \right)$$

$$= \sum_{\mathbf{t}_i=\pm 1} \prod_{i=0}^{N-1} e^{\beta J \mathbf{t}_i},$$

where the exponential has naturally factorized into a product over i .

⁷i.e. when $\tanh \beta J = 1$

The $\langle \sigma_j \sigma_i \rangle$ -correlator depends on just the difference in coordinates, ie. it presents translational invariance. This is not a generic result for a finite chain, but a peculiarity of this model in particular. In general, for an $N + 1$ -site chain, correlations between two spins will generally depend on where the two points are in relation to the ends. On the other hand, for $N \rightarrow \infty$ -models, translational invariance is expected to hold for correlations of spins far from the ends. In order for translational invariance to hold for a finite system, periodic boundary conditions must be imposed on the system. In said case, correlation functions will now only depend on the difference between the two coordinates, but will not decay monotonically with separation since as one point start moving away from one extreme, it starts approaching the extreme spin from the other side.

The scaling limit

The simplest way to understand the scaling limit is to first divide all lengths into "large" and "small" lengths. For the Ising chain,

- the correlation length ζ , the observation scale $J|i - j|$ and the system size $L = M|i - j|$ as large lengths,
- and the lattice spacing $|i - j|$ is the only small length.

Then, the scaling limit of an observable is then defined as its value when all corrections involving the ratio of small to large lengths are neglected.

There are two different, but equivalent, ways of thinking about the scaling limit. Particle physicists are inclined to send the small length $|i - j|$ to zero while keeping the large lengths fixed, while condensed matter physicists send all large lengths to infinity, while keeping $|i - j|$ fixed. The physics can only depend upon the ratio of lengths, thus making both approaches equivalent.

For the definition of the scaling limit to be complete, the manner in which βJ and the external magnetic field h are treated. From [equation \(2.1.2.29\)](#) it is clear that βJ can be thought as a ratio of lengths, $\frac{\zeta}{|i - j|}$, allowing for the elimination of the explicit dependence on the βJ -variable, by specifying $\frac{\zeta}{|i - j|} \rightarrow \infty$. In the βJ -limit, invariably, spins a few lattice spacing apart point in the same direction, and should be sensitive to the mean magnetic field h per unit length, measured by \tilde{h} ,

$$\tilde{h} = \frac{h}{|i - j|}.$$

So the scaling limit can be taken with $|i - j| \rightarrow 0$, while keeping \tilde{h} fixed.

2.1.2.3. **The Monte-Carlo method.** The partition function and two-point functions may be numerically calculated with the Monte-Carlo method. Consider the two point function,

$$(2.1.2.30) \quad \langle \sigma_i \sigma_j \rangle_{\Lambda, \beta, h} = \frac{\sum_{\sigma \in \Omega_\Lambda} \sigma_i \sigma_j e^{-\beta \mathbf{H}_{\Lambda, \beta, h}(\sigma)}}{\sum_{\sigma \in \Omega_\Lambda} e^{-\beta \mathbf{H}_{\Lambda, \beta, h}(\sigma)}} = \sum_{\sigma \in \Omega_\Lambda} \sigma_i \sigma_j(\sigma) \mathbb{P}_{\Lambda, \beta, h}(\sigma),$$

where $\sigma_i \sigma_j(\sigma)$ is the value of $\sigma_i \sigma_j$ at the $\sigma \in \Omega_\Lambda$ configuration. The Monte-Carlo method is useful for numerically computing multidimensional integrals or sums. Each configuration will occur at a rate proportional to its Boltzmann weight. Then, for each configuration, the value of $\sigma_i \sigma_j$ can be computed. Since the configurations are already weighted, the arithmetic average of these numbers will yield the weighted thermal average. From this correlation function, the correlation length ε can be calculated. By the same method, any other quantity of interest may be calculated eg. the magnetization. This can be done for any problem with a real, positive, Boltzmann weight, not just for Ising models. One specific algorithm for calculating these weighted configurations is the **Metropolis Method**, as follows

Data: System's parameters: J, h, β

Result:

```

Start with some configuration  $\sigma_i \in \Omega_\Lambda$  with energy  $E_i$ ; while not at end of this document do
  Consider another configuration  $\sigma_j \in \Omega_\Lambda$  with energy  $E_j$ , obtained by changing some degrees of
  freedom;
  if  $E_j < E_i$  then
    | jump to  $j$  with unit probability;
  else
    | jump to  $j$  with probability  $e^{-\beta(E_j - E_i)}$ 
  end
end

```

Algorithm 1: Metropolis Algorithm

After a great number of iterations, the configurations appear with probabilities obeying

$$\frac{p(i)}{p(j)} = e^{-\beta(E_i - E_j)}.$$

In effect, the rate equation reads

$$\frac{dp(i)}{dt} = -p(i) \sum_j R(i \rightarrow j) + \sum_j p(j) R(j \rightarrow i),$$

where $R(i \rightarrow j)$ is the rate of jumping from the i -th configuration to the j -th configuration and vice-versa. In equilibrium or steady state, $\dot{p}(i) = 0$. One way to impose such constrain on the right-hand side is by considering a detailed balance approach, to ensure that the contribution from every value j is separately zero, rather than just the sum. It is necessary but not sufficient. In this case, this implies

$$(2.1.2.31) \quad \frac{p(i)}{p(j)} = \frac{R(i \rightarrow j)}{R(j \rightarrow i)}.$$

2.1.2.4. Three kinds of phases. Even in simple interacting systems, the two point function will not have such simple values except in some extreme limits. However, knowing its dependence on the separation is valuable information, and it implies that the behaviour of the correlators is one of the best ways of understanding and characterising different phases of matter. Consider a correlator defined so that in the non-interacting limit, ie. the $T \rightarrow \infty$, it vanishes. Letting L be the system size, there are three main types of behaviour of such a two-point function as $L \gg |a - b| \gg 1$, as follows

- **ordered:**

$$\langle \sigma_a \sigma_b \rangle \sim C \neq 0,$$

In an ordered spin system, the spins are correlated arbitrarily far from each other: if a spin, at a given position, points in some direction, then another spin is going to most likely point in the same direction. For example, in an antiferromagnet, order typically means that spins at an odd number of sites apart, are antialigned, while those an even number of sites apart are aligned.

Another way to characterize order is in terms of a one-point function, of a given variable whose thermodynamic average vanishes in the non-interacting limit. In the Ising model, this is simply $\langle \sigma_a \rangle$, which is the probability that a given spin is either up or down. In principle, this is easier to compute than the two-point function but, in many cases including the Ising model, it vanishes for all values of a, b, J . Indeed, this follows from the Ising model's \mathbb{Z}_2 -symmetry $\sigma_i \rightarrow -\sigma_i$. Then, $E(\{\sigma_i\}_{i=1}^N) = E(-\{\sigma_i\}_{i=1}^N)$. Therefore, as long as the boundary conditions remain invariant under spin flip, it holds that

$$(2.1.2.32) \quad -\langle \sigma_a \rangle = -\sum_{\sigma_i = \pm 1} \sigma_i e^{-\beta E(\{\sigma_i\})} = \sum_{\sigma_i = \mp 1} s_a e^{-\beta E(\{s_i\})} \text{ where } s_i = -\sigma_i.$$

However, since the sum is performed over a finite number of lattice sites, the s_i -variables on the last equality may be renamed as σ_i . Thus, the last expression is simply $\langle \sigma_a \rangle = 0$.

- **disordered:**

$$\langle \sigma_a \sigma_b \rangle \sim e^{-\frac{|a-b|}{\epsilon}},$$

A more general way of characterizing order is to use the previous definition, ie. a non-vanishing value of a two-point function at arbitrarily large separation. Then, the order parameter can be defined as the square root of this absolute value.

- **critical:**

$$\langle \sigma_a \sigma_b \rangle \sim \frac{1}{|a - b|^{2x}}.$$

In a disordered phase, the two-point function class off to zero exponentially fast. Thus a convenient measure is the correlation length ϵ . A useful heuristic argument is to think of the degrees of freedom closer than ϵ as being correlated, and those further apart as being uncorrelated. One thing to bear in mind is that not all disordered systems behave in the same way. Some systems are disordered by definition, but possess a topological order, where an expectation value of some non-local quantities has a non-vanishing expectation value.

The system's criticality is the most difficult and physically meaningful case. The correlator is said to be algebraically decaying. In particular, in the algebraic decay, the spatial dependence does not involve in any way a dimensional parameter, but rather only a dimensionless parameter x . Thus, this behaviour is characteristic of a critical phase, where the system is scale-invariant, given that there is no parameter like a correlation length. The reason for the factor of 2 in the x -quantity's definition stems from this picture, in the continuum limit, x can be thought of as a "scaling dimension" of the spin field.

Note that, this picture is oversimplified since different types of correlators can occur in the same system.

A common case in statistical mechanics occurs when, at temperatures much larger than any parameter with dimensions of energy - like J in the Ising model-, one can effectively neglect the energy altogether. The partition function is then a sum over all allowed configurations with the same weight each. Then, the two-point function of any local operators vanishes quickly as they are brought apart (as long as non-local constraints are not present). The system is then in a disordered phase at sufficiently large temperatures. At low enough temperatures, the opposite occurs: since the Boltzmann weight is proportional to $e^{-E/T}$, the very low-energy configurations are much more likely. This can lead to order, the lowest-energy configurations have some regular pattern and so there is a non-vanishing local order parameter. When there is a disordered phase at high-temperature and an ordered phase at high temperature as well, there must be a phase transition in between. There are two categories of phase transitions, often referred to as first-order and continuous. In the latter case, correlators become algebraically decaying, and then the system is critical. In the former case, the system remains disordered at the phase transition and abruptly changes its behaviour.

2.2. Quantum Spin Systems. There is a deep correspondence between classical statistical mechanics and quantum statistical mechanics for spin- $\frac{1}{2}$ particles. In particular, the d -dimensional classical statistical mechanics problem may be mapped to a $d - 1$ -dimensional quantum statistical problem. The nature of the quantum variable will depend then on the nature of the classical variable. In general, the allowed values of the classical variable will correspond to the maximal set of simultaneous eigenvalues of the operators in the quantum problem (eg. σ_z).

The Hilbert space of a quantum spin is defined by choosing a representation for the spin operators. A representation of a Lie algebra is a set of three matrices satisfying the commutation relations of the $\mathfrak{su}(2)$ -algebra. An **irreducible representation** is a set of matrices such that no unitary transformation $U S^a U^\dagger$ block-diagonalizes all three matrices. It is known that for the $\mathfrak{su}(2)$ -Lie algebra there is exactly one set (up to unitary transformations) of irreducible complex $n \times n$ -matrices, for each integer n . It is customary to write $n = 2s + 1$ for all integers and half-integers s . A single spin- s quantum particle at a fixed point in space therefore has a Hilbert space \mathbb{C}^{2s+1} , so the matrices S^a are all $(2s + 1) \times (2s + 1)$. An orthonormal basis is given by the eigenstates of any one of the matrices.

- For $s = 0$, the matrices all consist of the number zero, thus this is the trivial representation.
- For $s = \frac{1}{2}$, the chosen basis is $S^a = \frac{\hbar}{2} \sigma^a$, this is the fundamental representation.
- For $s = 1$, the matrices can be written to have entries $(S^a)_{bc} = i\epsilon_{abc}$, yielding the adjoint representation.

In a given representation, an interesting invariant is given by the quadratic Casimir operator,

$$\mathbf{K} = \mathbf{S} \cdot \mathbf{S},$$

which commutes with each of the representation's generators. As a result, it must be proportional to the identity in a given irreducible representation. This is a fundamental consequence of Schur's lemma on theory of representations

Lemma 1. *Schur's lemma. Let \mathbb{V} be a \mathbb{C} -vector space, associated with a finite-dimensional irreducible representation of an algebra \mathfrak{A} over \mathbb{C} . Then, let $\phi : \mathbb{V} \rightarrow \mathbb{V}$ be a homomorphism ie. $\phi(av) = a\phi(v)$, $\forall a \in \mathfrak{A}, v \in \mathbb{V}$. Then, $\phi = \lambda \text{id}_{\mathbb{V}}$.*

2.2.0.1. Thermodynamics of the Quantum Ising model. Consider the following quantum Ising model

$$\mathbf{H}_I = -Jg \sum_i \sigma_i^x - J \sum_{\langle ij \rangle} \sigma_i^z \sigma_j^z, \quad \begin{array}{l} \text{where } J > 0 \text{ is an exchange constant which} \\ \text{sets the microscopic energy scale} \\ \text{and where } g > 0 \text{ is a dimensionless coupling.} \end{array}$$

When $g = 0$, wherein \mathbf{H}_I involves only σ^z , \mathbf{H}_I will be diagonal in the basis of eigenvalues of σ_i^z , and therefore it reduces simply to the familiar classical Ising model. However, when $g \neq 0$, the σ_i^x -terms are off-diagonal in the basis of these states, thus inducing quantum-mechanical tunnelling events which flip the orientation of the Ising spin on a site. The ground state of \mathbf{H}_I can only depend upon the value of the dimensionless coupling g , thus two extreme opposing limits swiftly arise,

- If $g \gg 1$, \mathbf{H}_I 's first term dominates and, upto $\mathcal{O}\left(\frac{1}{g}\right)$, the ground state is simply

$$(2.2.0.1) \quad |\Omega\rangle = \prod_i \frac{1}{\sqrt{2}}(|+\rangle + |-\rangle) \text{ where } \frac{1}{\sqrt{2}}(|+\rangle + |-\rangle) \text{ and } \frac{1}{\sqrt{2}}(|+\rangle - |-\rangle) \text{ are } \sigma_i^x \text{'s eigenvectors with } \pm 1 \text{ eigenvalues.}$$

Then, in this case, the values of σ_i^z on different sites are totally uncorrelated in the ground state, $\langle \Omega | \sigma_i^z \sigma_j^z | \Omega \rangle = \delta_{ij}$. Perturbative correction in $\frac{1}{g}$ will build in correlation in σ^z which increase in range at each order in $\frac{1}{g}$. For large enough g and large distances $|x_i - x_j|$, these correlations are expected to remain short-ranged, yielding

$$(2.2.0.2) \quad \langle 0 | \sigma_i^z \sigma_j^z | 0 \rangle \sim e^{-\frac{|x_i - x_j|}{\zeta}}, \quad |0\rangle \text{ is the exact ground state for large } g \text{ and where } \zeta \text{ is the correlation length}$$

- For $g \ll 1$, \mathbf{H}_I 's second term dominates. At $g = 0$, the spins are either all up or down. Thus, the ground state is degenerate, the spins are either all up or down,

$$(2.2.0.3) \quad |\Omega\rangle_+ = \prod_i |+\rangle_i \text{ or } |\Omega\rangle_- = \prod_i |-\rangle_i.$$

Turning on a small g will mix in a small fraction of spins of the opposite orientation. However, for an infinite system, the degeneracy will survive at any finite order in a perturbation theory in g . This is due to an exact global \mathbb{Z}_2 -symmetry transformation, generated by the unitary operator $\prod_i \sigma_i^x$, which maps the two ground states into each other and under which \mathbf{H}_I remains invariant:

$$\sigma_i^z \rightarrow -\sigma_i^z \text{ and } \sigma_i^x \rightarrow \sigma_i^x.$$

Therefore, there is no tunnelling matrix element between the majority up and down spin sectors of the infinite system at any finite order in g . Note that establishing this degeneracy to all order in g is not the same thing as establishing its existence for any small, non-zero g . A thermodynamic system will always choose one or the other of the states as its ground states (which may be preferred by some infinitesimal external perturbation) and this is commonly referred as a spontaneous breaking of the \mathbb{Z}_2 -symmetry. As in the large g -limit, the ground states can be characterized by the behaviour of the correlations on σ_i^z . The nature of the ($g = 0$)-ground states and the small g perturbation theory results suggest that

$$(2.2.0.4) \quad \pm \langle \Omega | \sigma_i^z \sigma_j^z | \Omega \rangle_{\pm} \underset{|x_i - x_j| \rightarrow \infty}{\sim} N_0^2,$$

where $N_0 \neq 0$ is the spontaneous magnetization of the ground state. The previous statement implies that

$$\pm \langle \Omega | \sigma_i^z | \Omega \rangle_{\pm} = \pm N_0.$$

For $g = 0$, $N_0 = 1$, but quantum fluctuations at small g reduce N_0 to a smaller, but non-zero, value.

It is straightforward to see that there is no analytical function on g which transforms the ground states for large and small g into each other. Furthermore, there must be a critical value $g = g_c$ at which the large $|x_i - x_j|$ limit of the two point correlators changes from [equation \(2.2.0.4\)](#) to [equation \(2.2.0.2\)](#). Note that systems may have more than one critical point, which does not happen for \mathbf{H}_I .

- For $g > g_c$, the ground state is a quantum paramagnet and [equation \(2.2.0.2\)](#) is obeyed. As $g \rightarrow g_c$, the correlation length diverges.
- Precisely at $g = g_c$, neither [equation \(2.2.0.2\)](#) nor [equation \(2.2.0.4\)](#) are obeyed. Instead, a power-law dependence on $|x_i - x_j|$ arises at large distances.
- [equation \(2.2.0.4\)](#) is obeyed for all $g < g_c$, where the ground state is magnetically ordered. The spontaneous magnetization of the ground state, N_0 , vanishes as a power law as $g \rightarrow g_c$.

Finally, note that in a finite lattice there is necessarily a non-zero energy gap between the ground state and the first excited state. However, this energy spacing can either remain finite or approach zero in the infinite lattice limit, the two cases being identified as having gapped or gapless energy spectrums, respectively. There is a non-zero energy gap Δ for all $g \neq g_c$, but it vanishes upon approaching g_c , producing a gapless spectrum at $g = g_c$.

3. ISING AND HEISENBERG MODELS

Ising and Heisenberg Spin Models

3.0.0.1. **Spin chain type.** Some of the most common examples of spin chains are treated as follows,

- The simplest example of a $SU(2)$ -symmetric spin Hamiltonian is therefore the XXX-nearest-neighbour Heisenberg model, where

$$(3.0.0.1) \quad \mathbf{H}_{\text{XXX}} = -J \sum_{\langle ij \rangle} \mathbf{S}_i \cdot \mathbf{S}_j = -J \sum_{i=1}^N \left(\frac{1}{2} (\sigma_i^+ \sigma_{i+1}^- + \sigma_i^- \sigma_{i+1}^+) + \frac{1}{4} \sigma_i^z \sigma_{i+1}^z \right),$$

where in the ferromagnetic case the diagonal terms in the Hamiltonian favour aligned spins, while in the antiferromagnetic case, the diagonal terms favour antialignment. The XXX-Hamiltonian commutes with the magnetization operator by construction. Note that Schur's lemma does not immediately apply since the magnetization operator is reducible.

- The XXZ model is a deformation of the Heisenberg model, breaking the $SU(2)$ -model symmetry down to a $U(1)$ -subgroup. The XXZ-hamiltonian reads

$$(3.0.0.2) \quad \mathbf{H}_{\text{XXZ}} = - \sum_{\langle jk \rangle} (J_{\perp} (\mathbf{S}_j^x \mathbf{S}_k^x + \mathbf{S}_j^y \mathbf{S}_k^y) + J_z \mathbf{S}_j^z \mathbf{S}_k^z) = - \sum_{j=1}^N \left(J_{\perp} (\sigma_j^+ \sigma_{j+1}^- + \sigma_j^- \sigma_{j+1}^+) + \frac{\Delta}{2} \sigma_j^z \sigma_{j+1}^z \right),$$

which is the most general $U(1)$ -symmetric nearest-neighbor interaction for spin- $\frac{1}{2}$ particles.

Although this model is no longer $SU(2)$ -symmetric, the magnetization operator commutes with the hamiltonian, thus generating a $U(1)$ -symmetry. Even though the full $SU(2)$ -symmetry is broken, the degrees of freedom are still typically referred to as spins. The sign of J_z alone determines whether the model is ferromagnetic, $J_z > 0$, or antiferromagnetic, $J_z < 0$, since the sign of J_{\perp} is unimportant in any bipartite lattice, redefining the states by changing the overall sign of the \mathbf{S}^x and \mathbf{S}^y -operators on half the lattice sites, leaves the algebra unchanged but flips the sign of J_{\perp} . Therefore, the physically meaningful coupling is $\Delta = \frac{J_z}{|J_{\perp}|}$, so that $\Delta = \pm 1$ for the ferromagnetic and antiferromagnetic spin models, respectively. In the $\Delta \rightarrow \pm\infty$ -limit, only the J_z remains, and the model is effectively classical. For the ferromagnetic case $J_z > 0$, all the spins simply line up with the maximum value of \mathbf{M}^z . In the antiferromagnetic case, $\Delta \rightarrow -\infty$ on a bipartite lattice, the spins take their maximum opposite values on every other site.

It is easy to check that this Hamiltonian commutes with the magnetization operator, and so preserves a $U(1) \times \mathbb{Z}_2$ -symmetry. In a classical notion, the $U(1)$ -symmetry corresponds to rotations around the z -axis, while the \mathbb{Z}_2 corresponds to flipping all the spins $\mathbf{S}_j^a \rightarrow -\mathbf{S}_j^a$.

- A still more general Hamiltonian is given by the XYZ-model, which only preserves the spin-flip symmetry. None of these Hamiltonians correspond to a quantum version of the Ising model.

3.0.1. Quantum Ising Model.

3.0.2. Bethe ansatz for the XXX-model.

3.0.2.1. **Heisenberg XXX-chain.** Consider a one-dimensional lattice with N lattice sites and a $\frac{1}{2}$ -spin particle positioned at every lattice site, which have a nearest neighbour spin-spin interaction. Each particles either has spin up or spin down, generating a two-dimensional local Hilbert space, \mathbb{V}_n . Strictly speaking there are two possible topologies for a one-dimensional chain, open or closed. Consider the XXX-Hamiltonian with closed topology, given by

$$(3.0.2.1) \quad \mathbf{H} = J \sum_{n=1}^N \left(\vec{\mathbf{S}}_n \cdot \vec{\mathbf{S}}_{n+1} - \frac{1}{4} \mathbb{1}^{\otimes N} \right) = \frac{J}{4} \sum_{n=1}^N \left(\vec{\sigma}_n \cdot \vec{\sigma}_{n+1} - \frac{1}{4} \mathbb{1}^{\otimes N} \right).$$

Note that $J < 0$ models a ferromagnetic state while $J > 0$ models an anti-ferromagnetic state. A ferromagnetic state refers to a state where all the spins are aligned while an anti-ferromagnetic state refers to a state where adjacent spins are anti-aligned within the domain.

Physical Context

The main point of the Bethe solution is that XXX Heisenberg Hamiltonian, which defines the model, can be expressed in terms of a complex-valued \mathcal{R} -matrix which is a solution of the Yang-Baxter equation, thus leading to the model's integrability. The most convenient approach is to construct a transfer \mathcal{T} -matrix -a one parameter commutative family of operators acting on the full state space of the Heisenberg spin chain.

3.0.2.2. **Some fundamentals of Matrix Theory.** Let \mathbb{V} be an n -dimensional \mathbb{C} -vector space and let $k \in \mathbb{1}_n$ such that $k \in \mathbb{N}_{[0,k]}$. Let $\mathfrak{J} = \mathbb{1}_n / \mathbb{1}_k = \{k+1, \dots, n\}$. Let $A \in \text{GL}(n, \mathbb{C})$, which can be written in block notation as

$$(3.0.2.2) \quad A = \begin{pmatrix} A_{\mathcal{I}} & A_{\mathcal{I}\mathfrak{J}} \\ A_{\mathfrak{J}\mathcal{I}} & A_{\mathfrak{J}} \end{pmatrix} \begin{matrix} \mathbb{C}^k \\ \mathbb{C}^{n-k} \end{matrix} \text{ where } \begin{matrix} A_{\mathcal{I}} \in \text{GL}(k, \mathbb{C}), & A_{\mathcal{I}\mathfrak{J}} \in \text{GL}(k \rightarrow n-k, \mathbb{C}), \\ A_{\mathfrak{J}\mathcal{I}} \in \text{GL}(n-k \rightarrow k, \mathbb{C}), & A_{\mathfrak{J}} \in \text{GL}(n-k, \mathbb{C}), \end{matrix}$$

This notation is consistent with the common matrix operations. The same idea holds for writing down a block matrix representation for a linear operator with respect to a given subspace. Let $\mathcal{S} \subset \mathbb{C}^n$ be a k -dimensional vector subspace, such that $\mathbb{C}^n = \mathcal{S} \oplus \mathcal{S}^\perp$. Then, every linear operator in $\text{End}(\mathbb{C}^n)$ can be rewritten as a 2×2 -block-matrix, with its entries being matrices of a given dimension ie. ⁸

$$A \in \text{End}(\mathbb{C}^n) \rightarrow \exists! \left\{ \begin{matrix} A_{\mathcal{S}} \in \text{GL}(k, \mathbb{C}), & A_{\mathcal{S}, \mathcal{S}^\perp} \in \text{GL}(k \rightarrow n-k, \mathbb{C}), \\ A_{\mathcal{S}^\perp, \mathcal{S}} \in \text{GL}(n-k \rightarrow k, \mathbb{C}), & A_{\mathcal{S}^\perp} \in \text{GL}(n-k, \mathbb{C}), \end{matrix} \right\} /$$

$$A = \begin{pmatrix} A_{\mathcal{S}} & A_{\mathcal{S}, \mathcal{S}^\perp} \\ A_{\mathcal{S}^\perp, \mathcal{S}} & A_{\mathcal{S}^\perp} \end{pmatrix} \begin{matrix} \mathcal{S} \\ \mathcal{S}^\perp \end{matrix}$$

The main idea behind the lax-operator and \mathcal{R} -operator methods is to find an operator, eg. a product-matrix, where each one of the matrices acts non-trivially only on a single Hilbert space $\mathbb{V} \simeq \mathbb{C}^2$. Said matrices can then be rewritten in the previous form.

3.0.2.3. **The Lax operator.** Consider an XXX-spin chain with N sites and a corresponding Hilbert space given by

$$(3.0.2.3) \quad \mathbb{H} \simeq \bigotimes_{n=1}^N \mathbb{V}_n \text{ where } \mathbb{V}_n \simeq \mathbb{C}^2 \quad \begin{matrix} \text{To these individual Hilbert spaces, additional} \\ \text{auxiliary, non-physical, spaces } \mathbb{V}_a \simeq \mathbb{C}^2, \text{ can be added.} \end{matrix}$$

⁸For example, consider the orthogonal projector over \mathcal{S} , labelled $P_{\mathcal{S}}$. In this notation, its matrix representation is given by

$$P_{\mathcal{S}} = \begin{pmatrix} \mathcal{I} & 0 \\ 0 & 0 \end{pmatrix} \begin{matrix} \mathcal{S} \\ \mathcal{S}^\perp \end{matrix}$$

The algebraic Bethe ansatz' basic tool is the lax operator \mathcal{L} , whose definition is as follows.

Definition 1. The lax operator \mathcal{L} can be defined as an operator which involves the local quantum space \mathbb{V}_n and the auxiliary space \mathbb{V}_a , as follows

(3.0.2.4)

$$\mathcal{L} : \mathbb{V}_n \otimes \mathbb{V}_a \rightarrow \mathbb{V}_n \otimes \mathbb{V}_a \quad \text{where } \mathbb{1}_n \text{ and } \mathbf{S}_n^\alpha \text{ act on } \mathbb{V}_n$$

$$\mathcal{L}_{n,a} = u(\mathbb{1}_n \otimes \mathbb{1}_a) + i \sum_{\alpha} \mathbf{S}_n^\alpha \otimes \sigma_a^\alpha, \quad \forall n \leq N \quad \text{while } \mathbb{1}_a \text{ and } \mathbf{S}_a^\alpha \text{ act upon } \mathbb{V}_a \text{ and where the}$$

complex-valued u -parameter is the spectral parameter.

Note that [equation \(3.0.2.4\)](#) can be rewritten as

(3.0.2.5)

$$\mathcal{L}_{n,a} = u(\mathbb{1}_n \otimes \mathbb{1}_a) + i \sum_{\alpha} \mathbf{S}_n^\alpha \otimes \sigma_a^\alpha = \begin{pmatrix} u\mathbb{1}_n + i\mathbf{S}_n^z & i\mathbf{S}_n^x + \mathbf{S}_n^y \\ i\mathbf{S}_n^x - \mathbf{S}_n^y & u\mathbb{1}_n - i\mathbf{S}_n^z \end{pmatrix}_a = \begin{pmatrix} u\mathbb{1}_n + i\mathbf{S}_n^z & i\mathbf{S}_n^- \\ i\mathbf{S}_n^+ & u\mathbb{1}_n - i\mathbf{S}_n^z \end{pmatrix}_a,$$

which is a 2×2 -matrix in the auxiliary space \mathbb{V}_a and the matrix entries are operators acting on the physical Hilbert space \mathbb{V}_n . Consider now the following mathematical object.

Definition 2. Let the permutation operator be

$$(3.0.2.6) \quad \mathcal{P}_{n,a} = \frac{1}{2} \left(\mathbb{1}_n \otimes \mathbb{1}_a + \sum_{\alpha} \sigma_n^\alpha \otimes \sigma_n^\alpha \right) = \frac{1}{2} \sigma^{n\mu} \sigma_\mu^\alpha$$

which ⁹ acts on both the physical and auxiliary systems.

Using this definition for the permutation operator, the lax operator can then be rewritten as

$$(3.0.2.8) \quad \begin{aligned} \mathcal{L}_{n,a} &= u(\mathbb{1}_n \otimes \mathbb{1}_a) + i \sum_{\alpha} \mathbf{S}_n^\alpha \otimes \sigma_a^\alpha = u(\mathbb{1}_n \otimes \mathbb{1}_a) + \frac{i}{2} \sum_{\alpha} \sigma_n^\alpha \otimes \sigma_a^\alpha \\ &= u(\mathbb{1}_n \otimes \mathbb{1}_a) - \frac{i}{2} (\mathbb{1}_n \otimes \mathbb{1}_a) + \frac{i}{2} (\mathbb{1}_n \otimes \mathbb{1}_a) + \frac{i}{2} \sum_{\alpha} \sigma_n^\alpha \otimes \sigma_a^\alpha \\ &= \left(u - \frac{i}{2} \right) (\mathbb{1}_n \otimes \mathbb{1}_a) + \frac{i}{2} \left(\mathbb{1}_n \otimes \mathbb{1}_a + \sum_{\alpha} \sigma_n^\alpha \otimes \sigma_n^\alpha \right) \\ &= \left(u - \frac{i}{2} \right) \mathbb{1}_{n,a} + i \mathcal{P}_{n,a}. \end{aligned}$$

3.0.2.4. **Fundamental commutation Relations and Monodromy Matrix.** Consider the following two lax operators

$$\mathcal{L}_{n,a_1}(u_1) : \mathbb{V}_n \otimes \mathbb{V}_{a_1} \rightarrow \mathbb{V}_n \otimes \mathbb{V}_{a_1} \quad \text{both of which act on the physical quantum}$$

$$\mathcal{L}_{n,a_2}(u_2) : \mathbb{V}_n \otimes \mathbb{V}_{a_2} \rightarrow \mathbb{V}_n \otimes \mathbb{V}_{a_2} \quad \text{state and on two different auxiliary spaces as well¹⁰}$$

The product of these two operator is then a triple tensor product in $\mathbb{V}_n \otimes \mathbb{V}_{a_1} \otimes \mathbb{V}_{a_2}$. Now, the following matrix arises

⁹Note that the permutation operator can be written as a 2×2 -matrix in the auxiliary space \mathbb{V}_a ,

$$(3.0.2.7) \quad \begin{aligned} \mathcal{P}_{n,a} &= \frac{1}{2} \left[\begin{pmatrix} \mathbb{1}_n & 0 \\ 0 & \mathbb{1}_a \end{pmatrix} + \sigma_n^x \otimes \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}_a + \sigma_n^y \otimes \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}_a + \sigma_n^z \otimes \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}_a \right] \\ &= \begin{pmatrix} \mathbb{1}_n + \sigma_n^z & \sigma_n^x - i\sigma_n^y \\ \sigma_n^x + i\sigma_n^y & \mathbb{1}_n - \sigma_n^z \end{pmatrix}_a = \end{aligned}$$

Definition 3. Consider an \mathcal{R} -matrix which relates the permutation and spectral parameters of two auxiliary spaces, which is given by

$$(3.0.2.10) \quad \mathcal{R} : \mathbb{V}_{a_1} \otimes \mathbb{V}_{a_2} \otimes \mathbb{V}_{a_1} \otimes \mathbb{V}_{a_2}$$

$$\mathcal{R}_{a_1, a_2}(u_1 - u_2) = (u_1 - u_2)\mathbb{1}_{a_1, a_2} + i\mathcal{P}_{a_1, a_2},$$

which doesn't act on the physical system at all, but on the auxiliary, phantom, systems¹¹.

An interesting relationship between the products of the two lax operators and the \mathcal{R} -matrix can then be found¹², consider

$$\begin{aligned} & \mathcal{R}_{a_1, a_2}(u_1 - u_2)\mathcal{L}_{n, a_1}(u_1)\mathcal{L}_{n, a_2}(u_2) \\ &= \left((u_1 - u_2)\mathbb{1}_{a_1, a_2} + i\mathcal{P}_{a_1, a_2} \right) \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} + i\mathcal{P}_{n, a_1} \right) \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i\mathcal{P}_{n, a_2} \right) \\ &= \left((u_1 - u_2)\mathbb{1}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} + (u_1 - u_2)\mathbb{1}_{a_1, a_2} i\mathcal{P}_{n, a_1} + i\mathcal{P}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} + i^2 \mathcal{P}_{a_1, a_2} \mathcal{P}_{n, a_1} \right) \\ &\quad \times \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i\mathcal{P}_{n, a_2} \right) \\ &= (u_1 - u_2)\mathbb{1}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + (u_1 - u_2)\mathbb{1}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} i\mathcal{P}_{n, a_2} \\ &\quad + (u_1 - u_2)\mathbb{1}_{a_1, a_2} i\mathcal{P}_{n, a_1} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + (u_1 - u_2)\mathbb{1}_{a_1, a_2} i\mathcal{P}_{n, a_1} i\mathcal{P}_{n, a_2} \\ &\quad + i\mathcal{P}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i\mathcal{P}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} i\mathcal{P}_{n, a_2} \\ &\quad - \mathcal{P}_{a_1, a_2} \mathcal{P}_{n, a_1} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} - \mathcal{P}_{a_1, a_2} \mathcal{P}_{n, a_1} i\mathcal{P}_{n, a_2} \\ &= (u_1 - u_2)\mathbb{1}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} \times \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} + (u_1 - u_2)\mathbb{1}_{a_1, a_2} i\mathcal{P}_{n, a_2} \times \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} \\ &\quad + (u_1 - u_2)\mathbb{1}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} \times i\mathcal{P}_{n, a_1} + (u_1 - u_2)\mathbb{1}_{a_1, a_2} i\mathcal{P}_{n, a_2} \times i\mathcal{P}_{n, a_1} \\ &\quad + i\mathcal{P}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i\mathcal{P}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} i\mathcal{P}_{n, a_2} \\ &\quad - \mathcal{P}_{a_1, a_2} \mathcal{P}_{n, a_1} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} - \mathcal{P}_{a_1, a_2} \mathcal{P}_{n, a_1} i\mathcal{P}_{n, a_2} \end{aligned}$$

¹¹Note that \mathcal{R} can be rewritten as a 2×2 -matrix in either of the auxiliary spaces, as follows

$$(3.0.2.11) \quad \mathcal{R}_{a_1, a_2} = \begin{pmatrix} \left(u + \frac{i}{2} \right) \mathbb{1}_{a_2} + i\mathbf{S}_{a_2}^z & i\mathbf{S}_{a_2}^- \\ i\mathbf{S}_{a_2}^+ & \left(u + \frac{i}{2} \right) \mathbb{1}_{a_2} - i\mathbf{S}_{a_2}^z \end{pmatrix}_{a_1} = \begin{pmatrix} \left(u + \frac{i}{2} \right) \mathbb{1}_{a_1} + i\mathbf{S}_{a_1}^z & i\mathbf{S}_{a_1}^- \\ i\mathbf{S}_{a_1}^+ & \left(u + \frac{i}{2} \right) \mathbb{1}_{a_1} - i\mathbf{S}_{a_1}^z \end{pmatrix}_{a_2}.$$

¹²Note that the following permutation properties hold

$$(3.0.2.12) \quad \mathcal{P}_{n, a_1} \mathcal{P}_{n, a_2} = \mathcal{P}_{a_1, a_2} \mathcal{P}_{n, a_1} = \mathcal{P}_{n, a_2} \mathcal{P}_{a_2, a_1} \text{ and } \mathcal{P}_{a, b} = \mathcal{P}_{b, a}$$

$$\begin{aligned}
&= \left((u_1 - u_2) \mathbb{1}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + (u_1 - u_2) \mathbb{1}_{a_1, a_2} i \mathcal{P}_{n, a_2} \right) \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} \\
&+ \left((u_1 - u_2) \mathbb{1}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + (u_1 - u_2) \mathbb{1}_{a_1, a_2} i \mathcal{P}_{n, a_2} \right) \times i \mathcal{P}_{n, a_1} \\
&+ i \mathcal{P}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i \mathcal{P}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} i \mathcal{P}_{n, a_2} \\
&- \mathcal{P}_{a_1, a_2} \mathcal{P}_{n, a_1} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} - \mathcal{P}_{a_1, a_2} \mathcal{P}_{n, a_1} i \mathcal{P}_{n, a_2} \\
&= \left((u_1 - u_2) \mathbb{1}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + (u_1 - u_2) \mathbb{1}_{a_1, a_2} i \mathcal{P}_{n, a_2} \right) \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i \mathcal{P}_{n, a_1} \right) \\
&+ i \mathcal{P}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i \mathcal{P}_{a_1, a_2} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_1} i \mathcal{P}_{n, a_2} \\
&- \mathcal{P}_{a_1, a_2} \mathcal{P}_{n, a_1} \left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} - \mathcal{P}_{a_1, a_2} \mathcal{P}_{n, a_1} i \mathcal{P}_{n, a_2} \\
&= \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i \mathcal{P}_{n, a_2} \right) \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i \mathcal{P}_{n, a_1} \right) (u_1 - u_2) \mathbb{1}_{a_1, a_2} \\
&+ \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i \mathcal{P}_{n, a_2} \right) \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i \mathcal{P}_{n, a_1} \right) i \mathcal{P}_{n, a_2} \\
&= \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i \mathcal{P}_{n, a_2} \right) \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i \mathcal{P}_{n, a_1} \right) \left(\left(u - \frac{i}{2} \right) \mathbb{1}_{n, a_2} + i \mathcal{P}_{n, a_2} \right),
\end{aligned}$$

which proves that

$$(3.0.2.13) \quad \mathcal{R}_{a_1, a_2}(u_1 - u_2) \mathcal{L}_{n, a_1}(u_1) \mathcal{L}_{n, a_2}(u_2) = \mathcal{L}_{n, a_1}(u_2) \mathcal{L}_{n, a_1}(u_1) \mathcal{R}_{a_1, a_2}(u_1 - u_2).$$

From the fundamental commutation relation it can then be shown that the \mathcal{R} -matrix follows the quantum Yang-Baxter equation. Consider then the product of the following \mathcal{L} -operators,

$$\begin{aligned}
(3.0.2.14) \quad \mathcal{L}_{n, 1} \mathcal{L}_{n, 2} \mathcal{L}_{n, 3} &= \mathcal{R}_{12}^{-1} \mathcal{L}_{n, 2} \mathcal{L}_{n, 1} \mathcal{L}_{n, 3} \mathcal{R}_{12} \\
&= \mathcal{R}_{12}^{-1} \mathcal{R}_{13}^{-1} \mathcal{L}_{n, 2} \mathcal{L}_{n, 3} \mathcal{L}_{n, 1} \mathcal{R}_{13} \mathcal{R}_{12} \\
&= \mathcal{R}_{12}^{-1} \mathcal{R}_{13}^{-1} \mathcal{R}_{23}^{-1} \mathcal{L}_{n, 3} \mathcal{L}_{n, 2} \mathcal{L}_{n, 1} \mathcal{R}_{23} \mathcal{R}_{13} \mathcal{R}_{12} \\
&= (\mathcal{R}_{23} \mathcal{R}_{13} \mathcal{R}_{12})^{-1} \mathcal{L}_{n, 3} \mathcal{L}_{n, 2} \mathcal{L}_{n, 1} \mathcal{R}_{23} \mathcal{R}_{13} \mathcal{R}_{12},
\end{aligned}$$

and similarly

$$\begin{aligned}
(3.0.2.15) \quad \mathcal{L}_{n, 1} \mathcal{L}_{n, 2} \mathcal{L}_{n, 3} &= \mathcal{R}_{23}^{-1} \mathcal{L}_{n, 1} \mathcal{L}_{n, 3} \mathcal{L}_{n, 2} \mathcal{R}_{23} \\
&= \mathcal{R}_{23}^{-1} \mathcal{R}_{13}^{-1} \mathcal{L}_{n, 3} \mathcal{L}_{n, 1} \mathcal{L}_{n, 2} \mathcal{R}_{13} \mathcal{R}_{23} \\
&= \mathcal{R}_{23}^{-1} \mathcal{R}_{13}^{-1} \mathcal{R}_{12}^{-1} \mathcal{L}_{n, 3} \mathcal{L}_{n, 2} \mathcal{L}_{n, 1} \mathcal{R}_{12} \mathcal{R}_{13} \mathcal{R}_{23} \\
&= (\mathcal{R}_{12} \mathcal{R}_{13} \mathcal{R}_{23})^{-1} \mathcal{L}_{n, 3} \mathcal{L}_{n, 2} \mathcal{L}_{n, 1} \mathcal{R}_{12} \mathcal{R}_{13} \mathcal{R}_{23},
\end{aligned}$$

both of which then yield an important constraint on the \mathcal{R} -matrix, namely

$$(3.0.2.16) \quad \mathcal{R}_{12} \mathcal{R}_{13} \mathcal{R}_{23} = \mathcal{R}_{23} \mathcal{R}_{13} \mathcal{R}_{12},$$

which is the quantum Yang-Baxter equation.

Definition 4. Let the monodromy matrix $\mathcal{T}_{N,a}$ be defined as a product of subsequent of \mathcal{L} -operators, ie.

$$(3.0.2.17) \quad \mathcal{T}_{N,a} = \prod_1^{j=N} \mathcal{L}_{j,a}$$

$$\mathcal{T}_{N,a} = \prod_1^{j=N} \begin{pmatrix} u\mathbb{1}_j + i\mathbf{S}_j^z & i\mathbf{S}_j^- \\ i\mathbf{S}_j^+ & u\mathbb{1}_j - i\mathbf{S}_j^z \end{pmatrix} \equiv \begin{pmatrix} A(u) & B(u) \\ C(u) & D(u) \end{pmatrix},$$

This matrix obeys the fundamental relation as well¹³, ie.

$$(3.0.2.19) \quad \mathcal{R}_{a_1,a_2}(u_1 - u_2) \mathcal{T}_{N,a_1}(u_1) \mathcal{T}_{N,a_2}(u_2) = \mathcal{T}_{N,a_2}(u_2) \mathcal{T}_{N,a_1}(u_1) \mathcal{R}_{a_1,a_2}(u_1 - u_2).$$

Definition 5. The transfer t-matrix is defined as the \mathcal{T} -matrix's partial trace over the auxiliary space \mathbb{V}_a , ie.

$$(3.0.2.20) \quad \mathfrak{t}(u) = \text{Tr}_{\mathbb{V}_a} \mathcal{T}_{N,a}(u) = \text{Tr}_{\mathbb{V}_a} \begin{pmatrix} A(u) & B(u) \\ C(u) & D(u) \end{pmatrix} = A(u) + D(u),$$

where \mathfrak{t} , A , B , C , D all are 2×2 -matrices in the physical system, the auxiliary systems being present no more.

With the previous definitions in mind, equation (3.0.2.19)'s double trace can then be written in terms of the transfer matrix, as follows

$$\begin{aligned} & \text{Tr}_{\mathbb{V}_{a_1}} \text{Tr}_{\mathbb{V}_{a_2}} \left[\mathcal{R}_{a_1,a_2}(u_1 - u_2) \mathcal{T}_{N,a_1}(u_1) \mathcal{T}_{N,a_2}(u_2) \right] \\ & \Rightarrow \text{Tr}_{\mathbb{V}_{a_1}} \text{Tr}_{\mathbb{V}_{a_2}} \left[\mathcal{R}_{a_1,a_2}(u_1 - u_2) \mathcal{T}_{N,a_1}(u_1) \mathcal{T}_{N,a_2}(u_2) - \mathcal{T}_{N,a_1}(u_2) \mathcal{T}_{N,a_1}(u_1) \mathcal{R}_{a_1,a_2}(u_1 - u_2) \right] = 0 \\ & \Rightarrow \text{Tr}_{\mathbb{V}_{a_1}} \text{Tr}_{\mathbb{V}_{a_2}} \left[\mathcal{R}_{a_1,a_2}(u_1 - u_2) \left(\mathcal{T}_{N,a_1}(u_1) \mathcal{T}_{N,a_2}(u_2) - \mathcal{T}_{N,a_2}(u_2) \mathcal{T}_{N,a_1}(u_1) \right) \right] = 0 \end{aligned}$$

which implies $\mathfrak{t}(u_1)\mathfrak{t}(u_2) = \mathfrak{t}(u_2)\mathfrak{t}(u_1)$.

Therefore the transfer matrix commutes with itself for different values of the spectral parameter. It turns out that the monodromy matrix can be expanded in a power series around any point $z_0 \in \mathbb{C}$, generating an infinite set of linearly independent commuting operators acting on the full quantum space. This formally leads to the model's integrability.

¹³For example, consider $N = 2$, then

$$(3.0.2.18) \quad \begin{aligned} \mathcal{R}_{a_1,a_2}(u_1 - u_2) \mathcal{T}_{N,a_1}(u_1) \mathcal{T}_{N,a_2}(u_2) &= \mathcal{R}_{a_1,a_2}(u_1 - u_2) \mathcal{T}_{2,a_1}(u_1) \mathcal{T}_{2,a_2}(u_2) \\ &= \mathcal{R}_{a_1,a_2}(u_1 - u_2) \mathcal{L}_{2,a_1}(u_1) \mathcal{L}_{2,a_1}(u_1) \mathcal{L}_{2,a_2}(u_2) \mathcal{L}_{1,a_2}(u_2) \\ &= \left(\mathcal{R}_{a_1,a_2}(u_1 - u_2) \mathcal{L}_{2,a_1}(u_1) \mathcal{L}_{2,a_2}(u_2) \right) \mathcal{L}_{1,a_1}(u_1) \mathcal{L}_{1,a_2}(u_2) \\ &= \mathcal{L}_{2,a_2}(u_2) \mathcal{L}_{2,a_1}(u_1) \left(\mathcal{R}_{a_1,a_2}(u_1 - u_2) \mathcal{L}_{1,a_1}(u_1) \mathcal{L}_{1,a_2}(u_2) \right) \\ &= \mathcal{L}_{2,a_2}(u_2) \mathcal{L}_{2,a_1}(u_1) \mathcal{L}_{1,a_2}(u_2) \mathcal{L}_{1,a_1}(u_1) \mathcal{R}_{a_1,a_2}(u_1 - u_2) \\ &= \mathcal{L}_{2,a_2}(u_2) \mathcal{L}_{1,a_2}(u_2) \mathcal{L}_{2,a_1}(u_1) \mathcal{L}_{1,a_1}(u_1) \mathcal{R}_{a_1,a_2}(u_1 - u_2) \\ &= \mathcal{T}_{n,a_1}(u_2) \mathcal{T}_{n,a_1}(u_1) \mathcal{R}_{a_1,a_2}(u_1 - u_2), \end{aligned}$$

where the fundamental commutation relation given by (3.0.2.13) was used. Note as well that \mathcal{L} -operators acting on different spaces commute as well, eg. $[\mathcal{L}_{n_1,a_1}, \mathcal{L}_{n_2,a_2}] \propto \delta_{n_1,n_2}$. The general fundamental relation can then be proved via mathematical induction for the general case.

3.0.2.5. Monodromy matrix and the Hamiltonian. The Hamiltonian operator belongs to the family of transfer matrices. Let $z_0 = \frac{i}{2}$, then the monodromy matrix can be expanded around said complex point yielding (up to first order)

$$\begin{aligned}
 \mathcal{T}_{N,a}(z_0) &= \prod_1^{j=N} \left[\cancel{\left(z_0 - \frac{i}{2}\right)} \mathbb{1}_{j,a} + i\mathcal{P}_{j,a} \right] = i^N \mathcal{P}_{N,a} \mathcal{P}_{N-1,a} \cdots \mathcal{P}_{1,a} = \\
 (3.0.2.21) \quad &= i^N \mathcal{P}_{1,2} \mathcal{P}_{2,3} \cdots \mathcal{P}_{N-1,N} \mathcal{P}_{N,a} \\
 &= i^N \prod_{j=1}^{N-1} \mathcal{P}_{j,j+1} \mathcal{P}_{N,a}
 \end{aligned}$$

where the second line holds given the properties of the permutation operators. Taking the partial trace over the auxiliary space yields the transfer matrix, as follows

$$\begin{aligned}
 \mathfrak{t}(z_0) &= \text{Tr}_{\mathbb{V}_a} \mathcal{T}_{N,a}(z_0) = \text{Tr}_{\mathbb{V}_a} i^N \mathcal{P}_{1,2} i^N \prod_{j=1}^{N-1} \mathcal{P}_{j,j+1} \mathcal{P}_{N,a} \\
 (3.0.2.22) \quad &= i^N \prod_{j=1}^{N-1} \mathcal{P}_{j,j+1} \text{Tr}_{\mathbb{V}_a} \mathcal{P}_{N,a} = i^N \prod_{j=1}^{N-1} \mathcal{P}_{j,j+1} \text{Tr}_{\mathbb{V}_a} \frac{1}{2} \begin{pmatrix} \mathbb{1}_n + \sigma_n^z & \sigma_n^x - i\sigma_n^y \\ \sigma_n^x + i\sigma_n^y & \mathbb{1}_n - \sigma_n^z \end{pmatrix}_a \\
 &= i^N \prod_{j=1}^{N-1} \mathcal{P}_{j,j+1} \mathbb{1}_N.
 \end{aligned}$$

Then $\mathcal{U} = i^{-N} \mathfrak{t}(z_0) = \prod_{j=1}^{N-1} \mathcal{P}_{j,j+1}$ is a shift operator in the full quantum Hilbert space \mathbb{H} . In effect, note that

$$\mathcal{P}_{n_1, n_2} \mathbf{X}_{n_2} \mathcal{P}_{n_1, n_2} = \mathbf{X}_{n_1},$$

ie. this permutation moves \mathbf{X} one step back. Furthermore, this operator is unitary (given that the permutations are, by definition, unitary operators). Then, according to Stone's theorem on one-parameter unitary groups, the shift operator is related to the momentum operator

$$(3.0.2.23) \quad \mathcal{U} = e^{iP}$$

The next order in the expansion of the transfer matrix can then be found via the derivative of the monodromy matrix at $z_0 = \frac{i}{2}$,

$$\begin{aligned}
\frac{d\mathcal{T}_{N,a}(u)}{du}\Big|_{u=z_0} &= \frac{d}{du} \prod_1^{j=N} \begin{pmatrix} u\mathbb{1}_j + i\mathbf{S}_j^z & i\mathbf{S}_j^- \\ i\mathbf{S}_j^+ & u\mathbb{1}_j - i\mathbf{S}_j^z \end{pmatrix} \Big|_{u=z_0} \\
&= \frac{d}{du} \left[\begin{pmatrix} u\mathbb{1}_N + i\mathbf{S}_N^z & i\mathbf{S}_N^- \\ i\mathbf{S}_N^+ & u\mathbb{1}_N - i\mathbf{S}_N^z \end{pmatrix} \right] \Big|_{u=z_0} \prod_1^{j=N-1} \begin{pmatrix} u\mathbb{1}_j + i\mathbf{S}_j^z & i\mathbf{S}_j^- \\ i\mathbf{S}_j^+ & u\mathbb{1}_j - i\mathbf{S}_j^z \end{pmatrix} \\
&\quad + \begin{pmatrix} u\mathbb{1}_N + i\mathbf{S}_N^z & i\mathbf{S}_N^- \\ i\mathbf{S}_N^+ & u\mathbb{1}_N - i\mathbf{S}_N^z \end{pmatrix} \frac{d}{du} \left[\begin{pmatrix} u\mathbb{1}_{N-1} + i\mathbf{S}_{N-1}^z & i\mathbf{S}_{N-1}^- \\ i\mathbf{S}_{N-1}^+ & u\mathbb{1}_{N-1} - i\mathbf{S}_{N-1}^z \end{pmatrix} \right] \Big|_{u=z_0} \\
&\quad \times \prod_1^{j=N-2} \begin{pmatrix} u\mathbb{1}_j + i\mathbf{S}_j^z & i\mathbf{S}_j^- \\ i\mathbf{S}_j^+ & u\mathbb{1}_j - i\mathbf{S}_j^z \end{pmatrix} \\
&\quad + \cdots + \prod_3^{j=N} \begin{pmatrix} u\mathbb{1}_j + i\mathbf{S}_j^z & i\mathbf{S}_j^- \\ i\mathbf{S}_j^+ & u\mathbb{1}_j - i\mathbf{S}_j^z \end{pmatrix} \frac{d}{du} \left[\begin{pmatrix} u\mathbb{1}_2 + i\mathbf{S}_2^z & i\mathbf{S}_2^- \\ i\mathbf{S}_2^+ & u\mathbb{1}_2 - i\mathbf{S}_2^z \end{pmatrix} \right] \Big|_{u=z_0} \\
&\quad \times \begin{pmatrix} u\mathbb{1}_1 + i\mathbf{S}_1^z & i\mathbf{S}_1^- \\ i\mathbf{S}_1^+ & u\mathbb{1}_1 - i\mathbf{S}_1^z \end{pmatrix} \\
&\quad + \prod_2^{j=N} \begin{pmatrix} u\mathbb{1}_j + i\mathbf{S}_j^z & i\mathbf{S}_j^- \\ i\mathbf{S}_j^+ & u\mathbb{1}_j - i\mathbf{S}_j^z \end{pmatrix} \frac{d}{du} \left[\begin{pmatrix} u\mathbb{1}_1 + i\mathbf{S}_1^z & i\mathbf{S}_1^- \\ i\mathbf{S}_1^+ & u\mathbb{1}_1 - i\mathbf{S}_1^z \end{pmatrix} \right] \Big|_{u=z_0} \\
&= i^{N-1} \left[\prod_1^{j=N-1} \mathcal{P}_j + \prod_N^{j=N} \mathcal{P}_j \times \prod_1^{j=N-2} \mathcal{P}_j + \cdots + \prod_k^{j=N} \mathcal{P}_j \times \prod_1^{k+2} \mathcal{P}_j \right. \\
&\quad \left. + \cdots + \prod_3^{j=N} \mathcal{P}_j \times \prod_1^{j=1} \mathcal{P}_j + \prod_2^{j=N} \mathcal{P}_j \right] \\
&= i^{N-1} \sum_{n \in \mathbb{N}} \mathcal{P}_{N,a} \cdots \mathcal{P}_{n-1,a} \mathcal{P}_{n+1,a} \cdots \mathcal{P}_{1,a} \\
&= i^{N-1} \sum_{n \in \mathbb{N}} \mathcal{P}_{1,2} \mathcal{P}_{2,3} \cdots \mathcal{P}_{n-1,n+1} \cdots \mathcal{P}_{N-1,N} \mathcal{P}_{N,a} \\
&= i^{N-1} \sum_{n \in \mathbb{N}} \prod_{j=1}^{n-2} \mathcal{P}_{j,j+1} \times \mathcal{P}_{n-1,n+1} \times \prod_{j=n+2}^{N-1} \mathcal{P}_{j,j+1} \mathcal{P}_{N,a}.
\end{aligned}$$

Therefore,

$$\begin{aligned}
(3.0.2.24) \quad \frac{d\mathbf{t}(u)}{du}\Big|_{u=z_0} &= \frac{d \operatorname{Tr}_{\mathbb{V}_a} \mathcal{T}_{N,a}(u)}{du}\Big|_{u=z_0} \\
&= i^{N-1} \sum_{n \in \mathbb{N}} \prod_{j=1}^{n-2} \mathcal{P}_{j,j+1} \times \mathcal{P}_{n-1,n+1} \times \prod_{j=n+2}^{N-1} \mathcal{P}_{j,j+1} \times \operatorname{Tr}_{\mathbb{V}_a} \mathcal{P}_{N,a} \\
&= i^{N-1} \sum_{n \in \mathbb{N}} \prod_{j=1}^{n-2} \mathcal{P}_{j,j+1} \times \mathcal{P}_{n-1,n+1} \times \prod_{j=n+2}^{N-1} \mathcal{P}_{j,j+1}.
\end{aligned}$$

Now, note that

(3.0.2.25)

$$\begin{aligned}
\frac{d}{du} \log(\mathfrak{t}(u)) \Big|_{u=z_0} &= \frac{d}{du} \log \operatorname{Tr}_{\mathbb{V}_a} \mathcal{T}(u) \Big|_{u=z_0} \\
&= \frac{d}{du} \left(\mathfrak{t}(u) \mathfrak{t}(u)^{-1} \right) \Big|_{u=z_0} \\
&= i^{N-1} \sum_{n \in \mathbb{N}} \left(\prod_{j=1}^{n-2} \mathcal{P}_{j,j+1} \times \mathcal{P}_{n-1,n+1} \times \prod_{j=n+2}^{N-1} \mathcal{P}_{j,j+1} \right) \times \left(i^N \prod_{j=1}^{N-1} \mathcal{P}_{j,j+1} \mathbb{1}_N \right)^{-1} \\
&= \frac{1}{i} \sum_{n \in \mathbb{N}} \left(\prod_{j=1}^{n-2} \mathcal{P}_{j,j+1} \times \mathcal{P}_{n-1,n+1} \times \prod_{j=n+2}^{N-1} \mathcal{P}_{j,j+1} \right) \times \prod_{j=1}^{N-1} \mathcal{P}_{N-1-j,N-j}^{-1} \\
&= \frac{1}{i} \sum_{n \in \mathbb{N}} \left(\prod_{j=1}^{n-2} \mathcal{P}_{j,j+1} \times \mathcal{P}_{n-1,n+1} \times \prod_{j=n+2}^{N-1} \mathcal{P}_{j,j+1} \right) \times \prod_{j=1}^{N-1} \mathcal{P}_{j,j+1} \\
&= \frac{1}{i} \sum_{n \in \mathbb{N}} \left(\cancel{\mathcal{P}_{1,2}} \cancel{\mathcal{P}_{2,3}} \cdots \mathcal{P}_{n-1,n+1} \cdots \cancel{\mathcal{P}_{N-1,N}} \times \cancel{\mathcal{P}_{N-1,N}} \cdots \mathcal{P}_{n,n+1} \cdots \cancel{\mathcal{P}_{2,3}} \cancel{\mathcal{P}_{1,2}} \right) \\
&= \frac{1}{i} \sum_{n \in \mathbb{N}} \mathcal{P}_{N-1,N}
\end{aligned}$$

Note that (3.0.2.1) can be rewritten in terms of the permutation operator as

$$(3.0.2.26) \quad \mathbf{H} = \frac{J}{2} \sum_n \left(\mathcal{P}_{n,n+1} - \mathbb{1}^{\otimes N} \right),$$

which in turn can be rewritten as

$$(3.0.2.27) \quad \mathbf{H} = \frac{J}{2} \left(i \frac{d}{du} \log(\mathfrak{t}(u)) \Big|_{u=z_0} - N \mathbb{1}^{\otimes N} \right),$$

which shows that the XXX Hamiltonian belongs to the family of $N - 1$ commuting operators generated by the trace of the monodromy matrix \mathcal{T} . As a result, the Hamiltonian commutes with the transfer matrix $[\mathbf{H}, \mathfrak{t}(u)] = 0$.

3.0.2.6. Diagonalizing the Hamiltonian. The only task left is to diagonalize the Hamiltonian, by diagonalizing the transfer matrix. The monodromy matrix $\mathfrak{t}_{N,a}$ is a 2×2 -matrix in the physical space \mathbb{V}_a . Recalling the fundamental commutation relation between the \mathcal{R} and \mathcal{T} -matrices, then the following theorem holds,

Theorem 1 *Let \mathcal{T} be the monodromy matrix for the Heisenberg model. Then, $\forall u_1, u_2 \in \mathbb{C}$ then*

- $[\mathcal{T}_{ij}(u_1), \mathcal{T}_{ij}(u_2)] = 0$, $i, j = 1, 2$, which entails that $[B(u_1), B(u_2)] = \cdots = [D(u_1), D(u_2)] = 0$
- $A(u_1)B(u_2) = f(u_1 - u_2)B(u_2)A(u_1) + g(u_1 - u_2)B(u_1)A(u_2)$
- $D(u_1)B(u_2) = h(u_1 - u_2)B(u_2)D(u_1) + k(u_1 - u_2)B(u_1)D(u_2)$,

where f, g, h , and k are given by

$$f(u) = \frac{u - i}{u}, \quad g(u) = \frac{i}{u}, \quad h(u) = \frac{u + i}{u}, \quad k(u) = -\frac{i}{u}.$$

Proof. Consider two different monodromy matrices, labelled $\mathcal{T}_{N,a_1}(u_1)$ and $\mathcal{T}_{N,a_2}(u_2)$, acting on two different auxiliary spaces, labelled \mathbb{V}_1 and \mathbb{V}_2 respectively. Then, these \mathcal{T} -matrices can be written as

$$\begin{aligned}
(3.0.2.28) \quad \mathcal{T}_{N,a_1}(u_1) &= \begin{pmatrix} A(u_1) & B(u_1) \\ C(u_1) & D(u_1) \end{pmatrix} \otimes \mathbb{1}_{V_2}, \quad \mathcal{T}_{N,a_2}(u_2) = \mathbb{1}_{V_1} \otimes \begin{pmatrix} A(u_2) & B(u_2) \\ C(u_2) & D(u_2) \end{pmatrix} \\
&= \begin{pmatrix} A(u_1) & 0 & B(u_1) & 0 \\ 0 & A(u_1) & 0 & B(u_1) \\ C(u_1) & 0 & D(u_1) & 0 \\ 0 & C(u_1) & 0 & D(u_1) \end{pmatrix}, \quad \begin{pmatrix} A(u_2) & B(u_2) & 0 & 0 \\ C(u_2) & D(u_2) & 0 & 0 \\ 0 & 0 & A(u_2) & B(u_2) \\ 0 & 0 & C(u_2) & D(u_2) \end{pmatrix}
\end{aligned}$$

Recall that the \mathcal{R} -matrix is given by (3.0.2.10), which can be rewritten as

$$\mathcal{R}_{a_1,a_2}(u_1 - u_2) = \begin{pmatrix} a(u_1 - u_2) & 0 & 0 & 0 \\ 0 & b(u_1 - u_2) & c(u_1 - u_2) & 0 \\ 0 & c(u_1 - u_2) & b(u_1 - u_2) & 0 \\ 0 & 0 & 0 & a(u_1 - u_2) \end{pmatrix} \text{ where } \begin{cases} a(\lambda) = \lambda + i \\ b(\lambda) = \lambda \\ c(\lambda) = i \end{cases},$$

Then the fundamental commutation relation (3.0.2.19) establishes a relationship between the monodromy \mathcal{T} -matrices and the \mathcal{R} -matrix. Then, plugging in the previous results yields

$$\begin{aligned}
& \begin{pmatrix} a_{12} & 0 & 0 & 0 \\ 0 & b_{12} & c_{12} & 0 \\ 0 & c_{12} & b_{12} & 0 \\ 0 & 0 & 0 & a_{12} \end{pmatrix} \begin{pmatrix} A(u_1)A(u_2) & A(u_1)B(u_2) & B(u_1)A(u_2) & B(u_1)B(u_2) \\ A(u_1)C(u_2) & A(u_1)D(u_2) & B(u_1)C(u_2) & B(u_1)D(u_2) \\ C(u_1)A(u_2) & C(u_1)B(u_2) & D(u_1)A(u_2) & D(u_1)B(u_2) \\ C(u_1)C(u_2) & C(u_1)D(u_2) & D(u_1)C(u_2) & D(u_1)D(u_2) \end{pmatrix} \\
&= \begin{pmatrix} A(u_2)A(u_1) & B(u_2)A(u_1) & A(u_2)B(u_1) & B(u_2)B(u_1) \\ C(u_2)A(u_1) & D(u_2)A(u_1) & C(u_2)B(u_1) & D(u_2)B(u_1) \\ A(u_2)C(u_1) & B(u_2)C(u_1) & A(u_2)D(u_1) & B(u_2)D(u_1) \\ C(u_2)C(u_1) & D(u_2)C(u_1) & C(u_2)D(u_1) & D(u_2)D(u_1) \end{pmatrix} \begin{pmatrix} a_{12} & 0 & 0 & 0 \\ 0 & b_{12} & c_{12} & 0 \\ 0 & c_{12} & b_{12} & 0 \\ 0 & 0 & 0 & a_{12} \end{pmatrix} \\
&\Rightarrow \\
&\Rightarrow \begin{pmatrix} a_{12}A_1A_2 & a_{12}A_1B_2 & a_{12}B_1A_2 & a_{12}B_1B_2 \\ b_{12}A_1C_2 + c_{12}C_1A_2 & b_{12}A_1D_2 + c_{12}C_1B_2 & b_{12}B_1C_2 + c_{12}D_1A_2 & b_{12}B_1D_2 + c_{12}D_1B_2 \\ c_{12}A_1C_2 + b_{12}C_1A_2 & b_{12}A_1D_2 + c_{12}C_1B_2 & c_{12}B_1C_2 + b_{12}D_1A_2 & c_{12}B_1D_2 + b_{12}D_1B_2 \\ a_{12}C_1C_2 & a_{12}C_1D_2 & a_{12}D_1C_2 & a_{12}D_1D_2 \end{pmatrix} \\
&= \begin{pmatrix} a_{12}A_2A_1 & b_{12}B_2A_1 + c_{12}A_2B_1 & c_{12}B_2A_1 + b_{12}A_2B_1 & a_{12}B_2B_1 \\ a_{12}C_2A_1 & b_{12}D_2A_1 + c_{12}C_2B_1 & c_{12}D_2A_1 + b_{12}C_2B_1 & a_{12}D_2B_1 \\ a_{12}A_2C_1 & b_{12}B_2C_1 + c_{12}A_2D_1 & c_{12}B_2C_1 + b_{12}A_2D_1 & a_{12}B_2D_1 \\ a_{12}C_2C_1 & b_{12}D_2C_1 + c_{12}C_2D_1 & c_{12}D_2C_1 + b_{12}C_2D_1 & a_{12}D_2D_1 \end{pmatrix} \\
&\Rightarrow \begin{pmatrix} a_{12}[A_1, A_2] & a_{12}A_1B_2 - (b_{12}B_2A_1 + c_{12}A_2B_1) & a_{12}B_1A_2 - (c_{12}B_2A_1 + c_{12}A_2B_1) & a_{12}[B_1, B_2] \\ (b_{12}A_1C_2 + c_{12}C_1A_2) - a_{12}C_2A_1 & b_{12}[A_1, D_2] + c_{12}[C_1B_2] & (b_{12}[B_1, C_2] + c_{12}[D_1, A_2]) & b_{12}B_1D_2 + c_{12}D_1B_2 - a_{12}D_2B_1 \\ c_{12}A_1C_2 + b_{12}C_1A_2 - a_{12}A_2C_1 & (b_{12}A_1D_2 + c_{12}C_1B_2) - b_{12}B_2C_1 + c_{12}A_2D_1 & c_{12}[B_1, C_2] + b_{12}[D_1, A_2] & c_{12}B_1D_2 + b_{12}D_1B_2 - a_{12}B_2D_1 \\ a_{12}[C_1, C_2] & a_{12}C_1D_2 - (b_{12}D_2C_1 + c_{12}C_2D_1) & a_{12}D_1C_2 - (c_{12}D_2C_1 + b_{12}C_2D_1) & a_{12}[D_1, D_2] \end{pmatrix} = \mathbf{0},
\end{aligned}$$

from which, it's been proven that $[\mathcal{T}_{ij}(u_1), \mathcal{T}_{ij}(u_2)] = 0$, $i, j = 1, 2$ as long as $a(u_1 - u_2) \neq 0$. From the other matrix-entries, after many long calculations, it can be seen that

$$a(u_1 - u_2)B(u_1)A(u_2) = c(u_1 - u_2)B(u_2)A(u_1) + b(u_1 - u_2)A(u_2)B(u_1).$$

Swapping $u_1 \rightarrow u_2$ implies $u_1 - u_2 \rightarrow -(u_1 - u_2)$. Now, note that $c(u_1)/b(u_1) = g(u_1) = g(-u_1)$ and similarly $a(-u_1)/b(-u_1) = h(-u_1) \equiv f(u_1)$, which yields the second line of theorem 1. The third claim can then be similarly derived.

3.0.2.7. Eigenvalues of the Bethe Ansatz. Theorem 1's relations shows the model's inner structure, identifying an underlying algebraic structure somewhat similar to that of the Harmonic oscillator, where $A + D$ are associated with the eigenenergies, and where B and C can be interpreted as ladder operators.

Then, a Fock-like state space can be constructed for the N -body system, with its ground state satisfying $C\Omega^{\otimes N} = 0$. Then, $\omega_+^i = \begin{pmatrix} 1 & 0 \end{pmatrix}_i^\dagger$ describes i -th spin's state, with up spin. Then,

$$\mathbf{S}^+ \omega_+^i = 0,$$

Therefore, consider the \mathcal{L} -matrix, defined by (3.0.2.8), its action on a tensor product of states can be written out as

$$(3.0.2.29) \quad \mathcal{L}_{N,a} v \otimes \omega_+^i = \begin{pmatrix} \lambda + \frac{i}{2} & \cdots \\ 0 & \lambda - \frac{i}{2} \end{pmatrix} v \otimes \omega_+^i, \quad \forall v \in \mathbb{V}_a,$$

since \mathbf{S}^+ annihilates ω_+^i . Thus, the ground state can be written as the physical ferromagnetic vacuum state with all its spins up. Said state is an eigenstate of $A(u)$ and $D(u)$ and is annihilated by $C(u)$, ie.

$$(3.0.2.30) \quad \Omega^{\otimes N} = \bigotimes_{i=1}^N \begin{pmatrix} 1 \\ 0 \end{pmatrix}, \quad \text{where} \quad \begin{aligned} A(u)\Omega^{\otimes N} &= \left(u + \frac{i}{2}\right)^N \Omega^{\otimes N}, \\ D(u)\Omega^{\otimes N} &= \left(u - \frac{i}{2}\right)^N \Omega^{\otimes N}, \\ C(u)\Omega^{\otimes N} &= 0. \end{aligned}$$

The $B(u)$ operators can then be used to construct the Bethe state

$$(3.0.2.31) \quad |u_1, \dots, u_M\rangle = \prod_{j=1}^{M < N} B(u_j) \Omega^{\otimes N}, \quad \text{for some set } \{u_j\}_{j=1}^{M < N} \subset \mathbb{C}.$$

A priori, not all complex-valued sequences are allowed since the Bethe states must be eigenvectors of $A(u) + D(u)$, which yields some algebraic conditions.

3.0.2.8. Conditions on the u -complex coefficients. Consider the Bethe state given by equation (3.0.2.31). Then,

$$A(v) |u_1, \dots, u_M\rangle = A(v) \prod_{j=1}^{M < N} B(u_j) \Omega^{\otimes N}.$$

Using the fundamental commutation relations given theorem 1,

$$A(v)B(u) = f(v-u)B(u)A(v) + g(v-u)B(v)A(u).$$

Therefore,

$$(3.0.2.32) \quad \begin{aligned} A(v) |u_1, \dots, u_M\rangle &= \left[f(v-u_1)B(u_1)A(v) + g(v-u_1)B(v)A(u_1) \right] \prod_{j=2}^{M < N} B(u_j) \Omega^{\otimes N} \\ &= \left[\prod_{k=1}^{\ell} f(v-u_k) \right] \alpha^N(v) |u_1, \dots, u_M\rangle + \mathcal{O}(A \times D), \end{aligned}$$

wherein, the first term is an scalar times the Bethe state, said scalar being the eigenvalue if and only if the second term, the term with products of A and D , cancels out. In that case, the Bethe state is an eigenvector of $A + D$. Then, continuing the previous calculation

(3.0.2.33)

$$\begin{aligned}
A(v) |u_1, \dots, u_M\rangle &= \left[f(v - u_1)B(u_1)A(v) + g(v - u_1)B(v)A(u_1) \right] B(u_2) \prod_{j=3}^{M < N} B(u_j) \Omega^{\otimes N} \\
&= \left[f(v - u_1)f(v - u_2)B(u_1)B(u_2)A(v) + f(v - u_1)g(v - u_2)B(u_1)B(v)A(u_2) \right. \\
&\quad \left. + g(v - u_1)f(u_1 - u_2)B(v)B(u_2)A(u_1) + g(v - u_1)g(v - u_2)B(v)B(u_2) \right] \prod_{j=3}^{M < N} B(u_j) \Omega^{\otimes N}.
\end{aligned}$$

Now, note that

$$\begin{aligned}
(A(v) + D(v)) |u_1, \dots, u_M\rangle &= \left[\left(\prod_{k=1}^{\ell} f(v - u_k) \right) \alpha^N(v) + \left(\prod_{k=1}^{\ell} h(v - u_k) \right) \delta^N(v) \right] |u_1, \dots, u_M\rangle \\
&\quad + \sum_{k=1}^{\ell} \left[\left(g(v - u_k) \left[\prod_{j \neq k}^{\ell} f(u_k - u_j) \right] \alpha^N(u_k) \right) + h(v - u_k) \left(\prod_{j \neq k}^{\ell} h(u_k - u_j) \right) \delta^N(u_j) \right. \\
&\quad \left. \times \prod_{\substack{j=1 \\ j \neq k}}^{\ell} \Omega^{\otimes N} B(u_i) \right] \\
&\equiv \left[\left(\prod_{k=1}^{\ell} f(v - u_k) \right) \alpha^N(v) + \left(\prod_{k=1}^{\ell} h(v - u_k) \right) \delta^N(v) \right] |u_1, \dots, u_M\rangle + 0 \\
&\Rightarrow \sum_{k=1}^{\ell} \left[\left(g(v - u_k) \left[\prod_{j \neq k}^{\ell} f(u_k - u_j) \right] \alpha^N(u_k) \right) + h(v - u_k) \left(\prod_{j \neq k}^{\ell} h(u_k - u_j) \right) \delta^N(u_j) \right] = 0 \\
&\quad \sum_{k=1}^{\ell} \left[\left(g(v - u_k) \left[\prod_{j \neq k}^{\ell} f(u_k - u_j) \right] \alpha^N(u_k) \right) - g(v - u_k) \left(\prod_{j \neq k}^{\ell} h(u_k - u_j) \right) \delta^N(u_j) \right] = 0 \\
&\quad \sum_{k=1}^{\ell} g(v - u_k) \left[\prod_{j \neq k}^{\ell} f(u_k - u_j) \right] \alpha^N(u_k) + \prod_{j \neq k}^{\ell} h(u_k - u_j) \delta^N(u_j) = 0,
\end{aligned}$$

which must be true for all k ,

$$(3.0.2.34) \quad \left[\prod_{j \neq k}^{\ell} f(v - u_k) \right] \alpha^N(u_j) = \left[\prod_{j \neq k}^{\ell} h(v - u_k) \right] \delta^N(u_j).$$

Now using **theorem 1** functions and the that $\alpha(u) = u + \frac{i}{2}$ and $\delta(u) = u - \frac{i}{2}$, the previous equation can be rewritten as

$$\begin{aligned}
&\left[\prod_{j \neq k}^{\ell} f(u_k - u_j) \right] \alpha^N(u_j) = \left[\prod_{j \neq k}^{\ell} h(u_k - u_j) \right] \delta^N(u_j) \\
(3.0.2.35) \quad &\prod_{j \neq k}^{\ell} \frac{u_k - u_j - i}{u_k - u_j} \left(u_j + \frac{i}{2} \right) = \prod_{j \neq k}^{\ell} \frac{u_k - u_j + i}{u_k - u_j} \left(u_j - \frac{i}{2} \right) \\
&0 = \prod_{j \neq k}^{\ell} \frac{u_k - u_j - i}{u_k - u_j} \left(u_j + \frac{i}{2} \right) - \prod_{j \neq k}^{\ell} \frac{u_k - u_j + i}{u_k - u_j} \left(u_j - \frac{i}{2} \right),
\end{aligned}$$

which can be rearranged to yield the Bethe ansatz equations,

$$(3.0.2.36) \quad \left(\frac{u_k + \frac{i}{2}}{u_k - \frac{i}{2}} \right)^N = \prod_{j \neq k}^{\ell} \frac{u_k - u_j + i}{u_k - u_j - i},$$

which is a set of ℓ non linear equations on the coefficients $u_k \in \mathbb{C}$.

3.0.2.9. Momentum and Energy. Let $\{u_k\}_{k=1}^{\ell < N} \subset \mathbb{C}$ satisfy the Bethe equations. Then,

(3.0.2.37)

$$\begin{aligned} (A(v) + D(v)) |u_1, \dots, u_M\rangle &= \left[\left(\prod_{k=1}^{\ell} f(v - u_k) \right) \alpha^N(v) + \left(\prod_{k=1}^{\ell} h(v - u_k) \right) \delta^N(v) \right] |u_1, \dots, u_M\rangle \\ (A(v) + D(v)) |u_1, \dots, u_M\rangle &= \left[\left(v + \frac{i}{2} \right)^N \prod_{k=1}^{\ell} \frac{v - u_k - i}{v - u_k} + \left(v - \frac{i}{2} \right)^N \prod_{k=1}^{\ell} \frac{v - u_k + i}{v - u_k} \right] |u_1, \dots, u_M\rangle \\ \left(A \left(\frac{i}{2} \right) + D \left(\frac{i}{2} \right) \right) |u_1, \dots, u_M\rangle &= i^N \prod_{k=1}^{\ell} \frac{u_k + \frac{i}{2}}{u_k - \frac{i}{2}} |u_1, \dots, u_M\rangle \equiv \lambda(\{u_k\}_{k=1}^{\ell < N}) |u_1, \dots, u_M\rangle \end{aligned}$$

Therefore, via the transfer matrix, it holds that $U = e^{iP} = i^{-N} \mathbf{t} \left(\frac{i}{2} \right) = i^{-N} \left(A \left(\frac{i}{2} \right) + D \left(\frac{i}{2} \right) \right)$, from which

$$(3.0.2.38) \quad P |u_1, \dots, u_{\ell}\rangle = \sum_{k=1}^{\ell} p(u_k) |u_1, \dots, u_{\ell}\rangle \text{ where } p(u_k) = \frac{1}{i} \log \frac{u_k + \frac{i}{2}}{u_k - \frac{i}{2}}.$$

Now, [equation \(3.0.2.27\)](#) establishes that the XXX Hamiltonian belongs to the monodromy matrices' family. Furthermore, it provides a link between the $A + D$'s eigenvalues and the eigenenergies. In effect, said eigenenergies may be found by differentiating the $A + D$'s eigenvalues over u and setting $u = z_0 = \frac{i}{2}$. Let the former

$$\Lambda(v) = \left(v + \frac{i}{2} \right)^N \prod_{k=1}^{\ell} \frac{v - u_k - i}{v - u_k} + \left(v - \frac{i}{2} \right)^N \prod_{k=1}^{\ell} \frac{v - u_k + i}{v - u_k},$$

then,

$$\begin{aligned} \left. \frac{d \log \Lambda(v)}{dv} \right|_{v=z_0} &= \frac{1}{\Lambda(v)} \left. \frac{d \Lambda(v)}{dv} \right|_{v=z_0} \\ &= \frac{1}{\Lambda(v) \Big|_{v=z_0}} \left[N i^{N-1} \prod_{k=1}^{\ell} \frac{-u_k - \frac{i}{2}}{\frac{i}{2} - u_k} + i^N \sum_k \frac{i}{(\frac{i}{2} - u_k)^2} \prod_{\substack{j=1 \\ j \neq k}}^{\ell} \frac{-u_j - \frac{i}{2}}{\frac{i}{2} - u_j} \right] \\ &= i \left(\sum_{k=1}^{\ell} \frac{1}{u_k^2 + \frac{1}{4}} - N \right). \end{aligned}$$

Then, using [equation \(3.0.2.27\)](#), yields the eigenenergies

$$(3.0.2.39) \quad H |u_1, \dots, u_M\rangle = \left[\frac{i}{2} i \left(\sum_{k=1}^{\ell} \frac{1}{u_k^2 + \frac{1}{4}} - N \right) - \frac{N}{2} \right] |u_1, \dots, u_M\rangle$$

or, equivalently,

$$(3.0.2.40) \quad \mathbf{H} |u_1, \dots, u_M\rangle = \sum_{k=1}^{\ell} \epsilon(u_k) |u_1, \dots, u_M\rangle \text{ where } \epsilon(u_k) = -\frac{1}{2} \frac{1}{u_k^2 + \frac{1}{4}},$$

which concludes the Bethe solution for the XXX-Heisenberg chain.

3.1. XX-model. Consider the XX-Heisenberg model, with its Hamiltonian given in terms of the traditional $\frac{1}{2}$ -spin operators ie.

$$(3.1.0.1) \quad \mathbf{H}_{\text{XX}} = J \sum_{i=1}^L (\mathbf{S}_i^x \mathbf{S}_{i+1}^x + \mathbf{S}_i^y \mathbf{S}_{i+1}^y) - \lambda \sum_{j=1}^L \mathbf{S}_j^z,$$

which describes interacting spins in a one-dimensional chain, with periodic boundary conditions. (3.1.0.1)'s first terms represents nearest neighbour interactions in the x and y -directions interactions, with J being either positive or negative and quantifying the strength and type of interactions, while the second term represents a magnetic field of strength λ , applied in the z -direction of the spins.

Exact solution

In order to solve this problem, it is necessary to rewrite (3.1.0.1) and apply a Jordan-Wigner transformation, mapping the spin problem into a fermionic problem. But first, it is convenient to write the spin-operators in terms of the raising and lowering $\mathfrak{su}(2)$ -operators, ie.

$$\mathbf{S}_j^\pm = \mathbf{S}_j^x \pm i\mathbf{S}_j^y \Rightarrow \begin{aligned} \mathbf{S}_j^x &= \frac{1}{2}(\mathbf{S}_j^+ + \mathbf{S}_j^-), \\ \mathbf{S}_j^y &= \frac{1}{2i}(\mathbf{S}_j^+ - \mathbf{S}_j^-). \end{aligned} \Rightarrow \mathbf{H}_{\text{XXZ}} = \frac{J}{2} \sum_{i=1}^L (\mathbf{S}_i^+ \mathbf{S}_{i+1}^- + \mathbf{S}_i^- \mathbf{S}_{i+1}^+) - \lambda \sum_{j=1}^L \mathbf{S}_j^z$$

where the interacting terms in the first summation, flip neighboring spins if said spins are anti-aligned¹⁴. In addition, the XX-Hamiltonian has a total magnetization symmetry, since the Hamiltonian given by (3.1) commutes with the magnetization operator.

Under the Jordan-Wigner map, nearest-neighbours spin flipping is translated into to nearest-neighbours fermionic hopping, ie. $\mathbf{S}_j^+ \mathbf{S}_{j+1}^- = f_j^\dagger f_{j+1}$ and $\mathbf{S}_j^- \mathbf{S}_{j+1}^+ = f_{j+1}^\dagger f_j$. However, due to the boundary conditions' periodicity, the XX-Hamiltonian cannot be rewritten as a fermionic model yet since (it will contain an additional boundary term), for example, the fermionic counterparts to the $\mathbf{S}_L^+ \mathbf{S}_1^-$ interaction are highly non-local operators and are not desirable. Indeed, under the Jordan-Wigner mapping

$$\mathbf{S}_L^+ \mathbf{S}_1^- = f_L^\dagger \exp\left(-i\pi \sum_{\ell=1}^{L-1} f_\ell^\dagger f_\ell\right) f_1 \Rightarrow \begin{aligned} \mathbf{S}_L^+ \mathbf{S}_1^- &= \mathcal{Q} f_L^\dagger f_1 \\ \mathbf{S}_L^- \mathbf{S}_1^+ &= \mathcal{Q} f_1^\dagger f_L, \end{aligned}$$

then (3.1) can be rewritten as

$$(3.1.0.3) \quad \mathbf{H}_{\text{XX}} = \frac{J}{2} \sum_{i=1}^{L-1} \left(f_i^\dagger f_{i+1} + f_{i+1}^\dagger f_i \right) - \lambda \sum_{j=1}^L \left(f_j^\dagger f_j - \frac{1}{2} \right) + \frac{J}{2} \mathcal{Q} (f_L^\dagger f_1 + f_1^\dagger f_L),$$

where the first term accounts for fermionic nearest-neighbour hopping, the second term accounts for the magnetic field, and the third term being the non-local boundary term. Note that this fermionic Hamiltonian hasn't got any type of boundary conditions, since the L -lattice site is disconnected in any way whatsoever from the

¹⁴In effect, consider for example, a two-spin problem. Then, the interaction term is given by

$$\mathbf{S}_1^+ \mathbf{S}_2^- + \mathbf{S}_1^- \mathbf{S}_2^+,$$

and consider an anti-aligned state $|\downarrow\uparrow\rangle$. Then, the action of the previous two-spin operator over this state yields

$$(\mathbf{S}_1^+ \mathbf{S}_2^- + \mathbf{S}_1^- \mathbf{S}_2^+) |\downarrow\uparrow\rangle = |\uparrow\downarrow\rangle + 0,$$

since $\mathbf{S}_1^- \mathbf{S}_2^+$ destroys the state. Similarly, $(\mathbf{S}_1^+ \mathbf{S}_2^- + \mathbf{S}_1^- \mathbf{S}_2^+) |\uparrow\downarrow\rangle = |\downarrow\uparrow\rangle$. However, note that, should both spins be either up or down, the state remain invariant under the action of the two-spin operator.

$$(3.1.0.2) \quad (\mathbf{S}_1^+ \mathbf{S}_2^- + \mathbf{S}_1^- \mathbf{S}_2^+) |\downarrow\downarrow\rangle = |\downarrow\downarrow\rangle \text{ and } (\mathbf{S}_1^+ \mathbf{S}_2^- + \mathbf{S}_1^- \mathbf{S}_2^+) |\uparrow\uparrow\rangle = |\uparrow\uparrow\rangle$$

first lattice site. Then, the standard procedure is to add and subtract terms from the Hamiltonian, so that the nearest-neighbour hopping term in (3.1.0.3) can also have periodic boundary conditions, thus yielding

$$(3.1.0.4) \quad \mathbf{H}_{\text{XX}} = \frac{J}{2} \sum_{i=1}^L \left(f_j^\dagger f_{j+1} + f_{j+1}^\dagger f_j \right) - \lambda \sum_{j=1}^L \left(f_j^\dagger f_j - \frac{1}{2} \right) + \frac{J}{2} (\mathcal{Q} - 1) (f_L^\dagger f_1 + f_1^\dagger f_L),$$

where now the fermionic hopping term has the standard boundary conditions. The third term, since it does not involve any type of summation over lattice sites, only contributes at $\mathcal{O}\left(\frac{1}{L}\right)$ -order to any microscopic quantity. In the thermodynamic limit, this non-local term can be dropped, thus yielding an $\mathcal{O}(L)$ -Hamiltonian given by

$$(3.1.0.5) \quad \mathbf{H}_{\text{XX}} = \frac{J}{2} \sum_{i=1}^L \left(f_j^\dagger f_{j+1} + f_{j+1}^\dagger f_j - \lambda f_j^\dagger f_j \right) + \frac{\lambda L}{2},$$

which is now fully cyclic and where its operators obey fermionic algebras. This Hamiltonian can then be diagonalized via a discrete Fourier transform on the fermionic operators

$$(3.1.0.6) \quad f_j = \frac{1}{\sqrt{L}} \sum_{\substack{k=2\pi m/L \\ m \in \mathbb{Z}_{[1,L]}}} e^{ijk} d_k, \quad f_j^\dagger = \frac{1}{\sqrt{L}} \sum_{\substack{k=2\pi m/L \\ m \in \mathbb{Z}_{[1,L]}}} e^{-ijk} d_k^\dagger,$$

with the inverse transformation given by

$$(3.1.0.7) \quad d_k = \frac{1}{\sqrt{L}} \sum_{j=1}^L e^{-ikj} f_j, \quad d_k^\dagger = \frac{1}{\sqrt{L}} \sum_{j=1}^L e^{ikj} f_j^\dagger.$$

Note that the d_k -operators follow the standard fermionic anticommutation algebra. Note as well, that the f_j -vacuum state, defined such that $f_j |0\rangle_f = 0$, $\forall j$, is the same as the d_k -vacuum state, defined such that $d_j |0\rangle_d = 0$, $\forall k$, ie. $|0\rangle_f = |0\rangle_d$. Another important relationship is the Fourier transform's consistency condition, ie.

$$\sum_{j=1}^L e^{i(k-q)j} = L \delta_{kq}.$$

Under the Fourier transform, (3.1.0.5)'s terms are mapped as follows

$$\begin{aligned} \sum_{j=1}^L f_j^\dagger f_{j+1} &= \sum_{j=1}^L \frac{1}{L} \sum_{k,q} e^{-ikj} e^{iq(j+1)} d_k^\dagger d_q = \sum_{j=1}^L \frac{1}{L} \sum_{k,q} e^{i(q-k)j} e^{iq} d_k^\dagger d_q \\ &= \sum_{k,q} \frac{1}{L} e^{iq} \delta_{qk} d_k^\dagger d_q = \sum_k e^{ik} d_k^\dagger d_k \\ \sum_{j=1}^L f_{j+1}^\dagger f_j &= \sum_{j=1}^L \frac{1}{L} \sum_{k,q} e^{-ik(j+1)} e^{iqj} d_k^\dagger d_q = \sum_{j=1}^L \frac{1}{L} \sum_{k,q} e^{i(q-k)j} e^{-ik} d_k^\dagger d_q \\ &= \sum_{k,q} \frac{1}{L} e^{-ik} \delta_{qk} d_k^\dagger d_q = \sum_k e^{-ik} d_k^\dagger d_k \\ \sum_{j=1}^L f_j^\dagger f_j &= \sum_{j=1}^L \frac{1}{L} \sum_{k,q} e^{-ikj} e^{iqj} d_k^\dagger d_q = \sum_{j=1}^L \frac{1}{L} \sum_{k,q} e^{i(q-k)j} d_k^\dagger d_q \\ &= \sum_{k,q} \frac{1}{L} \delta_{qk} d_k^\dagger d_q = \sum_k d_k^\dagger d_k \\ \Rightarrow \mathbf{H}_{\text{XX}} &= \frac{J}{2} \sum_k \left(e^{ik} d_k^\dagger d_k + e^{-ik} d_k^\dagger d_k - \lambda d_k^\dagger d_k \right) + \frac{\lambda L}{2} = \sum_k \left(J \cos k - \lambda \right) d_k^\dagger d_k + \frac{\lambda L}{2}, \end{aligned}$$

which can be rewritten as

$$(3.1.0.8) \quad \mathbf{H}_{\text{XX}} = \sum_k \epsilon_k d_k^\dagger d_k + \frac{\lambda L}{2} \quad \begin{array}{l} \text{with the eigenvalues being given by } \epsilon_k = J \cos k - \lambda + \frac{\lambda L}{2} \\ \text{and the eigenvectors being given by } |E_n\rangle = \prod_n (d_k^\dagger)^n, \\ \text{with eigenvalue } E_n = \sum_n \epsilon_n \end{array}$$

where $k = \frac{2\pi m}{L}$, $m = -\frac{L}{2}+1, \dots, \frac{L}{2}$, ie. $k \in (-\pi, \pi]$, this is called the first Brillouin zone. Thus the problem has been solved. As for its thermal properties, since this is a fermionic model, the fermions will obey the Fermi-Dirac distribution, ie.

$$(3.1.0.9) \quad \mathcal{N}_{jk} = \langle d_j^\dagger d_k \rangle_{\text{th}} = \frac{1}{1 + e^{\beta \epsilon_k + \mu}} \delta_{jk}$$

Quantum Phase Transition. Note that each of the fermion states can be either occupied or vacant, corresponding to the dimension 2^N of the Hilbert space for N spins with $S = \frac{1}{2}$. The ground state, the state with the lowest energy, has all levels with $\epsilon(k) \leq 0$ and can then be defined such that

$$(3.1.0.10) \quad \langle E_G | d_k^\dagger d_k | E_G \rangle = 1 \text{ if } \epsilon_k \leq 0 \text{ and } \langle E_G | d_k^\dagger d_k | E_G \rangle = 0 \text{ if } \epsilon_k > 0.$$

Then the groundstate energy is given by

$$(3.1.0.11) \quad E_G = \sum_k \epsilon_k \langle E_G | d_k^\dagger d_k | E_G \rangle.$$

Now, there exists a momenta k_F such that $\epsilon_{k_F} = 0$, with k_F being the Fermi momentum, ie. $\epsilon_{k_f} = J \cos k_f - \lambda$. The previous equation depends both on the coupling strength J and the magnetic field strength λ . From the previous equation some special cases arise, namely

- if $\lambda > J$, all fermions levels are occupied, which lead to maximum positive magnetization. This entails that $k_F = 0$.
- If $\lambda < -J$, all fermion levels are vacant, leading to maximum negative magnetization. This implies a $k_F = \pi$ Fermi momentum.
- If $|\lambda| \leq J$, the Fermi momentum is $k_F = \arccos\left(\frac{\lambda}{J}\right)$.

Therefore, the system groundstate will be defined by the following relations,

$$(3.1.0.12) \quad \begin{array}{ll} \langle E_G | d_k^\dagger d_k | E_G \rangle = 1 & \langle E_G | d_k^\dagger d_k | E_G \rangle = 0 \\ \text{if } k \in [-\pi, -k_F] \cup [k_F, \pi], & \text{if } k \in [-k_F, k_F]. \end{array}$$

Note that the periodic boundary conditions in spin space are modified by the Jordan-Wigner transformation, the boundary term in the Hamiltonian explicitly depends on the fermion number L and leads to different Hamiltonians for the two subspaces of even or odd fermion number. For fixed fermion number, this reduces to different sets of allowed fermion momenta k , if the total number of spins L is even, the allowed values of fermion momenta are given by $k_n = \frac{2\pi m}{L}$, where the m -numbers are integer/half-odd integers if the number of fermions $L_f = M^z + \frac{N}{2}$ is odd/even. The total momentum for the ground states is thus $p_F = L_f \pi$ ¹⁵.

¹⁵The same two sets of k -values are found via the Bethe ansatz solution of the XXZ-chain. The complication of two different Hilbert spaces is avoided entirely with free boundar conditions, giving up translational symmetry.

The total magnetization operator is defined as the sum of the S_z operators, along the entire chain ie. $M^z = \sum_{j=1} S_j^z$. Under the Jordan-Wigner transformation, $S_j^z = f_j^\dagger f_j - \frac{1}{2}$. In turn, these fermionic number operators can be mapped to fermionic number operators acting on momentum space, ie. $\sum_{j=1}^L f_j^\dagger f_j = \sum_k d_k^\dagger d_k$. This entails

$$(3.1.0.13) \quad M^z = \sum_k \left(d_k^\dagger d_k - \frac{1}{2} \right).$$

Then the total magnetization operator's expectation value for the groundstate can be calculated. Given the system's fermionic nature, momentum eigenmodes are either occupied or unoccupied, contributing only one or zeros respectively. Then, this expectation value yields

$$(3.1.0.14) \quad \langle E_G | M^z | E_G \rangle = \sum_{k=-\pi}^{-k_F} 1 + \sum_{k=k_F}^{\pi} 1 - \frac{1}{2} \sum_{k=-\pi}^{\pi} 1 = \frac{1}{2} \left(\sum_{k=-\pi}^{-k_F} 1 + \sum_{k=k_F}^{\pi} 1 - \sum_{k=-k_F}^{k_F} 1 \right),$$

where the total magnetization expectation value is just half the difference of two competing summations, related to the negative and positive eigenmodes. Defining the total magnetization per site operator, $m^z = \frac{1}{L} M^z$. Now, as the chain gets progressively bigger, with more and more lattice sites, the summations in (3.1.0.14) can be approximated by Riemannian integrals, ie.

$$\sum_k f(k) = \frac{1}{\Delta k} \sum_k f(k) \Delta k \xrightarrow{L \rightarrow \infty} \frac{L}{2\pi} \int_{k \in \text{BZ}} dk f(k) \text{ where } \Delta k = \frac{2\pi}{L}.$$

Using this trick, (3.1.0.14) can be rewritten as

$$(3.1.0.15) \quad \langle m^z \rangle_G = \frac{1}{4\pi} \left(\int_{-\pi}^{-k_F} dk + \int_{k_F}^{\pi} dk - \int_{-k_F}^{k_F} dk \right) = \frac{1}{4\pi} \left(-k_F + \pi + \pi - k_F - k_F + k_F \right)$$

$$(3.1.0.16) \quad = \frac{1}{2} - \frac{k_F}{\pi},$$

which yields the total magnetization per lattice site. Note that at $\lambda < J$, $k_F = \pi$ whereby there is total polarization ie. $\langle m^z \rangle_G = -\frac{1}{2}$. Similarly, for $\lambda > J$, $k_F = 0$ and then $\langle m^z \rangle_G = \frac{1}{2}$. However, if $|\lambda| \leq J$, then $k_F = \arccos\left(\frac{\lambda}{J}\right)$, which then yields a total magnetization per lattice site given by

$$(3.1.0.17) \quad \langle m^z \rangle_G = \begin{cases} \frac{1}{2} - \frac{1}{\pi} \arccos\left(\frac{\lambda}{J}\right) & \text{if } |\lambda| \leq J \\ -\frac{1}{2} & \text{if } \lambda < -J \\ \frac{1}{2} & \text{if } \lambda > J \end{cases}$$

Note that for $\lambda < J$, maximum negative magnetization is obtained, while for $\lambda > J$ maximum positive magnetization is obtained. Another interesting quantity is the total magnetic susceptibility, which can be written in terms of the total magnetization per lattice site as follows

$$(3.1.0.18) \quad \chi = \left(\frac{\partial \langle m^z \rangle_G}{\partial \lambda} \right) \Rightarrow \chi = \begin{cases} \frac{1}{J\pi\sqrt{1-\frac{\lambda^2}{J^2}}} & \text{if } |\lambda| \leq J \\ 0 & \text{if } \lambda < -J \\ 0 & \text{if } \lambda > J \end{cases},$$

which is divergent if $\lambda = J$. This shows that the XX-Heisenberg model has a second order (since there is no latent heat involved it can't be a first order phase transition) quantum phase transition, with an associated power law with a critical exponent.

Static correlation functions for the XX-model can be calculated for the discrete system, without taking the continuum limit. The longitudinal correlation function in the ground state is given by

$$\langle 0 | \mathbf{S}_n^z \mathbf{S}_0^z | 0 \rangle = \begin{cases} -\frac{1}{4} \left(\frac{2}{\pi n} \right)^2 & \text{for } n \in 2\mathbb{N} + 1 \\ 0 & \text{for } n \in 2\mathbb{N} \end{cases} \quad \text{and} \quad \langle 0 | \mathbf{S}_n^x \mathbf{S}_0^x | 0 \rangle = \langle 0 | \mathbf{S}_n^y \mathbf{S}_0^y | 0 \rangle \sim \frac{C}{\sqrt{n}}, \quad C \approx 0.5884....$$

The transverse correlation function is expressed as a product of two $\frac{n}{2} \times \frac{n}{2}$ determinants with an explicit expression is available only for the asymptotic behaviour. Dynamic correlation functions cannot be obtained at the same level of rigour as static ones since they involve transitions between states in different Hilbert spaces (one for even and another for odd fermion number). Quantities of experimental relevance can be easily calculated from the exact expression for the free energy in terms of the basic fermion dispersion,

$$(3.1.0.19) \quad F = -Lk_B T \left[\log 2 + \frac{2}{\pi} \int_{\mathbb{R}[0, \frac{\pi}{2}]} dk \log \cosh \left(\frac{\epsilon(k)}{2k_B T} \right) \right].$$

An important quantity is the specific heat whose low-temperature behaviour is linear in T :

$$C(T) \approx \frac{\pi T}{6v_F}, \quad \text{where } v_F = \left. \frac{\partial \epsilon(k)}{\partial k} \right|_{k=k_F} = J \text{ is the Fermi velocity.}$$

Low-lying excitations are also simply described in the fermion picture: they are either obtained by adding or removing fermions, thus changing the total magnetization by one unity and adding or removing the energy $\epsilon(k)$, or particle-hole excitations which do not change the total magnetization. Creating a particle-hole excitation involves moving a fermion with momentum k_i inside the Fermi sea at some momentum k_f , outside of the Fermi sea. It is clear that moving a fermion just across the Fermi point costs arbitrarily low energy since the excitation spectrum is gapless. It is easily seen that for a given total momentum $q = k_f - k_i$, a finite range of excitation energies is possible, thus the spectrum of particle-hole excitations is a continuum with the initial momentum $k = k_i$ as an internal degree of freedom,

$$\omega(q, k) = \epsilon(k + q) - \epsilon(k).$$

For $\Delta \neq 0$ the interacting fermion Hamiltonian can be treated in perturbation theory. From this approach, and more generally from the Bethe ansatz and field-theoretic methods, it is established that the behaviour for $-1 < \Delta < 1$ is qualitatively the same as the free fermion limit $\Delta = 0$: the excitation spectrum is gapless, a Fermi point exists and correlation functions show a power-law behaviour. The Heisenberg chain is in the XY regime, and thus is in a critical phase. The phase is equivalent to the so-called Tomonaga-Luttinger liquid. The fermion dispersion, upto first order in Δ , is obtained by direct perturbation theory starting from the free fermion limit (in unit of J),

$$\epsilon(k) = \Delta - \lambda + \cos q - \frac{2\Delta}{\pi} \Theta(1 - \lambda) \left[\arccos \lambda - \frac{\cos q}{(1 - \lambda^2)^{\frac{1}{2}}} \right].$$

For the more general XYZ-chain, wherein the Hamiltonian reads

$$(3.1.0.20) \quad \mathbf{H}_{\text{XYZ}} = J \sum_{n=1}^N \left((1 + \gamma) \mathbf{S}_n^x \mathbf{S}_{n+1}^x + (1 - \gamma) \mathbf{S}_n^y \mathbf{S}_{n+1}^y + \Delta \mathbf{S}_n^z \mathbf{S}_{n+1}^z \right) - \lambda \sum_{n=1}^N \mathbf{S}_n,$$

the rotational symmetry in the xy -plane is broken and a unique preferred direction in spin space exists: $\Delta = 0$ continues to result in a free fermion system, but the basic fermion dispersion acquires a gap and the ground state correlation functions $\langle 0 | \mathbf{S}_n^x \mathbf{S}_0^x | 0 \rangle$ develops long range order.

3.2. **XXZ-model.** Consider the $d = 1$ -XXZ-Heisenberg model, with its Hamiltonian (written in its longitudinal and transverse terms) as

$$(3.2.0.1) \quad H_{\text{XXZ}} = J \sum_{n=1}^N \left(\frac{1}{2} (\mathbf{S}_n^+ \mathbf{S}_{n+1}^- + \mathbf{S}_n^- \mathbf{S}_{n+1}^+) + \Delta \mathbf{S}_n^z \mathbf{S}_{n+1}^z \right),$$

where the effect of the transverse part, for spin- $\frac{1}{2}$ particles, is nothing but to interchange up and down spins, $|\uparrow\downarrow\rangle \Leftrightarrow |\downarrow\uparrow\rangle$, upto a factor of $\frac{1}{2}$. This model has a number of symmetries, leading to good quantum numbers such as wave vector \mathbf{q} (translation), \mathbf{M}^z (rotation about z -axis), \mathbf{S}_{tot} (for general rotations, $|\Delta| = 1$) and parity (spin inversion). The most important link between theory and experiment are the spin correlation functions or response dynamical structure factors, which for a spin chain are defined as follows

$$(3.2.0.2) \quad \begin{aligned} \mathcal{S}^{\alpha,\alpha}(q, \omega) &= \sum_n \int dt e^{i(qn - \omega t)} \langle \mathbf{S}_n^\alpha \mathbf{S}_0^\alpha(t=0) \rangle, \\ \mathcal{S}^{\alpha,\alpha}(q) &= \sum_n e^{iqn} \langle \mathbf{S}_n^\alpha \mathbf{S}_0^\alpha \rangle = \frac{1}{2\pi} \int d\omega \mathcal{S}^{\alpha,\alpha}(q, \omega), \end{aligned}$$

where $\mathcal{S}(q, \omega)$ determines the cross-section for scattering experiments as well as line shapes in NMR and ESR experiments. A useful sum rule is the total intensity, obtained by integrating $\mathcal{S}(q)$ over frequency and wave vector,

$$(3.2.0.3) \quad \frac{1}{4\pi^2} \int d\omega \mathcal{S}^{\alpha,\alpha}(q, \omega) = \frac{1}{2\pi} \int dq \mathcal{S}^{\alpha,\alpha}(q) = \langle (\mathbf{S}_0^\alpha)^2 \rangle,$$

which is simply equal to $\frac{1}{3} \mathcal{S}(\mathcal{S} + 1)$ in the isotropic case.

Physical Context

The XXZ-Heisenberg chain is an important model to describe real material and is the most important paradigm of low-dimensional quantum magnetism, as it allows to introduce many scenarios, such as broken symmetry, the gapless Luttinger liquid, the Kosterlitz-Thouless phase transition, gapped and gapless excitation spectra.

3.2.1. **Ferromagnetic phase:** Consider a Heisenberg XXZ chain supplemented by an external magnetic field. For $\Delta < -1$, the XXZ-chain is in the ferromagnetic Ising phase: the ground state is the saturated state, with all its spins aligned in either the z -direction or $-z$ -direction. This entails that the classical ground state has magnetization $\mathbf{M}^z = \pm \frac{N}{2}$, where N is the number of lattice sites. This is a phase with broken symmetry since the ground state does not exhibit the discrete symmetry of spin reflection $\mathbf{S}^z \rightarrow -\mathbf{S}^z$, under which the Hamiltonian is invariant. Furthermore, when an external magnetic field in the z -axis is considered, and additional term is included into the Hamiltonian, namely a Zeeman interaction

$$H_{\text{Zeeman}} = -g\mu_B H \sum_n \mathbf{S}_n^z.$$

Since the XXZ-Hamiltonian commutes with the magnetization operator, the external magnetic field results in an additional energy contribution $-g\mu_B H \mathbf{M}^z$, without affecting the wave-functions. The symmetry under spin reflection is lifted and the saturated ground state is stabilized. The low-lying excited states in the ferromagnetic phase are magnons with the total spin quantum number $\mathbf{M}_{\text{tot}}^z = \frac{1}{2}N - 1$ and the dispersion law, which is valid for general spin S

$$(3.2.1.1) \quad \epsilon(q) = 2JS(1 - \cos(q) - (\Delta + 1)) + 2g\mu_B HS.$$

These states are exact eigenstates of the XXZ-Hamiltonian. In zero field, the excitation spectrum has a gap at $q = 0$ of magnitude $|\Delta| - 1$ for $\Delta < -1$. At $\Delta = -1$, the discrete spin-reflection symmetry generalizes to a continuous rotation symmetry and the spectrum thus becomes gapless. This is a consequence of Goldstone's theorem: the breaking of a continuous symmetry in the ground state results in the emergence of a gapless excitation mode. Whereas the ground state exhibit long range order, the large phase space available to the low-lying excitations in 1D leads to exponential decay of correlations at arbitrarily small finite temperatures following the Theorem of Mermin and Wagner.

Physical Context

Two interesting consequences of [equation \(3.2.1.1\)](#) quickly arise, which can be generalized to any XXZ-type Hamiltonian conserving the total magnetization:

- 1) For sufficiently strong external magnetic field, the classical saturated state is forced to be the ground state for arbitrary value of Δ and the lowest excitations are exactly known. If the necessary magnetic fields are within experimentally accessible range, this can be used for an experimental determination of the exchange constants from the magnon dispersions (an example in 2D are recent neutron scattering experiments on Cs_2CuCl_4).
- 2) The ferromagnetic ground state becomes unstable when the low spin wave frequency becomes negative. This allows to determine eg. the boundary of the ferromagnetic phase for $\Delta > -1$ in an external field as $H = H_c$ with $g\mu_B H_c = \Delta + 1$.

3.2.2. Néel phase: For $\Delta > +1$, the XXZ-chain is in the antiferromagnetic Ising or Néel phase with, in the thermodynamic limit, broken symmetry and one from 2 degenerate ground states, the $S = \frac{1}{2}$ remnants of the classical Néel states. The spatial period is $2a$, and the states are described in the reduced Brillouin zone with wave vectors $0 \leq q \leq \frac{\pi}{a}$. The ground states have zero magnetisation but finite sublattice magnetization,

$$N^z = \sum_n (-1)^n S_n^z,$$

and long range order in the corresponding correlation function. In contrast to the ferromagnet, however, quantum fluctuations prevent the order from being complete since the sublattice magnetization does not commute with the XXZ-Hamiltonian. For periodic boundary conditions and large but finite N , as in the situation in numerical approaches, the two ground states mix with energy separation $\propto e^{-\text{const} \times N}$, as $N \rightarrow \infty$. Then invariance under translation by the original lattice constant a is restored and the original Brillouin zone can be used, $BZ = \{q | 0 \leq q \leq \frac{2\pi}{a}\}$.

The elementary excitations in the antiferromagnetic Ising phase are described most clearly close to the Ising limit $\delta \rightarrow \infty$, starting from one of the two ideal Néel states: turning around one spin breaks two bonds and leads to a state with energy Δ , degenerate with all state resulting from turning around an arbitrary number of subsequent spins. These states have total magnetization $M^z = \pm 1$, response 0 for an odd, and response even number of turned spins. They are appropriately called two-domain wall states since each of the two broken bonds mediates between two different Néel states. They are appropriately called two-domain wall states, since each of the two broken bonds mediates between two different Néel states. The total number of these states is $N(N - 1)$,

- there are $\frac{N^2}{4}$ states with total magnetization $M^z = 1$ and $M^z = -1$ (number of turned spins odd),
- and $\frac{N^2}{2} - N$ states with null magnetization (number of spins turned even).

These states are no more eigenstates when Δ^{-1} is finite, but for $\Delta^{-1} \ll 1$, they can be dealt with in perturbation theory, leading to the excitation spectrum in the first order in $\frac{1}{\Delta}$

$$\omega(q, k) = \Delta + 2 \cos q \cos 2\Phi = \epsilon \left(\frac{\epsilon}{2} + \Phi \right) + \epsilon \left(\frac{\epsilon}{2} - \Phi \right),$$

with $\epsilon(k) = \frac{\Delta}{2} + \cos 2k$, where q is the total momentum and takes values $q = \frac{2\pi\ell}{N}$, $\ell \in \mathbb{Z}_{[1, \frac{N}{2}]}$, Φ is the wave vector related to the superposition of domain walls with different distances and for $\mathbf{M}^z = \pm 1$ takes values $\Phi = \frac{m\pi}{N+2}$ with $m \in \mathbb{Z}_{[1, \frac{N}{2}]}$. Φ is, in essence, a relative momentum.

3.2.3. XY phase. For $-1 < \Delta < 1$ and zero external field, the XXZ-chain undergoes an XY phase, characterized by uniaxial symmetry of the easy-plane type and a gapless excitation continuum. Whereas the full analysis of this phase for general Δ , requires the use of bosonization and Bethe ansatz, a somewhat simpler approach is based on the mapping of $S = \frac{1}{2}$ spin operators in 1D to spinless fermions via the non-local Jordan Wigner transformations, wherein the mapping reads:

$$(3.2.3.1) \quad \mathbf{S}_n^+ = f_n^\dagger e^{i\pi \sum_{j=1}^n f_j^\dagger f_j}, \quad \mathbf{S}_n^z = f_n^\dagger f_n - \frac{1}{2}.$$

When a fermion is at site n , the spin projection is thus $\mathbf{S}_n^z = \pm \frac{1}{2}$. Thus, in terms of the fermionic operators, the XXZ-Hamiltonian reads

$$(3.2.3.2) \quad \mathbf{H}_{\text{XXZ}} = J \sum_n \left[\frac{1}{2} (f_n^\dagger f_{n+1} + f_{n+1}^\dagger f_n) + \Delta \left(f_n^\dagger f_n - \frac{1}{2} \right) \left(f_{n+1}^\dagger f_{n+1} - \frac{1}{2} \right) - g\mu_B H \left(f_n^\dagger f_n - \frac{1}{2} \right) \right].$$

For general Δ , the XXZ-chain is thus equivalent to an interacting 1D fermion system. In the $\Delta = 0$ case, this model simplifies itself to an integrable XX-chain, wherein the fermion chain becomes noninteracting and has a simpler analytical solution.

XY-Model

3.2.3.1. Setting of the problem and the Jordan-Wigner Transformation. The XY-Heisenberg chain Hamiltonian is given by

$$(3.2.3.3) \quad \mathbf{H}_{\text{XY}} = \sum_{j \in \Lambda} \left[(1 + \gamma) \mathbf{S}_j^x \mathbf{S}_{j+1}^x + (1 - \gamma) \mathbf{S}_j^y \mathbf{S}_{j+1}^y \right],$$

where γ is a parameter characterizing the degree of anisotropy in the xy -plane. The endpoint may be treated in two differently, but physically reasonable ways,

- 1) as free ends, in which case the range of the summation index is $1 \leq i \leq N - 1$, a situation convenient for discussing long-range order,
- 2) or as a cyclic chain, in which $1 \leq i \leq N$ and $\vec{\mathbf{S}}_{N+1} = \vec{\mathbf{S}}_1$, being the most convenient choice for calculating physical quantities other than the long-range order.

The XY-model is exactly soluble for all values of γ . In particular, consider $\gamma \in \mathbb{R}_{[-1,1]}$. As $\gamma \rightarrow 1$, both models tend to the Ising model in which the x -component of spin are completely ordered and the y -component and z -components are completely disordered. For $|\gamma| \neq 1$ both models are genuinely quantum mechanical because different components of $\vec{\mathbf{S}}$, present in the Hamiltonian, do not commute. The effect of the transverse terms (those with a $(1 - \gamma)$ -weight) is opposing the ordering of the x -components and to favour the ordering of the y -components.

Hitherto, it has not been clear for the Heisenberg model, just how strong these two effects are. Thus, the following question arises,

For any particular positive value of γ , do the transverse terms either establish any long-range order among the y - and z -components or destroy the long-range order of the x -components (which would imply absence of long range order for the y - and z -components as well)?

It turns out that for the XY-model, the transverse terms do neither until the limiting case $\gamma \rightarrow 0$, when these succeed in destroying the order of the x -components.

To solve the XY-model, consider the following raising and lowering operators

$$\begin{aligned}
\mathbf{S}_i^x &= \frac{1}{2}(\mathbf{a}_i + \mathbf{a}_i^\dagger) \\
(3.2.3.4) \quad \begin{aligned} \mathbf{a}_i^\dagger &= \mathbf{S}_i^x + i\mathbf{S}_i^y \\ \mathbf{a}_i &= \mathbf{S}_i^x - i\mathbf{S}_i^y \end{aligned} \Leftrightarrow \mathbf{S}_i^y = \frac{1}{2i}(\mathbf{a}_i - \mathbf{a}_i^\dagger) \Rightarrow \mathbf{H}_\gamma = \frac{1}{2} \sum_i \left(\mathbf{a}_i^\dagger \mathbf{a}_{i+1} + \gamma \mathbf{a}_i^\dagger \mathbf{a}_{i+1}^\dagger + \mathbf{a}_{i+1}^\dagger \mathbf{a}_i + \gamma \mathbf{a}_{i+1} \mathbf{a}_i \right). \\
\mathbf{S}_i^z &= \mathbf{a}_i^\dagger \mathbf{a}_i - \frac{1}{2}
\end{aligned}$$

These operators partly resemble Fermi operators: $\{\mathbf{a}_i, \mathbf{a}_i^\dagger\} = 1, \mathbf{a}_i^2 = (\mathbf{a}_i^\dagger)^2 = 0$,
and partly resemble bosonic operators: $[\mathbf{a}_i^\dagger, \mathbf{a}_j] = [\mathbf{a}_i^\dagger, \mathbf{a}_j^\dagger] = [\mathbf{a}_i, \mathbf{a}_j] = 0, i \neq j$.

It is therefore not possible to diagonalize the quadratic form appearing in [equation \(3.2.3.4\)](#) directly with a canonical transformation, since principal axis transformations of the \mathbf{a} - and \mathbf{a}^\dagger -operators do not preserve this mixed set of canonical rules. However, it is possible to transform to a new set of variables that are strictly fermionic operators, and in terms of which the Hamiltonian is just as simple. This new operators are given by the Jordan-Wigner transformation

$$(3.2.3.5) \quad \begin{aligned} \mathbf{c}_i &= e^{i\pi \sum_{j=1}^{i-1} \mathbf{a}_j^\dagger \mathbf{a}_j} \mathbf{a}_i & \Leftrightarrow \mathbf{a}_i &= e^{i\pi \sum_{j=1}^{i-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \mathbf{c}_i \\ \mathbf{c}_i^\dagger &= \mathbf{a}_i^\dagger e^{-i\pi \sum_{j=1}^{i-1} \mathbf{a}_j^\dagger \mathbf{a}_j} & \Leftrightarrow \mathbf{a}_i^\dagger &= \mathbf{c}_i^\dagger e^{-i\pi \sum_{j=1}^{i-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \end{aligned}, \text{ such that } \mathbf{c}_i^\dagger \mathbf{c}_i = \mathbf{a}_i^\dagger \mathbf{a}_i$$

The \mathbf{c} - and \mathbf{c}^\dagger -operators are truly fermionic operators, this is they follow fermionic anti-commutation relations,

$$(3.2.3.6) \quad \{\mathbf{c}_i, \mathbf{c}_j^\dagger\} = \delta_{ij}, \quad \{\mathbf{c}_i, \mathbf{c}_j\} = 0 \quad \{\mathbf{c}_i^\dagger, \mathbf{c}_j^\dagger\} = 0.$$

Note that

$$(3.2.3.7) \quad e^{\pm i\pi \sum_{j=1}^m \mathbf{a}_j^\dagger \mathbf{a}_j} = \prod_{j=1}^m e^{\pm i\pi \mathbf{a}_j^\dagger \mathbf{a}_j}.$$

Furthermore, given that $\mathbf{c}_i^\dagger \mathbf{c}_i$ is a well-defined hermitian occupation number having values 0 or 1, it follows

$$(3.2.3.8) \quad e^{i\pi \mathbf{c}_j^\dagger \mathbf{c}_j} = e^{-i\pi \mathbf{c}_j^\dagger \mathbf{c}_j} \Rightarrow e^{\pm i\pi \mathbf{a}_i^\dagger \mathbf{a}_i} = \sum_{\ell=0}^{\infty} \frac{1}{\ell!} (\pi i)^\ell (\mathbf{a}_i^\dagger \mathbf{a}_i)^\ell = 1 + \sum_{\ell=1}^{\infty} \frac{1}{\ell!} (\pm i\pi)^\ell \mathbf{a}_i^\dagger \mathbf{a}_i = 1 + (e^{\pm \pi i} - 1) \mathbf{a}_i^\dagger \mathbf{a}_i = 1 - 2\mathbf{a}_i^\dagger \mathbf{a}_i.$$

Note as well that

$$\begin{aligned}
(3.2.3.9) \quad \begin{aligned} \{\mathbf{a}_i, 1 - 2\mathbf{a}_i^\dagger \mathbf{a}_i\} &= \mathbf{a}_i(1 - 2\mathbf{a}_i^\dagger \mathbf{a}_i) + (1 - 2\mathbf{a}_i^\dagger \mathbf{a}_i)\mathbf{a}_i \\ &= \mathbf{a}_i - 2\mathbf{a}_i \mathbf{a}_i^\dagger \mathbf{a}_i + \mathbf{a}_i - 2\mathbf{a}_i^\dagger \mathbf{a}_i \mathbf{a}_i \\ &= 2\mathbf{a}_i - 2\left(\cancel{\mathbf{a}_i \mathbf{a}_i^\dagger} + \cancel{\mathbf{a}_i^\dagger \mathbf{a}_i}\right) \mathbf{a}_i = 0 \end{aligned} & \quad \begin{aligned} \{\mathbf{a}_i^\dagger, 1 - 2\mathbf{a}_i^\dagger \mathbf{a}_i\} &= \mathbf{a}_i^\dagger(1 - 2\mathbf{a}_i^\dagger \mathbf{a}_i) + (1 - 2\mathbf{a}_i^\dagger \mathbf{a}_i)\mathbf{a}_i^\dagger \\ &= \mathbf{a}_i^\dagger - 2\mathbf{a}_i^\dagger \mathbf{a}_i^\dagger \mathbf{a}_i + \mathbf{a}_i^\dagger - 2\mathbf{a}_i^\dagger \mathbf{a}_i \mathbf{a}_i^\dagger \\ &= 2\mathbf{a}_i^\dagger - 2\mathbf{a}_i^\dagger \left(\cancel{\mathbf{a}_i \mathbf{a}_i^\dagger} + \cancel{\mathbf{a}_i^\dagger \mathbf{a}_i}\right) = 0 \end{aligned} \end{aligned}$$

$$\Rightarrow \left[e^{\pm i\pi \sum_{\ell=n}^m \mathbf{a}_j^\dagger \mathbf{a}_j}, \mathbf{a}_i \right] = \left[e^{\pm i\pi \sum_{\ell=n}^m \mathbf{a}_j^\dagger \mathbf{a}_j}, \mathbf{a}_i^\dagger \right] = 0, \text{ if } i \notin [n, m]$$

$$\Rightarrow \left\{ e^{\pm i\pi \sum_{\ell=n}^m \mathbf{a}_j^\dagger \mathbf{a}_j}, \mathbf{a}_i \right\} = \left\{ e^{\pm i\pi \sum_{\ell=n}^m \mathbf{a}_j^\dagger \mathbf{a}_j}, \mathbf{a}_i^\dagger \right\} = 0, \text{ if } i \in [n, m]$$

For $i \leq i \leq N$,

$$\begin{aligned}
\mathbf{a}_i \mathbf{a}_{i+1}^\dagger &= e^{i\pi \sum_{j=1}^{i-1} \mathbf{a}_j^\dagger \mathbf{c}_j} \mathbf{c}_i \mathbf{c}_{i+1}^\dagger e^{-i\pi \sum_{j=1}^i \mathbf{c}_j^\dagger \mathbf{a}_j} = \mathbf{c}_i \mathbf{c}_{i+1}^\dagger e^{i\pi \sum_{j=1}^{i-1} \mathbf{c}_j^\dagger \mathbf{c}_j} e^{-i\pi \sum_{j=1}^i \mathbf{c}_j^\dagger \mathbf{c}_j} \\
(3.2.3.10) \quad &= \mathbf{c}_i e^{i\pi \mathbf{a}_i^\dagger \mathbf{a}_i} \mathbf{c}_{i+1}^\dagger = \mathbf{c}_i (1 - 2\mathbf{c}_i^\dagger \mathbf{c}_i) \mathbf{c}_{i+1}^\dagger = -(1 - 2\mathbf{c}_i^\dagger \mathbf{c}_i) \mathbf{c}_i \mathbf{c}_{i+1}^\dagger = -\mathbf{c}_i^\dagger \mathbf{c}_{i+1} + 2\mathbf{c}_i^\dagger \mathbf{c}_i \mathbf{c}_{i+1}^\dagger \\
&= -\mathbf{c}_i^\dagger \mathbf{c}_{i+1} = \mathbf{c}_{i+1}^\dagger \mathbf{c}_i \\
\mathbf{a}_{i+1} \mathbf{a}_i^\dagger &= (\mathbf{a}_i \mathbf{a}_{i+1}^\dagger)^\dagger = \mathbf{c}_i^\dagger \mathbf{c}_{i+1}.
\end{aligned}$$

Then, two possible Hamiltonians arise, depending on which boundary condition has been imposed,

- For free boundary conditions, the Hamiltonian reads

$$(3.2.3.11) \quad \mathbf{H}_\gamma = \frac{1}{2} \sum_{j=1}^{N-1} \left[\mathbf{c}_j^\dagger \mathbf{c}_{j+1} + \mathbf{c}_{j+1}^\dagger \mathbf{c}_j + \gamma (\mathbf{c}_j^\dagger \mathbf{c}_{j+1}^\dagger + \mathbf{c}_{j+1} \mathbf{c}_j) \right],$$

- whereas, for the cyclic chain, there are additional terms which couple the N -th spin to the first spin, namely

$$\begin{aligned}
\mathbf{a}_N^\dagger \mathbf{a}_1 &= -\mathbf{c}_N^\dagger \mathbf{c}_1 e^{i\pi \mathfrak{F}} \neq \mathbf{c}_N^\dagger \mathbf{c}_1 \longleftrightarrow \mathbf{a}_1^\dagger \mathbf{a}_N = -\mathbf{c}_1^\dagger \mathbf{c}_N e^{-i\pi \mathfrak{F}} \neq \mathbf{c}_1^\dagger \mathbf{c}_N \\
\mathbf{a}_N^\dagger \mathbf{a}_1^\dagger &= -\mathbf{c}_1^\dagger \mathbf{c}_N e^{i\pi \mathfrak{F}} \neq \mathbf{c}_1^\dagger \mathbf{c}_N^\dagger \longleftrightarrow \mathbf{a}_1 \mathbf{a}_N = -\mathbf{c}_1 \mathbf{c}_N e^{-i\pi \mathfrak{F}} \neq \mathbf{c}_1 \mathbf{c}_N
\end{aligned}$$

where $\mathfrak{F} = \sum_{j=1}^N \mathbf{c}_j^\dagger \mathbf{c}_j = \sum_{i=1}^N \left(\mathbf{S}_i^z + \frac{1}{2} \right)$,

yielding the following Hamiltonian

$$(3.2.3.13) \quad \mathbf{H}_\gamma = \frac{1}{2} \sum_{j=1}^{N-1} \left[\mathbf{c}_j^\dagger \mathbf{c}_{j+1} + \mathbf{c}_{j+1}^\dagger \mathbf{c}_j + \gamma (\mathbf{c}_j^\dagger \mathbf{c}_{j+1}^\dagger + \mathbf{c}_{j+1} \mathbf{c}_j) \right] - \frac{1}{2} \left[\mathbf{c}_N^\dagger \mathbf{c}_1 + \mathbf{c}_1^\dagger \mathbf{c}_N + \gamma \mathbf{c}_N^\dagger \mathbf{c}_1^\dagger + \gamma \mathbf{c}_1 \mathbf{c}_N \right] (e^{i\pi \mathfrak{F}} + 1),$$

in which case the Hamiltonian, written in terms of the Fermi operators ($\mathbf{c}, \mathbf{c}^\dagger$), no longer has a simple cyclic structure. For large-enough systems, the last term may be neglected, in which case the problem is known as "c-cyclic". Otherwise, as it stands now, the problem is called the "a-cyclic" problem. Both problems can be analytically solved, though the former is simple.

Physical Context

Observe that the original Hamiltonian conserves total spin along the z -axis, for $[\mathbf{H}, \mathbf{S}^z] = 0$. After the Jordan-Wigner transformation, this symmetry is manifested as conservation of the total fermion number, $[\mathbf{H}, \sum_i \mathbf{c}_i^\dagger \mathbf{c}_i] = 0$. In effect,

$$(3.2.3.14) \quad \mathbf{S}_i^z = \frac{1}{2} [\mathbf{a}_i^\dagger, \mathbf{a}_i] = \mathbf{a}_i^\dagger \mathbf{a}_i - \frac{1}{2} = \mathbf{c}_i^\dagger \mathbf{c}_i - \frac{1}{2}, \quad \mathbf{S}^z = \sum_i \mathbf{S}_i^z = \sum_i \mathbf{c}_i^\dagger \mathbf{c}_i - \frac{N}{2}$$

Thus, each "spinless" fermion created by \mathbf{c}^\dagger carries $\mathbf{S}^z = 1$. Furthermore, unlike the XY model, the Heisenberg Hamiltonian is not a free theory under a Jordan-Wigner transformation, but a theory of spinless fermions on a lattice with nearest-neighbour hopping and nearest-neighbour interactions

In all cases, the Hamiltonian is a simple quadratic form in Fermi operators and can be exactly diagonalized. The particular simplicity of \mathbf{H}_γ depends on the fact that the spins can be arranged in a definite order, that interacting occur only between neighbouring spins in this ordering, and that the z -components of spin do not enter. If the interactions were to extend to n -nearest neighbours, \mathbf{H}_γ would involve a polynomial of $2n$ -order on the \mathbf{c} -operators. In two-dimensional models, it can be seen that any ordering and non-trivial scheme of interactions must lead to a Hamiltonian involving a polynomial roughly of $2N$ -order for a system with N^2 spins.

3.2.3.2. Ground State Energy, Elementary Excitations and Free Energy.

- Antiferromagnetic case and anisotropic general case:

Consider the following quadratic form,

$$(3.2.3.15) \quad \mathbf{H} = \mathbf{H}_b + \sum_{i,j} \left[\mathbf{c}_i^\dagger \mathcal{A}_{ij} \mathbf{c}_j + \frac{1}{2} \left(\mathbf{c}_i^\dagger \mathcal{B}_{ij} \mathbf{c}_j^\dagger + \mathbf{c}_i \mathcal{B}_{ij}^* \mathbf{c}_j \right) \right], \quad \mathbf{H}_b = -\frac{1}{2} (\mathbf{c}_1^\dagger \mathbf{c}_N + \mathbf{c}_N^\dagger \mathbf{c}_1) e^{i\pi \mathcal{M}},$$

$$\mathcal{A}_{ij} \sim \delta_{i,j+1} + \delta_{i+1,j}, \quad \mathcal{B}_{ij} \sim \delta_{i,j}$$

where the $(\mathbf{c}, \mathbf{c}^\dagger)$ -operators are fermionic in nature and where \mathbf{H} is hermitian, which induces \mathcal{A} 's hermiticity and \mathcal{B} 's antisymmetric nature. For this section's purpose, they can be treated as real-valued matrices.

Let $\{\phi_k, \psi_k\}_k$ be a complete and orthonormal set of eigenvectors of the second term in the previous Hamiltonian, with $\{\Lambda_k\}$ -eigenvalues. Then, let

$$(3.2.3.16) \quad \begin{aligned} \eta_k &= \sum_i \phi_{ki} \mathbf{c}_i + \psi_{ki} \mathbf{c}_i^\dagger, & \text{which imply canonical anticommutation} & & \mathbf{c}_i &= \sum_k \phi_{ki} \eta_k + \psi_{ki} \eta_k^\dagger \\ \eta_k^\dagger &= \sum_i \psi_{ki} \mathbf{c}_i + \phi_{ki} \mathbf{c}_i^\dagger, & \text{relations of the } \eta\text{'s, and conversely} & & \mathbf{c}_i^\dagger &= \sum_k \psi_{ki} \eta_k + \phi_{ki} \eta_k^\dagger. \end{aligned}$$

Therefore, the second term in the Hamiltonian can be written as

$$(3.2.3.17) \quad \mathbf{H} = c + \sum_k \eta_k^\dagger \Lambda_k \eta_k.$$

If this is possible, then

$$(3.2.3.18) \quad \begin{aligned} [\eta_k, \mathbf{H}] - \Lambda_k \eta_k &= 0 = \left[\sum_i (\phi_{ki} \mathbf{c}_i + \psi_{ki} \mathbf{c}_i^\dagger), \mathbf{H} \right] - \Lambda_k \sum_i (\phi_{ki} \mathbf{c}_i + \psi_{ki} \mathbf{c}_i^\dagger) \\ &= \sum_i \phi_{ki} [\mathbf{c}_i, \mathbf{H}] - \Lambda_k \psi_{ki} \mathbf{c}_i^\dagger + \psi_{ki} [\mathbf{c}_i^\dagger, \mathbf{H}] - \Lambda_k \phi_{ki} \mathbf{c}_i \\ &= \sum_{i,m,n} \left[\mathbf{c}_i, \mathbf{c}_m^\dagger \mathcal{A}_{mn} \mathbf{c}_n + \frac{1}{2} \left(\mathbf{c}_m^\dagger \mathcal{B}_{mn} \mathbf{c}_n^\dagger + \mathbf{c}_m \mathcal{B}_{mn}^* \mathbf{c}_n \right) \right] - \Lambda_k \psi_{ki} \mathbf{c}_i^\dagger \\ &\quad + \psi_{ki} \left[\mathbf{c}_i^\dagger, \mathbf{c}_m^\dagger \mathcal{A}_{mn} \mathbf{c}_n + \frac{1}{2} \left(\mathbf{c}_m^\dagger \mathcal{B}_{mn} \mathbf{c}_n^\dagger + \mathbf{c}_m \mathcal{B}_{mn}^* \mathbf{c}_n \right) \right] - \Lambda_k \phi_{ki} \mathbf{c}_i^\dagger \\ &= \sum_{i,m,n} \phi_{ki} \left\{ \mathcal{A}_{mn} \left(\cancel{[\mathbf{c}_i^\dagger, \mathbf{c}_m^\dagger] \mathbf{c}_n}^{\delta_{im}} + \mathbf{c}_m^\dagger \cancel{[\mathbf{c}_i, \mathbf{c}_n]}^{\delta_{in}} \right) + \frac{1}{2} \mathcal{B}_{mn} \left(\cancel{[\mathbf{c}_i, \mathbf{c}_m^\dagger] \mathbf{c}_n^\dagger}^{\delta_{im}} + \mathbf{c}_m^\dagger \cancel{[\mathbf{c}_i^\dagger, \mathbf{c}_n]}^{\delta_{in}} \right) \right. \\ &\quad \left. + \frac{1}{2} \mathcal{B}_{mn}^* \left(\cancel{[\mathbf{c}_i, \mathbf{c}_m] \mathbf{c}_n} + \mathbf{c}_m \cancel{[\mathbf{c}_i^\dagger, \mathbf{c}_n]} \right) \right\} - \Lambda_k \psi_{ki} \mathbf{c}_i^\dagger \\ &\quad + \psi_{ki} \left\{ \mathcal{A}_{mn} \left(\cancel{[\mathbf{c}_i^\dagger, \mathbf{c}_m^\dagger] \mathbf{c}_n} + \mathbf{c}_m^\dagger \cancel{[\mathbf{c}_i^\dagger, \mathbf{c}_n]}^{-\delta_{in}} \right) + \frac{1}{2} \mathcal{B}_{mn} \left(\cancel{[\mathbf{c}_i, \mathbf{c}_m^\dagger] \mathbf{c}_n^\dagger} + \mathbf{c}_m^\dagger \cancel{[\mathbf{c}_i^\dagger, \mathbf{c}_n]} \right) \right. \\ &\quad \left. + \frac{1}{2} \mathcal{B}_{mn}^* \left(\cancel{[\mathbf{c}_i^\dagger, \mathbf{c}_m] \mathbf{c}_n}^{-\delta_{im}} + \mathbf{c}_m \cancel{[\mathbf{c}_i, \mathbf{c}_n]}^{-i\delta_{in}} \right) \right\} - \Lambda_k \phi_{ki} \mathbf{c}_i \\ &= \sum_{i,n} \phi_{ki} \mathcal{A}_{in} \mathbf{c}_n + \mathcal{B}_{in} \mathbf{c}_n^\dagger - \Lambda_k \phi_{ki} \mathbf{c}_i^\dagger \\ &\quad - \psi_{ki} \mathcal{A}_{ni} \mathbf{c}_n^\dagger - \psi_{ki} \mathcal{B}_{ni}^* \mathbf{c}_n - \Lambda_k \psi_{ki} \mathbf{c}_i^\dagger = 0 \Rightarrow \begin{cases} \Lambda_k \phi_{ki} = \sum_n \phi_{kn} \mathcal{A}_{ni} - \psi_{kn} \mathcal{B}_{ni}^* \\ \Lambda_k \phi_{ki} = \sum_n \phi_{kn} \mathcal{B}_{ni} - \psi_{kn} \mathcal{A}_{ni} \end{cases} \end{aligned}$$

These equations can be further simplified by introducing the linear combinations

$$(3.2.3.19) \quad \Phi_{kn} = \phi_{kn} + \psi_{kn} \quad \Psi_{kn} = \phi_{kn} - \psi_{kn} \quad \Rightarrow \quad \begin{aligned} \Phi_k(\mathcal{A} + \mathcal{B}) &= \Lambda \Psi_{ki} \\ \Psi_k(\mathcal{A} - \mathcal{B}) &= \Lambda \Phi_{ki} \end{aligned}$$

Note that the relative sign between Φ_k and Ψ_k is arbitrary, having no effect on the energy of the system, since a relative sign simply interchanges $\psi \leftrightarrow \phi$ and interchanges the definitions of occupied and unoccupied zero-energy modes. Eliminating the first of these variables yields

$$(3.2.3.20) \quad \Phi_k(\mathcal{A} + \mathcal{B})(\mathcal{A} - \mathcal{B}) = \Lambda_k^2 \Phi_k.$$

Given that \mathcal{A} is symmetric and that \mathcal{B} is antisymmetric, it follows that $(\mathcal{A} + \mathcal{B})(\mathcal{A} - \mathcal{B}) = (\mathcal{A} - \mathcal{B})(\mathcal{A} + \mathcal{B})$ and that said product is positive semi-definite. Thus, all the Λ_k s are real and it is possible to choose all the Φ_k and Ψ_k to be real as well as orthogonal. Moreover, if the Φ_k s are normalized, then the Ψ_k s are automatically normalized when $\Lambda_k \neq 0$ or can be chosen when $\Lambda_k = 0$. This ensures that

$$(3.2.3.21) \quad \begin{aligned} \delta_{kk'} &= \sum_n \phi_{kn} \phi_{k'n} - \psi_{kn} \psi_{k'n} \\ 0 &= \sum_n \phi_{kn} \psi_{k'n} - \psi_{kn} \phi_{k'n}, \end{aligned}$$

which are the necessary and sufficient conditions for the (η, η^\dagger) -operators to be canonical fermionic operators. Furthermore, the constant in the Hamiltonian can be determined from the trace of [equation \(3.2.3.15\)](#), for the trace is invariant under canonical transformations, thus yielding

$$(3.2.3.22)$$

$$\begin{aligned} \text{Tr}\{\mathbf{H}\} &= 2^{N-1} \sum_i \mathcal{A}_{ii} \\ \text{Tr}\{\mathbf{H}\} &= 2^{N-1} \sum_k \Lambda_k + 2^N \times c \end{aligned} \Rightarrow c = \frac{1}{2} \sum_i \mathcal{A}_{ii} - \sum_k \Lambda_k,$$

$$\mathbf{H} = \sum_k \Lambda_k \left(\eta_k^\dagger \eta_k - \frac{1}{2} \right) + \frac{1}{2} \sum_i \mathcal{A}_{ii}, \quad \text{with the complete set of solutions and eigenvalues given by}$$

$$\phi_{kj} = \begin{cases} \sqrt{\frac{2}{N}} \sin kj \\ \sqrt{\frac{2}{N}} \cos kj \end{cases} \quad \Lambda_k^2 = 1 - (1 - \gamma^2) \sin^2 k.$$

where $k = \frac{2\pi m}{N}$ and $m = -\frac{N}{2}, \dots, 0, 1, \dots, \frac{N}{2} - 1$.

- If $k > 0$, the upper solution for ϕ_{kj} is taken and the lower solution for $k \leq 0$.
- The ϕ_{kj} -coefficients can be calculated from [equation \(3.2.3.19\)](#), yielding

$$(3.2.3.23) \quad \psi_{kj} = \begin{cases} \frac{1}{\Lambda_k} \left(\phi_{kj} \cos k + \gamma \phi_{-kj} \sin k \right) & \text{if } \Lambda_k \neq 0 \\ \pm \phi_{kj} & \text{if } \Lambda_k = 0 \end{cases}$$

Furthermore, $\Lambda_k = 0$ is an eigenvalue only if $\gamma = 0$ and $\frac{N}{4} \in \mathbb{Z}$.

Consider N such that $\frac{N}{4} \notin \mathbb{Z}$, hence $\Lambda_k \neq 0$. The sign of Λ_k is arbitrary and can be taken to be positive. This corresponds to a particle-hole picture for the η -particles, where the ground states has no elementary fermions and the fermion excitations, both above and below the Fermi surface,¹⁶ have positive energies. The energy excitations for this model can be visualized in [figure 1-\(a\)](#). Only in the isotropic case ($\gamma = 0$) there is no energy gap.

The ground state is defined as the state with no elementary excitations, thus

Physical Context

The Fermi surface consists of two points, $k = \frac{\pi}{2}$. The alternate picture in which all excitations are particles, which in the isotropic case is an up-spin $|\uparrow\rangle$, can be obtained by letting Λ_k have the sign of Λ_k . The Fermi sea would then be defined by $|k| > \frac{\pi}{2}$. However, the particle-hole picture is preferred.

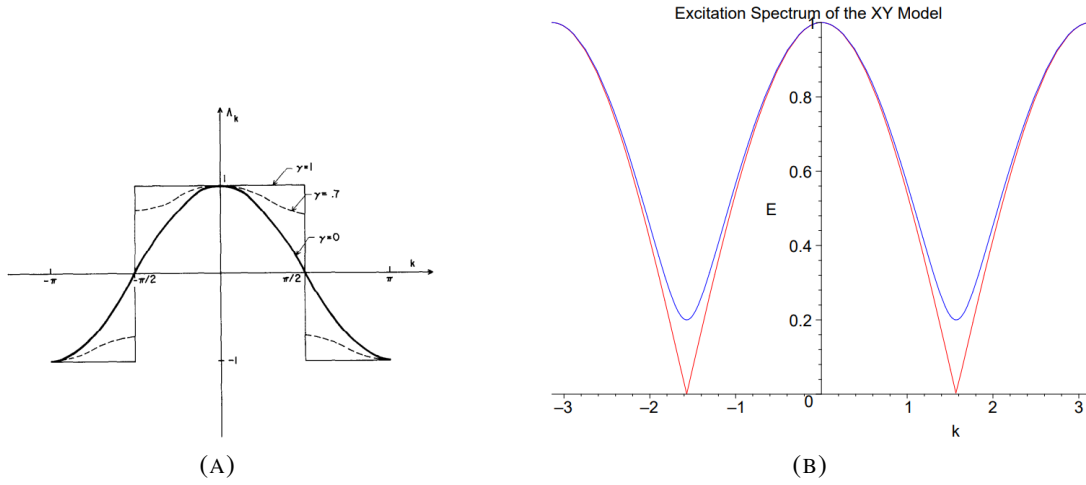


FIGURE 1. Fig. A: Energies of the elementary excitations in the $XY_{S=\frac{1}{2}}$ Heisenberg model, as function of the wave vector for three different degrees of anisotropy. Retrieved from [?]. Fig. B: Excitation spectra for the XY model for different anisotropy values γ . The **red** curve corresponds to the isotropic case $\gamma = 0$ - where the spectrum is gapless - and the **blue** curve corresponds to the anisotropic case of $\gamma = 0.2$, where a gap appears.

(3.2.3.24)

$$\begin{aligned} \eta_k |\Phi_0\rangle &= 0, \quad \forall k. \\ E_0 &= -\frac{1}{2} \sum_k \Lambda_k, \end{aligned}$$

where E_0 is the ground-state energy. Notice that the ground state has zero total spin. This zero-total-spin ground state is non-degenerate, which is the same result as in the full Heisenberg model. In the thermodynamic limit, the preceding sum can be replaced by an integral, yielding

$$\begin{aligned} \frac{E_0}{N} &= -\frac{1}{4\pi} \int_{-\pi}^{\pi} dk \left[1 - (1 - \gamma^2) \sin k \right] \\ (3.2.3.25) \quad &= -\frac{1}{\pi} \mathcal{E}(1 - \gamma^2), \quad \text{which yields the following limiting cases} \end{aligned}$$

$\frac{E_0}{N} = -\frac{1}{\pi}$ for the isotropic case, $\gamma = 0$

$\frac{E_0}{N} = -\frac{1}{2}$ for the Ising case, $\gamma = 1$.

where $\mathcal{E}(\cdot)$ is one of the complete elliptic integrals. The $\frac{E_0}{N}$ -quantity is analytic and goes smoothly between the $(\gamma = 0)$ - and $(\gamma = 1)$ -limiting cases. In the thermodynamic limit $N \rightarrow \infty$, there are always gapless excitations near $k = \pm\frac{\pi}{2}$. These have a dispersion, $k = \pm\frac{\pi}{2} + q$, $\Lambda_k = |\sin q| \approx q$ as $q \rightarrow 0$. In **figure 1**-(b) plots of the excitation spectra are shown for different values of γ .

Physical Context

The previous relations hold, in principle, only for the c -cyclic problem, rather than the full a -problem. The Hamiltonian in the a -cyclic term has an additional non-local term

$$-\frac{1}{2} \left[\mathbf{c}_N^\dagger \mathbf{c}_1 + \mathbf{c}_1^\dagger \mathbf{c}_N + \gamma \mathbf{c}_N^\dagger \mathbf{c}_1^\dagger + \gamma \mathbf{c}_1 \mathbf{c}_N \right] (e^{i\pi \mathfrak{F}} + 1).$$

The \mathfrak{F} is not invariant under the transformations shown in [equation \(3.2.3.5\)](#), but its evenness or oddness does remain invariant, with the same being true for $e^{i\pi \mathfrak{F}}$.

- Now, in the ground state of the c -cyclic problem, given by [equation \(3.2.3.24\)](#), and in all states with an even number of excitations, the number of c -particles is odd (check the definition of the c -particles in terms of the η -particles, given by [equation \(3.2.3.16\)](#), and consider that the ground state is defined in terms of the η -particles)^a. Hence an additional term gives zero net difference, when acting on said states and thus, for an even number of excitations, these states remain eigenstates of the a -cyclic problem.
- For states with an odd number of excitations, on the other hand, \mathfrak{F} is even, giving rise to an additional term,

$$-\frac{1}{2} \left(\mathbf{c}_N^\dagger \mathbf{c}_1 + \gamma \mathbf{c}_N^\dagger \mathbf{c}_1 + \text{h.c.} \right),$$

in the Hamiltonian. This induces corrections of $\mathcal{O}(N^{-1})$ on the k s, the Φ_k 's and the Ψ_k 's, all of which can be calculated exactly but are negligible in the calculation of physical quantities. Strictly speaking, the elementary excitations are not independent for the a -cyclic problem, since there is a dependence on the evenness or oddness of the total number of excitations present.

^aassuming $\frac{N}{4} \notin \mathbb{Z}$, the k s are occupied symmetrically around $K = 0$, except that $k = \pi$, but no $k = -\pi$, is occupied.

The free energy is given by the grand potential of such a system of non-interacting fermions with zero chemical potential,

(3.2.3.26)

$$\frac{F_\gamma}{N} = -\beta^{-1} \left[\alpha + \frac{2}{\pi} \int_0^{\frac{\pi}{2}} dk \log \cosh \frac{\beta \Lambda_k}{2} \right], \quad \text{which yields the following limiting cases}$$

$$\frac{F_{\text{isotrop.}}}{N} = -\beta^{-1} \left[\alpha + \frac{2}{\pi} \int_0^{\frac{\pi}{2}} dk \log \cosh \frac{\beta \cos k}{2} \right]$$

$$\frac{F_{\text{isotrop.}}}{N} = -\frac{\alpha + \log \cosh \frac{\beta}{2}}{\beta}$$

where $\alpha = \log 2$. Neither case exhibits any singular behaviour as a function of the temperature, a result to be expected in view of the one-dimensional nature of the model.

- Antiferromagnetic case, isotropic case:

In the isotropic case, the mathematical treatment is simpler than that of the general anisotropic case. Since some of the eigenenergies will be negative, it is convenient to introduce an additional transformation, which reads

$$\begin{aligned} \zeta_k = \eta_k & \quad \text{if } \Lambda_k \geq 0, \\ \zeta_k = -\eta_k & \quad \text{if } \Lambda_k < 0. \end{aligned} \Rightarrow \mathbf{H}_c = \sum_{k, \Lambda_k \geq 0} \Lambda_k \zeta_k^\dagger \zeta_k + \sum_{k, \Lambda_k < 0} \Lambda_k \zeta_k \zeta_k^\dagger$$

$$= \sum_k |\Lambda_k| \zeta_k^\dagger \zeta_k - \sum_{k, \Lambda_k < 0} |\Lambda_k|$$

$$= \sum_k |\Lambda_k| \left(\zeta_k^\dagger \zeta_k - \frac{1}{2} \right),$$

(3.2.3.27)

where $\text{Tr}\{\mathbf{A}\} = \sum_k \Lambda_k = 0$ was used. Here, the ζ_k -particles are fermionic in nature, thus the ground state satisfies $\zeta_k |0\rangle = 0$, $\forall k$, with the ζ_k^\dagger -operator generating elementary fermionic excitations with $|\Lambda_k|$ -energy above the ground state. Using translational invariance, the eigenvectors and eigenvalues of the \mathcal{A} -matrix are

$$(3.2.3.28) \quad \phi_{kj} = \frac{1}{\sqrt{N}} e^{ikj}, \quad \Lambda_k = \cos k, \quad \text{with } k = \frac{2\pi n}{N} \text{ for } -\frac{N}{2} \leq n \leq \frac{N}{2} - 1.$$

Hence, in the thermodynamic limit $N \rightarrow \infty$ there are always gapless excitations near the Fermi sphere, $k = \pm \frac{\pi}{2}$. These have a dispersion $k = \pm \frac{\pi}{2} + q$ with $\epsilon(k) = |\sin q| \approx |q|$ as $q \rightarrow 0$. The ground state energy per spin becomes

$$\frac{U}{N} = -\frac{1}{N} \sum_k \frac{|\Lambda_k|}{2} \rightarrow -\int_{-\pi}^{\pi} \frac{dk}{4\pi} |\cos k| = -\frac{1}{\pi}.$$

Consider the $(N \in 2\mathbb{N} \wedge \frac{N}{4} \notin \mathbb{Z})$ -case, so that the ground state is non-degenerate. In said case, the ground state of \mathbf{H}_c is an exact eigenstate of \mathbf{H} with the same eigenvalue. At any rate, the ground state is - at most - four-times degenerate. Then, the ground state spin properties read

$$(3.2.3.29) \quad \begin{aligned} \mathbf{S}^z &= \sum_i \mathbf{c}_i^\dagger \mathbf{c}_i - \frac{N}{2} = \sum_k \eta_k^\dagger \eta_k - \frac{N}{2} = \sum_{k, \Lambda_k \geq 0} \zeta_k^\dagger \zeta_k + \sum_{k, \Lambda_k < 0} \zeta_k \zeta_k^\dagger - \frac{N}{2} \\ &= \sum_{k, \Lambda_k \geq 0} \zeta_k^\dagger \zeta_k + \sum_{k, \Lambda_k < 0} (1 - \zeta_k^\dagger \zeta_k) - \frac{N}{2} \\ &= \sum_k \text{sgn}(\Lambda_k) \zeta_k^\dagger \zeta_k + \sum_{k, \Lambda_k < 0} 1 - \frac{N}{2} \end{aligned}$$

Thus, for excitations with $|k| < \frac{\pi}{2}$ carry $\mathbf{S}^z = 1$ and for excitations with $|k| > \frac{\pi}{2}$ carry $\mathbf{S}^z = -1$

$$(3.2.3.30) \quad \Rightarrow \quad \text{The total spin of the ground state is } \mathbf{S}^z |\Omega\rangle = \left(\sum_{k, \Lambda_k < 0} 1 - \frac{N}{2} \right) |\Omega\rangle = 0,$$

ie. the ground state is non-degenerate and carries zero total spin, as in the full Heisenberg model.

- Ferromagnetic case

The previous results can be trivially generalized to the case of ferromagnetic coupling, wherein the Hamiltonian reads

$$(3.2.3.31) \quad \mathbf{H}_{\text{ferro}} = -\mathbf{H} \approx -\mathbf{H}_c = \sum_k -|\Lambda_k| \left(\zeta_k^\dagger \zeta_k - \frac{1}{2} \right) = \sum_k |\Lambda_k| \left(\zeta_k \zeta_k^\dagger - \frac{1}{2} \right).$$

Notice that the new, ferromagnetic, ground state now satisfies $\zeta_k^\dagger |0_F\rangle = 0$, $\forall k$. Hence, the ferromagnetic ground state can be obtained from the antiferromagnetic ground state by turning on all excitations, and the ferromagnetic excitations are obtained by removing the antiferromagnetic ones. In particular, the ground state energy and the excitation spectrum are the same as in the ferromagnetic case. The ground state is still non-degenerate and has zero total spin.

$$(3.2.3.32) \quad \begin{aligned} \mathbf{S}^z |0_F\rangle &= \left(\sum_k \text{sgn}(\Lambda_k) \zeta_k^\dagger \zeta_k + \sum_{k, \Lambda_k < 0} 1 - \frac{N}{2} \right) |0_F\rangle = \left(\sum_k \text{sgn}(\Lambda_k) \zeta_k^\dagger \zeta_k + \sum_{k, \Lambda_k < 0} 1 - \frac{N}{2} \right) |0_F\rangle \\ &= \left(\sum_{k, \Lambda_k \geq 0} 1 - \frac{N}{2} |0_F\rangle \right) = 0. \end{aligned}$$

This is in clear contrast to the full ferromagnetic Heisenberg model, where the ground state is greatly degenerate, with one of the ground states carrying $S^z = \frac{N}{2}$.

3.2.3.3. Short and long-range order in the Ground state. The long-range order for the Heisenberg model is often defined in terms of two sublattices¹⁷. It is taken to be the preponderance of spins up to spins down on one of the sublattices or viceversa.

Physical Context

For the ($\gamma = 1$)-case, the ground state has long range Néel order, ie. $S_i^x |\Psi_0\rangle = (-1)^i |\Psi_0\rangle$ for the antiferromagnetic case, while in the ferromagnetic case $S_i^x |\Psi_0\rangle = |\Psi_0\rangle$. Therefore it is interesting to understand what happens in the intermediate values of γ . It is clear that states with spin- x - and spin- y -ordered components start to compete, yet it is not obvious whether the ground state has any long and/or short range order at $\gamma = 0$.

- Since the Heisenberg Hamiltonian is $SU(2)$ -invariant, in particular it is invariant under a rotation in $\theta^i = \pi$, $i = x, y, z$,
- is translationally invariant,
- and the ground state is non-degenerate,

then, according to this definition for long-range order, it must be zero even if the state itself was ordered. However note that the completely ferromagnetic states can have long-range order by this definition only because these are highly degenerate. This definition for long-range order, though somewhat lacking, is useful since the approximate states considered have not always had the full symmetries of the Hamiltonian.

A much better measure for the long-range order is the time dependent dynamical spin correlation tensor,

$$(3.2.3.33) \quad \mathcal{S}^{mn}(t) = \langle \Psi_0 | S_n(0) S_m(t) | \Psi_0 \rangle \rightarrow \mathcal{S}^{mn} = \langle \Psi_0 | S_n(0) S_m | \Psi_0 \rangle,$$

where the second expression is the ($t = 0$)-contraction of the dynamical spin correlation tensor, which enters in the calculation of any process conceived to directly measure order. Then,

$$(3.2.3.34)$$

$$\begin{aligned} \mathcal{S}_{nm}^x &= \langle \Psi_0 | S_n^x S_m^x | \Psi_0 \rangle = \frac{1}{4} \langle \Psi_0 | (\mathbf{c}_n^\dagger + \mathbf{a}_n)(\mathbf{c}_m^\dagger + \mathbf{a}_m) | \Psi_0 \rangle \\ &= \frac{1}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger e^{-i\pi \sum_{j=1}^{n-1} \mathbf{c}_j^\dagger \mathbf{c}_j} + \mathbf{c}_n e^{i\pi \sum_{j=1}^{n-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \right) \left(\mathbf{c}_m^\dagger e^{-i\pi \sum_{j=1}^{m-1} \mathbf{c}_j^\dagger \mathbf{c}_j} + \mathbf{c}_m e^{i\pi \sum_{j=1}^{m-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \right) | \Psi_0 \rangle \\ &= \frac{1}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger + \mathbf{c}_n e^{2i\pi \sum_{j=1}^{n-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \right) e^{-i\pi \sum_{j=1}^{n-1} \mathbf{c}_j^\dagger \mathbf{c}_j} e^{i\pi \sum_{j=1}^{m-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \left(\mathbf{c}_m^\dagger e^{-2i\pi \sum_{j=1}^{m-1} \mathbf{c}_j^\dagger \mathbf{c}_j} + \mathbf{c}_m \right) | \Psi_0 \rangle \\ &= \frac{1}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger + \mathbf{c}_n \right) e^{i\pi \sum_{j=n}^{m-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \left(\mathbf{c}_m^\dagger + \mathbf{c}_m \right) | \Psi_0 \rangle \\ &= \frac{1}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger + \mathbf{c}_n \right) \prod_{j=n}^{m-1} (1 - 2\mathbf{c}_j^\dagger \mathbf{c}_j) \left(\mathbf{c}_m^\dagger + \mathbf{c}_m \right) | \Psi_0 \rangle \\ &= \frac{1}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger + \mathbf{c}_n \right) (1 - 2\mathbf{c}_n^\dagger \mathbf{c}_n) \prod_{j=n+1}^{m-1} (1 - 2\mathbf{c}_j^\dagger \mathbf{c}_j) \left(\mathbf{c}_m^\dagger + \mathbf{c}_m \right) | \Psi_0 \rangle \\ &= \frac{1}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger - \mathbf{c}_n \right) \prod_{j=n+1}^{m-1} (\mathbf{c}_j^\dagger + \mathbf{c}_j)(\mathbf{c}_j^\dagger - \mathbf{c}_j) \left(\mathbf{c}_m^\dagger + \mathbf{c}_m \right) | \Psi_0 \rangle. \end{aligned}$$

Now, let $\mathbf{A}_n = \mathbf{c}_n^\dagger + \mathbf{c}_n$ and $\mathbf{B}_n = \mathbf{c}_n^\dagger - \mathbf{c}_n$, such that

¹⁷In the case of a linear one-dimensional chain, these sublattices may be defined to be, one corresponding to all even sites and the other one corresponding to all odd sites.

(3.2.3.35)

$$\begin{aligned}
\mathbf{A}_n = \mathbf{c}_n^\dagger + \mathbf{c}_n \quad \mathbf{B}_n = \mathbf{c}_n^\dagger - \mathbf{c}_n &\Rightarrow \begin{cases} \{\mathbf{A}_n, \mathbf{A}_m\} = \delta_{ij} \\ \{\mathbf{B}_n, \mathbf{B}_m\} = -2\delta_{ij} \\ \{\mathbf{A}_n, \mathbf{B}_m\} = 0 \end{cases} \Rightarrow \mathcal{S}_{nm}^x = \frac{1}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger - \mathbf{c}_n \right) \prod_{j=n+1}^{m-1} (\mathbf{c}_j^\dagger + \mathbf{c}_j)(\mathbf{c}_j^\dagger - \mathbf{c}_j) (\mathbf{c}_m^\dagger + \mathbf{c}_m) | \Psi_0 \rangle \\
&= \frac{1}{4} \langle \Psi_0 | \mathbf{B}_n \prod_{j=n+1}^{m-1} \mathbf{A}_j \mathbf{B}_j \mathbf{A}_j | \Psi_0 \rangle \\
&= \frac{1}{4} \langle \Psi_0 | \mathbf{B}_n \mathbf{A}_{n+1} \mathbf{B}_{n+1} \mathbf{A}_{n+2} \cdots \mathbf{A}_{m-1} \mathbf{B}_{m-1} \mathbf{A}_m | \Psi_0 \rangle
\end{aligned}$$

Similarly, for the other two correlators

(3.2.3.36)

$$\begin{aligned}
\mathcal{S}_{nm}^y &= \langle \Psi_0 | \mathbf{S}_n^y \mathbf{S}_m^y | \Psi_0 \rangle = \frac{1}{4} \langle \Psi_0 | (\mathbf{a}_n^\dagger - \mathbf{a}_n)(\mathbf{a}_m^\dagger - \mathbf{a}_m) | \Psi_0 \rangle \\
&= \frac{1}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger e^{-i\pi \sum_{j=1}^{n-1} \mathbf{c}_j^\dagger \mathbf{c}_j} - \mathbf{c}_n e^{i\pi \sum_{j=1}^{n-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \right) \left(\mathbf{c}_m^\dagger e^{-i\pi \sum_{j=1}^{m-1} \mathbf{c}_j^\dagger \mathbf{c}_j} - \mathbf{c}_m e^{i\pi \sum_{j=1}^{m-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \right) | \Psi_0 \rangle \\
&= \frac{1}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger - \mathbf{c}_n e^{2i\pi \sum_{j=1}^{n-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \right) e^{-i\pi \sum_{j=1}^{n-1} \mathbf{c}_j^\dagger \mathbf{c}_j} e^{i\pi \sum_{j=1}^{m-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \left(\mathbf{c}_m^\dagger e^{-2i\pi \sum_{j=1}^{m-1} \mathbf{c}_j^\dagger \mathbf{c}_j} - \mathbf{c}_m \right) | \Psi_0 \rangle \\
&= \frac{(-1)^{m-n}}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger - \mathbf{c}_n \right) e^{i\pi \sum_{j=n}^{m-1} \mathbf{c}_j^\dagger \mathbf{c}_j} \left(\mathbf{c}_m^\dagger - \mathbf{c}_m \right) | \Psi_0 \rangle \\
&= \frac{(-1)^{m-n}}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger - \mathbf{c}_n \right) \prod_{j=n}^{m-1} (1 - 2\mathbf{c}_j^\dagger \mathbf{c}_j) \left(\mathbf{c}_m^\dagger - \mathbf{c}_m \right) | \Psi_0 \rangle \\
&= \frac{(-1)^{m-n}}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger - \mathbf{c}_n \right) (1 - 2\mathbf{c}_n^\dagger \mathbf{c}_n) \prod_{j=n+1}^{m-1} (1 - 2\mathbf{c}_j^\dagger \mathbf{c}_j) \left(\mathbf{c}_m^\dagger - \mathbf{c}_m \right) | \Psi_0 \rangle \\
&= \frac{(-1)^{m-n}}{4} \langle \Psi_0 | \left(\mathbf{c}_n^\dagger + \mathbf{c}_n \right) \prod_{j=n+1}^{m-1} (\mathbf{c}_j^\dagger - \mathbf{c}_j)(\mathbf{c}_j^\dagger + \mathbf{c}_j) \left(\mathbf{c}_m^\dagger - \mathbf{c}_m \right) | \Psi_0 \rangle \\
&= \frac{(-1)^{m-n}}{4} \langle \Psi_0 | \mathbf{A}_n \prod_{j=n+1}^{m-1} \mathbf{B}_j \mathbf{A}_j \mathbf{B}_j | \Psi_0 \rangle \\
&= \frac{(-1)^{m-n}}{4} \langle \Psi_0 | \mathbf{A}_n \mathbf{B}_{n+1} \mathbf{A}_{n+1} \mathbf{B}_{n+2} \cdots \mathbf{B}_{m-1} \mathbf{A}_{m-1} \mathbf{B}_m | \Psi_0 \rangle \\
\mathcal{S}_{nm}^z &= \langle \Psi_0 | \mathbf{S}_n^z \mathbf{S}_m^z | \Psi_0 \rangle = \langle \Psi_0 | \left(\mathbf{a}_n^\dagger \mathbf{a}_n - \frac{1}{2} \right) \left(\mathbf{a}_m^\dagger \mathbf{a}_m - \frac{1}{2} \right) | \Psi_0 \rangle = \frac{1}{4} \langle \Psi_0 | \mathbf{A}_i \mathbf{B}_i \mathbf{B}_i \mathbf{A}_i | \Psi_0 \rangle.
\end{aligned}$$

These expressions may be evaluated using Wick's theorem, which allows to express the vacuum expectation value of a product of operators, all of which obey anticommutation rules, in terms of contractions of pairs. A particular simplification occurs in evaluating the previous set of expectation values. The basic contractions may be readily calculated,

$$(3.2.3.37) \quad \langle \mathbf{A}_i \mathbf{A}_j \rangle = \sum_k \phi_{ki} \phi_{kj} = \delta_{ij} \quad \langle \mathbf{B}_i \mathbf{B}_j \rangle = - \sum_k \psi_{ki} \psi_{kj} = -\delta_{ij} \quad \langle \mathbf{B}_i \mathbf{A}_j \rangle = - \sum_k \phi_{ki} \psi_{kj} = \mathfrak{G}_{ij}.$$

Thus, the most straightforward pairing contribution to \mathcal{S}_{nm}^x is

$$(3.2.3.38) \quad \langle \mathbf{B}_n \mathbf{A}_{n+1} \rangle \langle \mathbf{B}_{n+1} \mathbf{A}_{n+2} \rangle \cdots \langle \mathbf{B}_{m-1} \mathbf{A}_m \rangle,$$

since all other pairings can be obtained from this one via permutations of the \mathbf{A} -operators, keeping the \mathbf{B} -operators fixed. Since the number of crossings of \mathbf{B} -operators by \mathbf{A} -operators is always even, the sign associated with a given permutation is $(-1)^{-\text{sgn}(\sigma)}$, with $\text{sgn}(\sigma)$ the signature of the permutation of the \mathbf{A} -operators. Thus, the following identities hold

$$\begin{aligned}
\mathcal{S}_{nm}^x &= \frac{1}{4} \sum_{\sigma \in \mathcal{S}[n+1, m]} (-1)^{\text{sgn}(\sigma)} \mathfrak{G}_{n, \sigma(n+1)} \cdots \mathfrak{G}_{m-1, \sigma(m)} = \frac{1}{4} \begin{vmatrix} \mathfrak{G}_{n, n+1} & \mathfrak{G}_{n, n+2} & \cdots & \mathfrak{G}_{nm} \\ \vdots & & & \vdots \\ \mathfrak{G}_{m-1, n+1} & \mathfrak{G}_{m-1, n+2} & \cdots & \mathfrak{G}_{m-1, m} \end{vmatrix} \\
\mathcal{S}_{nm}^y &= \frac{1}{4} \sum_{\sigma \in \mathcal{S}[n, m-1]} (-1)^{\text{sgn}(\sigma)} \mathfrak{G}_{n+1, \sigma(n)} \cdots \mathfrak{G}_{m, \sigma(m-1)} = \frac{1}{4} \begin{vmatrix} \mathfrak{G}_{n+1, n} & \mathfrak{G}_{n+1, n+1} & \cdots & \mathfrak{G}_{n+1, m-1} \\ \vdots & & & \vdots \\ \mathfrak{G}_{m, n} & \mathfrak{G}_{m, n+1} & \cdots & \mathfrak{G}_{m, m-1} \end{vmatrix} \\
\mathcal{S}_{nm}^z &= \frac{1}{4} \left(\langle \mathbf{A}_n \mathbf{B}_m \rangle \langle \mathbf{A}_m \mathbf{B}_m \rangle - \langle \mathbf{A}_n \mathbf{B}_m \rangle \langle \mathbf{B}_m \mathbf{A}_n \rangle \right) = \frac{1}{4} (\mathfrak{G}_{nn} \mathfrak{G}_{m, m} - \mathfrak{G}_{m, n} \mathfrak{G}_{n, m}), \text{ for } n < m,
\end{aligned}
\tag{3.2.3.39}$$

and where $\mathcal{S}[a, b]$ is the group of permutations of integers $\mathbb{Z}_{[a, b]}$. Note that the first two submatrices of [equation \(3.2.3.39\)](#) can be written as

$$\begin{aligned}
(\mathcal{S}_{nm}^x)_{ij} &= \frac{1}{4} (\mathfrak{G}_{nm}^x)_{ij} & (\mathcal{S}_{nm}^y)_{ij} &= \frac{1}{4} (\mathfrak{G}_{nm}^y)_{ij} \\
&= \frac{1}{4} \mathfrak{G}_{n+i-1, j+n} & &= \frac{1}{4} \mathfrak{G}_{n+i, n+j-1}
\end{aligned}$$

Thus, both \mathcal{S}_{nm}^x and \mathcal{S}_{nm}^y are particular subdeterminants of $\det \mathfrak{G}$, where \mathfrak{G} is a matrix with entries

$$(\mathfrak{G})_{ij} = -\langle c \rangle_{ij}, \quad \begin{array}{l} \text{where } \Phi \text{ are the matrices of } \phi_{k,i} \\ \text{where } \Psi \text{ are the matrices of } \psi_{k,i}. \end{array}$$

It is obvious that this \mathfrak{G} -matrix is unitary, since both Ψ and Φ are unitary as well, and its determinant is then ± 1 . Its actual sign can be found to be $\det \mathfrak{G} = (-1)^N \det \mathfrak{G}$.

In order to calculate this Green function matrix, defined in [equation \(3.2.3.37\)](#), the thermodynamical limit is taken, yielding

$$\begin{aligned}
\mathfrak{G}_{ij} &= - \sum_k \psi_{ki} \phi_{kj} = - \sum_k \frac{1}{\Lambda_k} \left(\phi_{ki} \cos k + \gamma \phi_{-ki} \sin k \right) \phi_{k,j} \\
&= \frac{1 + (-1)^{i-j+1}}{2} \left(-\frac{2}{\pi} \right) \int_0^{\frac{\pi}{2}} dk \frac{1}{\Lambda_k} \left(\cos k \cos k(i-j) - \gamma \sin k \sin k(i-j) \right),
\end{aligned}
\tag{3.2.3.40}$$

where [equation \(3.2.3.23\)](#) was used in the second equality in the first line. Notice that \mathfrak{G} depends only on the difference $i - j$, a consequence of the cyclic-ness of the spin chain and its translational invariance. Hence, $\mathfrak{G}_{ij} = \mathfrak{G}_{i-j}$, being zero if $i - j$ is even. In the limiting cases, closed expressions can be obtained

$$\begin{aligned}
\mathfrak{G}_{ij} &= \begin{cases} -\delta_{i-j, -1} & \text{for } i - j \text{ odd} \\ 0 & \text{for } i - j \text{ even} \end{cases} & \text{in the isotropic limit, } \gamma = 0, \\
\mathfrak{G}_{ij} &= \begin{cases} (-1)^{\frac{i-j+1}{2}} \frac{2}{\pi(i-j)} & \text{for } i - j \text{ odd} \\ 0 & \text{for } i - j \text{ even} \end{cases} & \text{in the Ising limit, } \gamma = 1.
\end{aligned}
\tag{3.2.3.41}$$

With these elements, the ground state order can be evaluated. In particular, the ground state correlations for the XY-model between nearest-neighbours at finite anisotropy,

$$\begin{aligned}
\mathcal{S}_1^x &= \mathcal{S}_{i, i+1}^x = \frac{1}{4} \mathfrak{G}_{i, i+1} \equiv \frac{1}{4} \mathfrak{G}_{-1} \\
\mathcal{S}_1^y &= \mathcal{S}_{i, i+1}^y = \frac{1}{4} \mathfrak{G}_{i+1, i} \equiv \frac{1}{4} \mathfrak{G}_1 \text{ where } \mathfrak{G}_{\pm 1} = -\frac{2}{\pi} \frac{1}{1 \mp \gamma} \left(\mathcal{E}(\sqrt{1 - \gamma^2}) \mp \gamma \mathcal{K}(\sqrt{1 - \gamma^2}) \right), \\
\mathcal{S}_1^z &= \mathcal{S}_{i, i+1}^z = \frac{1}{4} \left(\mathfrak{G}_{i, i} \mathfrak{G}_{i+1, i+1} - \mathfrak{G}_{i, i+1} \mathfrak{G}_{i+1, i} \right) \equiv \frac{1}{4} \mathfrak{G}_{-1} \mathfrak{G}_1
\end{aligned}
\tag{3.2.3.42}$$

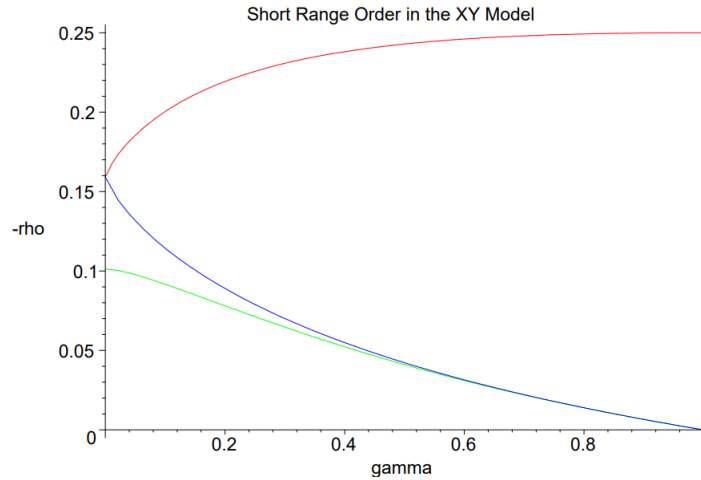


FIGURE 2. Nearest-neighbour correlators plotted as function of the anisotropy γ , where, the red curve corresponds to S_1^x , the blue curve corresponds to S_1^y , and where the green curve corresponds to S_1^z . From these plots it is clear that the x - and y -correlations compete as $\gamma \rightarrow 0$.

where the elliptic E and K functions were introduced. These correlators are plotted in figure 2. For any value of γ , all of these correlations are negative in value, and hence the nearest-neighbours display antiferromagnetic correlations.

- For the Ising limit, $\gamma = 1$, $S_1^x = \frac{1}{4}$ and $S_1^y = S_1^z = 0$.
- As γ decreases, the antiferromagnetic correlation between the y -components induce an anti-ferromagnetic correlation between the z -components.
- Finally, in the isotropic limit $\gamma = 0$, the x - and y -correlations between nearest-neighbours spins become equal (as expected for symmetry reasons), and $S_1^x = S_1^y = -\frac{1}{2\pi}$, $S_1^z = -\frac{1}{\pi^2}$. So, short-range order persists in the isotropic limit.

3.2.3.4. $\Delta \approx 1$: **The isotropic Heisenberg Antiferromagnet.** The most interesting regime for the XXZ chain is $\Delta \approx 1$, ie. the vicinity of the Isotropic Heisenberg antiferromagnet (HAF), wherein the ground state energy is given by

$$E_0 = -NJ \log 2.$$

3.3. General considerations of the Heisenberg model.

3.3.1. *Antiferromagnetic Ground States.* Consider an N -lattice system with spin- S , with Hamiltonian

$$(3.3.1.1) \quad \begin{aligned} \mathbf{H} &= \mathbf{H}^{zz} + \mathbf{H}^{xy}, \\ \mathbf{H}^{zz} &= \frac{1}{2} \sum_{ij \in \Lambda \subset \mathbb{Z}^d} J_{ij} \mathbf{S}_i^z \mathbf{S}_j^z, \\ \mathbf{H}^{xy} &= \frac{1}{4} \sum_{ij \in \Lambda \subset \mathbb{Z}^d} J_{ij} (\mathbf{S}_i^+ \mathbf{S}_j^- + \mathbf{S}_i^- \mathbf{S}_j^+), \end{aligned}$$

with $J_{ij} = J_{ji}$ having translational symmetry. Furthermore, this Hamiltonian is rotationally invariant, having an $SU(2)$ -symmetry, since it commutes with all three components of the total spin. Thus, the eigenstates may be labelled as

$$(3.3.1.2) \quad \begin{aligned} |\Phi\rangle &= |\mathbf{S}_{\text{tot}}, M, \dots\rangle, \\ M &= -\mathbf{S}_{\text{tot}}, -\mathbf{S}_{\text{tot}} + 1, \dots, \mathbf{S}_{\text{tot}}, \\ \mathbf{S}_{\text{tot}} &\leq NS, \end{aligned} \quad \begin{aligned} &\text{where } M \text{ is the eigenvalue of the total magnetization } \mathbf{S}_{\text{tot}}^z, \\ &\text{and where } \mathbf{S}_{\text{tot}}^2 = \mathbf{S}_{\text{tot}}(\mathbf{S}_{\text{tot}} + 1) \text{ is the } SU(2)\text{-Casimir invariant.} \\ &\text{Note that, if the model is translationally invariant} \\ &\text{the k-momentum is also a good quantum number.} \end{aligned}$$

The antiferromagnetic ground state is, in general, far more complicated than the ferromagnetic ground state. There are a series of theorems on the ground states, which hold for bipartite antiferromagnetic Hamiltonians. Consider two lattices, A and B , then the staggered magnetization operator on these sublattices is given by

$$(3.3.1.3) \quad \mathbf{S}^{\text{stagg}} = \sum_{i \in A} \mathbf{S}_i^z - \sum_{i \in B} \mathbf{S}_i^z.$$

The Ising configurations form the basis set

$$\prod_{i=1}^N |S, m_i^\alpha\rangle_i \text{ where } |S, m_i\rangle_i \text{ is an eigenstate of } \mathbf{S}_i^2, \mathbf{S}_i^z \text{ with eigenvalues } S(S+1), m_i \text{ respectively.}$$

In this bipartite system, the Néel state maximizes the staggered magnetization $\mathbf{S}^{\text{stagg}}$ and is given by the Ising configuration $|\Psi^{\text{Néel}}\rangle = \prod_{i=1}^N |S, \eta_i S_i\rangle$, $\eta_i = \begin{cases} 1 & i \in A \\ -1 & i \in B \end{cases}$

Generally, this state is not an eigenstate of the Hamiltonian, since the spin flip terms $\mathbf{S}_i^+ \mathbf{S}_j^-$ connect the Néel state to other (non-Néel) Ising spin configurations. Now, it is convenient to rotate the spin axes on the B -sublattice about the z -axis, which amounts to the unitary transformation

$$(3.3.1.4) \quad \begin{aligned} \mathbf{S}_i^+ &\rightarrow -\tilde{\mathbf{S}}_i^+ \\ \mathbf{S}_i^- &\rightarrow -\tilde{\mathbf{S}}_i^- \\ \mathbf{S}_i^z &\rightarrow +\tilde{\mathbf{S}}_i^z, \end{aligned} \quad |S, m_i\rangle_i \rightarrow |S, \tilde{m}_i\rangle_i = \begin{cases} |S, m_i\rangle_i & i \in A \\ (1-)^{S+m_i} |S, m_i\rangle_i & i \in B \end{cases},$$

where \tilde{m}_i are eigenvalues of \mathbf{S}_i^z and $\tilde{\mathbf{S}}_i^z$ on A and B sublattices, respectively. By restricting the problem to an M -subspace, all wave functions may be rewritten as

$$(3.3.1.5) \quad |\Psi^M\rangle = \sum_{\alpha} f_{\alpha}^M |\tilde{\Psi}_{\alpha}^M\rangle, \text{ where } |\tilde{\Psi}_{\alpha}^M\rangle = \prod_{\sum_i \tilde{m}_i^{\alpha} = M} |S, \tilde{m}_i^{\alpha}\rangle.$$

Then, the following theorems, first presented by Marshall and then extended by Lieb and Mattis, arise

Theorem 2 Consider the Hamiltonian *equation (3.3.1.8)* of a bipartite lattice, with antiferromagnetic exchanges $J_{ij} \geq 0$ which connect two sublattices A and B . Said antiferromagnetic exchanges are such that any two lattice sites are connected by a sequence of finite exchanges between intermediary sites. In any allowed M -sector, the lowest energy state $|\Psi_0^M\rangle$ can be chosen to have positive-definite coefficients in the rotated Ising basis $|\tilde{\Psi}_\alpha^M\rangle$ ie.

$$(3.3.1.6) \quad |\Psi_0^M\rangle = \sum_{\alpha} f_{\alpha}^M |\tilde{\Psi}_{\alpha}^M\rangle, \quad \{f_{\alpha}^M\}_{\alpha} > 0.$$

Therefore, the coefficients of $|\Psi_0^M\rangle$, in terms of the unrotated Ising configurations $|\Psi_{\alpha}\rangle$ obey the **Marshall sign criterion**,

$$(3.3.1.7) \quad |\Psi_0^M\rangle = \sum_{\alpha} (-1)^{\Gamma(\alpha)} f_{\alpha}^M |\Psi_{\alpha}^M\rangle, \quad \Gamma(\alpha) = \sum_{i \in B} (S + m_i^{\alpha}).$$

Theorem 3 The absolute ground state $|\Omega\rangle$, for equal size sublattices A and B , is a singlet of zero total spin, this is

$$\mathbf{S}_{\text{tot}} |\Psi_0^M\rangle = 0.$$

Note that not all total singlets obey the so-called Marshall sign criterion and, conversely, not all Marshall sign-obeying states are singlets of the total spin. The ground state of the Heisenberg Hamiltonian must, however, obey both conditions.

Proof. Consider an $SU(2)$ -rotation, acting trivially on the first lattice and non-trivially in the second system. This is,

$$U \in SU(2) | U : (\mathfrak{su}(2)_A) \times (\mathfrak{su}(2)_B) \rightarrow (\mathfrak{su}(2)_A) \times (\mathfrak{su}(2)_B) \wedge \quad ()$$

Consider an $SU(2)$ -rotation, under which the spin operators and Ising basis transform as *equation (3.3.1.4)*. The Heisenberg Hamiltonian in this new representation then reads

$$(3.3.1.8) \quad \begin{aligned} \tilde{\mathbf{H}} &= \tilde{\mathbf{H}}^{zz} + \tilde{\mathbf{H}}^{xy}, \\ \tilde{\mathbf{H}}^{zz} &= \sum_{ij \in \Lambda \subset \mathbb{Z}^d} |J_{ij}| \mathbf{S}_i^z \tilde{\mathbf{S}}_j^z, \\ \tilde{\mathbf{H}}^{xy} &= -\frac{1}{2} \sum_{ij \in \Lambda \subset \mathbb{Z}^d} |J_{ij}| (\mathbf{S}_i^+ \tilde{\mathbf{S}}_j^- + \mathbf{S}_i^- \tilde{\mathbf{S}}_j^+), \end{aligned}$$

wherein, in this new sublattice-rotated representation,

the longitudinal term, $\tilde{\mathbf{H}}^{zz}$, is diagonal in the rotated Ising configuration: $\tilde{\mathbf{H}}^{zz} |\tilde{\Psi}_{\alpha}\rangle = e_{\alpha} |\tilde{\Psi}_{\alpha}\rangle$ and where the transverse $\tilde{\mathbf{H}}^{xy}$ has non-positive matrix elements: $\langle \tilde{\Psi}_{\alpha} | \tilde{\mathbf{H}}^{xy} | \tilde{\Psi}_{\beta} \rangle = -|K_{\alpha\beta}|$.

Both of these matrix representations induce the following eigenvalue equation on the f -coefficients,

$$\tilde{\mathbf{H}} |\tilde{\Psi}_{\alpha}\rangle = E |\tilde{\Psi}_{\alpha}\rangle \Rightarrow - \sum_{\beta} |K_{\alpha\beta}| f_{\beta} + e_{\alpha} f_{\alpha} = E f_{\alpha}.$$

Consider then the trial function $|\tilde{\Omega}\rangle = \sum_{\alpha} |f_{\alpha}| |\tilde{\Psi}_{\alpha}\rangle$ whose energy is $\langle \tilde{\Omega} | \tilde{\mathbf{H}} | \tilde{\Omega} \rangle = \sum_{\alpha} e_{\alpha} |f_{\alpha}|^2 - \sum_{\alpha\beta} |K_{\alpha\beta}| |f_{\alpha}| |f_{\beta}|$

$$\leq \sum_{\alpha} e_{\alpha} |f_{\alpha}|^2 - \sum_{\alpha\beta} |K_{\alpha\beta}| f_{\alpha} f_{\beta} = E,$$

where the last equality holds if and only if all of the f_α -coefficients are non-negative or all of the f_α -coefficients are negative. This state must, therefore, also be a ground state and satisfy the eigenvalue equation

$$-\sum_{\beta} |K_{\alpha\beta}| |f_{\beta}| + e_{\alpha} |f_{\alpha}| = E |f_{\alpha}|.$$

In general, the $|\tilde{\Psi}\rangle$ -states are not eigenstates of the rotated Hamiltonian, except in the case in which $M = NS$ which is the case of a one-state space. Therefore, $e_{\alpha} > E$, $\forall \alpha$. Then,

$$(3.3.1.9) \quad \begin{aligned} (e_{\alpha} - E) f_{\alpha} &= \sum_{\beta} |K_{\alpha\beta}| f_{\beta} \\ (e_{\alpha} - E) |f_{\alpha}| &= \sum_{\beta} |K_{\alpha\beta}| |f_{\beta}| \Rightarrow \left| \sum_{\beta} |K_{\alpha\beta}| f_{\beta} \right| = \sum_{\beta} |K_{\alpha\beta}| |f_{\beta}| \Rightarrow f_{\beta} \geq 0. \end{aligned}$$

Thus, the trial state and the ground states are the same, upto phase factors: $|\tilde{\Omega}\rangle = \pm |\Psi_0^M\rangle$. An even stronger condition on the f_{β} s can be imposed.

Lemma 2. $f_{\beta} > 0$, $\forall \beta$.

Proof. In effect, any Ising configuration with total magnetization M is connected to all other configurations in the same sector by successive application of the pairwise spin flip operator $\mathbf{K}_{\alpha\beta} \approx (\sum_{i,j} \mathbf{S}_i^+ \tilde{\mathbf{S}}_j^- + \mathbf{S}_i^- \tilde{\mathbf{S}}_j^+)_{\alpha\beta}$. Thus, if $\exists \alpha : f_{\alpha} = 0$, then f_{β} must vanish for all other β in the same M -sector. Since, by design, the ground state is such that in its expansions there is at least one non-zero coefficient, it follows that $f_{\beta} > 0$, $\forall \beta$, thus proving **theorem 2**.

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Corollary 1. For any fixed M , $|\Psi_0^M\rangle$ is non-degenerate.

Proof. In effect, since all f_{β} s are non-zero, there cannot be an orthogonal state to $|\Omega\rangle$ that has only positive definite coefficients in said M subspace.

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For **theorem 3** to hold, the ground state of the M -sector must have minimally possible total spin,

Lemma 3. For a bipartite antiferromagnetic Heisenberg model,

$$(\mathbf{S}_{\text{tot}})^2 |\Psi^M\rangle = M(M+1) |\Psi^M\rangle$$

Proof. Consider the infinite range Hamiltonian on a bipartite lattice with equal number of sites on the two sublattices,

$$\mathbf{H}^{\infty} = J \sum_{\substack{i \in A \\ j \in B}} \vec{\mathbf{S}}_i \cdot \vec{\mathbf{S}}_j = J \mathbf{S}_{\text{tot}}^A \cdot \mathbf{S}_{\text{tot}}^B$$

This problem is thus trivially solvable, since it is a two spin problem using $\mathbf{S}_{\text{tot}} = \mathbf{S}_{\text{tot}}^A + \mathbf{S}_{\text{tot}}^B$, since $\mathbf{S}_{\text{tot}}^A \cdot \mathbf{S}_{\text{tot}}^B = \frac{1}{2} (\mathbf{S}_{\text{tot}}^2 - (\mathbf{S}_{\text{tot}}^A)^2 - (\mathbf{S}_{\text{tot}}^B)^2)$. Therefore, the possible values of the total spin operators are

$$\mathbf{S}_{\text{tot}}^A \cdot \mathbf{S}_{\text{tot}}^B = 0, 1, \dots, \frac{NS}{2},$$

which for equal size sublattice means that

$$(3.3.1.10) \quad 0 \leq \mathbf{S}_{\text{tot}} \leq NS.$$

This infinite-range Hamiltonian's eigenvalues are then

$$(3.3.1.11) \quad E^\infty(\mathbf{S}_{\text{tot}}) = \frac{J}{2} \left[\mathbf{S}_{\text{tot}}(\mathbf{S}_{\text{tot}} + 1) - \mathbf{S}_{\text{tot}}^A(\mathbf{S}_{\text{tot}}^A + 1) - \mathbf{S}_{\text{tot}}^B(\mathbf{S}_{\text{tot}}^B) \right],$$

which are monotonically increasing with \mathbf{S}_{tot} , and since $\mathbf{S}_{\text{tot}} > M$, the ground state in a given M -sector has

$$\mathbf{S}_{\text{tot}} = M.$$

Now, note that both the Heisenberg Hamiltonian and the infinite range Hamiltonian satisfy the requirements for the first of Marshall's theorems, [theorem 2](#), and thus their ground states have Marshall signs. Therefore, their overlap is non-zero, since it involves a sum over positive number, and these must have the same total spin quantum numbers. Therefore, the lowest energy state of \mathbf{H} in the M -sector has spin $\mathbf{S}_{\text{tot}} = M$. Since all allowed values of total spin $\mathbf{S}'_{\text{tot}} > M$ have members in the sector with magnetization M , then the energy in terms of the total spin obeys

$$\forall \mathbf{S}'_{\text{tot}} > \mathbf{S}_{\text{tot}} \Rightarrow E(\mathbf{S}'_{\text{tot}}) > E(\mathbf{S}_{\text{tot}}),$$

ie. the energy is a monotonically increasing function of the total spin. Now, consider the $(M = 0)$ -sector. This inequality proves that the ground state must have minimal possible total spin \mathbf{S}_{tot} which is zero, according to [equation \(3.3.2.4\)](#). Thus concluding the proof of [theorem 3](#).

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3.3.2. Ferromagnetic Ground States. For the ferromagnetic Heisenberg chains, a similar result can be proved.

Corollary 2. *For the ferromagnetic model with non-positive charges,*

$$\mathbf{H} = -\frac{1}{2} \sum_{ij} |J_{ij}| \vec{\mathbf{S}}_i \cdot \vec{\mathbf{S}}_j,$$

the fully ferromagnetic state $|\Psi^{\text{FM}}\rangle$ is a member of the ground state multiplet, where

$$|\Psi^{\text{FM}}\rangle = \prod_{i=1}^N |S, S\rangle_i.$$

Proof. Consider the ferromagnetic Heisenberg model, whose Hamiltonian may be explicitly written as

$$(3.3.2.1) \quad \begin{aligned} \mathbf{H} &= \mathbf{H}^{zz} + \mathbf{H}^{xy}, \\ \mathbf{H}^{zz} &= -\frac{1}{2} \sum_{ij \in \Lambda \subset \mathbb{Z}^2} |J_{ij}| \mathbf{S}_i^z \mathbf{S}_j^z, \\ \mathbf{H}^{xy} &= -\frac{1}{4} \sum_{ij \in \Lambda \subset \mathbb{Z}^2} |J_{ij}| (\mathbf{S}_i^+ \mathbf{S}_j^- + \mathbf{S}_i^- \mathbf{S}_j^+). \end{aligned}$$

For a fixed M -subspace, all wave functions may be written in terms of Ising configurations which lie in said M -subspace, namely

$$(3.3.2.2) \quad |\Psi^M\rangle = \sum_{\alpha} f_{\alpha}^M |\Psi_{\alpha}^M\rangle, \text{ where } |\Psi_{\alpha}^M\rangle = \prod_{\sum_i m_i^{\alpha} = M} |S, m_i^{\alpha}\rangle.$$

Now,

- the longitudinal term, \mathbf{H}^{zz} , is diagonal in the Ising configurations, $\mathbf{H}^{zz} |\Psi_{\alpha}\rangle = e_{\alpha} |\Psi_{\alpha}\rangle$,
- and where the transverse \mathbf{H}^{xy} has non-negative $\langle \Psi_{\alpha} | \mathbf{H}^{xy} | \Psi_{\beta} \rangle = |K_{\alpha\beta}|$.

Consider the following ansatz, $|\Omega\rangle = \sum_{\alpha} |f_{\alpha}| |\Psi\rangle_{\alpha}$, whose energy is

$$(3.3.2.3) \quad \begin{aligned} \langle \Omega | H | \Omega \rangle &= \sum_{\alpha} e_{\alpha} |f_{\alpha}|^2 + \sum_{\alpha\beta} |K_{\alpha\beta}| |f_{\alpha}| |f_{\beta}| \\ &\leq \sum_{\alpha} e_{\alpha} |f_{\alpha}|^2 + \sum_{\alpha\beta} |K_{\alpha\beta}| f_{\alpha} f_{\beta} = E \Rightarrow \text{sgn}(f_{\alpha}) = \text{sgn}(f_{\beta}), \forall \alpha, \beta. \end{aligned}$$

Without loss of generality, the f_α -coefficients may be taken to be positive, with the extra, overall, sign to be associated with the Ising configurations. In effect, in general, the $|\Psi\rangle$ -states are not eigenstates of the Hamiltonian, except in the case in which $M = NS$ which is the case of a one-state space. Therefore, $e_\alpha > E$, $\forall \alpha$. Consider the infinite range ferromagnetic Hamiltonian on a bipartite lattice with equal number of sites on the two sublattices,

$$\mathbf{H}^\infty = -J \sum_{\substack{i \in A \\ j \in B}} \vec{\mathbf{S}}_i \cdot \vec{\mathbf{S}}_j = -J \mathbf{S}_{\text{tot}}^A \cdot \mathbf{S}_{\text{tot}}^B$$

Again, this problem is trivially solvable, since it is a two spin problem using $\mathbf{S}_{\text{tot}} = \mathbf{S}_{\text{tot}}^A + \mathbf{S}_{\text{tot}}^B$, since $\mathbf{S}_{\text{tot}}^A \cdot \mathbf{S}_{\text{tot}}^B = \frac{1}{2} (\mathbf{S}_{\text{tot}}^2 - (\mathbf{S}_{\text{tot}}^A)^2 - (\mathbf{S}_{\text{tot}}^B)^2)$. Therefore, the possible values of the total spin operators are

$$\mathbf{S}_{\text{tot}}^A \cdot \mathbf{S}_{\text{tot}}^B = 0, 1, \dots, \frac{NS}{2},$$

which for equal size sublattice means that

$$(3.3.2.4) \quad 0 \leq \mathbf{S}_{\text{tot}} \leq NS.$$

This infinite-range Hamiltonian's eigenvalues are then

$$(3.3.2.5) \quad E^\infty(\mathbf{S}_{\text{tot}}) = -\frac{J}{2} \left[\mathbf{S}_{\text{tot}}(\mathbf{S}_{\text{tot}} + 1) - \mathbf{S}_{\text{tot}}^A(\mathbf{S}_{\text{tot}}^A + 1) - \mathbf{S}_{\text{tot}}^B(\mathbf{S}_{\text{tot}}^B) \right],$$

which are monotonically increasing with \mathbf{S}_{tot} , and since $\mathbf{S}_{\text{tot}} > M$, the ground state in a given M -sector has

$$\mathbf{S}_{\text{tot}} = M.$$

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As a member of a degenerate multiplet, this state spontaneously breaks the rotational symmetry of the Hamiltonian. This symmetry breaking is special in that the operator $\mathbf{S}_{\text{tot}}^z$ commutes with the Hamiltonian and, therefore, M is a good quantum number (ie. it labels the eigenstates of \mathbf{H}). In this respect, the antiferromagnet differs from the ferromagnet. The staggered magnetization does not commute in general with the Hamiltonian. Spontaneous symmetry breaking, however, is still possible for the antiferromagnet in the strict thermodynamic limit.

3.3.3. Half-Odd Integer Spin Chains. As a Lie group, $SU(2)$ is the simply-convex three-sphere \mathcal{S}^3 . Thus, particle with half-odd integer spin, eg. electrons, have a peculiar quantum property: their wave functions acquires a minus sign under a 2π -rotation about any spin axis. For localized systems, any particular direction in spin space, this effect is not noticeable. However, for low-dimensional systems, this zero-point fluctuations are large. Thus, interference effects due to these negative signs may become important. In fact, the antiferromagnetic one-dimensional Heisenberg model has different spectra for integer and half-odd integer spins. Thus the following theorem, due to Lieb, Schultz and Mattis, arises:

Theorem 4 Lieb-Schultz-Mattis Theorem on Half-odd Integer Spin Chains. Consider the antiferromagnetic Heisenberg spin chain, with $J > 0$ on a one dimensional chain with closed topology

$$(3.3.3.1) \quad \mathbf{H} = J \sum_{j=1}^N \mathbf{S}_j \cdot \mathbf{S}_{j+1} + J \mathbf{S}_N \cdot \mathbf{S}_1,$$

where N is even. For half-odd integer spin, there exists an excited state whose energy vanishes in the thermodynamic limit, $N \rightarrow \infty$.

Let $|\Psi_0\rangle$ be the system's ground state and let the twist operator be

$$\mathcal{O}^k = \exp \left(ik \sum_{j=1}^N j \mathbf{S}_j^z \right),$$

which incrementally rotates each spin in the xy -plane, such that between the first and last spin, the spin coordinates are rotated by 2π about the z -axis. Thus, the twisted state is constructed as

$$|\Psi_k\rangle = \mathcal{O}^k |\Psi_0\rangle$$

Two statements swiftly arise,

Lemma 4.

$$(3.3.3.2) \quad \langle \Psi_k | \Psi_0 \rangle = 0 \text{ and } \left[\langle \Psi_k | \mathbf{H} | \Psi_k \rangle - E_0 \right] \xrightarrow{N \rightarrow \infty} 0$$

Proof. In effect, consider a translation operator, defined by cyclically shifting the spins by one site modulo N , this is

$$(3.3.3.3) \quad \begin{aligned} \mathcal{T} : \mathfrak{A}_i &\rightarrow \mathfrak{A}_j \\ \mathcal{T}^{-1} \vec{S}_i \mathcal{T} &= \vec{S}_{i+1} \end{aligned} \quad \text{where } \mathfrak{A}_i, \mathfrak{A}_j \simeq \mathfrak{su}(2) \ \forall i, j.$$

By construction, the Heisenberg Hamiltonian with closed boundary conditions commutes with the translation operator, $[\mathbf{H}, \mathcal{T}] = 0$. Then, if $|\Psi_0\rangle$ is an eigenstate of the Hamiltonian, so is $\mathcal{T}|\Psi_0\rangle$. Given $|\Psi_0\rangle$'s non-degeneracy, and guaranteed by [corollary 1](#), and letting k_0 be the ground state lattice momentum, the following holds

$$(3.3.3.4) \quad \begin{aligned} \mathcal{T}|\Psi_0\rangle &= e^{ik_0} |\Psi_0\rangle \Rightarrow \langle \Psi_0 | \Psi_k \rangle = \langle \Psi_0 | \mathcal{O}^k | \Psi_0 \rangle = \langle \Psi_0 | \mathcal{T}^{-1} \mathcal{O}^k \mathcal{T} | \Psi_0 \rangle \\ &= \langle \Psi_0 | \exp \left[-ikN \sum_{j=1}^N j \mathcal{T}^{-1} \mathbf{S}_j^z \mathcal{T} \right] | \Psi_0 \rangle \\ &= \langle \Psi_0 | \exp \left[-ikN \left(\sum_{j=1}^{N-1} j \mathbf{S}_{j+1}^z + N \mathbf{S}_1^z \right) \right] | \Psi_0 \rangle \\ &= \langle \Psi_0 | e^{ikN \mathbf{S}_1^z} \exp \left[-ikN \left(\sum_{j=1}^{N-1} (j+1) \mathbf{S}_{j+1}^z + \mathbf{S}_1^z - \sum_{j=1}^N \mathbf{S}_{j+1}^z \right) \right] | \Psi_0 \rangle \\ &= \langle \Psi_0 | e^{ikN \mathbf{S}_1^z} \exp \left(-ikN \left(\sum_{j=1}^N (j+1) \mathbf{S}_{j+1}^z - \sum_{j=1}^N \mathbf{S}_{j+1}^z \right) \right) | \Psi_0 \rangle \\ &= \langle \Psi_0 | e^{i2\pi \mathbf{S}_1^z} \mathcal{O}^k \exp \left(-ikN \mathbf{S}_{\text{tot}}^z \right) | \Psi_0 \rangle \\ &= \langle \Psi_0 | \mathcal{O}^k e^{ik \mathbf{S}_1^z} \exp \left(-ikN \sum_{j=1}^N j \mathbf{S}_j^z \right) | \Psi_0 \rangle \end{aligned}$$

This last equality is, in fact, zero by [theorem 3](#), the ground state has zero-total magnetization. Given that $|\Psi_0\rangle$ is a singlet, note that

$$(3.3.3.5) \quad \mathbf{S}_1^z = \begin{pmatrix} \frac{1}{2} & 0 \\ 0 & -\frac{1}{2} \end{pmatrix} \Rightarrow e^{ik \mathbf{S}_1^z} = \begin{cases} 1 & S = 0, 1, 2, \dots \\ -1 & S = \frac{1}{2}, \frac{3}{2}, \frac{5}{2}, \end{cases} \Rightarrow \mathcal{T}^{-1} \mathcal{O}^k \mathcal{T} = \begin{cases} \mathcal{O}^k & S = 0, 1, 2, \dots \\ -\mathcal{O}^k & S = \frac{1}{2}, \frac{3}{2}, \frac{5}{2}, \end{cases}$$

which then implies that $\langle \Psi_0 | \Psi_k \rangle = -\langle \Psi_0 | \Psi_k \rangle = 0$, thus proving the first proposition. The $|\Psi_k\rangle$ -state energies are also readily calculated. In effect, given that the twist operator is a unitary rotation on the xy -plane, the following results arise

$$\begin{aligned}
\mathcal{O}^{k-1} \mathbf{S}_n^x \mathcal{O}^k &= \cos kn \mathbf{S}_n^x + \sin kn \mathbf{S}_n^y \\
\mathcal{O}^{k-1} \mathbf{S}_n^y \mathcal{O}^k &= \cos kn \mathbf{S}_n^y - \sin kn \mathbf{S}_n^x \\
\mathcal{O}^{k-1} \mathbf{S}_n^z \mathcal{O}^k &= \mathbf{S}_n^z \Rightarrow \langle \Psi_k | \mathbf{H} | \Psi_k \rangle = \langle \Psi_0 | \mathcal{O}^{k-1} \mathbf{H} \mathcal{O}^k | \Psi_0 \rangle = \langle \Psi_0 | \mathcal{O}^{k-1} \mathbf{H} \mathcal{O}^k | \Psi_0 \rangle \\
(3.3.3.6) \quad &= \langle \Psi_0 | \mathcal{O}^{k-1} \left(\frac{J}{2} \sum_{j \in \Lambda} (\mathbf{S}_j^x \mathbf{S}_{j+1}^x + \mathbf{S}_j^y \mathbf{S}_{j+1}^y + \mathbf{S}_j^z \mathbf{S}_{j+1}^z) \right) \mathcal{O}^k | \Psi_0 \rangle \\
&= E_0 + \langle \Psi_0 | [\cos k - 1] \sum_{j=1}^N (\mathbf{S}_j^x \mathbf{S}_{j+1}^x + \mathbf{S}_j^y \mathbf{S}_{j+1}^y) + \\
&\quad \sin k \sum_{j=1}^N (\mathbf{S}_j^x \mathbf{S}_{j+1}^y - \mathbf{S}_j^y \mathbf{S}_{j+1}^x) | \Psi_0 \rangle,
\end{aligned}$$

examining each term in the previous equation yields the following approximations

$$\begin{aligned}
(3.3.3.7) \quad &\Rightarrow [\cos k - 1] \langle \Psi_0 | \sum_{j=1}^N (\mathbf{S}_j^x \mathbf{S}_{j+1}^x + \mathbf{S}_j^y \mathbf{S}_{j+1}^y) | \Psi_0 \rangle + \sin k \sum_{j=1}^N (\mathbf{S}_j^x \mathbf{S}_{j+1}^y - \mathbf{S}_j^y \mathbf{S}_{j+1}^x) \\
&= \left[-\frac{1}{2} \left(\frac{2\pi}{N} \right)^2 - \mathcal{O}(N^{-4}) \right] \langle \Psi_0 | \sum_{j=1}^N (\mathbf{S}_j^x \mathbf{S}_{j+1}^x + \mathbf{S}_j^y \mathbf{S}_{j+1}^y) | \Psi_0 \rangle = -i \sin k \langle \Psi_0 | [\sum_j j \mathbf{S}_j^z, \mathbf{H}] | \Psi_0 \rangle \\
&\leq \left(\frac{2\pi}{N} \right)^2 \frac{N}{2} + \mathcal{O}(N^{-3}) = 0
\end{aligned}$$

Thus, for $k = \frac{2\pi}{N}$,

$$\langle \Psi_k | \mathbf{H} | \Psi_k \rangle \leq E_0 + \frac{2\pi^2}{N},$$

and then, there is no energy gap in the thermodynamic limit. Thus concluding [lemma 4](#)'s proof.

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However, note that in the thermodynamic limit, the admixture of the ground state and this low-lying excited state may break a symmetry of the Hamiltonian. This is the case for the ground state and the first excited state of the Majumdar-Gosh model, which are superpositions of dimer configurations, break lattice symmetry. An alternative scenario occurs in the nearest-neighbour Heisenberg model of spin half. [des Cloizeaux and Pearson](#) have found that magnon excitations are gapless, said excitations not being spin waves since they are not small fluctuations about a ground state with broken spin symmetry. In contrast to spin-half chains, integer spin chains have a gap in the excitation spectrum, the Haldane gap.

The previous results can be readily extended to two and three-dimensional lattices. In effect, for the former, consider a square lattice of N sites in the x -direction and $M \sim \mathcal{O}(N^\nu)$ sites in the y -direction, where $\nu \in \mathbb{R}_{[0,1]}$, with a cyclic Hamiltonian, ie.

$$\mathbf{S}_{n,M+1} = \mathbf{S}_{n,1} \text{ and } \mathbf{S}_{N+1,m} = \mathbf{S}_{1,m},$$

ie. the lattice is wrapped on a two-dimensional torus \mathbb{T}^2 . Let the two-dimensional twisted operator \mathcal{O}^k be

$$\mathcal{O}^k = \exp \left(ik \sum_{n,m} n \mathbf{S}_{n,m}^z \right),$$

which twists the x -direction sites by the same amount and the y -direction sites by another amount. The $|\Psi_k\rangle$ -states construction and orthonormality follow directly from the one-dimensional case. In this case, there is no energy gap in the continuum limit,

$$\langle \Psi_k | \mathbf{H} | \Psi_k \rangle \leq E_0 + \frac{2\pi^2}{N^{1-\nu}}.$$

Since the excitation energy of exact low-lying states should not depend on the shape of the entire lattice, there should be no energy gap for an $(N \times N)$ -lattice either. The $|\Psi_k\rangle$ state is, unfortunately, not sufficiently like an exact low-lying excited state to give this result. A similar extension to three-dimensional lattices swiftly follows.

3.4. Numerical solution to Fermionic models. Consider a Hamiltonian describing a fermionic system, given by

$$(3.4.0.1) \quad \mathbf{H} = J \sum_{j=1}^{L-1} (f_j^\dagger f_{j+1} + f_{j+1}^\dagger f_j) + \sum_{j=1} \lambda_j f_j^\dagger f_j, \quad \text{with the usual commutation rules} \quad \begin{cases} \{f_j, f_k\} = \{f_j^\dagger, f_k^\dagger\} = 0 \\ \{f_j, f_k^\dagger\} = \delta_{jk} \end{cases}$$

where L indicates the number of lattice sites, J is the hopping strength, which could be either positive or negative, and where λ_j is the on-site potential strength¹⁸. Said Hamiltonian has open boundaries conditions since there is no hopping term across the boundary. Note that we can rewrite (3.4.0.1) as

$$(3.4.0.2) \quad \mathbf{H} = \sum_{i,j=1}^L \mathcal{M}_{ij} f_i^\dagger f_j \quad \text{with} \quad \mathcal{M}_{ij} = \begin{cases} \lambda_i & \text{if } i = j \\ J & \text{if } j = i + 1 \text{ or } i = j + 1 \\ 0 & \text{otherwise} \end{cases}, \quad \mathcal{M} \in \text{GL}(L, \mathbb{R}),$$

which is a positive-defined tri-diagonal matrix. Let $\mathbf{f} = (f_1 \ f_2 \ \cdots \ f_L)^\mathbf{T}$ be a vector of the L fermionic operators. Then, (3.4.0.1) can be rewritten as

$$(3.4.0.3) \quad \mathbf{H} = \mathbf{f}^\dagger \mathcal{M} \mathbf{f}.$$

Since \mathcal{M} is symmetric, then it can be diagonalized $\mathcal{M} = A \mathcal{D} A^\mathbf{T}$, where $A \in \mathbb{R}^{L \times L}$ is a real orthogonal matrix and with $\mathcal{D}_{ij} \in \mathbb{R}^{L \times L} \mid \mathcal{D}_{ij} = \epsilon_i \delta_{ij}$. In this context, the A -matrix acts on the fermionic operator as a Bogoliubov transformation, allowing for (3.4.0.1) to be rewritten as

$$(3.4.0.4) \quad \mathbf{H} = \mathbf{f}^\dagger A \mathcal{D} A^\mathbf{T} \mathbf{f} = \mathbf{d}^\dagger \mathcal{D} \mathbf{d}$$

where $\mathbf{d} = A^\mathbf{T} \mathbf{f}$. Since the A -matrix is orthogonal, the new d_k -operators are fermionic operators as well, satisfying (3.4.0.1)'s anti-commutation rules. Then, the new fermionic operators are

$$\begin{aligned} d_k &= \sum_{j=1}^L A_{jk} f_j \\ f_j &= \sum_{k=1}^L A_{jk} d_k \end{aligned} \quad \text{since } A^\mathbf{T} A = \sum_{j,k=1}^L A_{jk} A_{kj} = \mathbb{1}_L.$$

Then, we can expand (3.4.0.4) in terms of the lattices, as follows

$$(3.4.0.5) \quad \mathbf{H} = \sum_{k=1}^L \epsilon_k d_k^\dagger d_k,$$

which is a sum of number operators with potentials. The eigenstates can then be constructed from the theory's vacuum state, by applying the d_k^\dagger -fermionic operators. In the Heisenberg-picture, the d_k -operators can be evolved via the Heisenberg equation of motion

$$(3.4.0.6) \quad \frac{d}{dt} d_k = i[\mathbf{H}, d_k],$$

¹⁸The λ_j -term frequently appears in many condensed matter models, with different numerical values and interpretations, eg.

- In the XX model, $\lambda_j = \lambda \forall j$.
- While for the Anderson model $\lambda_j \in \mathcal{U}_{\mathbb{R}[-W, W]}$, a uniform random variable, with W being the disorder strength.
- In the Aubry-André model, $\lambda_j = \lambda \cos(2\pi\sigma j)$, with $\sigma \in \mathbb{I}$ and λ quantifying the disorder strength.

and using that $d_k^2 = 0$, it turns out that (3.4.0.6)'s solution is simply $d_k(t) = e^{-i\epsilon_k t} d_k$. The system's correlation can be easily found by analyzing the following matrix. Let $\mathcal{N}_{jk} = \langle d_j^\dagger d_k \rangle$, where the expectation value is taken via calculating the operator's trace along the Fock space, which takes the following values

$$(3.4.0.7) \quad \mathcal{N}_{jk} = \langle d_j^\dagger d_k \rangle = \begin{cases} 0 \text{ or } 1 & \text{if } j = k \\ 0 & \text{if } j \neq k \end{cases},$$

ie. different lattice-sites are not correlated and there can only be a single fermion at most per lattice site, in accordance with Pauli's principle. A ground state, for example, would choose to turn on all fermions in the eigenmode d -space such that $\epsilon_k < 0$. If instead, the expectation value is taken with thermal states, the Fermi-Dirac distribution is returned,

$$(3.4.0.8) \quad \mathcal{N}_{jk} = \langle d_j^\dagger d_k \rangle_{\text{th}} = \frac{1}{1 + e^{\beta\epsilon_k + \mu}} \delta_{jk}.$$

Another interesting quality is a system with an initial configuration where the system's initial state, in real space, is known. In this setting, \mathcal{N}_{jk} is known for all lattices. Consider for example the Anderson model, where the system's initial state is given by a single tensor product of n -fermionic states in real space, with $n < L$. Then, for all lattice sites, we have that \mathcal{N}_{jj} is either zero or one. The \mathcal{N}_{jk} -matrix entries can then be evaluated as

$$\langle d_j^\dagger d_k \rangle = \sum_{i,j=1}^{n < L} A_{ik} A_{jl} \langle f_i^\dagger f_j \rangle = \sum_{j=1} A_{jk} A_{jl} \langle f_j^\dagger f_j \rangle,$$

which can then be numerically computed to obtain the LHS expectation value. In general, this \mathcal{N}_{jk} -matrix will not be diagonal, which is reasonable since the system's real configuration is not an eigenstate. In principle and in practice, by inverting (3.4.0.9), we can evolve any number operator or two-body correlation operator, ie.

$$(3.4.0.9) \quad \langle f_j^\dagger f_k \rangle = \sum_{k,l=1}^n A_{jk} A_{jl} \langle d_k^\dagger d_l \rangle.$$

This quantities' time evolution can then be found out to be

$$(3.4.0.10) \quad \langle f_j^\dagger(t) f_k(t) \rangle = \sum_{k,l=1}^L e^{i(\epsilon_k - \epsilon_l)t} A_{jk} A_{jl} \langle d_k^\dagger d_l \rangle,$$

which can then be numerically solved.

3.4.0.1. Ferromagnet and antiferromagnet. Unless geometric frustration is present, in classical systems, there is no significant difference between ferromagnets and antiferromagnets. Eg, with NN-interaction, geometric frustration occurs for lattices that are not bipartite. In a bipartite lattice, the sites can be divided into two sub-lattices such that nearest neighbours always belong to different sub-lattices. A nearest-neighbour antiferromagnetic interaction in a classical model on a bipartite lattice can typically be changed into a ferromagnetic one, by redefining the spin via a flip on all sites on one of the sub-lattices but not on the other (eg. $\uparrow \rightarrow \downarrow$ in the Ising model). The physics of such classical antiferromagnets is therefore essentially equivalent to that of the ferromagnets.

Antiferromagnetic quantum systems on non-bipartite lattices also exhibit interesting behaviour. But the interesting thing is that on bipartite lattices, there are a number of important differences between quantum ferromagnets and antiferromagnets.

3.4.0.2. Quantum Ferromagnet. Consider the Heisenberg interaction $-J\vec{S}_1 \cdot \vec{S}_2$ across a single bond. For spin- $\frac{1}{2}$ particles, this is a simple 4×4 -matrix acting on the computational basis $\mathcal{B} = \{|\uparrow\uparrow\rangle, |\uparrow\downarrow\rangle, |\downarrow\uparrow\rangle, |\downarrow\downarrow\rangle\}$, yielding

$$\vec{S}_1 \cdot \vec{S}_2 = \frac{1}{4} \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & -1 & 4 & 0 \\ 0 & 4 & -1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}$$

Diagonalizing this matrix will yield its eigenvalues and its eigenvectors. Given the Clebsh-Gordan decomposition rule $\frac{1}{2} \otimes \frac{1}{2} = 0 \oplus 1$, these eigenvectors can be grouped into the $s = 1$ -triplet representation and the $s = 0$ -singlet representation. In effect,

$$(3.4.0.11) \quad \begin{aligned} \text{triplet} : & |\uparrow\uparrow\rangle, \frac{1}{\sqrt{2}}(|\downarrow\uparrow\rangle + |\uparrow\downarrow\rangle), |\downarrow\downarrow\rangle, \text{ with } \lambda_i = -\frac{J}{4} \\ \text{singlet} : & \frac{1}{\sqrt{2}}(|\downarrow\uparrow\rangle - |\uparrow\downarrow\rangle), \text{ with } \mu_i = +\frac{J}{4} \end{aligned}$$

It is simple to check that $\mathbf{K} = s(s+1)$ in both cases, and that while the singlet is annihilated by all three generators \vec{S} , acting with S^+ and S^- takes members of the triplet to each other. One important difference between the ferromagnetic ($J > 0$) and the antiferromagnetic ($J < 0$) models now arises,

- if $J > 0$, there are multiple ground states for the ferromagnet, each member of the triplet having minimum energy $-J$,
- but if $J < 0$, the antiferromagnetic ground state, the singlet, is unique. Moreover, the antiferromagnetic ground state is invariant under the $SU(2)$ -symmetry, whereas the ferromagnetic ground states are not.

Additional insight comes from solving the Heisenberg model on a four-site chain, with its Hilbert space being 16-dimensional, by exploiting the model's inner symmetries. The magnetization operator M^z commutes with the Hamiltonian, which makes them simultaneously diagonalizable. For periodic boundary conditions, the translational invariance plays a powerful role. The translation operator is defined by shifting the spins by one site modulo N , this is

$$(3.4.0.12) \quad \begin{aligned} \mathcal{T} : \mathfrak{A}_i &\rightarrow \mathfrak{A}_j \\ \mathcal{T}^{-1}\vec{S}_i\mathcal{T} &= \vec{S}_{i+1} \end{aligned} \quad \text{where } \mathfrak{A}_i, \mathfrak{A}_j \simeq \mathfrak{su}(2) \ \forall i, j.$$

The Hamiltonian commutes with the translation operator when the boundary conditions are periodic, and commutes with the magnetization as well. Thus the Hamiltonian can be broken into blocks acting on states with a fixed eigenvalue of the translation operator and of the magnetization.

Since $\mathcal{T}^N = \mathbb{1}$ for an N -site chain, the eigenvalues of \mathcal{T} are $\lambda_i = e^{\frac{2\pi i n}{N}}, \forall n \in \mathbb{Z}$, from which the corresponding momentum can be defined as $k = \frac{2\pi n}{N}$. Then, the following holds

$$\sum_{n \in \mathbb{Z}_{[0, N-1]}} e^{-\frac{2\pi i n}{N}} \mathcal{T}^n |s\rangle = |s\rangle + e^{-\frac{2\pi i}{N}} |s\rangle + \cdots + e^{-\frac{2\pi i (N-1)}{N}} \mathcal{T}^{N-1} |s\rangle.$$

For example, for a four-site chain, the translational invariance diagonalizes the entire Hamiltonian save for $m = 0$ and $k = 0$ or $k = \pi$, since the corresponding eigenstates are

$$(3.4.0.13) \quad \begin{aligned} |A\rangle &= \frac{1}{\sqrt{2}} \left(|\uparrow\downarrow\uparrow\downarrow\rangle + e^{-i\pi k} |\downarrow\uparrow\downarrow\uparrow\rangle \right), & \text{with the hamiltonian} \\ |B\rangle &= \frac{1}{2} \left(|\uparrow\uparrow\downarrow\downarrow\rangle + |\downarrow\downarrow\uparrow\uparrow\rangle + \cdots \right) & \text{on these} \\ & & k = 0 \text{ and } k = \pi \\ & & \text{states being} \end{aligned} \quad -J \begin{pmatrix} -1 & \sqrt{2} \cos\left(\frac{\pi k}{2}\right) \\ \sqrt{2} \cos\left(\frac{\pi k}{2}\right) & 0 \end{pmatrix}.$$

These sector's eigenvalues are therefore $-J$ and $2J$ for $k = 0$ and $J = 0$ for $k = \pi$. Again organising the eigenstates into $\mathfrak{su}(2)$ -multiplets yield the energy levels divided by $-J$, as follows

$$(3.4.0.14) \quad \begin{aligned} &\text{quintuplet : } 1 \\ &\text{triplets : } \cos(\pi k) (k \neq 0) \\ &\text{singlets : } -2, 0 \text{ for } k = 0, \pi \text{ respectively.} \end{aligned}$$

This $\mathfrak{su}(2)$ -invariance allows for the construction of the entire multiplet once one of the states is known.

As with the two-site chain, the ferromagnetic ground state is a multiplet, whereas the antiferromagnetic is a singlet. For a general N -site ferromagnet, the completely ferromagnetic states (all spins up or all spins down) are exact ground states of the Hamiltonian. This suggest using an order parameter for ferromagnetism. A ferromagnetic order parameter is simply the quantum analog of the magnetization, the expectation value of the z -component of the total spin is

$$\langle \text{g.s.} | \mathbf{M}^z | \text{g.s.} \rangle = \sum_{i=1}^N \text{g.s.} \langle \text{g.s.} | \mathbf{S}_i^z | \text{g.s.} \rangle,$$

which is not a particularly great order parameter, since it vanishes if there is symmetry under spin flip, as with the classical case. A better example is $\langle \text{g.s.} | (\mathbf{M}^z)^2 | \text{g.s.} \rangle$.

The completely ferromagnetically ordered states, ie. all spins up or all spins down, are both \mathbf{M}^z -eigenvalues with maximum magnitude of eigenvalues, namely N^2 , and are ground states of the Hamiltonian as well. The antiferromagnetic situation is different. The staggered magnetization, or the **Néel order parameter**, is commonly used to understand antiferromagnetic order. It can only be defined on bipartite lattices. For example, consider the odd-sites on a sublattice, and the even-sites on the other, then the Néel operator is

$$\mathcal{N} = \sum_{i=1}^N (-1)^i \mathbf{S}_i^z.$$

A Néel state has N^2 -eigenvalues under \mathcal{N}^2 , ie. the spins are spin up on one sublattice and spin down on the other, like the state $|A\rangle$ which has four sites. Note, however, that the Néel operator does not commute with Hamiltonian. Then, a Néel state is not an eigenstate of the Heisenberg Hamiltonian. This is a huge difference between ferromagnets and antiferromagnets, even on bipartite lattices, not just on geometrically frustrated ones.

For example, for a four-site chain and $J < 0$, the ground state is a spin singlet and translationally invariant. The properties are thus quite typical of antiferromagnetic ground states. The diagonal terms are J and 0 for the translationally-invariant Néel state $|A\rangle$ and the other state, $|B\rangle$, respectively. This does indeed favour the Néel state. However, the off-diagonal terms in the Hamiltonian, mean

3.4.0.3. Hamiltonian in terms of Projectors. It is often useful to write a Hamiltonian as a sum of projection operators, especially if a form can be found where all the coefficients are positive. This can be done for both the Heisenberg ferromagnet and antiferromagnets, and greatly illuminates the distinction between the two.

A projector is an idempotent operator, thus its eigenvalues can only be either zero or one. Any Hamiltonian which can be written in such a form, will non-negative eigenvalues. Let $\mathbf{P}^{(s)}$ be a projector onto spin- s such that

- it gives 1 when acting on a collection of spins whose total spin is s ,
- and 0 otherwise.

A projector onto m spins can be built by taking product of the operators $(\sum_{i=1}^m \mathbf{S}_i) \cdot (\sum_{i=1}^m \mathbf{S}_i) - j(j+1)$. For example, there are two projectors, $\mathbf{P}^{(0)}$ and $\mathbf{P}^{(1)}$, acting on two spin- $\frac{1}{2}$ particles, since the following $SU(2)$ representations holds $\frac{1}{2} \otimes \frac{1}{2} = 0 \oplus 1$. Said projectors are given by

$$(3.4.0.15) \quad \begin{aligned} \mathbf{P}^{(1)}(\mathbf{S}_1 + \mathbf{S}_2) &= \frac{1}{2} \left((\mathbf{S}_1 + \mathbf{S}_2) \cdot (\mathbf{S}_1 + \mathbf{S}_2) \right) \\ \mathbf{P}^{(0)}(\mathbf{S}_1 + \mathbf{S}_2) &= -\frac{1}{2} \left((\mathbf{S}_1 + \mathbf{S}_2) \cdot (\mathbf{S}_1 + \mathbf{S}_2) - 2 \right). \end{aligned}$$

Note that $\mathbf{P}^{(1)} + \mathbf{P}^{(0)} = 1$, since all states belong to some $\mathfrak{su}(2)$ -multiplet. Since $\mathbf{S} \cdot \mathbf{S} = \frac{3}{4}$ for a spin- $\frac{1}{2}$ particle,

$$\mathbf{P}^{(1)}(\mathbf{S}_i + \mathbf{S}_{i+1}) = \frac{3}{4} + \mathbf{S}_i \cdot \mathbf{S}_{i+1}, \text{ and } \mathbf{P}^{(0)}(\mathbf{S}_i + \mathbf{S}_{i+1}) = \frac{1}{4} - \mathbf{S}_i \cdot \mathbf{S}_{i+1}.$$

Therefore, the spin- $\frac{1}{2}$ Heisenberg Hamiltonian can therefore be rewritten as a sum of projectors with positive weights in both the ferromagnetic and antiferromagnetic cases.

- In the ferromagnetic case, the ferromagnetic hamiltonian is a sum of projectors with positive coefficients,

$$\mathbf{H} = J \sum_{i=1}^N \mathbf{P}^{(0)}(\mathbf{S}_i + \mathbf{S}_{i+1}).$$

Each of these projectors annihilates any ferromagnetic ground state. This is obvious for the completely aligned states, since $\mathbf{P}^{(0)}$ annihilates any pair of spins that are the same. It is not as obviously true for other parts of the ferromagnetic ground-state multiplet, but follows from the facts that \mathbf{S}^\pm commutes with the Hamiltonian and act on the completely aligned state to give the rest of the multiplet. Thus, the ferromagnetic Heisenberg model is frustration free.

- The antiferromagnetic spin- $\frac{1}{2}$ Heisenberg Hamiltonian can also be written as a sum of projectors with positive coefficients. In effect,

$$\mathbf{H} = -J \sum_{i=1}^N \mathbf{P}^{(1)}(\mathbf{S}_i + \mathbf{S}_{i+1}).$$

Here, however, there is no state annihilated by all the operators. This is simple to see by checking explicitly just for three spins in a row. Therefore, the Heisenberg Hamiltonian is not frustration free.

3.5. Disorder in Low Dimensions. Spontaneously broken symmetry is a general phenomenon in statistical mechanics and particle physics. It can occur at finite temperatures and when the order parameter does not commute with the Hamiltonian. For example, consider the spin density wave operator in the z -direction

$$(3.5.0.1) \quad \mathbf{S}_{\mathbf{q}}^z = \sum_{i \in \Lambda} e^{i\mathbf{q} \cdot \mathbf{x}_i} \mathbf{S}_i^z,$$

where \mathbf{q} is the ordering wave vector. A symmetry breaking term is added to the rotationally invariant Hamiltonian \mathbf{H}^0 using an ordering field h ,

$$(3.5.0.2) \quad \mathbf{H} = \mathbf{H}^0 - h\mathbf{S}_{\mathbf{q}}^z.$$

Then, the magnetization per site is given by

$$m_{\mathbf{q}}(h) = \frac{1}{N\mathcal{Z}} \text{Tr} \left[e^{-\beta\mathbf{H}} \mathbf{S}_{\mathbf{q}}^z \right], \text{ where } \mathcal{Z} = \text{Tr} \left[e^{-\beta\mathbf{H}} \right].$$

In particular, this discussion is restricted to symmetric lattices under reflection about the origin, inducing $\mathbf{S}_{\mathbf{q}}^z$'s hermiticity and $m_{\mathbf{q}}$ -realness. At finite fields, $m_{\mathbf{q}}(h) > 0$, since the magnetization is induced by the ordering field.

Definition 1.

The system has spontaneously broken symmetry if it sustains a finite magnetization in the thermodynamic limit even if the ordering field is taken to zero from above, this is

$$(3.5.0.3) \quad \lim_{h \rightarrow 0^+} \lim_{N \rightarrow \infty} m_{\mathbf{q}}(h, N, T) \neq 0.$$

The order of limits is crucial in the previous definition.

Spontaneously broken symmetries can be found in the absence of an ordering field by examining the two point correlation function,

$$(3.5.0.4) \quad S^{\alpha\alpha}(\mathbf{q}) = \lim_{h \rightarrow 0^+} \frac{1}{N\mathcal{Z}} \text{Tr} \left[e^{-\beta\mathbf{H}} \mathbf{S}_{\mathbf{q}}^{\alpha} \mathbf{S}_{-\mathbf{q}}^{\alpha} \right], \text{ with } \alpha = \{x, y, z\}.$$

In the absence of the ordering field, the correlations are isotropic in spin space, i.e., independent on the α -direction. It can be shown that spontaneously broken symmetry at \mathbf{q} -momentum implies true long-range order in the two point correlation function,

Definition 2.

True-Long range order

$$(3.5.0.5) \quad \lim_{N \rightarrow \infty} \frac{S^{\alpha\alpha}(\mathbf{q})}{N} > 0 \Leftrightarrow \lim_{|\mathbf{x}_i - \mathbf{x}_j| \rightarrow \infty} \lim_{N \rightarrow \infty} \langle \mathbf{S}_i \cdot \mathbf{S}_j \rangle \neq 0.$$

This definition thus distinguishes the true-long range order from a quasi-long-range order, in which there is a power law decay of correlations at large distances.

Spontaneously broken symmetry and true-long range order

Spontaneously broken symmetry implies true-long range order.

Proof. Given two operators acting on the Hilbert space of a quantum system, a scalar product between two operators \mathbf{A} and \mathbf{B} can be defined as

$$(\mathbf{A}, \mathbf{B}) = \frac{1}{\mathcal{Z}} \sum_{n,m, E_n \neq E_m} \langle n | \mathbf{A}^\dagger | m \rangle \langle m | \mathbf{B} | n \rangle \left(\frac{e^{-\beta E_m} - e^{-\beta E_n}}{E_n - E_m} \right),$$

where $\{E_n, |n\rangle\}$ are the energies and eigenstates of the Hamiltonian. This is a well-defined real-valued non-negative proper scalar product in the product Hilbert space, spanned by $\{|n\rangle\} \times \{|n\rangle\}$. Now, the Cauchy-Schwartz inequality establishes that

$$|(\mathbf{A}, \mathbf{B})|^2 \leq (\mathbf{A}, \mathbf{A})(\mathbf{B}, \mathbf{B}), \text{ or equivalently } (\mathbf{A}, \mathbf{B}) \leq |\mathbf{A}||\mathbf{B}|,$$

where $|\mathbf{A}|$ is to be understood as $\sqrt{(\mathbf{A}, \mathbf{A})}$. But then,

$$(3.5.0.6) \quad \begin{aligned} |(\vec{\mathbf{S}}_{\mathbf{q}}, \vec{\mathbf{S}}_{-\mathbf{q}})|^2 &\leq \frac{1}{2T} \langle \vec{\mathbf{S}}_{\mathbf{q}} \cdot \vec{\mathbf{S}}_{-\mathbf{q}} + \vec{\mathbf{S}}_{-\mathbf{q}} \cdot \vec{\mathbf{S}}_{\mathbf{q}} \rangle (\vec{\mathbf{S}}_{-\mathbf{q}}, \vec{\mathbf{S}}_{-\mathbf{q}}) \\ &\Rightarrow |(\vec{\mathbf{S}}_{\mathbf{q}}, \vec{\mathbf{S}}_{-\mathbf{q}})|^2 \leq \frac{1}{T} \langle \vec{\mathbf{S}}_{\mathbf{q}} \cdot \vec{\mathbf{S}}_{-\mathbf{q}} \rangle (\vec{\mathbf{S}}_{-\mathbf{q}}, \vec{\mathbf{S}}_{-\mathbf{q}}) \Rightarrow \langle \vec{\mathbf{S}}_{\mathbf{q}} \cdot \vec{\mathbf{S}}_{-\mathbf{q}} \rangle \geq T \frac{|(\vec{\mathbf{S}}_{\mathbf{q}}, \vec{\mathbf{S}}_{-\mathbf{q}})|^2}{(\vec{\mathbf{S}}_{-\mathbf{q}}, \vec{\mathbf{S}}_{-\mathbf{q}})} \geq N^2 |m_{\mathbf{q}}|^2. \end{aligned}$$

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Now, consider a rotationally symmetric Hamiltonian. Taking the $\hbar \rightarrow 0^+$ limit yields

$$(3.5.0.7) \quad ss$$

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3.5.0.1. Mervin-Wagner theorem. For a system with an ordered ground state, thermally excited states reduce the spin correlations at finite temperatures, for both classical and quantum systems alike. When the temperature is much higher than the typical coupling energy scale J , the spins are expected to be uncorrelated at large distances and the magnetization $m_{\mathbf{q}}$ to vanish in the absence of an ordering field. This, thus, requires a phase transition at some temperature T_c between the ordered and disordered phases.

The ordered ground state of the Heisenberg model breaks a continuous $O(3)$ symmetry in spin-space. The spontaneous magnetization could be made to point anywhere on the S^2 -sphere by choosing the orientation of the ordering field in that direction. The consequences of this continuous symmetry are especially pronounced in lattices of low-dimensionality. In fact, it turns out, that for $d = 1$ and $d = 2$ dimensions, the phase transition occurs exactly at $T = 0$, i.e. thermal excitations disorder the spins at infinitesimally low temperatures.

Historical Note

In 1966, Hohenberg used the Bogoliubov's inequality to prove the absence of superfluidity at finite temperatures in one and two dimensions. Following the same approach, Mervin and Wagner showed that, for short-range spin models in one and two dimensions, there cannot be spontaneous ordering at any finite temperature. Said proof was initially valid for Heisenberg models only, but was then readily generalized to any system with a larger class of models with continuous symmetries.

Theorem 5 *Mervin-Wagner's Theorem*
For the quantum Heisenberg model,

$$(3.5.0.8) \quad \mathbf{H} = \frac{1}{2} \sum_{ij} J_{ij} \mathbf{S}_i \cdot \mathbf{S}_j - h \mathbf{S}_q^z, \quad \begin{array}{l} \text{with short-range} \\ \text{interaction that obey} \end{array} \quad \bar{J} = \frac{1}{2N} \sum_{ij} |J_{ij}| |\mathbf{x}_i - \mathbf{x}_j| < \infty,$$

there cannot be spontaneously broken spin symmetry at finite temperatures in one and two dimensions.

Proof. For $\mathbf{A} = \mathbf{B}$, this yields the square norm of \mathbf{A} and yields \mathbf{A} 's susceptibility of the \mathbf{A} ,

$$(\mathbf{A}, \mathbf{A})_{\text{th}} = \Re R_{\mathbf{A}\mathbf{A}}(0) = \chi_{\mathbf{A}\mathbf{A}}.$$

Let \mathbf{C} be an operator such that

$$(3.5.0.9) \quad \mathbf{B} = [\mathbf{C}^\dagger, \mathbf{H}] \Rightarrow \begin{array}{l} (\mathbf{A}, \mathbf{B}) = \langle [\mathbf{C}^\dagger, \mathbf{A}^\dagger] \rangle \\ (\mathbf{A}, \mathbf{B}) = \langle [\mathbf{C}^\dagger, [\mathbf{H}, \mathbf{C}]] \rangle, \end{array}$$

and given that Since $\tanh x < x$, this yields the following inequality

$$\begin{aligned} 0 < \frac{e^{-\beta E_m} - e^{-\beta E_n}}{E_n - E_m} &\leq \frac{1}{2T} (e^{-\beta E_m} + e^{-\beta E_n}) \Rightarrow (\mathbf{A}, \mathbf{A})_{\text{th}} \leq \frac{1}{2T} \langle \mathbf{A}^\dagger \mathbf{A} + \mathbf{A} \mathbf{A}^\dagger \rangle \\ &\Rightarrow |(\mathbf{A}, \mathbf{B})|^2 \leq (\mathbf{A}, \mathbf{A})(\mathbf{B}, \mathbf{B}) \\ &\leq \frac{1}{2T} \langle \mathbf{A}^\dagger \mathbf{A} + \mathbf{A} \mathbf{A}^\dagger \rangle (\mathbf{B}, \mathbf{B}) \end{aligned}$$

$$(3.5.0.10) \quad \Rightarrow |\langle [\mathbf{C}^\dagger, \mathbf{A}^\dagger] \rangle|^2 \leq \frac{1}{2T} \langle \mathbf{A}^\dagger \mathbf{A} + \mathbf{A} \mathbf{A}^\dagger \rangle \langle [\mathbf{C}^\dagger, [\mathbf{H}, \mathbf{C}]] \rangle,$$

the last relation is Bogoliubov's inequality. In particular, consider the following two operators, $\mathbf{C} = \mathbf{S}_{\mathbf{k}}^x$ and $\mathbf{A} = \mathbf{S}_{-\mathbf{k}-\mathbf{q}}^y$, then

$$(3.5.0.11) \quad \begin{aligned} \frac{1}{2} \langle \mathbf{A}^\dagger \mathbf{A} + \mathbf{A} \mathbf{A}^\dagger \rangle &= \langle \mathbf{S}_{\mathbf{k}+\mathbf{q}}^y \mathbf{S}_{-\mathbf{k}-\mathbf{q}}^y \rangle = N \mathbf{S}^{yy}(\mathbf{k} + \mathbf{q}), \text{ where } \mathbf{S}^{yy} \text{ is the equal-time correlation function.} \\ \langle [\mathbf{C}^\dagger, \mathbf{A}^\dagger] \rangle &= \langle [\mathbf{S}_{-\mathbf{k}}^x, \mathbf{S}_{\mathbf{k}+\mathbf{q}}^y] \rangle = i \langle \mathbf{S}_{\mathbf{q}}^z \rangle = i N m_{\mathbf{q}} \\ \langle [\mathbf{C}^\dagger, [\mathbf{H}, \mathbf{C}]] \rangle &= N^{-1} \left\langle \langle [\mathbf{S}_{-\mathbf{k}}^x, [\mathbf{H}, \mathbf{S}_{\mathbf{k}}^x]] \rangle \right\rangle. \end{aligned}$$

Now, the Bogoliubov inequality reads

$$(3.5.0.12) \quad m_{\mathbf{q}}^2 \leq \frac{1}{T} \mathbf{S}^{yy}(\mathbf{k} + \mathbf{q}) \left\langle [\mathbf{S}_{-\mathbf{k}}^x, [\mathbf{H}, \mathbf{S}_{\mathbf{k}}^x]] \right\rangle.$$

The double Hamiltonian can be bounded using the Heisenberg Hamiltonian as follows,

(3.5.0.13)

$$\begin{aligned}
\left\langle [\mathbf{S}_{-\mathbf{k}}^x, [\mathbf{H}, \mathbf{S}_{\mathbf{k}}^x]] \right\rangle &= hm_{\mathbf{q}} + \frac{i}{N} \sum_{ij\ell \in \Lambda} \left\langle \left[\mathbf{S}_i^x, [e^{i\mathbf{k} \cdot \mathbf{x}_\ell - \mathbf{x}_i} J_{j\ell} (\mathbf{S}_j^z \mathbf{S}_\ell^y - \mathbf{S}_\ell^y \mathbf{S}_j^z)] \right] \right\rangle \\
&= hm_{\mathbf{q}} + \frac{1}{N} \sum_{j\ell \in \Lambda} J_{j\ell} \left(\cos \mathbf{k}(\mathbf{x}_j - \mathbf{x}_\ell) - 1 \right) \langle \mathbf{S}_j^y \mathbf{S}_\ell^y + \mathbf{S}_j^z \mathbf{S}_\ell^z \rangle \\
&\leq hm_{\mathbf{q}} + \frac{|\mathbf{k}|^2}{2N} \sum_{j\ell \in \Lambda} |J_{j\ell}| |\mathbf{x}_j - \mathbf{x}_\ell|^2 |\langle \mathbf{S}_j^y \mathbf{S}_\ell^y + \mathbf{S}_j^z \mathbf{S}_\ell^z \rangle|
\end{aligned}$$

where an upper bound can be established $|\langle \mathbf{S}_j^y \mathbf{S}_\ell^y + \mathbf{S}_j^z \mathbf{S}_\ell^z \rangle| \leq |\langle \mathbf{S}_i \cdot \mathbf{S}_\ell \rangle| \leq S(S+1)$ by considering the maximal eigenvalues of $(\mathbf{S}_j + \mathbf{S}_\ell)^2$ and \mathbf{S}_i^2 .

$$\Rightarrow \left\langle [\mathbf{S}_{-\mathbf{k}}^x, [\mathbf{H}, \mathbf{S}_{\mathbf{k}}^x]] \right\rangle \leq hm_{\mathbf{q}} + S(S+1)\bar{J}|\mathbf{k}|^2.$$

This double commutator contains dynamical information on the system, since it explicitly depends on the Hamiltonian. The previous finite bound thus holds only if \bar{J} is itself finite, ie. the interactions decay sufficiently rapidly at large distances. All in all, combining all the previous approximations, the bound on $m_{\mathbf{q}}^2$ reads

$$\begin{aligned}
(3.5.0.14) \quad m_{\mathbf{q}}^2 &\leq \frac{1}{T} \mathbf{S}^{yy}(\mathbf{k} + \mathbf{q}) \left(hm_{\mathbf{q}} + S(S+1)\bar{J}|\mathbf{k}|^2 \right) \Rightarrow \mathbf{S}^{yy}(\mathbf{k} + \mathbf{q}) \geq \frac{Tm_{\mathbf{q}}^2}{hm_{\mathbf{q}} + S(S+1)\bar{J}|\mathbf{k}|^2} \\
\frac{1}{N} \sum_{\mathbf{k}} \mathbf{S}^{yy}(\mathbf{k} + \mathbf{q}) &\geq \frac{Tm_{\mathbf{q}}^2}{hm_{\mathbf{q}} + S(S+1)\bar{J}|\mathbf{k}|^2} \\
\frac{1}{N} \sum_i \langle (\mathbf{S}_i^y)^2 \rangle &\geq \frac{Tm_{\mathbf{q}}^2}{hm_{\mathbf{q}} + S(S+1)\bar{J}|\mathbf{k}|^2} \\
S(S+1) &\geq \frac{Tm_{\mathbf{q}}^2}{(2\pi)^d} \int_0^{\bar{k}} dk k^{d-1} \frac{1}{hm_{\mathbf{q}} + S(S+1)\bar{J}|\mathbf{k}|^2},
\end{aligned}$$

where, in the last equality, the sum over all momenta has been transformed into an integral and a lower bound has been estimated by choosing \bar{k} to be smaller than the Brillouin zone edge wave vector. This yields an analytical estimate for the right-hand side of the previous relation

(3.5.0.15)

$$S(S+1) \geq \begin{cases} \frac{Tm_{\mathbf{q}}}{2\pi\sqrt{\bar{J}S(S+1)hm_{\mathbf{q}}}} \arctan \left[\bar{k} \sqrt{\frac{\bar{J}S(S+1)}{hm_{\mathbf{q}}}} \right] & \text{if } d = 1 \\ \frac{Tm_{\mathbf{q}}}{4\pi S(S+1)\bar{J}} \log \left[1 + \frac{\bar{J}S(S+1)\bar{k}^2}{hm_{\mathbf{q}}} \right] & \text{if } d = 2 \end{cases} \Rightarrow m_{\mathbf{q}} \underset{h \rightarrow 0^+}{\leq} \begin{cases} c \frac{S(S+1)\bar{J}^{\frac{1}{3}}}{T^{\frac{2}{3}}} h^{\frac{1}{3}} & \text{if } d = 1 \\ \leq c \frac{S(S+1)\sqrt{\bar{J}}}{\sqrt{T}} \frac{1}{\sqrt{|\log h|}} & \text{if } d = 2 \end{cases}$$

These two equations thus prove that the magnetization must vanish with the ordering field for all finite values of T, S, \bar{J} . Now, these bounds do not depend on the lattice size. Then, these results hold in the thermodynamic limit.

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This result holds for all S and it holds for the classical S -independent limit of the Heisenberg model

$$(3.5.0.16) \quad \mathbf{H}^{\text{cl}} = \frac{1}{2} \sum_{ij \in \Lambda} J_{ij}^{\text{cl}} \hat{\Omega}_i \cdot \hat{\Omega}_j - h^{\text{cl}} \hat{\Omega}_{\mathbf{q}}^z,$$

where $\hat{\Omega}$ are c-number unit vectors. The correspondence between the quantum and classical Heisenberg models is as follows

$$J_{ij}S(S+1) \rightarrow J^{\text{cl}} \quad hS \rightarrow h^{\text{cl}} \quad \frac{m_{\mathbf{q}}}{S} \rightarrow m_{\mathbf{q}}^{\text{cl}}.$$

Using these parameters eliminates S from inequalities, thus the theorem holds for the classical model as well.

In $d = 3$ -dimensions, the integral does not diverge for $h \rightarrow 0^+$. Therefore, it is possible to satisfy the inequality at some finite temperature. Therefore, spontaneous symmetry breaking and an ordered phase is expected at low temperatures in the three dimensions.

3.5.0.2. Quantum Disorder at $T = 0$. Mervin-Wagner's theorem does not hold at zero temperature. Consider a quantum system in the presence of an infinitesimal ordering field, such that there is a unique ground state $|0\rangle$. At zero temperature, the expectation value for any operator is simply $\langle \mathbf{A} \rangle = \langle 0 | \mathbf{A} | 0 \rangle$. The scalar product is then given by

$$(\mathbf{A}, \mathbf{B}) = \Re \sum_{m \neq 0} \frac{\langle 0 | \mathbf{A}^\dagger | m \rangle \langle m | \mathbf{B} | 0 \rangle}{E_m - E_0}.$$

Then, consider the previous choices for the \mathbf{A} , \mathbf{B} and \mathbf{C} . The susceptibility is given by

$$(3.5.0.17) \quad \chi^{yy}(\mathbf{k}) = \frac{1}{N} (\mathbf{S}_{\mathbf{k}}^y, \mathbf{S}_{-\mathbf{k}}^y) = \frac{1}{N} \sum_{m \neq 0} \frac{|\langle 0 | \mathbf{S}_{\mathbf{k}}^y | m \rangle|^2 + |\langle 0 | \mathbf{S}_{-\mathbf{k}}^y | m \rangle|^2}{E_m - E_0}$$

$$\Rightarrow m_{\mathbf{q}}^2 \leq \chi^{yy}(\mathbf{k} + \mathbf{q}) \left[\mathbf{S}_{-\mathbf{k}}^x, [\mathbf{H}, \mathbf{S}_{\mathbf{k}}^x] \right]$$

$$\leq \chi^{yy} [h m_{\mathbf{q}} + S(S+1) \bar{J} |\mathbf{k}|^2]$$

where, in the last line, [equation \(3.5.0.10\)](#) and [equation \(3.5.0.13\)](#) have been used, thus yielding the zero-temperature Bogoliubov inequality. Inverting this inequality and summing over the Brillouin zone reads

$$(3.5.0.18) \quad \frac{m_{\mathbf{q}}^2}{(2\pi)^d} \int_0^{\bar{k}} dk \frac{k^{d-1}}{h m_{\mathbf{q}} + S(S+1) \bar{J} |\mathbf{k}|^2} \leq \frac{1}{N} \sum_{\mathbf{k}} \chi^{yy}(\mathbf{k}).$$

Therefore, if the right hand side is finite, $\frac{1}{N} \sum_{\mathbf{k}} \chi^{yy}(\mathbf{k}) = C < \infty$, and $d = 1$ or $d = 2$, $m_{\mathbf{q}}$ must vanish with h in order for the last equation to hold. In that case, there will be no long-range order at zero temperature.

Furthermore, it can be proved that if there is a gap in the excitation spectrum,

$$E_m - E_0 > \Delta, \quad \forall m \neq 0,$$

where Δ is independent of h , N , the ground state of Heisenberg model must be disordered. This gap induces a bound on the susceptibility, said bound given by the correlation function

$$(3.5.0.19) \quad \chi^{yy}(\mathbf{k}) \leq \frac{2}{N\Delta} \sum_m \left| \langle 0 | \mathbf{S}_{\mathbf{k}}^y | m \rangle \right|^2 = \frac{2\mathbf{S}^{yy}(\mathbf{k})}{\Delta},$$

and given that $\frac{1}{N} \sum_{\mathbf{k}} \mathbf{S}^{yy}(\mathbf{k}) \leq S(S+1) < \infty$, the "finite-ness" condition of the susceptibility is satisfied. Here, the energy gap plays the role of temperature in the previous example.

Physical Context

As an example, the antiferromagnetic integer spin chains are known to exhibit the so-called Haldane gap in their excitation spectrum. The Mermin-Wagner theorem at zero temperature implies that these models do not possess true long-range order in their ground states. In fact, the non-linear sigma model analysis predicts that their correlations exponentially decay at large distances. The converse statement (gapless excitations imply long-range order) is false. For example, the spin-half Heisenberg antiferromagnet in one dimension has gapless excitations but no long-range order at zero temperature.

Some interesting implications:

Long-range order and magnetic susceptibility

Given that

$$\lim_{h \rightarrow 0^+} m_{\mathbf{q}} = - \lim_{h \rightarrow 0^-} m_{-\mathbf{q}},$$

long-range order requires the divergence of the zero-frequency susceptibility

$$\lim_{h, \omega \rightarrow 0^+} \chi^{zz}(\mathbf{q}, \omega) = \infty.$$

Proof. In effect, the zero-frequency susceptibility is defined as

$$\begin{aligned} \chi^{zz}(\mathbf{q}, \omega) &= \frac{1}{N} \langle \mathbf{S}_{\mathbf{q}}^z, \mathbf{S}_{\mathbf{q}}^z \rangle \Rightarrow \langle \mathbf{S}_{\mathbf{q}}^z, \mathbf{S}_{-\mathbf{q}}^z \rangle \leq \frac{1}{2T} \langle (\mathbf{S}_{\mathbf{q}}^z \mathbf{S}_{-\mathbf{q}}^z) + (\mathbf{S}_{-\mathbf{q}}^z, \mathbf{S}_{\mathbf{q}}^z) \rangle = \frac{1}{T} \langle \mathbf{S}_{\mathbf{q}}^z \mathbf{S}_{\mathbf{q}}^z \rangle \\ (3.5.0.20) \quad &\Rightarrow \langle \mathbf{S}_{\mathbf{q}}^z \mathbf{S}_{-\mathbf{q}}^z \rangle \geq T \langle \mathbf{S}_{\mathbf{q}}^z, \mathbf{S}_{-\mathbf{q}}^z \rangle \\ &\Rightarrow \frac{\langle \mathbf{S}_{\mathbf{q}}^z \mathbf{S}_{-\mathbf{q}}^z \rangle}{N} \geq T \chi^{zz}(\mathbf{q}, \omega). \end{aligned}$$

If the system has true-long range order, then it follows that the two-point correlator has a non-zero limit, in the thermodynamic limit, this is

$$\lim_{N \rightarrow \infty} \frac{S^{\alpha\alpha}(\mathbf{q})}{N} > 0.$$

In particular, choose $\alpha = z$ -direction, then the $S^{zz}(\mathbf{q})$ correlator is given by

$$S^{zz}(\mathbf{q}) = \lim_{h \rightarrow 0^+} \frac{1}{N\mathcal{Z}} \text{Tr} \left[e^{-\beta \mathbf{H}} \mathbf{S}_{\mathbf{q}}^z \mathbf{S}_{-\mathbf{q}}^z \right], \text{ with } \alpha = \{x, y, z\}.$$

Then, consider [equation \(3.5.0.20\)](#)

$$\begin{aligned} \frac{\langle \mathbf{S}_{\mathbf{q}}^z \mathbf{S}_{-\mathbf{q}}^z \rangle}{N} &\geq T \chi^{zz}(\mathbf{q}, \omega) \Rightarrow e^{-\beta \mathbf{H}} \frac{\langle \mathbf{S}_{\mathbf{q}}^z \mathbf{S}_{-\mathbf{q}}^z \rangle}{N} \geq e^{-\beta \mathbf{H}} T \chi^{zz}(\mathbf{q}, \omega) \\ (3.5.0.21) \quad &\Rightarrow \frac{1}{\mathcal{Z}} \text{Tr} \left[e^{-\beta \mathbf{H}} \frac{\langle \mathbf{S}_{\mathbf{q}}^z \mathbf{S}_{-\mathbf{q}}^z \rangle}{N} \right] \geq \frac{1}{\mathcal{Z}} \text{Tr} \left[e^{-\beta \mathbf{H}} T \chi^{zz}(\mathbf{q}, \omega) \right] \\ &\Rightarrow \lim_{N \rightarrow \infty} \lim_{h \rightarrow 0^+} \frac{1}{\mathcal{Z}} \text{Tr} \left[\frac{\langle \mathbf{S}_{\mathbf{q}}^z \mathbf{S}_{-\mathbf{q}}^z \rangle}{N} \right] \geq \lim_{N \rightarrow \infty} \lim_{h \rightarrow 0^+} T \chi^{zz}(\mathbf{q}, \omega) \\ &\Rightarrow \lim_{h \rightarrow 0^+} \frac{S^{zz}(\mathbf{q})}{N} \geq \lim_{N \rightarrow \infty} \lim_{h \rightarrow 0^+} T \chi^{zz}(\mathbf{q}, \omega) \end{aligned}$$

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No broken spin symmetry for the Two-dimensional antiferromagnetic Heisenberg model

Marshall's theorem implies that there cannot be any broken spin symmetry for the Heisenberg antiferromagnet on a finite bipartite lattice with two equal size sublattices.

Proof.

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3.6. Spin Representations.

3.6.0.1. Holstein-Primakoff Bosons. The spin is a vector operator such that, in a broken symmetry phase, the expectation value of at least one of its components is non-zero. Thus, it is natural to describe the ordered phase in terms of small spin fluctuations about their expectation values. This is the main idea behind spin wave theory. Consider now the Holstein-Primakoff bosonic operator \mathbf{b} , which represents the three spin component operators, as follows

$$(3.6.0.1) \quad \mathbf{S}^+ = \left(\sqrt{2S - \mathbf{b}^\dagger \mathbf{b}}\right) \mathbf{b} \quad \mathbf{S}^- = \mathbf{b}^\dagger \left(\sqrt{2S - \mathbf{b}^\dagger \mathbf{b}}\right) \quad \mathbf{S}^z = S - \mathbf{b}^\dagger \mathbf{b},$$

such that $[\mathbf{b}, \mathbf{b}^\dagger] = 1$. These spin operators, for their part, obey the standard $\mathfrak{su}(2)$ -spin commutation relations. In this mapping, the vacuum states has spin of $+S$ in the z -direction and each Holstein-Primakoff boson represents a spin-1 moment in the $-z$ -direction, thereby they represent a perturbation from the classical ferromagnetic ground state, with the antiferromagnetic case being more complicated. With this physical picture in mind, the $(\sqrt{2S - \mathbf{b}^\dagger \mathbf{b}})$ -factor -present in both the raising and lowering spin operators- limits the number of bosons on a given site to $2S$, since the z -projection of the spin moment at a given site must be between $-S$ and $+S$.

Spurious Fock states, ie. those with $\langle \mathbf{b}^\dagger \mathbf{b} \rangle_\Psi > 2S$, are eliminated by a projector \mathbf{P}_S . In this subspace, the $\mathfrak{su}(2)$ -Casimir invariant holds, $\mathbf{S}^2 = S(S+1)$. Furthermore, note that these spin operators do not connect the physical and unphysical subspaces. Furthermore, at low temperatures ($\beta^{-1} \ll J$, with J being the exchange energy), the number of perturbations about the classical ground state is very small $\langle \mathbf{n} \rangle \ll S$, a linear approximation can be performed

$$(3.6.0.2) \quad \mathbf{S}^z = S - \mathbf{b}^\dagger \mathbf{b}, \quad \mathbf{S}^+ \approx \mathbf{b} \sqrt{2S} \quad \mathbf{S}^- \approx \mathbf{b}^\dagger \sqrt{2S}.$$

The Holstein-Primakoff representation is useful to describe broken symmetry phases of the quantum Heisenberg model. The square root factors in [equation \(3.6.0.1\)](#) can be expanded in powers of S^{-1} ,

$$\sqrt{2S - \mathbf{b}^\dagger \mathbf{b}} = \sqrt{2S} \left(1 - \frac{(\mathbf{b}^\dagger \mathbf{b})^2}{4S} - \frac{n^2}{32S^2} \right),$$

which is a semiclassical expansion of the spin fluctuations about the z -direction. Using this expansion into the Hamiltonian and preserving only quadratic terms, yields a non-interacting spin wave Hamiltonian, where the higher order terms introduce interactions between the spin waves. The truncation will, however, couple the physical and unphysical subspaces ¹⁹.

¹⁹

Physical Context

The low-order spin wave approximation has been found to be remarkably succesful in certain ordered phases of the Heisenberg ferromagnet and antiferromagnet.

3.6.0.2. **Schwinger Bosons.** The symmetric phases of the Heisenberg model are easier to describe using representations in which the rotational invariance of the Hamiltonian is manifested. Two Schwinger bosons, \mathbf{a} , \mathbf{b} , represent the spin operators as follows

$$(3.6.0.3) \quad \mathbf{S}^x + i\mathbf{S}^y = \mathbf{a}^\dagger \mathbf{b} \quad \mathbf{S}^x - i\mathbf{S}^y = \mathbf{b}^\dagger \mathbf{a} \quad \mathbf{S}^z = \frac{1}{2}(\mathbf{a}^\dagger \mathbf{a} - \mathbf{b}^\dagger \mathbf{b}).$$

These spin components satisfy the $\mathfrak{su}(2)$ -commutation relations as well. The spin magnitude S defines the physical subspace,

$$(3.6.0.4) \quad \left\{ |n_{\mathbf{a}}, n_{\mathbf{b}}\rangle \mid n_{\mathbf{a}} + n_{\mathbf{b}} = 2S \right\}, \quad \text{with the physical subspace given by a projector } \mathbf{P}_S \text{ such that} \quad \mathbf{P}_S(\mathbf{a}^\dagger \mathbf{a} + \mathbf{b}^\dagger \mathbf{b} - 2S) = 0.$$

$$\mathbf{S}^2 \mathbf{P}_S = S(S+1) \mathbf{P}_S.$$

where the last relation indicates that the spin magnitude is well defined in the projected (physical) subspace. The spin states are given by

$$(3.6.0.5) \quad |S, m\rangle = \frac{(\mathbf{a}^\dagger)^{S+m}}{\sqrt{(S+m)!}} \frac{(\mathbf{b}^\dagger)^{S-m}}{\sqrt{(S-m)!}} |0\rangle,$$

where $|S, m\rangle$ labels the Casimir invariant's and the z -spin component eigenvalues and where $|0\rangle$ is the Schwinger bosons vacuum. For example, the spin- $\frac{1}{2}$ states are given by $|\uparrow\rangle = \mathbf{a}^\dagger |0\rangle$ and $|\downarrow\rangle = \mathbf{b}^\dagger |0\rangle$.

Physical Context

Schwinger bosons are useful for calculating matrix elements of spin operators. Since they do not contain square roots (in contrast to Holstein-Primakoff bosons), the matrix elements of spin operators between Fock states factorize as for free bosons. However, this does not necessarily simplify the calculation of spin correlations in non-Fock wavefunctions. The local constraints on the Hilbert space introduce correlations between the a and b occupation numbers.

A generalization to the spin- $\frac{1}{2}$ representation to N flavours defines the generators of $SU(N)$ as generalized spins. The large- N generalizations are amenable to simple mean field theories which are controlled by the small parameter $\frac{1}{N}$. The large- N mean field theories are described by effectively non-interacting Bose quasiparticles, where correlations between different flavours are ignored.

The Schwinger bosons and the Holstein-Primakoff bosons are closely related. In effect, by eliminating the a -boson using the constraint, the correspondence is found to be

$$(3.6.0.6) \quad \begin{array}{ccc} \text{Schwinger} & \leftrightarrow & \text{Holstein-Primakoff} \\ \mathbf{b} & \leftrightarrow & \mathbf{b} \\ \mathbf{a} & \leftrightarrow & \sqrt{2S - \mathbf{b}^\dagger \mathbf{b}}. \end{array}$$

While Schwinger bosons provide a symmetric representation in spin space, Holstein-Primakoff bosons single out the \mathbf{S}^z direction, which defines their vacuum. Therefore, these two representations lend themselves to different approximation schemes: Holstein-Primakoff for the broken symmetry phases and Schwinger for the symmetry phases.

Furthermore, Schwinger boson creation operators transform as vectors under $SU(2)$ global rotations,

$$(3.6.0.7) \quad \begin{pmatrix} \mathbf{a}^\dagger \\ \mathbf{b}^\dagger \end{pmatrix} = \mathcal{R} \begin{pmatrix} \mathbf{a}^\dagger \\ \mathbf{b}^\dagger \end{pmatrix} \mathcal{R}^{-1} = e^{i\frac{1}{2}\chi\sigma^z} e^{i\frac{1}{2}\theta\sigma^y} e^{i\frac{1}{2}\phi\sigma^x} \begin{pmatrix} \mathbf{a}^\dagger \\ \mathbf{b}^\dagger \end{pmatrix}$$

3.6.0.3. Spin Coherent States. Spin coherent states are a family of spin states created by applying the rotation operator \mathcal{R} to the maximally polarized state $|S, S\rangle$,

$$(3.6.0.8) \quad \begin{aligned} |\hat{\Omega}\rangle &= \mathcal{R}(\chi, \theta, \varphi) |S, S\rangle \\ &= e^{i\mathbf{S}^z \varphi} e^{i\mathbf{S}^y \theta} e^{i\mathbf{S}^x \chi} |S, S\rangle, \end{aligned}$$

where the unit vector

$$(3.6.0.9) \quad \hat{\Omega} = (\sin \theta \cos \phi, \sin \theta \sin \phi, \cos \theta),$$

parametrizes the spin coherent state. Note that these \mathcal{R} -transformations form a group, $SU(2)$ precisely, and are parametrized by the Euler angles φ, θ, χ . There is a gauge freedom, since χ can be arbitrarily defined. The two independent variables are thus the θ -angle, the latitude, and φ -angle, the longitude, which are within the domains

$$\{\theta \in [0, \pi], \phi \in [-\pi, \pi)\}.$$

The spin coherent states are not orthogonal, since

$$\langle \Omega | \Omega' \rangle = \left(\frac{1 + \hat{\Omega} \cdot \hat{\Omega}'}{2} \right)^S e^{-iS\phi},$$

where ψ is a function on the six Euler angles. For references, check [?].

Spin coherent states are not orthogonal, as seen by the preceding relation, but they do span the space of states of spin S . The measure of integration over the group parameters is defined as

$$(3.6.0.10) \quad \frac{2S+1}{4\pi} \hat{\Omega} = \frac{2S+1}{4\pi} d\theta \sin \theta d\phi.$$

This is the Haar measure of the $SU(2)$ -Lie group. Using this measure, the resolution of identity is provided by the integral

$$(3.6.0.11) \quad \frac{2S+1}{4\pi} \int_{S^3} d\hat{\Omega} |\hat{\Omega}\rangle \langle \hat{\Omega}| = \sum_m |S, m\rangle \langle S, m| = \text{id}_{SU(2)},$$

thus the coherent states form an overcomplete basis. Another useful relation can be proven in the same fashion

$$(3.6.0.12) \quad \frac{(S+1)(2S+1)}{4\pi} \int_{S^3} d\hat{\Omega} \hat{\Omega}^\alpha |\hat{\Omega}\rangle \langle \hat{\Omega}| = \mathbf{S}^\alpha, \quad \alpha = x, y, z.$$

The previous relations can be readily extended to many spins. In effect, consider an N -site lattice. The many spin coherent states are products of single spin states,

$$(3.6.0.13)$$

$$|\hat{\Omega}\rangle = \prod_{i=1}^N |\hat{\Omega}_i\rangle, \quad \text{whose overlap is now given by } \langle \hat{\Omega} | \hat{\Omega}' \rangle = \prod_i \left(\frac{1 + \hat{\Omega}_i \cdot \hat{\Omega}'_i}{2} \right)^S e^{-iS \sum_i \psi[\hat{\Omega}_i, \hat{\Omega}'_i]}.$$

$$\text{The resolution of the identity in the product space is } \int_{\otimes SU(2)} d\hat{\Omega}_i \frac{2S+1}{4\pi} \langle \hat{\Omega} | \hat{\Omega} \rangle = \text{id}_{\otimes SU(2)}.$$

$$\text{The spin correlations of any wave function } |\Psi\rangle \text{ can be computed via the integral } \frac{\langle \Psi | \mathbf{S}_i \cdot \mathbf{S}_j | \Psi \rangle}{\langle \Psi | \Psi \rangle} = \frac{(S+1 - \delta_{ij})(S+1)}{\mathcal{Z}} \int_{\otimes SU(2)} d\hat{\Omega}_i |\Psi[\hat{\Omega}]|^2 \hat{\Omega}_i \cdot \hat{\Omega}_j,$$

$$\text{with } \mathcal{Z} = \int_{\otimes SU(2)} d\hat{\Omega}_i |\Psi[\hat{\Omega}]|^2,$$

where note that the representation $\Psi[\hat{\Omega}] = \langle \Psi | \hat{\Omega} \rangle$ is continuous. Furthermore, all spin operators can be represented by differential operators. Spin coherent states can then be used to evaluate the trace of any operator. If $|n\rangle$ is any orthonormal basis, then it

$$\begin{aligned}
 \text{Tr}\{\mathbf{O}\} &= \sum_n \langle n | \mathbf{O} | n \rangle = \left(\frac{2S+1}{4\pi} \right)^2 \int_{S^3} d\hat{\Omega} \int_{S^3} d\hat{\Omega}' \sum_n \langle n | \hat{\Omega} \rangle \langle \hat{\Omega} | \mathbf{O} | \hat{\Omega}' \rangle \langle \hat{\Omega}' | n \rangle \\
 (3.6.0.14) \quad &= \frac{2S+1}{4\pi} \int_{S^3} d\hat{\Omega} \langle \hat{\Omega} | \mathbf{O} | \hat{\Omega} \rangle.
 \end{aligned}$$

The coherent states elucidate the correspondence between classical and quantum spins. The classical limit is achieved when $S \rightarrow \infty$. In that limit, the overlap of different coherent states vanishes exponentially with S , as can be seen from the first identity of [equation \(3.6.0.14\)](#). The expectation values of spin operators are functions of unit vectors, exactly like classical spins. Quantum effects are therefore associated with the non-orthogonality of coherent states, which also implies a finite width of $\Psi[\hat{\Omega}]$ in $\hat{\Omega}$ -space.

4. SPIN WAVE THEORY

Spin Wave Theory Spin Wave Theory

4.1. Spin Path Integrals. Path integrals provide formal expressions which, in general without a known closed analytic form, lead to useful approximation schemes, such as generating asymptotic expansions and to formulate mean field theories.

The spin coherent states path integral describes quantum spins in terms of time-dependent histories of unit vectors, thus providing a connection between the classical and quantum phenomena.

4.1.1. Construction of the path-integrals. Spin coherent states can be used to construct a path integral representation of the Heisenberg model, with its generating functional in imaginary time being

$$(4.1.1.1) \quad \mathcal{Z}[j] = \text{Tr } \mathcal{T}_\tau \exp \left(- \int_0^\beta d\tau \mathbf{H}(\tau) \right) \quad \text{where } \mathcal{T}_\tau \text{ is the time-ordering operator, } \epsilon = \frac{\beta}{N_\epsilon},$$

$$= \lim_{N \rightarrow \infty} \text{Tr } \mathcal{T}_\tau \prod_{n=0}^{N_\epsilon-1} [1 - \epsilon \mathbf{H}(\tau_n)], \quad \text{and where } \tau_n = n\epsilon \text{ is the discrete imaginary time}$$

$$\quad \text{and where the generating Hamiltonian includes source currents } \mathbf{H}(\tau) = \mathbf{H} - \sum_{i,\alpha} j_i^\alpha(\tau) \mathbf{S}_i^\alpha, \text{ for } \tau \in [0, \beta).$$

The spin coherent resolution identity is given by [equation \(3.6.0.11\)](#), wherein the spin operators in the spin coherent states are given by [equation \(3.6.0.12\)](#). Using these identities, the generating functional may be rewritten as a tridimensional integral

$$(4.1.1.2) \quad \mathcal{Z}[j] = \lim_{N \rightarrow \infty} \text{Tr } \mathcal{T}_\tau \prod_{n=0}^{N_\epsilon-1} [1 - \epsilon \mathbf{H}(\tau_n)]$$

$$= \int \left(d\hat{\Omega}_0 d\hat{\Omega}_\epsilon \cdots d\hat{\Omega}_{N_\epsilon-1} \right) \left\langle \hat{\Omega}_0 \hat{\Omega}_\epsilon \cdots \hat{\Omega}_{N_\epsilon-1} \left| \mathcal{T}_\tau [1 - \epsilon \mathbf{H}(\tau_0)] [1 - \epsilon \mathbf{H}(\tau_\epsilon)] \cdots [1 - \epsilon \mathbf{H}(\tau_{N_\epsilon-1})] \right| \hat{\Omega}_0 \hat{\Omega}_\epsilon \cdots \hat{\Omega}_{N_\epsilon-1} \right\rangle$$

$$= \int \left(d\hat{\Omega}_0 d\hat{\Omega}_\epsilon \cdots d\hat{\Omega}_{N_\epsilon-1} \right) \left\langle \hat{\Omega}_0 \hat{\Omega}_\epsilon \cdots \hat{\Omega}_{N_\epsilon-1} \left| \right. \right.$$

$$\quad \times \mathcal{T}_\tau [1 - \epsilon \mathbf{H}(\tau_0)] \text{id}_{SU(2)} [1 - \epsilon \mathbf{H}(\tau_\epsilon)] \text{id}_{SU(2)} \cdots \text{id}_{SU(2)} [1 - \epsilon \mathbf{H}(\tau_{N_\epsilon-1})] \text{id}_{SU(2)} \left. \right| \hat{\Omega}_0 \hat{\Omega}_\epsilon \cdots \hat{\Omega}_{N_\epsilon-1} \rangle$$

$$= \int \left(d\hat{\Omega}_0 d\hat{\Omega}_\epsilon \cdots d\hat{\Omega}_{N_\epsilon-1} \right) \left\langle \hat{\Omega}_0 \hat{\Omega}_\epsilon \cdots \hat{\Omega}_{N_\epsilon-1} \left| \right. \right.$$

$$\quad \times \mathcal{T}_\tau \int d\hat{\Omega}'_0 \left| \hat{\Omega}'_0 \right\rangle \langle \hat{\Omega}'_0 | [1 - \epsilon \mathbf{H}(\tau_0)] \int d\hat{\Omega}'_1 \left| \hat{\Omega}'_1 \right\rangle \langle \hat{\Omega}'_1 | [1 - \epsilon \mathbf{H}(\tau_\epsilon)] \int d\hat{\Omega}'_2 \left| \hat{\Omega}'_2 \right\rangle \langle \hat{\Omega}'_2 | \cdots$$

$$\quad \times \int d\hat{\Omega}'_{N_\epsilon-1} \left| \hat{\Omega}'_{N_\epsilon-1} \right\rangle \langle \hat{\Omega}'_{N_\epsilon-1} | [1 - \epsilon \mathbf{H}(\tau_{N_\epsilon})] \int d\hat{\Omega}'_{N_\epsilon} \left| \hat{\Omega}'_{N_\epsilon} \right\rangle \langle \hat{\Omega}'_{N_\epsilon} | \hat{\Omega}_0 \hat{\Omega}_\epsilon \cdots \hat{\Omega}_{N_\epsilon-1} \rangle$$

$$\quad \hat{\Omega} = (\hat{\Omega}_1, \cdots, \hat{\Omega}_N)$$

$$= \lim_{N_\epsilon \rightarrow \infty} \int \left(\prod_{r,i} \Omega_{i\tau} \right) \prod_{\tau=\epsilon}^\beta \left\langle \hat{\Omega}(\tau) \left| \hat{\Omega}(\tau - \epsilon) \right. \right\rangle [1 - \epsilon \mathcal{H}(\tau)], \text{ where}$$

$$\mathcal{H}(\tau) = \frac{\langle \hat{\Omega}(\tau) | \mathbf{H}(\tau) | \hat{\Omega}(\tau - \epsilon) \rangle}{\langle \hat{\Omega}(\tau) | \hat{\Omega}(\tau - \epsilon) \rangle}$$

where, in the second line a total of N_ϵ [equation \(3.6.0.11\)](#)- identities were introduced and labelled as $\hat{\Omega}'_i$, where $\hat{\Omega}_\tau = (\sin \theta_\tau \cos \phi_\tau, \sin \theta_\tau \sin \phi_\tau, \cos \theta_\tau) = \Omega_{i\tau}$ and where $\mathcal{H}(\tau)$ is a "classical Hamiltonian". In the last line, the $\hat{\Omega}'_i$ s were identified with $\hat{\Omega}_{i-\epsilon}$. Let $\hat{\Omega}(\beta) = \hat{\Omega}(0)$.

A priori, $\hat{\Omega}(\tau)$ is an arbitrary discrete function on τ . Though unjustified, it may be taken to be a continuous and differentiable function in the $(N_\epsilon \rightarrow \infty)$ -limit, in such a way that

$$(4.1.1.3) \quad \frac{\hat{\Omega}_i(\tau + \epsilon) - \hat{\Omega}_i(\tau)}{\epsilon} = \dot{\hat{\Omega}}_i(\tau) + \mathcal{O}(\epsilon),$$

holds. The implicit assumption is that the generating functional is dominated by smooth paths, ie. $|\dot{\hat{\Omega}}| < \infty$. This turns out to be unjustified. By ignoring discontinuous paths, information is lost about the ordering of operators in the quantum Hamiltonian. For this reason, path integral results must be checked against operator methods, which do not suffer from ordering ambiguities. The overlap between coherent states, given in the first line of [equation \(3.6.0.13\)](#), induces the following representation for the overlap between coherent states at nearby timesteps, to leading order in ϵ , as

$$(4.1.1.4) \quad \left\langle \hat{\Omega}(\tau) \left| \hat{\Omega}(\tau - \epsilon) \right. \right\rangle = \exp \left(-iS\epsilon \sum_i \dot{\phi}_i \cos(\theta_i(\tau)) + \dot{\chi}_i \right),$$

where $\chi(\tau)$ is an arbitrary gauge convention. Furthermore, the classical Hamiltonian, which is multiplied by ϵ , can be evaluated at equal times, $\mathcal{H}[\hat{\Omega}(\tau)] \rightarrow \langle \hat{\Omega}(\tau) | \mathbf{H}(\tau) | \hat{\Omega}(\tau) \rangle$. The $(N_\epsilon \rightarrow \infty)$ -limit will transform [equation \(4.1.1.2\)](#) into a path integral, as follow

$$(4.1.1.5) \quad \begin{aligned} \mathcal{Z}[j] &= \oint \mathcal{D}\hat{\Omega}(\tau) e^{-\tilde{\mathcal{S}}[\hat{\Omega}]} \\ \tilde{\mathcal{S}}[\hat{\Omega}] &= -iS \sum_i \omega[\hat{\Omega}_i] + \int_0^\beta \mathcal{H}[\hat{\Omega}(\tau)] \end{aligned} \quad \text{with the integration measure } \mathcal{D}\hat{\Omega}(\tau) = \lim_{N_\epsilon \rightarrow \infty} \prod_{i,n} d\hat{\Omega}_i(\tau_n),$$

where, in the first and second lines, the Hamiltonian was exponentiated and higher-order terms in ϵ were discarded and the formal continuum limit was taken. The \oint -notation reflects the system's periodic boundary conditions. For the $\chi_i(\tau) = 0$, the ω -functional depends on the history of a single spin as follows

$$(4.1.1.6) \quad \omega[\hat{\Omega}] = - \int_0^\beta d\tau \dot{\phi} \cos \theta = - \int_{\phi_0}^{\phi_0} d\phi \cos \theta_\phi,$$

which implies that ω is geometric: it depends on the trajectory of $\hat{\Omega}_i(\tau)$ on the 3-sphere and not on its explicit time dependence. The $(S\omega)$ -functional is called the Berry phase of the spin history, since it describes the phase acquired by a spin that aligns with an adiabatically rotating external magnetic field which is parallel to $\hat{\Omega}_i(\tau)$ ²⁰.

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The Berry phase measures the area enclosed by the $\hat{\Omega}(\tau)$ -path on the unit three-sphere. A path increment $d\hat{\Omega}$ is connected to the north pole by longitudes ϕ and $\phi + d\phi$. This area increment is a triangle on the sphere whose area is given by $d\omega' = (1 - \cos \theta)d\phi$. Thus, the area enclosed by this closed orbit, parametrized by $\theta(\tau), \phi(\tau)$ is given by

$$\omega' = \int_0^\beta d\tau \dot{\phi} (1 - \cos \theta).$$

Therefore, ω' is the area enclosed on the left on the counterclockwise orbit $\{\hat{\Omega}(\tau)\}_0^\beta$. The $(\omega' = \omega)$ -identity holds if $\phi(\beta)$ doesn't cross the "date line boundary" $\pm\pi$.

It is useful to express the ω -Berry phase in a gauge invariant form, ie. without specifying any parametrization on the sphere. Consider then a vector potential $\mathbf{A}(\hat{\Omega})$ such that

$$\omega = \int_0^\beta d\tau \mathbf{A}(\hat{\Omega}) \cdot \dot{\hat{\Omega}},$$

where $\hat{\Omega}$ is a unit magnetic monopole vector potential, whose line integral over the $\{\hat{\Omega}(\tau)\}_0^\beta$ -orbit is equal to the solid angle ω subtended by that orbit. By Stokes' theorem, this vector potential satisfies

$$(4.1.1.7) \quad \nabla \times \mathbf{A} \cdot \hat{\Omega} = \epsilon^{\alpha\beta\gamma} \partial_{\hat{\Omega}_\alpha} A^\beta \hat{\Omega}_\gamma = 1,$$

4.1.2. *The Green Function.* The Green function, which describes the real-time evolution at zero-temperature, is given by the matrix element of the evolution operator between two coherent states,

$$(4.1.2.1) \quad \mathfrak{G}(\hat{\Omega}_0, \hat{\Omega}_t; t) = \langle \hat{\Omega}_t | \mathcal{T}_t \exp \left(-i \int_0^t dt' \mathbf{H}(t') \right) | \hat{\Omega}_0 \rangle,$$

whose path integral representation can be found in close analogy to that obtained for the generating \mathcal{Z} -functional by replacing the imaginary time τ -variable with the real time t ,

$$\tau \rightarrow it' \text{ and } t' \in [0, t],$$

and the $[0, t]$ -interval is discretized into N_ϵ -timesteps of ϵ -width. This procedure yields the formal expression,

$$(4.1.2.2) \quad \mathfrak{G}(t) = \int_{\hat{\Omega}_0}^{\hat{\Omega}_t} \mathcal{D}\hat{\Omega}(t') e^{i\mathcal{S}[\hat{\Omega}]}, \text{ where } \mathcal{S}[\hat{\Omega}] = \int_0^t dt' \left(S \sum_i \mathbf{A} \cdot \dot{\hat{\Omega}} - \mathcal{H}[\hat{\Omega}(t'), t'] \right).$$

Here $\mathcal{S}[\hat{\Omega}]$ is the real-time action.

4.1.3. *The large- S expansion.* The spin coherent path integrals are useful starting points for deriving semiclassical approximations, wherein coordinates are simply unit vectors ie. classical spins. The quantum effects enter through their real or imaginary time dependence. Consider the classical Hamiltonian \mathcal{H} , which is the result obtained by scaling the parameters of $\mathbf{H}[\hat{\Omega}]$ such that these are independent of S . By considering the $(S \rightarrow \infty)$ -limit, all time-dependent paths with $\dot{\hat{\Omega}} \neq 0$ are suppressed by the rapid oscillations of the Berry phase factor. Thus, a classical partition function is obtained

$$(4.1.3.1) \quad \mathcal{Z} \underset{S \rightarrow \infty}{\sim} \mathcal{Z}' \int \mathcal{D}\hat{\Omega} e^{-\beta \mathcal{H}[\hat{\Omega}]},$$

where \mathcal{Z}' includes normalization factors and higher-order quantum corrections and where the integral represents the generating functional for the classical Hamiltonian. This allows for the use of S as the control parameter for a systematic asymptotic expansion for the Green function. This semiclassical expansion introduces quantum corrections to the classical theory. The first step is to rescale the time variable

$$(4.1.3.2) \quad \tau \rightarrow S\tau = \bar{\tau} \quad \beta \rightarrow S\beta = \bar{\beta}.$$

The classical inverse temperature β is taken to be independent of S , allowing for scaling S out of action, yielding

$$(4.1.3.3) \quad \mathcal{Z}(\bar{\beta}) = \oint \mathcal{D}\hat{\Omega}[\tau] \exp \left(-S\mathcal{S}^{\text{cl}}[\hat{\Omega}(\bar{\tau}), \bar{\beta}] \right), \text{ with } \mathcal{S}^{\text{cl}} = \int_0^{\bar{\beta}} d\bar{\tau} \left(i \sum_i \mathbf{A} \cdot \dot{\hat{\Omega}} + \mathbf{H}^{\text{cl}}[\hat{\Omega}(\tau)] \right).$$

Using the method of steepest descent on \mathcal{Z} , using S as a large parameter, yields

$$(4.1.3.4) \quad \mathcal{Z} \sim \sum_{\alpha} \exp \left(-S\mathcal{S}^{\text{cl}}[\hat{\Omega}(\bar{\tau}), \bar{\beta}] \right) \mathcal{Z}'_{\alpha}, \quad \text{with the saddle point equations } \left. \frac{\delta \mathcal{S}^{\text{cl}}}{\delta \hat{\Omega}} \right|_{\hat{\Omega}=\hat{\Omega}^{\text{cl},\alpha}}.$$

These saddle point equations yield the prefactors \mathcal{Z}'_{α} , which contain subdominant corrections in powers of S^{-1} . These can be evaluated by expanding the fluctuation integrals,

$$(4.1.3.5) \quad \mathcal{Z}'_{\alpha} = \oint \mathcal{D}\delta\hat{\Omega} \exp \left[-S \left(\mathcal{S}^{\text{cl}}[\hat{\Omega}^{\text{cl}}] - \mathcal{S}^{\text{cl}}[\hat{\Omega}^{\text{cl},\alpha}] \right) \right], \text{ where } \delta\hat{\Omega} = \hat{\Omega} - \hat{\Omega}^{\text{cl},\alpha}.$$

allowing for different gauge choices for the vector potential. For further discussion on this point and the Dirac string, see [?] page 105.

4.1.3.1. Semiclassical Dynamics. The large- S expansion of the Green function requires a rescaling $t \rightarrow St = \bar{t}$, which yields

$$(4.1.3.6) \quad \mathfrak{G}(\bar{t}) = \int_{\hat{\Omega}_0}^{\hat{\Omega}_{\bar{t}}} \mathcal{D}\hat{\Omega}(\bar{t}') \exp \left(iSS^{\text{cl}}\hat{\Omega} \right), \text{ with } S^{\text{cl}} = \int_0^{\bar{t}} d\bar{t}' \left(\sum_i \mathbf{A} \cdot \dot{\hat{\Omega}} - \mathcal{H}[\hat{\Omega}(\bar{t}')] \right),$$

The Green function is dominated by time-dependent paths $\hat{\Omega}_{\text{cl}}(t')$, which extremize the action. Henceforth, let $\bar{t} \rightarrow t$ and employ the method of steepest descents, which yields a sum over saddle points

$$(4.1.3.7) \quad \mathfrak{G} \sim \sum_{\alpha} \exp \left(iSS[\hat{\Omega}^{\text{cl},\alpha}, \bar{t}] \right) \mathfrak{G}'_{\alpha}, \quad \begin{array}{l} \text{where the } \hat{\Omega}_{\text{cl}}(t')\text{-paths} \\ \text{are determined by} \\ \text{the saddle-point equations} \end{array} \quad \frac{\delta}{\delta \hat{\Omega}} \mathcal{S}[\hat{\Omega}] \Big|_{\hat{\Omega}^{\text{cl},\alpha}} = 0 \quad \begin{array}{l} \text{subject to boundary conditions} \\ \hat{\Omega}^{\text{cl},\alpha} = \hat{\Omega}_0 \\ \hat{\Omega}^{\text{cl},\alpha}(t) = \hat{\Omega}_t. \end{array}$$

The variation of the Berry phase part of the action is given by

$$(4.1.3.8) \quad \begin{aligned} \delta\omega[\hat{\Omega}] &= \int_0^t dt' \delta(\mathbf{A} \cdot \dot{\hat{\Omega}}) \\ &= \int_0^t dt' \left[\frac{\partial A^{\alpha}}{\partial \hat{\Omega}^{\beta}} \delta \hat{\Omega}^{\beta} \dot{\hat{\Omega}}^{\alpha} + A^{\alpha} \frac{d}{dt} \delta \hat{\Omega}^{\alpha} + \frac{\partial A^{\alpha}}{\partial \hat{\Omega}^{\beta}} \dot{\hat{\Omega}}^{\beta} \delta \hat{\Omega}^{\alpha} - \frac{\partial A^{\alpha}}{\partial \hat{\Omega}^{\beta}} \dot{\hat{\Omega}}^{\beta} \delta \hat{\Omega}^{\alpha} \right] \\ &= \int_0^t dt' \frac{\partial A^{\alpha}}{\partial \hat{\Omega}^{\beta}} \epsilon^{\alpha\beta\gamma} (\dot{\hat{\Omega}} \times \delta \hat{\Omega})_{\gamma} + \int_0^t dt' \frac{d}{dt'} (\mathbf{A} \cdot \delta \hat{\Omega}) \xrightarrow{0} 0 \\ &= \int_0^t dt' \hat{\Omega} \cdot (\dot{\hat{\Omega}} \times \delta \hat{\Omega}), \end{aligned}$$

where the integral over the total time derivative cancels itself out since the endpoints are fixed. Given that $\delta \hat{\Omega} \cdot \hat{\Omega} = \dot{\hat{\Omega}} \cdot \hat{\Omega} = 0$ and $\dot{\hat{\Omega}} \times \delta \hat{\Omega} \parallel \hat{\Omega}$. From these, the Euler-Lagrange equations of motion are thus obtained

$$(4.1.3.9) \quad \begin{aligned} \frac{\delta}{\delta \hat{\Omega}} \mathcal{S}[\hat{\Omega}] \Big|_{\hat{\Omega}^{\text{cl},\alpha}} &= 0 = \int_0^t \frac{\delta}{\delta \hat{\Omega}} \left[\sum_i \mathbf{A} \cdot \dot{\hat{\Omega}} - \mathcal{H}[\hat{\Omega}(\bar{t}')] \right] \\ &= \delta\omega[\hat{\Omega}] - \int_0^t dt' \frac{\delta \mathcal{H}[\hat{\Omega}(t')]}{\delta \hat{\Omega}} \\ &= \int_0^t dt' \left[\hat{\Omega} \cdot (\dot{\hat{\Omega}} \times \delta \hat{\Omega}) - \frac{\partial \mathcal{H}[\hat{\Omega}(t')]}{\partial \hat{\Omega}} \delta \hat{\Omega} \right] \\ &= \int_0^t dt' \left[\left(\hat{\Omega} \times \dot{\hat{\Omega}} - \frac{\partial \mathcal{H}[\hat{\Omega}(t')]}{\partial \hat{\Omega}} \right) \delta \hat{\Omega} \right] \Rightarrow \hat{\Omega}^{\text{cl}} \times \dot{\hat{\Omega}}^{\text{cl}} - \frac{\partial \mathcal{H}[\hat{\Omega}^{\text{cl}}]}{\partial \hat{\Omega}} = 0 \\ &\Rightarrow \dot{\hat{\Omega}}_i^{\text{cl}}(t') = \hat{\Omega}^{\text{cl}} \times \frac{\partial \mathcal{H}[\hat{\Omega}^{\text{cl}}]}{\partial \hat{\Omega}_i(t')} \Big|_{\hat{\Omega}^{\text{cl}}}. \end{aligned}$$

The last equation describes a system of classical rotators in the fast-top-limit, ie. when the rotator's internal rotational energy is much larger than the typical inter-rotator interaction energy. The right-hand side of the second row is the torque applied to rotator $\hat{\Omega}_i$, which changes its direction but not its magnitude. Using this equation of motion, it is clear that the Hamiltonian is a constant of motion on the classical path

$$(4.1.3.10) \quad \begin{aligned} \frac{d\mathcal{H}[\hat{\Omega}^{\text{cl}}(t')]}{dt} &= \sum_i \frac{\partial \mathcal{H}}{\partial \hat{\Omega}_i^{\text{cl}}(t')} \cdot \dot{\hat{\Omega}}_i^{\text{cl}}(t') \\ &= \sum_i \frac{\partial \mathcal{H}}{\partial \hat{\Omega}_i^{\text{cl}}(t')} \cdot \left[\hat{\Omega}^{\text{cl}} \times \frac{\partial \mathcal{H}[\hat{\Omega}^{\text{cl}}]}{\partial \hat{\Omega}_i(t')} \Big|_{\hat{\Omega}^{\text{cl}}} \right] = 0, \end{aligned}$$

thus ensuring conservation of energy $\mathcal{H}[\hat{\Omega}^{\text{cl}}(t')] = \mathcal{H}[\hat{\Omega}_0^{\text{cl}}]$ along the classical trajectory. A problem arises however, in that the equations of motion are first order in time, but the solution must satisfy two boundary conditions at $t' = 0$ and at $t' = t$. This is impossible for almost all boundary conditions (eg. when $\mathcal{H}[\hat{\Omega}_t] = \mathcal{H}[\hat{\Omega}_0]$). Klauder has suggested a way to overcome this problem by including second-order transistent terms of $\mathcal{O}(\epsilon)$ -order in the classical equations of motion. This allows one to solve a second-order equation, with fixed boundary conditions, and calculate the classical action and at the end, take the $(\epsilon \rightarrow 0)$ -limit.

4.1.3.2. Semiclassical spectrum. The energy dependent Green's function is defined as

$$\mathfrak{G}(\hat{\Omega}_0, \hat{\Omega}_t; E) = i \int_{\mathbb{R}_+} dt \mathfrak{G}(\hat{\Omega}_0, \hat{\Omega}_t; t) e^{i(E+i0^+)t}, \quad \text{with its spectral function } \Gamma(E) \text{ being } \Gamma(E) = \int d\hat{\Omega}_0 \mathfrak{G}(\hat{\Omega}_0, \hat{\Omega}_t; E) \\ = \sum_{\alpha} \frac{1}{E - E_{\alpha} + i0^+},$$

where the poles of $\Gamma(E)$ are the eigenenergies $\{E_{\alpha}\}_{\alpha}$ of the Hamiltonian. The path representation²¹ of the spectral function has the additional $d\hat{\Omega}_0$ - and dt -integrations. The classical paths have \square^E -duration, which is determined by the saddle-point approximation for the dt -integral,

$$(4.1.3.11) \quad \left. \frac{\partial \mathcal{S}[\hat{\Omega}]}{\partial \hat{\Omega}} \right|_{\hat{\Omega}^{\text{cl}}} \Big|_{t=t^E} + E = 0.$$

Since the Berry phase term ω is geometric and doesn't depend on t^E , it follows that

$$(4.1.3.12) \quad \mathcal{H}[\hat{\Omega}^{\text{cl}}] = E, \quad \mathcal{S}(t^E) + Et^E = \sum_i \omega[\hat{\Omega}_i].$$

Consider then a periodic orbit $\hat{\Omega}^{E,\alpha}$ of E -energy which is traversed once. The saddle point equation for the trace over $\hat{\Omega}_0$ implies that $\hat{\Omega} = \hat{\Omega}(t^E)$. Therefore, Γ is given by summing over all repetitions of this periodic orbit $\hat{\Omega}^{E,\alpha}$,

$$(4.1.3.13) \quad \Gamma \sim \sum_{\alpha} \sum_n \exp \left(inS \sum_i \omega[\hat{\Omega}^{E,\alpha}] \right) \\ = \sum_{\alpha} \frac{\exp \left(inS \sum_i \omega[\hat{\Omega}^{E,\alpha}] \right)}{1 - \exp \left(inS \sum_i \omega[\hat{\Omega}^{E,\alpha}] \right)}$$

By comparing this expression to the resolvent function's one, the semiclassical spectrum can be obtained from its poles, E_{α} , which are determined by the Bohr-Sommerfeld quantization condition

$$(4.1.3.14) \quad \sum_i \omega[\hat{\Omega}_i^{E,\alpha}] = \frac{2n\pi}{S},$$

yielding the semiclassical spectrum $\{E_{\alpha}^{\text{sc}}\}$. Since two functions are the same, upto a constant, if they share the same poles and residues, then for energies E_{α}^{sc} that are order one (ie. $n = \mathcal{O}(S)$), then the quantum spectrum may be approximated by its semiclassical counterpart,

$$(4.1.3.15) \quad E_{\alpha} = E_{\alpha}^{\text{sc}} + \mathcal{O}\left(\frac{1}{S}\right).$$

²¹This present derivation is shaky on mathematical grounds. A much more robust formalization using the path integral method was derived by Gutzwiller, and his formula is widely used in quantizing Hamiltonians with chaotic classical dynamics.

4.2. Spin Wave Theory.

4.2.1. Spin Waves: Path Integral Approachs.

Physical Context

The low-order harmonic spin wave expansion of the quantum Heisenberg model allows for a physically appealing treatment of the **long-range ordered phases**, since the quantum effects enter as real or imaginary time dependent fluctuations about the classical ground state. The spin wave modes and dispersions are obtained from the linearized classical equations of motion. To lowest order in $\frac{1}{S}$, spin waves are independent harmonic oscillators which can be semiclassically quantized. A great advantage of this approach is that it can be followed for any frustrated Heisenberg model with a complicated ground state. The path-integral, however, suffers from the ambiguity in operator ordering due to the continuous definition of the time derivative in the kinetic term.

Consider the quantum Heisenberg Hamiltonian

$$\mathbf{H} = \frac{1}{2} \sum_{i,j} J_{ij} \mathbf{S}_i \cdot \mathbf{S}_j.$$

The expectation value of this Hamiltonian in the coherent state representation,

$$(4.2.1.1) \quad \mathbf{H}[\hat{\Omega}] = \langle \hat{\Omega} | \mathbf{H} | \hat{\Omega} \rangle = \frac{S^2}{2} \sum_{ij} J_{ij} \hat{\Omega}_i \cdot \hat{\Omega}_j.$$

The classical ground states of the Hamiltonian depend on the details of the coupling constants J_{ij} . Given the model's $O(3)$ -symmetry, global spin rotations generate a continuous manifold of degenerate ground states. Certain models have additional ground state degeneracies due to frustration. Frustration occurs when no configuration can minimize all individual bond interactions simultaneously. Degeneracies that are not related to symmetries of the Hamiltonian are usually lifted by thermal and quantum fluctuations. In the literature, this is often referred to order due to disorder.

The basic assumption of spin wave theory is that, for any chosen classical ground state manifold $\hat{\Omega}^{\text{cl}}$, the partition function and Green functions can be expanded about it using the saddle-point expansion, controlled by the size of the spin S .

Consider two transverse unit vectors at each site of the lattice, $\hat{\phi}_i, \hat{\theta}_i$, which describe the azimuthal and longitudinal directions at $\hat{\Omega}^{\text{cl}}$,

$$\hat{\Omega}_i^{\text{cl}} = \hat{\phi}_i \times \hat{\theta}_i.$$

The spin fluctuations $\delta\hat{\Omega} = \hat{\Omega} - \hat{\Omega}^{\text{cl}}$ are parametrized by two sets of variables given by the projections,

$$(4.2.1.2) \quad \begin{aligned} \mathbf{q} &= \{q_i\}_{i=1}^N \equiv \{\delta\hat{\Omega}_i \cdot \hat{\phi}_i\}_{i=1}^N \\ \mathbf{p} &= \{p_i\}_{i=1}^N \equiv \{S\delta\hat{\Omega}_i \cdot \hat{\theta}_i\}_{i=1}^N. \end{aligned}$$

If the path is dominated by small fluctuations $|\delta\hat{\Omega}| < 1$, the spin-state measure may be replaced by the phase space measure,

$$(4.2.1.3) \quad \mathcal{D}\delta\hat{\Omega} \rightarrow \nu \mathcal{D}\mathbf{q} \mathcal{D}\mathbf{p}, \text{ with } \nu = \left(\frac{2S+1}{4\pi S} \right)^{N_\epsilon N},$$

where N_ϵ is the number of timesteps and thus, ν is formally infinite normalization constant. Thus, a phase-space path integral is yielded

$$(4.2.1.4) \quad \mathfrak{G} \propto e^{iS[\hat{\Omega}^{\text{cl}}]} \int \mathcal{D}\mathbf{q} \mathcal{D}\mathbf{p} \exp \left[\frac{iS}{2} \int_0^t dt' \begin{pmatrix} \mathbf{q} & \mathbf{p} \end{pmatrix} \mathcal{L}^{(2)} \begin{pmatrix} \mathbf{q} \\ \mathbf{p} \end{pmatrix} \right] [1 + \mathcal{O}(\mathbf{p}, \mathbf{q})^3],$$

where $\mathcal{L}^{(2)}$ is the spin wave lagrangian matrix, expanded about a fixed ground state configuration $\hat{\Omega}^{\text{cl}}$. The cubic and higher-order fluctuations introduce corrections, which are higher in order in S^{-1} . Since the classical ground state configuration is, by definition, stationary $\frac{d\hat{\Omega}^{\text{cl}}}{dt} = 0$, the lowest order contribution to the Berry phase is

$$(4.2.1.5) \quad \begin{aligned} S \sum_i \omega_i &\approx S \sum_i \int_0^t dt' \left(\hat{\Omega}^{\text{cl}} \cdot \delta \dot{\Omega} \times \delta \hat{\Omega} \right) \\ &= \frac{1}{2} \int_0^t dt (\mathbf{p} \cdot \dot{\mathbf{q}} - \mathbf{q} \cdot \dot{\mathbf{p}}), \end{aligned}$$

where [equation \(4.1.3.8\)](#) was used in the second equality in the first line. This implies that \mathbf{q} and \mathbf{p} play the role of coordinates and canonical momenta, respectively. The second-order expansion of the Hamiltonian yields the spin wave Hamiltonian matrix

$$(4.2.1.6) \quad \begin{aligned} \mathbf{H} - \mathbf{H}[\hat{\Omega}^{\text{cl}}] &\approx \frac{1}{2} \begin{pmatrix} \mathbf{q} & \mathbf{p} \end{pmatrix} \mathcal{H}^{(2)} \begin{pmatrix} \mathbf{q} \\ \mathbf{p} \end{pmatrix}, \\ \text{where } \mathcal{H}^{(2)} &= \begin{pmatrix} K & P \\ P^T & M^{-1} \end{pmatrix} \\ \text{with } K &= \left. \frac{\partial^2 \mathcal{H}^{(2)}}{\partial \mathbf{q} \partial \mathbf{q}} \right|_{\mathbf{q}=\mathbf{p}=0}, \quad M^{-1} = \left. \frac{\partial^2 \mathcal{H}^{(2)}}{\partial \mathbf{p} \partial \mathbf{p}} \right|_{\mathbf{q}=\mathbf{p}=0}, \quad P = \left. \frac{\partial^2 \mathcal{H}^{(2)}}{\partial \mathbf{p} \partial \mathbf{q}} \right|_{\mathbf{q}=\mathbf{p}=0}. \end{aligned}$$

Here $\mathcal{H}^{(2)}$ is a dynamical matrix of coupled harmonic oscillators, where K and M are the mass and force constant matrices and where P couples coordinates and momenta. Combining the previous expressions, the small coefficient contribution to the action is yielded,

$$(4.2.1.7) \quad \mathcal{S}^{(2)} \approx \mathcal{S}_0 + \int_0^t dt' \begin{pmatrix} \mathbf{q} & \mathbf{p} \end{pmatrix} \mathcal{L}^{(2)} \begin{pmatrix} \mathbf{q} \\ \mathbf{p} \end{pmatrix}, \text{ where } \mathcal{L}^{(2)} = - \begin{pmatrix} K & \partial_t + P \\ -\partial_t + P^T & M^{-1} \end{pmatrix}.$$

If $\hat{\Omega}^{\text{cl}}$, and thus $\mathcal{L}^{(2)}$, are periodic on the lattice, the spin-wave can be labelled by \mathbf{k} -momentum and band index α . The spin-wave modes $(q_{\mathbf{k},\alpha}, p_{\mathbf{k},\alpha})$ and spin wave frequencies can be obtained by solving the classical equations of motion for small oscillations, namely

$$(4.2.1.8) \quad \frac{\delta \mathcal{S}^{(2)}(\mathbf{p}, \mathbf{q})}{\delta \hat{\Omega}} = \mathcal{L}^{(2)} \begin{pmatrix} q_{\mathbf{k},\alpha} \\ p_{\mathbf{k},\alpha} \end{pmatrix} e^{-i\omega_{\mathbf{k},\alpha} t}, \quad \det \begin{pmatrix} K & i\omega + P \\ -i\omega + P^T & M^{-1} \end{pmatrix} \Big|_{\omega=\omega_{\mathbf{k},\alpha}}.$$

The harmonic spin waves are non-interacting bosons. Their eigenmodes and energies allows for the calculation of all thermodynamic averages and dynamical response functions. In particular, the spin-wave correction to the free energy is given by

$$(4.2.1.9) \quad F^{(2)} = -\frac{1}{\beta} \log \det \mathcal{L}^{(2)} = T \sum_{\mathbf{k},\alpha} \log \sinh \frac{\omega_{\mathbf{k},\alpha}}{2T}.$$

The spin-wave variables can be mapped to Bose coherent $(\mathbf{z}, \mathbf{z}^*)$ variables, which transform the phase-space path integral [equation \(4.2.1.4\)](#) into a bosonic coherent state path integral, as follows

$$(4.2.1.10) \quad \begin{aligned} \begin{pmatrix} \mathbf{z}_i(\tau) \\ \mathbf{z}_i^*(\tau) \end{pmatrix} &= \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & i \\ 1 & -i \end{pmatrix} \begin{pmatrix} q_i(\tau) \\ p_i(\tau) \end{pmatrix}, \quad \mathcal{D}\mathbf{p} \mathcal{D}\mathbf{q} = \mathcal{D}^2 \mathbf{z}, \\ \text{The kinetic term is } &\frac{1}{2} \int_0^t dt (\mathbf{p} \cdot \dot{\mathbf{q}} - \mathbf{q} \cdot \dot{\mathbf{p}}) = \frac{1}{2} \int_0^t dt (\mathbf{z}^* \dot{\mathbf{z}} - \dot{\mathbf{z}} \mathbf{z}^*). \end{aligned}$$

Then, the Green function in bosonic coherent representation is given by

(4.2.1.11)

$$\mathfrak{G} = \int \mathcal{D}^2 \mathbf{z} \exp \left[\frac{i}{2} \int_0^t dt' (\mathbf{z}^* \dot{\mathbf{z}} - \mathbf{z} \dot{\mathbf{z}}^*) - \mathbf{H}[\mathbf{z}^*, \mathbf{z}] \right], \quad \text{where the bosonic Hamiltonian is } \mathbf{H}[\mathbf{z}^*, \mathbf{z}] \approx \frac{1}{2} \begin{pmatrix} \mathbf{z} & \mathbf{z}^* \end{pmatrix} \tilde{\mathcal{H}}^{(2)} \begin{pmatrix} \mathbf{z} \\ \mathbf{z}^* \end{pmatrix}$$

$$\tilde{\mathcal{H}}^{(2)} = \begin{pmatrix} \frac{1}{2}(K + M^{-1}) & \frac{1}{2}(K + M^{-1}) + iP \\ \frac{1}{2}(K - M^{-1}) - iP & \frac{1}{2}(K + M^{-1}). \end{pmatrix}$$

This Hamiltonian contains both normal (ie. $z^* z$) terms as well as anomalous (eg. $z^* z^*$ and $z z$) terms. Using the real-to-imaginary-time transformation, $t' \rightarrow -i\tau$, $t \rightarrow -i\beta$, the path-integral representation for the Green function is transformed into a partition function, as was anticipated in [equation \(4.1.3.1\)](#). The spin-wave contribution to the partition function, defined in [equation \(4.1.1.5\)](#), is

$$\mathcal{Z}' \approx \oint \mathcal{D}^2 \mathbf{z}_\omega \exp \left\{ -\beta \sum_{\omega_n} \left[i\omega_n \mathbf{z}^* \mathbf{z} - \frac{1}{2} \begin{pmatrix} \mathbf{z} & \mathbf{z}^* \end{pmatrix} \tilde{\mathcal{H}}^{(2)} \begin{pmatrix} \mathbf{z} \\ \mathbf{z}^* \end{pmatrix} \right] \right\}$$

$$= \prod_n \det \left[\begin{pmatrix} i\omega_n & 0 \\ 0 & -i\omega_n \end{pmatrix} - \tilde{\mathcal{H}}^{(2)} \right]^{-1}, \quad \omega_n = \frac{2n\pi}{\beta}.$$

Wherein, ω_n are the Bose-Matsubara frequencies. A second-quantized bosonic Hamiltonian can be obtained from [equation \(4.2.1.11\)](#), yielding

$$\mathcal{H}^{(2)} = \frac{1}{2} \begin{pmatrix} \mathbf{b}^\dagger & \mathbf{b} \end{pmatrix} \mathbf{H} \begin{pmatrix} \frac{1}{2}(K + M^{-1}) & \frac{1}{2}(K + M^{-1}) + iP \\ \frac{1}{2}(K - M^{-1}) - iP & \frac{1}{2}(K + M^{-1}). \end{pmatrix} \begin{pmatrix} \mathbf{b}^\dagger \\ \mathbf{b} \end{pmatrix} + E'_0,$$

which contains both normal and anomalous terms, with E'_0 an unknown constat. The problem with [equation \(4.2.1.12\)](#) is that there is no prescription for the ordering of the $(\mathbf{b}, \mathbf{b}^\dagger)$ -operators in the $\mathcal{H}^{(2)}$ -quadratic form, hence the ambiguity of the constant E'_0 .

The techniques developed in [equation \(4.2.1.8\)](#)-[equation \(4.2.1.13\)](#) are of general use, since these determine the spin wave modes of any Heisenberg model once the classical ground state $\hat{\Omega}^{\text{cl}}$ has been specified. In the following sections, the simplest models of nearest-neighbour interaction Heisenberg ferromagnets and antiferromagnets are considered for $d = 1, 2, 3$ -dimensional lattices are to be considered:

$$\mathbf{H} = \pm |J| S^2 \sum_{\langle ij \rangle} \hat{\Omega}_i \cdot \hat{\Omega}_j$$

$$= JS^2 \sum_{\langle ij \rangle} \left[\cos \theta_i \cos \theta_j + \sin \theta_i \sin \theta_j \cos(\phi_i - \phi_j) \right].$$

The spin-wave expansions of these Hamiltonians are simple since their classical ground states are either uniform ferromagnets or Néel antiferromagnets.

4.2.1.1. The Ferromagnet. Assuming a spin-wave expansion of the path-integral is valid for the ordered classical ground state of the Heisenberg ferromagnet, the spin can be polarized (either with a small magnetic field or with the appropriate boundary conditions) to point in the x -direction, ie.

$$\hat{\Omega}^{\text{cl}} = (\theta^{\text{cl}}, \phi^{\text{cl}}) = \left(\frac{\pi}{2}, 0 \right), \quad \text{with the fluctuations parameterized by } \begin{matrix} q_i = \phi_i \\ p_i = S \cos \theta_i. \end{matrix}$$

The ferromagnetic Hamiltonian, with $-|J|$, written in [equation \(4.2.1.14\)](#), can be expanded for small θ, ϕ , yielding a second order quadratic form on the (q, p) -coordinates as follows

(4.2.1.16)

$$\mathbf{H} \approx -\frac{Nz}{2}|J|S^2 + \frac{1}{2}|J|S^2 \sum_{\langle ij \rangle} \left[\frac{(p_i - p_j)^2}{S^2} + (q_i - q_j)^2 \right], \text{ where}$$

this expression is valid for $|\cos \theta| \ll 1$
 $|\phi_i| \ll \pi$,
 N is the number of sites and where
 $z = 2d$ is the coordination number in d -dimensions.

From [equation \(4.2.1.6\)](#)'s nested relations, the K , M^{-1} , P submatrices can be easily found. In particular for a one-dimensional system, it yields

(4.2.1.17)

$$\begin{aligned} K_{ij} &= \frac{\partial^2 \mathcal{H}^{(2)}}{\partial q_i \partial q_{i+1}} \Big|_{\mathbf{q}=\mathbf{p}=0} = \frac{1}{2} |J| S^2 \frac{\partial^2 \mathcal{H}^{(2)}}{\partial q_i \partial q_j} (q_i - q_{i+1})^2 \Big|_{\mathbf{q}=\mathbf{p}=0} \\ &= \frac{1}{2} |J| S^2 \frac{\partial}{\partial \mathbf{q}_j} \left(2q_i \delta_{ij} - 2(q_i \delta_{j,i+1} + q_{i+1} \delta_{ij}) + 2q_{i+1} \delta_{j,i+1} \right) \Big|_{\mathbf{q}=\mathbf{p}=0} \\ &= \frac{1}{2} |J| S^2 \left(2q \delta_{ij} - 2(\delta_{j,i+1} + \cancel{\delta_{i,i+1} \delta_{ij}}) + 2\cancel{\delta_{i,i+1} \delta_{j,i+1}} \right) = -z |J| S^2 \nabla_{i,i}^2, \quad \text{where } \nabla_{i,i}^2 = \frac{1}{z} (\delta_{i,i+1} - \delta_{ij}) \end{aligned}$$

$$\begin{aligned} M_{ij}^{-1} &= \frac{\partial^2 \mathcal{H}^{(2)}}{\partial p_i \partial p_{i+1}} \Big|_{\mathbf{q}=\mathbf{p}=0} = \frac{1}{2} |J| \frac{\partial^2 \mathcal{H}^{(2)}}{\partial p_i \partial p_j} (p_i - p_{i+1})^2 \Big|_{\mathbf{q}=\mathbf{p}=0} \\ &= \frac{1}{2} |J| \frac{\partial}{\partial \mathbf{q}_j} \left(2q_i \delta_{ij} - 2(p_i \delta_{j,i+1} + p_{i+1} \delta_{ij}) + 2p_{i+1} \delta_{j,i+1} \right) \Big|_{\mathbf{q}=\mathbf{p}=0} \\ &= \frac{1}{2} |J| S^2 \left(2q \delta_{ij} - 2(\delta_{j,i+1} + \cancel{\delta_{i,i+1} \delta_{ij}}) + 2\cancel{\delta_{i,i+1} \delta_{j,i+1}} \right) = -z |J| \nabla_{i,i}^2, \quad \text{where } \nabla_{i,i}^2 = \frac{1}{z} (\delta_{i,i+1} - \delta_{ij}) \end{aligned}$$

$$P = 0.$$

These relations hold only for the one-dimensional lattice, but they are easily generalized to $d > 1$ -dimensional lattices, yielding

$$(4.2.1.18) \quad \begin{aligned} K_{ij} &= -z |J| S^2 \nabla_{ij}^2, \\ M_{ij}^{-1} &= -z |J| \nabla_{ij}^2, \quad \nabla_{ij}^2 = \frac{1}{z} \sum_{\eta} (\delta_{i+\eta,j} - \delta_{ij}), \\ P &= 0 \end{aligned}$$

where η are the nearest-neighbour vectors and where ∇_{ij}^2 is the lattice laplacian, its eigenvalues given by

$$(4.2.1.19) \quad \begin{aligned} \nabla_{\mathbf{k}}^2 &= \frac{1}{N} \sum_{ij} e^{-i(\mathbf{k}(\mathbf{x}_i - \mathbf{x}_j))} \nabla_{ij}^2 \\ &= \frac{1}{z} \sum_{\eta} (e^{i\mathbf{k}\eta} - 1) \equiv \gamma_{\mathbf{k}} - 1, \end{aligned}$$

which defines a tight-binding function $\gamma_{\mathbf{k}}$. Using [equation \(4.2.1.18\)](#)'s relations in [equation \(4.2.1.8\)](#) yields the following matrix

$$(4.2.1.20) \quad \det \begin{pmatrix} -z |J| S^2 (\gamma_{\mathbf{k}} - 1) & i\omega_{\mathbf{k}} \\ -i\omega_{\mathbf{k}} & -z |J| (\gamma_{\mathbf{k}} - 1) \end{pmatrix} = 0 \Rightarrow \omega_{\mathbf{k}} \equiv z |J| S (1 - \gamma_{\mathbf{k}}).$$

Thus, the spin wave is described by a time dependent precession of all spins about the x -direction with $\omega_{\mathbf{k}}$ -angular frequency.

4.2.1.2. The Antiferromagnet. In the antiferromagnetic case, the path integral can be expanded about the classical Néel state, which can be made to point in the $\pm x$ -direction for the A, B sublattices respectively, without loss of generality, with coordinates

$$(4.2.1.21) \quad (\theta_i^{\text{cl}}, \phi_i^{\text{cl}}) = \begin{cases} \left(\frac{\pi}{2}, 0\right) & i \in A \\ \left(\frac{\pi}{2}, -\pi\right) & i \in B \end{cases}$$

wherein the singularity of the vector potential is assumed to be far from both the $+x$ - and $-x$ -directions, eg. only at the south pole. The small fluctuations are then

$$(4.2.1.22) \quad q_i = \begin{cases} \phi_i & i \in A \\ \phi_i + \pi & i \in B \end{cases}, \quad p_i = S \cos \theta_i.$$

The second order expansion of the antiferromagnetic ($+|J|$) Hamiltonian [equation \(4.2.1.14\)](#) is given by

$$(4.2.1.23) \quad \mathbf{H} \approx -\frac{Nz}{2}|J|S^2 + \frac{1}{2}|J|S^2 \sum_{\langle ij \rangle} \left[\frac{(p_i + p_j)^2}{S^2} + (q_i - q_j)^2 \right] + \mathcal{O}(p^3, q^3).$$

The mass and force constants matrices, which will be simultaneously diagonalized in Fourier space and allow for the determination of the spin wave spectrum of the spin-wave spectrum, are given by

$$(4.2.1.24) \quad \begin{aligned} M_{ij}^{-1} &= zJ(\nabla_{ij}^2 + 2\delta_{ij}) \quad , K_{ij} = -zJS^2\nabla_{ij}^2 \quad , P = 0. \\ M_{\mathbf{k}}^{-1} &= zJ(1 + \gamma_{\mathbf{k}}), \quad K_{\mathbf{k}} = zJS^2(1 - \gamma_{\mathbf{k}}). \end{aligned}$$

$$\Rightarrow \det \begin{pmatrix} -z|J|S^2(1 - \gamma_{\mathbf{k}}) & i\omega_{\mathbf{k}} \\ -i\omega_{\mathbf{k}} & z|J|(1 + \gamma_{\mathbf{k}}) \end{pmatrix} = 0 \rightarrow \omega_{\mathbf{k}} = zJS\sqrt{1 - \gamma_{\mathbf{k}}^2} \sim c|\mathbf{k} - \mathbf{k}_c|,$$

where $\mathbf{k}_c = 0, \pi\mathbb{1}_N$. The spin-wave velocity for a cubic lattice in d dimensions is $c = \sqrt{d}JS$. The antiferromagnetic spin waves are unit vectors which precess about the classical directions $\pm x$. Furthermore, the spin precess in opposite directions on the two sublattices. Unlike ferromagnetic spin waves, given by [equation \(4.2.1.20\)](#), there are two degenerate antiferromagnetic spin wave modes whose frequencies vanish as $\mathbf{k} \rightarrow 0$ and $\mathbf{k} \rightarrow \pi\mathbb{1}_N$.

4.2.2. *Spin Waves: Holstein-Primakoff Approach.* Low-order spin waves can be written as non-interacting bosons. For a quantitative quantum theory, the correct order of bosonic operators in the Hamiltonian is needed. To that end, generalized Holstein-Primakoff bosons can be used to represent the quantum model. Consider an arbitrary classical configuration $\hat{\Omega}^{\text{cl}}$ which minimizes $\mathbf{H}[\hat{\Omega}]$. Using $\hat{\Omega}^{\text{cl}}$, three basis vectors can be defined at every site, $(\mathbf{e}^1, \mathbf{e}^2, \hat{\Omega}^{\text{cl}})$, such that

$$\mathbf{e}^1 \times \mathbf{e}^2 = \hat{\Omega}_i^{\text{cl}}.$$

In this coordinate frame, the raising and lowering operators are defined

$$(4.2.2.1) \quad \mathbf{S}^{\pm} = \mathbf{S} \cdot \mathbf{e}^1 \pm i\mathbf{S} \cdot \mathbf{e}^2.$$

With respect to this triad of basis vectors, the spin components can be represented by HP-bosons

$$(4.2.2.2) \quad \mathbf{S}^+ = \sqrt{2S - \mathbf{n}_b} \mathbf{b}, \quad \mathbf{S}^- = \mathbf{b}^\dagger \sqrt{2S - \mathbf{n}_b}, \quad \mathbf{S} \cdot \hat{\Omega}^{\text{cl}} = -\mathbf{n}_b + S.$$

This representation is exact in the Hilbert subspace of $\mathbf{n}_b \leq S$. However, the square-root function represents an infinite power series of number operators multiplied by $\frac{1}{S}$ -factors. The truncation of the series to low order can be justified if one can show, a posteriori, that $\langle \mathbf{n}_b \rangle \ll 2S$, ie. that the spin fluctuations about the classical

directions are small. Next, the spin operator in the Hamiltonian are substituted by the spin operators written in [equation \(4.2.2.2\)](#). Terms upto quadratic order in the bosons constitute the spin wave Hamiltonian.

4.2.2.1. The Ferromagnet. Consider a Heisenberg ferromagnet whose classical direction, without loss of generality, is chosen to be the z -direction and substitute the HP-operators into the ferromagnetic Hamiltonian,

(4.2.2.3)

$$\begin{aligned} \mathbf{H} &= -|J| \sum_{\langle ij \rangle} \mathbf{S}_i \cdot \mathbf{S}_j \\ &= -S^2 |J| N \frac{z}{2} - |J| \sum_{\langle ij \rangle} \left[S \mathbf{b}_i^\dagger \sqrt{1 - \frac{\mathbf{n}_i}{2S}} \sqrt{1 - \frac{\mathbf{n}_j}{2S}} \mathbf{b}_j - \frac{1}{2} S (\mathbf{n}_i + \mathbf{n}_j) + \frac{1}{2} \mathbf{n}_i \cdot \mathbf{n}_j \right]. \\ &= -S^2 |J| N \frac{z}{2} + \mathbf{H}_1 + \mathbf{H}_2 + \mathcal{O}\left(\frac{1}{S}\right), \text{ where} \end{aligned}$$

$$\begin{aligned} \mathbf{H}_1 &= \sum_{\mathbf{k}} \omega_{\mathbf{k}} \mathbf{b}_{\mathbf{k}}^\dagger \mathbf{b}_{\mathbf{k}} \text{ where } \mathbf{b}_{\mathbf{k}} = \frac{1}{\sqrt{N}} \sum_i e^{-i\mathbf{k} \cdot \mathbf{x}_i} \mathbf{b}_i \\ \mathbf{H}_2 &= \frac{|J|}{4} \sum_{\langle ij \rangle} \left[\mathbf{b}_i \mathbf{b}_j (\mathbf{b}_i - \mathbf{b}_j)^2 + (\mathbf{b}_i^\dagger - \mathbf{b}_j^\dagger)^2 \mathbf{b}_i \mathbf{b}_j \right] \end{aligned}$$

with

$$\omega_{\mathbf{k}} = S|J|z \left(1 - \frac{1}{z} \sum_{j, \langle ij \rangle} e^{i(\mathbf{x}_j - \mathbf{x}_i) \cdot \mathbf{k}} \right) = S|J|z(1 - \gamma_{\mathbf{k}}),$$

which agrees with [equation \(4.2.1.20\)](#). In the previous expression for the Hamiltonian, [equation \(4.2.2.3\)](#), \mathbf{H}_1 describes non-interacting spin waves with dispersion $\omega_{\mathbf{k}}$, \mathbf{H}_2 and the higher-order terms describe interactions between spin waves, which can be treated by perturbation theory or by mean field approximations. At $T = 0$, it is clear from [equation \(4.2.2.3\)](#) that the energy is equal to the classical energy

$$E_0 = -S^2 |J| N \frac{z}{2}.$$

This is in agreement with [corollary 2](#), which states that the ground state of the quantum ferromagnet is the classical ground state.

The long-wavelength limit of the ferromagnetic spin wave dispersion vanishes as

$$(4.2.2.4) \quad \omega_{\mathbf{k}} \sim S|J||\mathbf{k}|^2.$$

This is the **gapless Goldstone mode**, which is a consequence of the broken symmetry of the ferromagnetic ground state, as stated in [theorem 6](#). The Goldstone mode dominates the low-temperature and long-wavelength correlations of the ferromagnet. The lowest-order corrections to the ground state magnetization at finite temperatures is given by

$$\begin{aligned} (4.2.2.5) \quad \Delta m_0 &= \frac{1}{N} \langle \mathbf{S}_{\text{tot}}^z - S \rangle \\ &= -\langle \mathbf{n}_i \rangle = -\frac{1}{N} \sum_{\mathbf{k}} \mathbf{n}_{\mathbf{k}}, \text{ where the Bose-Einstein occupation number is } \mathbf{n}_{\mathbf{k}} = \frac{1}{e^{\beta \omega_{\mathbf{k}}} - 1}. \end{aligned}$$

The asymptotic low- T behaviour of the sum given in [equation \(4.2.2.5\)](#) can be found by introducing a small infrared cutoff k_0 . Consider an additional, small but, finite momentum $\bar{k} > k_0$ as well, which is well inside the region where [equation \(4.2.2.4\)](#) holds, ie.

$$\omega_{\bar{k}} \ll T \ll |J|S.$$

The sum in [equation \(4.2.2.5\)](#) can be broken into two regions, one with $k_0 < |\mathbf{k}| < \bar{k}$ and $|\mathbf{k}| \geq \bar{k}$ and find that

$$(4.2.2.6) \quad \Delta m_0 \approx - \int_{k_0}^{\bar{k}} \frac{dk k^{d-1}}{(2\pi)^d} \frac{T}{JSk^2} - \frac{1}{N} \sum_{\mathbf{k} > \bar{k}} \frac{1}{e^{\beta\omega_{\mathbf{k}}} - 1}.$$

For $T > 0$ and $d = 1, 2$, the first integral diverges at low k_0 as

$$(4.2.2.7) \quad \Delta m_0 \approx \begin{cases} -\frac{t}{k_0} + \dots & d = 1 \\ t \log k_0 + \dots & d = 2 \end{cases}, \quad t = \frac{T}{JS},$$

but with the second sum being finite, having no infrared singularity. The most important conclusion emerging from the previous equation's result is that the magnetization correction due to fluctuations diverges as the infrared cutoff k_0 vanishes in one and two dimensions. Therefore, for the $d = 1, 2$ -dimensional cases, the basic assumption that the spin fluctuations $\langle (\mathbf{S}_i^z - S)^2 \rangle$ are small, turns out to be wrong. Thus, the truncation of the expansion of the HP-operators at low orders is unjustified. The breakdown of spin wave theory is consistent with Mermin-Wagner's theorem, [theorem 5](#), which rules out a finite spontaneous magnetization in $(d = 1, 2)$ -dimensions at all non-zero temperatures.

In $d = 3$ dimensions, there are no infrared divergences. The leading temperature dependent of M can be calculated by writing the Bose function as a geometric sum and integrating over momenta up to infinity as follows,

$$(4.2.2.8) \quad \begin{aligned} \Delta m_0^{d=3} &= - \int_{k_0}^{\bar{k}} \frac{dk k^2}{2\pi^2} \sum_{n=1}^{\infty} e^{-\frac{nk^2}{t}} [1 + \mathcal{O}(t)] \\ &= -\frac{1}{8} \left(\frac{t}{\pi}\right)^{\frac{3}{2}} \sum_{n=1}^{\infty} n^{-\frac{3}{2}} = -\frac{1}{8} \left(\frac{t}{\pi}\right)^{\frac{3}{2}} \zeta\left(\frac{3}{2}\right), \end{aligned}$$

where ζ is the Riemann zeta function. Thus, in three-dimensions the leading temperature correction to the ordered moment is proportional to $-T^{\frac{3}{2}}$.

4.2.2.2. The Antiferromagnet. Consider the classical Néel state, which is chosen to be in the z - and $-z$ -directions on the A and B sublattices, respectively. The rotated spin $\tilde{\mathbf{S}}$ are defined as

$$(4.2.2.9) \quad j \in B, \quad \begin{cases} \tilde{\mathbf{S}}_j^z = -\mathbf{S}_j^z \\ \tilde{\mathbf{S}}_j^x = \mathbf{S}_j^x \\ \tilde{\mathbf{S}}_j^y = -\mathbf{S}_j^y \end{cases} \Rightarrow \begin{cases} \text{Sublattice A : } \begin{aligned} \mathbf{S}_i^z &= S - \mathbf{b}_i^\dagger \mathbf{b}_i \\ \mathbf{S}_i^+ &= \sqrt{2S - \mathbf{b}_i^\dagger \mathbf{b}_i} \mathbf{b}_i \\ \mathbf{S}_i^- &= \mathbf{b}_i^\dagger \sqrt{2S - \mathbf{b}_i^\dagger \mathbf{b}_i} \end{aligned} \\ \text{Sublattice B : } \begin{aligned} \mathbf{S}_i^z &= \mathbf{b}_i^\dagger \mathbf{b}_i - S \\ \mathbf{S}_i^+ &= \mathbf{b}_i^\dagger \sqrt{2S - \mathbf{b}_i^\dagger \mathbf{b}_i} \\ \mathbf{S}_i^- &= \sqrt{2S - \mathbf{b}_i^\dagger \mathbf{b}_i} \mathbf{b}_i \end{aligned} \end{cases}$$

It is clear that the $\tilde{\mathbf{S}}^\alpha$ obey the same commutation relations as \mathbf{S}^α and therefore can be represented by Holstein-Primakoff bosons. Using the sublattice-rotated representation, the antiferromagnetic Hamiltonian reads

$$(4.2.2.10) \quad \begin{aligned} \mathbf{H} &= + |J| \sum_{\langle ij \rangle} \mathbf{S}_i^z \cdot \tilde{\mathbf{S}}_j^z + \frac{|J|}{2} \sum_{\langle ij \rangle} (\mathbf{S}_i^+ \tilde{\mathbf{S}}_j^+ + \mathbf{S}_i^- \tilde{\mathbf{S}}_j^-) \\ &= + |J| \sum_{\langle ij \rangle} \mathbf{S}_i^z \cdot \tilde{\mathbf{S}}_j^z + \frac{|J|}{2} \sum_{\langle ij \rangle} (\mathbf{S}_i^+ \mathbf{S}_j^- + \mathbf{S}_i^- \mathbf{S}_j^+) \end{aligned}$$

Using the Holstein-Primakoff representation given by [equation \(4.2.2.2\)](#) for \mathbf{S} on the A -sublattice and $\tilde{\mathbf{S}}$ on the B -sublattice, one obtains

(4.2.2.11)

$$\begin{aligned}
\mathbf{H} &= +|J| \sum_{\langle ij \rangle} \left[(S - \mathbf{b}_i^\dagger \mathbf{b}_i)(\mathbf{b}_j^\dagger \mathbf{b}_j - S) \right] \\
&\quad + \frac{|J|}{2} \sum_{\langle ij \rangle} \left[(\sqrt{2S - \mathbf{b}_i^\dagger \mathbf{b}_i} \mathbf{b}_i)(\sqrt{2S - \mathbf{b}_j^\dagger \mathbf{b}_j} \mathbf{b}_j) + (\mathbf{b}_i^\dagger \sqrt{2S - \mathbf{b}_i^\dagger \mathbf{b}_i})(\mathbf{b}_j^\dagger \sqrt{2S - \mathbf{b}_j^\dagger \mathbf{b}_j}) \right] \\
&= +|J| \sum_{\langle ij \rangle} \left[S \mathbf{b}_j^\dagger \mathbf{b}_j - S^2 + S \mathbf{b}_i^\dagger \mathbf{b}_i + \mathbf{b}_i^\dagger \mathbf{b}_i \mathbf{b}_j^\dagger \mathbf{b}_j \right] \\
&\quad + S|J| \sum_{\langle ij \rangle} \left[\left(1 - \frac{\mathbf{n}_i}{4S}\right) \mathbf{b}_i \left(1 - \frac{\mathbf{n}_j}{4S}\right) \mathbf{b}_j + \mathbf{b}_i^\dagger \left(1 - \frac{\mathbf{n}_i}{4S}\right) \mathbf{b}_j^\dagger \left(1 - \frac{\mathbf{n}_j}{4S}\right) \right] \\
&= +|J| \sum_{\langle ij \rangle} \left[S \mathbf{n}_j + S \mathbf{n}_i - S^2 + \mathbf{n}_i \mathbf{n}_j \right] \\
&\quad + S|J| \sum_{\langle ij \rangle} \left[\mathbf{b}_i \mathbf{b}_j - \mathbf{b}_i \frac{\mathbf{n}_j}{4S} \mathbf{b}_j - \frac{\mathbf{n}_i}{4S} \mathbf{b}_i \mathbf{b}_j + \frac{\mathbf{n}_i}{4S} \mathbf{b}_i \frac{\mathbf{n}_j}{4S} \mathbf{b}_j + \mathbf{b}_i^\dagger \mathbf{b}_j^\dagger - \mathbf{b}_i^\dagger \mathbf{b}_j^\dagger \frac{\mathbf{n}_j}{4S} - \mathbf{b}_i^\dagger \frac{\mathbf{n}_j}{4S} \mathbf{b}_j^\dagger + \mathbf{b}_i^\dagger \frac{\mathbf{n}_i}{4S} \mathbf{b}_j^\dagger \frac{\mathbf{n}_j}{4S} \right] \\
&= -\frac{S^2|J|Nz}{2} + JSz \sum_{\langle ij \rangle} \left[\mathbf{n}_i + \frac{1}{2} (\mathbf{b}_i \mathbf{b}_i + \mathbf{b}_i^\dagger \mathbf{b}_i^\dagger) \right] \\
&\quad + \frac{-|J|}{4} \sum_{\langle ij \rangle} \left[\mathbf{b}_i \mathbf{b}_j^\dagger \mathbf{b}_j \mathbf{b}_j - \mathbf{b}_i^\dagger \mathbf{b}_i \mathbf{b}_i \mathbf{b}_j + \mathbf{b}_i^\dagger \mathbf{b}_j^\dagger \mathbf{b}_j^\dagger \mathbf{b}_j - \mathbf{b}_i^\dagger \mathbf{b}_j^\dagger \mathbf{b}_j \mathbf{b}_j^\dagger - \frac{\mathbf{n}_i}{4S} \mathbf{b}_i \frac{\mathbf{n}_j}{4S} \mathbf{b}_j - \mathbf{b}_i^\dagger \frac{\mathbf{n}_i}{4S} \mathbf{b}_j^\dagger \frac{\mathbf{n}_j}{4S} \right] \mathcal{O}(S^{-2})
\end{aligned}$$

the first term is the classical ground-state energy of the antiferromagnet and contains the antiferromagnetic magnon dispersion relation,

$$(4.2.2.12) \quad \mathbf{H}_1 = JSz \sum_{\langle ij \rangle} \left[\mathbf{n}_i + \frac{1}{2} (\mathbf{b}_i \mathbf{b}_i + \mathbf{b}_i^\dagger \mathbf{b}_i^\dagger) \right]$$

The sum over nearest neighbours can be rewritten as a sum over the displacement vector δ_i between neighbouring spin, such that the $\mathbf{b}_{i+\delta_i}$ annihilates a boson at the B -site with position $\mathbf{x}_i + \delta_i$, thus yielding the following Hamiltonian in momentum space:

(4.2.2.13)

$$\begin{aligned}
\mathbf{H}_1 &= JSz \frac{1}{N/2} \frac{1}{2} \sum_{i \in A} \sum_{\delta, \mathbf{k}, \mathbf{k}'} \left[\mathbf{n}_i + \left(e^{i\mathbf{k} \cdot \mathbf{x}_i} e^{i\mathbf{k}' \cdot (\mathbf{x}_i + \delta_i)} \mathbf{b}_{\mathbf{k}} \mathbf{b}_{\mathbf{k}'} + e^{-i\mathbf{k} \cdot \mathbf{x}_i} e^{-i\mathbf{k}' \cdot (\mathbf{x}_i + \delta_i)} \mathbf{b}_{\mathbf{k}}^\dagger \mathbf{b}_{\mathbf{k}'}^\dagger \right) \right] \\
&= JSz \frac{1}{2} \sum_i \sum_{\delta, \mathbf{k}, \mathbf{k}'} \left[\mathbf{n}_i + \left(\sum_{i \in A} e^{i(\mathbf{k} + \mathbf{k}') \cdot \mathbf{x}_i} e^{i\mathbf{k}' \cdot \delta_i} \mathbf{b}_{\mathbf{k}} \mathbf{b}_{\mathbf{k}'} + \sum_{i \in A} e^{-i(\mathbf{k} + \mathbf{k}') \cdot \mathbf{x}_i} e^{-i\mathbf{k}' \cdot \delta_i} \mathbf{b}_{\mathbf{k}}^\dagger \mathbf{b}_{\mathbf{k}'}^\dagger \right) \right] \\
&= JSz \sum_{\mathbf{k}} \left[\mathbf{b}_{\mathbf{k}}^\dagger \mathbf{b}_{\mathbf{k}} + \frac{\gamma_{\mathbf{k}}}{2} \left(e^{-i\mathbf{k}' \cdot \delta_i} \mathbf{b}_{\mathbf{k}}^\dagger \mathbf{b}_{-\mathbf{k}}^\dagger + e^{i\mathbf{k}' \cdot \delta_i} \mathbf{b}_{\mathbf{k}} \mathbf{b}_{-\mathbf{k}} \right) \right], \quad \text{with } \gamma_{\mathbf{k}} = \frac{1}{z} \sum_{\delta} e^{i\mathbf{k} \cdot \delta}. \\
\Rightarrow \mathbf{H} &= -\frac{S^2|J|Nz}{2} + \mathbf{H}_1 + \mathbf{H}_2 + \mathcal{O}\left(\frac{1}{S}\right), \quad \begin{aligned} \mathbf{H}_1 &= JSz \sum_{\mathbf{k}} \left[\mathbf{b}_{\mathbf{k}}^\dagger \mathbf{b}_{\mathbf{k}} + \frac{\gamma_{\mathbf{k}}}{2} (\mathbf{b}_{\mathbf{k}}^\dagger \mathbf{b}_{-\mathbf{k}}^\dagger + \mathbf{b}_{\mathbf{k}} \mathbf{b}_{-\mathbf{k}}) \right] \\ \mathbf{H}_2 &= -\frac{J}{4} \sum_{\langle ij \rangle} \left[\mathbf{b}_i^\dagger (\mathbf{b}_j^\dagger + \mathbf{b}_i)^2 \mathbf{b}_j + \mathbf{b}_i (\mathbf{b}_j^\dagger + \mathbf{b}_i)^2 \mathbf{b}_j^\dagger \right]. \end{aligned}
\end{aligned}$$

where $N/2$ is the number of sublattice A and sublattice B sites, and with $\sum_i e^{i(\mathbf{k} + \mathbf{k}') \cdot \mathbf{x}_i} = \frac{N}{2} \delta_{\mathbf{k}', -\mathbf{k}}$. Now \mathbf{H}_1 is a quadratic Hamiltonian which includes normal and anomalous terms, but can be diagonalized via a Bogoliubov transformation, as follows

(4.2.2.14)

$$\begin{pmatrix} \alpha_{\mathbf{k}} \\ \beta_{\mathbf{k}}^\dagger \end{pmatrix} = \begin{pmatrix} u_{\mathbf{k}} & -v_{\mathbf{k}} \\ -v_{\mathbf{k}} & u_{\mathbf{k}} \end{pmatrix} \begin{pmatrix} \mathbf{b}_{\mathbf{k}} \\ \mathbf{b}_{-\mathbf{k}}^\dagger \end{pmatrix}, \quad u_{\mathbf{k}}^2 - v_{\mathbf{k}}^2 = 1, \quad \text{which preserves the commutation relations}$$

$$\begin{aligned} [\alpha_{\mathbf{k}}, \alpha_{\mathbf{k}'}^\dagger] &= [\beta_{\mathbf{k}}, \beta_{\mathbf{k}'}^\dagger] = \delta_{\mathbf{k}\mathbf{k}'} \\ [\alpha_{\mathbf{k}}, \alpha_{\mathbf{k}'}] &= [\beta_{\mathbf{k}}, \beta_{\mathbf{k}'}] = 0 \\ [\alpha_{\mathbf{k}}, \beta_{\mathbf{k}'}] &= [\alpha_{\mathbf{k}}, \alpha_{\mathbf{k}'}] = 0 \end{aligned}$$

With the inverse transformation given by

$$(4.2.2.15) \quad \begin{pmatrix} \mathbf{b}_{\mathbf{k}} \\ \mathbf{b}_{-\mathbf{k}}^\dagger \end{pmatrix} = \begin{pmatrix} u_{\mathbf{k}} & v_{\mathbf{k}} \\ v_{\mathbf{k}} & u_{\mathbf{k}} \end{pmatrix} \begin{pmatrix} \alpha_{\mathbf{k}} \\ \beta_{\mathbf{k}}^\dagger \end{pmatrix} \Rightarrow \begin{aligned} \mathbf{b}_{\mathbf{k}} &= u_{\mathbf{k}}\alpha_{\mathbf{k}} + v_{\mathbf{k}}\beta_{\mathbf{k}}^\dagger, & \mathbf{b}_{\mathbf{k}}^\dagger &= u_{\mathbf{k}}\alpha_{\mathbf{k}}^\dagger + v_{\mathbf{k}}\beta_{\mathbf{k}} \\ \mathbf{b}_{-\mathbf{k}}^\dagger &= v_{\mathbf{k}}\alpha_{\mathbf{k}} + u_{\mathbf{k}}\beta_{\mathbf{k}}^\dagger, & \mathbf{b}_{-\mathbf{k}} &= v_{\mathbf{k}}\alpha_{\mathbf{k}}^\dagger + u_{\mathbf{k}}\beta_{\mathbf{k}} \end{aligned}$$

Then,

(4.2.2.16)

$$\begin{aligned} \mathbf{H}_1 &= JSz \sum_{\mathbf{k}} \left[\mathbf{b}_{\mathbf{k}}^\dagger \mathbf{b}_{\mathbf{k}} + \frac{\gamma_{\mathbf{k}}}{2} (\mathbf{b}_{\mathbf{k}}^\dagger \mathbf{b}_{-\mathbf{k}}^\dagger + \mathbf{b}_{\mathbf{k}} \mathbf{b}_{-\mathbf{k}}) \right] \\ &= JSz \sum_{\mathbf{k}} \left[(u_{\mathbf{k}}\alpha_{\mathbf{k}}^\dagger + v_{\mathbf{k}}\beta_{\mathbf{k}})(u_{\mathbf{k}}\alpha_{\mathbf{k}} + v_{\mathbf{k}}\beta_{\mathbf{k}}^\dagger) + v_{\mathbf{k}}^2 + \frac{\gamma_{\mathbf{k}}}{2} \begin{pmatrix} (u_{\mathbf{k}}\alpha_{\mathbf{k}}^\dagger + v_{\mathbf{k}}\beta_{\mathbf{k}})(v_{\mathbf{k}}\alpha_{\mathbf{k}} + u_{\mathbf{k}}\beta_{\mathbf{k}}^\dagger) \\ + (u_{\mathbf{k}}\alpha_{\mathbf{k}} + v_{\mathbf{k}}\beta_{\mathbf{k}}^\dagger)(v_{\mathbf{k}}\alpha_{\mathbf{k}}^\dagger + u_{\mathbf{k}}\beta_{\mathbf{k}}) \end{pmatrix} \right] \\ &= JSz \sum_{\mathbf{k}} \left[u_{\mathbf{k}}^2 \alpha_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}} + v_{\mathbf{k}}^2 \beta_{\mathbf{k}}^\dagger \beta_{\mathbf{k}} + u_{\mathbf{k}} v_{\mathbf{k}} (\alpha_{\mathbf{k}}^\dagger \beta_{\mathbf{k}}^\dagger + \beta_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}}^\dagger) + v_{\mathbf{k}}^2 \right. \\ &\quad \left. + \frac{\gamma_{\mathbf{k}}}{2} \begin{pmatrix} u_{\mathbf{k}} v_{\mathbf{k}} \alpha_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}} + u_{\mathbf{k}}^2 \alpha_{\mathbf{k}}^\dagger \beta_{\mathbf{k}}^\dagger + v_{\mathbf{k}}^2 \beta_{\mathbf{k}} \alpha_{\mathbf{k}} + u_{\mathbf{k}} v_{\mathbf{k}} \beta_{\mathbf{k}} \beta_{\mathbf{k}}^\dagger \\ + u_{\mathbf{k}} v_{\mathbf{k}} \alpha_{\mathbf{k}} \alpha_{\mathbf{k}}^\dagger + u_{\mathbf{k}}^2 \alpha_{\mathbf{k}} \beta_{\mathbf{k}} + v_{\mathbf{k}}^2 \beta_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}}^\dagger + u_{\mathbf{k}} v_{\mathbf{k}} \beta_{\mathbf{k}}^\dagger \beta_{\mathbf{k}} \end{pmatrix} \right] \\ &= JSz \sum_{\mathbf{k}} \left[\begin{aligned} &(u_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}}) \alpha_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}} + (v_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}}) \beta_{\mathbf{k}}^\dagger \beta_{\mathbf{k}} \\ &+ u_{\mathbf{k}} v_{\mathbf{k}} (\alpha_{\mathbf{k}}^\dagger \beta_{\mathbf{k}}^\dagger + \beta_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}}^\dagger) + 2\gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}} + v_{\mathbf{k}}^2 + \frac{\gamma_{\mathbf{k}}}{2} (u_{\mathbf{k}}^2 \alpha_{\mathbf{k}}^\dagger \beta_{\mathbf{k}}^\dagger + v_{\mathbf{k}}^2 \beta_{\mathbf{k}} \alpha_{\mathbf{k}} + u_{\mathbf{k}}^2 \alpha_{\mathbf{k}} \beta_{\mathbf{k}} + v_{\mathbf{k}}^2 \beta_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}}^\dagger) \end{aligned} \right] \\ &= JSz \sum_{\mathbf{k}} \left[\begin{aligned} &(u_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}}) \alpha_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}} + (v_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}}) \beta_{\mathbf{k}}^\dagger \beta_{\mathbf{k}} \\ &+ (u_{\mathbf{k}} v_{\mathbf{k}} + \gamma_{\mathbf{k}} u_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} v_{\mathbf{k}}^2) \alpha_{\mathbf{k}}^\dagger \beta_{\mathbf{k}}^\dagger + (u_{\mathbf{k}} v_{\mathbf{k}} + \gamma_{\mathbf{k}} u_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} v_{\mathbf{k}}^2) \alpha_{\mathbf{k}} \beta_{\mathbf{k}} + v_{\mathbf{k}}^2 + 2\gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}} \end{aligned} \right] \\ &= JSz \sum_{\mathbf{k}} \left[\begin{aligned} &(u_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}}) \alpha_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}} + (v_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}}) \beta_{\mathbf{k}}^\dagger \beta_{\mathbf{k}} \\ &+ (u_{\mathbf{k}} v_{\mathbf{k}} + \gamma_{\mathbf{k}} u_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} v_{\mathbf{k}}^2) (\alpha_{\mathbf{k}}^\dagger \beta_{\mathbf{k}}^\dagger + \alpha_{\mathbf{k}} \beta_{\mathbf{k}}) + 2\gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}} + v_{\mathbf{k}}^2 \end{aligned} \right] \end{aligned}$$

Hence, in order for this Hamiltonian to be diagonal, the u, v -coefficients must be such that the net antidiagonal coefficient is zero, ie. such that

$$(4.2.2.17) \quad \begin{aligned} u_{\mathbf{k}}^2 &= \frac{1}{2} \left(\frac{1}{\sqrt{1-\gamma_{\mathbf{k}}^2}} + 1 \right) \\ u_{\mathbf{k}} v_{\mathbf{k}} + \gamma_{\mathbf{k}} u_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} v_{\mathbf{k}}^2 &= 0, \quad \Rightarrow \quad v_{\mathbf{k}}^2 = \frac{1}{2} \left(\frac{1}{\sqrt{1-\gamma_{\mathbf{k}}^2}} - 1 \right) \\ u_{\mathbf{k}} v_{\mathbf{k}} &= -\frac{1}{2} \frac{\gamma_{\mathbf{k}}}{\sqrt{1-\gamma_{\mathbf{k}}^2}}, \end{aligned}$$

which leaves the classical Hamiltonian in diagonal form²², ie.

$$(4.2.2.18) \quad \begin{aligned} \mathbf{H}_1 &= JSz \sum_{\mathbf{k}} \left[(u_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}}) \alpha_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}} + (v_{\mathbf{k}}^2 + \gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}}) \beta_{\mathbf{k}}^\dagger \beta_{\mathbf{k}} + 2\gamma_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}} + v_{\mathbf{k}}^2 \right] \\ &= JSz \sum_{\mathbf{k}} \sqrt{1-\gamma_{\mathbf{k}}^2} (\alpha_{\mathbf{k}}^\dagger \alpha_{\mathbf{k}} + \beta_{\mathbf{k}}^\dagger \beta_{\mathbf{k}}) - Jz \sum_{\mathbf{k}} \left(1 - \sqrt{1-\gamma_{\mathbf{k}}^2} \right), \end{aligned}$$

which agrees with [equation \(4.2.1.24\)](#). The previous result for the antiferromagnet, unlike the ferromagnet, states that \mathbf{H}_1 has a quantum zero-point energy of size

²²Evidentemente tengo un error acá en la cuenta porque lo que realmente da es simétrico en los betas y alpha

$$(4.2.2.19) \quad E'_0 = E_0 - \left(-\frac{z}{2} N |J| S^2 \right) = \frac{1}{2} \sum_{\mathbf{k}} |J| S z \left(1 - \sqrt{1 - \gamma_{\mathbf{k}}^2} \right),$$

wherein $E'_0 < 0$, i.e. quantum fluctuations reduce the energy of the antiferromagnet. This entails that the classical Néel state is not the quantum antiferromagnet's true ground state. This can be understood by considering the simplest case of a Heisenberg antiferromagnetic spin chain, a one-dimensional two-site lattice. While the classical ground state has $(-JS^2)$ -energy, the quantum energy is considerably lower, standing at $-JS(S+1)$.

The ground state of the quantum antiferromagnet is the vacuum of the (α, β) -bosons. The Bogoliubov transformation thus introduces arbitrary numbers of Holstein-Primakoff bosons in the ground state. Directly in terms of spins, this means that the ground state admixes configurations with arbitrary numbers of spin-flips relative to the Néel state. This reduces the long-range Néel order at zero temperature.

Near $\mathbf{k} = 0$ and near $\mathbf{k} = \vec{\pi} = \pi \mathbb{1}_L$, the spin-wave spectrum vanishes as

$$(4.2.2.20) \quad \omega_{\mathbf{k}} \sim \begin{cases} JS\sqrt{2z}|\mathbf{k}| & |\mathbf{k}| \simeq 0 \\ JS\sqrt{2z}|\mathbf{k} - \vec{\pi}| & |\mathbf{k} - \vec{\pi}| \simeq 0 \end{cases},$$

with the order parameter being the staggered (Néel) magnetization. Its low-order corrections are given by \mathbf{H}_1 and read,

$$(4.2.2.21)$$

$$\Delta m_0^s = \frac{1}{N} \langle \sum_i e^{i\vec{\pi} \cdot \mathbf{x}_i} \mathbf{S}_i^z \rangle - S = -\frac{1}{N} \langle \sum_i \mathbf{b}_i^\dagger \mathbf{b}_i \rangle$$

$$= \frac{1}{2} - \frac{1}{N} \sum_{\mathbf{k}} \left(\mathbf{n}_{\mathbf{k}} + \frac{1}{2} \right) \frac{1}{\sqrt{1 - \gamma_{\mathbf{k}}^2}} \Rightarrow \Delta m_0 \sim \begin{cases} \frac{c_1}{k_0} & d = 1 \\ -c_2 + t \log k_0 & d = 2 \\ -c_3 + c'_3 t^2 & d = 3 \end{cases}, \text{ where } \begin{cases} c_d = \frac{1}{2N} \sum_{\mathbf{k}} \frac{1}{\sqrt{1 - \gamma_{\mathbf{k}}^2}} - 1, \\ c'_3 = 6^{-5/2}/2 \\ t = \frac{T}{|J|S\sqrt{d}}, \end{cases} d$$

where, in the last line, a truncation of the Holstein-Primakoff expansion upto quadratic order was employed and is justified only when $\Delta m_0^s \ll S$. The leading infrared singularities in Δm_0^s at low-temperatures is evaluated by expanding the Bose function near $\mathbf{n}_{\mathbf{k}} \approx 0$ and by using k_0 as a momentum cut-off, truncating the sum near $\mathbf{k} \approx 0, \vec{\pi}$.

As in the case of the Heisenberg ferromagnet, the truncation of the Holstein-Primakoff expansion is only justified when true-long-range order, see [definition 2](#) for the definition of true long-range order, is present and when $\Delta m_0^s \ll S$. Note that the staggered magnetization correction diverges, in the $k_0 \rightarrow 0$ limit for a one-dimensional system, signalling the failure of spin-wave theory and no long-range order in the ground state. Moreover, note that the existence of long-range order in the ground state for the $(d = 2, 3)$ -dimensional models depends on the relative sizes of c' and S . A three-dimensional model is the lowest dimension where long-range spin order can possibly exist at finite temperatures.

Physical Context

It turns out that the absence of long-range order in the d -dimensional quantum model is related to the Mermin-Wagner theorem for the classical Heisenberg model in $d + 1$ dimensions at finite temperatures.

4.2.3. The Continuum Approximation. In the previous sections, the spin-wave theory was derived in a semiclassical expansion of the spin path integral. Spin wave theory, however, is not applicable to the disordered (or "rotationally symmetric") phases of the Heisenberg model, since it assumes that the $O(3)$ symmetry is spontaneously broken. However, Mermin-Wagner's **theorem 5** states that there can be no spontaneously broken symmetry at finite temperatures in one and two dimensions. Moreover, for the Heisenberg antiferromagnet, there can be no long-range spin order even at zero temperature.

The first interesting question to arise is: can one use the semiclassical approximation in the absence of spontaneously broken symmetries?

- The answer is yes. A short-range classical Hamiltonian is sensitive mostly to short-range correlations. Thus, for the Heisenberg antiferromagnet, the important configurations in the semiclassical (large- S) limit (ie. the semiclassical configurations) have at least short-range antiferromagnetic order. At longer length scales, the semiclassical configurations can deviated largely from the Néel state. Thus, instead of breaking the rotational symmetry of the path integral as in the "naïve" spin-wave theory (which could yield incorrect results), the short-length scale fluctuations can be eliminated and thus keep the full rotational symmetry of the long-wavelength semiclassical modes.

Consider the imaginary-time path integral given by **equation (4.1.1.5)** and consider the antiferromagnetic Hamiltonian,

$$\mathbf{H}[\hat{\Omega}] = \langle \hat{\Omega} | \mathbf{H} | \hat{\Omega} \rangle = \frac{S^2}{2} \sum_{ij} J_{ij} \hat{\Omega}_i \cdot \hat{\Omega}_j.$$

Without loss of generality, one can restrict the discussion to cubic d -dimensional lattices with lattice-constant a , number of sites N and even number of sites in each dimension, wherein the J_{ij} -constants have full lattice symmetry. For the classical Hamiltonian $\mathbf{H}[\hat{\Omega}]$, this gives rise to a Néel ground state. Furthermore, the J_{ij} -constants are assumed to be short-range, see **theorem 5** for further reference.

4.2.3.1. Haldane's Mapping.

Physical Context

The Haldane's mapping is the correspondence between the effective long-wavelength action of the quantum d -dimensional Heisenberg antiferromagnet into the $(d + 1)$ -dimensional nonlinear sigma model (NLSM).

The essence of the Haldane map is the separation between the short- and long-length scale fluctuations via suitable choice of coordinates. In effect, consider two continuous vector fields $\hat{\mathbf{n}}$, \mathbf{L} which parameterize the spins as follows,

$$(4.2.3.1) \quad \hat{\Omega}_i = \eta_i \hat{\mathbf{n}}(\mathbf{x}_i) \sqrt{1 - \left| \frac{\mathbf{L}(\mathbf{x}_i)}{S} \right|^2} + \frac{\mathbf{L}(\mathbf{x}_i)}{S}, \text{ where } \eta_i = e^{i\mathbf{X}_i \cdot \vec{\pi}} \text{ has opposite signs on the sublattices}$$

and where $\hat{\mathbf{n}}$ is the unimodular Néel field, such that $|\hat{\mathbf{n}}(\mathbf{X}_i)| = 1$, and where \mathbf{L} is the transverse canting field, chosen to obey $\mathbf{L}(\mathbf{x}_i) \cdot \hat{\mathbf{n}}(\mathbf{x}_i) = 0$.

In the previous definition, at first glance, it may seem that the two independent degrees of freedom per site, (θ_i, ϕ_i) , have been replaced with four variables, six degrees degrees of freedom for $(\mathbf{n}_i, \mathbf{L}_i)$ minus the two constraints. This discrepancy is readily resolved by fixing the total number of Fourier components in the integration measure

(4.2.3.2)

$$\mathcal{D}\hat{\Omega} = \prod_{\mathbf{q} \leq \Lambda_{\text{BZ}}} d\hat{\mathbf{n}}_{\mathbf{q}} d\mathbf{L}_{\mathbf{q}} \Delta(\mathbf{L} \cdot \hat{\mathbf{n}}) \mathcal{J}[\hat{\mathbf{n}}, \mathbf{L}],$$

$$\text{where } \left\{ \begin{array}{l} J \text{ is the Jacobian of the transformation given by equation (4.2.3.1)} \\ \text{where the Fourier transform reads } \mathbf{X}_{\mathbf{q}} = \sum_i e^{-i\mathbf{q} \cdot \mathbf{x}_i} \mathbf{X}(\mathbf{x}_i), \mathbf{X} = \hat{\mathbf{n}}, \mathbf{L}, \dots, \\ \text{and where } \Lambda_{\text{BZ}} \text{ is the spherical Brillouin zone radius, chosen such that } 2N = 4 \sum \mathbf{q} \leq \Lambda_{\text{BZ}}. \end{array} \right.$$

The left- and right-hand sides of the definition of the spherical Brillouin one radius count the effective degrees of freedom, for the left- and right-hand sides of equation (4.2.3.2) respectively.

In the short-range ordered phases, \mathcal{Z} is dominated by configurations with sizeable antiferromagnetic correlations for distances below the ε -correlation length. For large $\frac{\varepsilon}{a}$, it can be assumed that the intermediate momentum cutoff scale Λ is much smaller than the microscopic momenta, but also much larger than the inverse correlation length; in more mundane terms, that

$$(4.2.3.3) \quad \zeta^{-1} \ll \Lambda \ll \frac{\Lambda_{\text{BZ}}}{2\pi/R_J},$$

where R_J is the characteristic range of J_{ij} . The previous requirements are equivalent to assuming that the dominant configurations in the path integral are *slowly varying* on the scale of Λ^{-1} , or that these configurations have negligible Fourier components for $|\hat{\mathbf{n}}_{|\mathbf{q}|>\Lambda}| \ll 1$. In these regimes, the cubic Brillouin zone of the lattice can be replaced by the spherical zone written in equation (4.2.3.2).

Moreover, $|\hat{\mathbf{n}}_{|\mathbf{q}|>\Lambda}| \ll 1$ is consistent with assuming that the canting field is small, namely that

$$\left| \frac{\mathbf{L}_i}{S} \right| \Rightarrow \begin{array}{l} \text{which implies that, upto leading order } \mathcal{O}\left(\frac{\mathbf{L}}{S}\right) \quad \mathcal{J} \approx S^{-N}. \\ \text{the Jacobian in equation (4.2.3.2) is constant} \end{array}$$

The assumption of a slowly varying \mathbf{L} -field in space is not needed. The main assumption is encapsulated in the existence of a wave vector scale Λ . The theory's self-consistency, in particular that of this assumption, must be checked by evaluating $\varepsilon(\Lambda)$ and by verifying that the inequality, written in equation (4.2.3.3), holds.

4.2.3.2. The Continuum Hamiltonian. The Heisenberg Hamiltonian can then be written in terms of the $(\hat{\mathbf{n}}, \mathbf{L})$ -fields. Upto quadratic order $\mathcal{O}\left(\frac{\mathbf{L}}{S}\right)^2$, the interactions read

$$(4.2.3.4) \quad \begin{aligned} \hat{\Omega}_i \cdot \hat{\Omega}_j &= \left[\eta_i \hat{\mathbf{n}}(\mathbf{x}_i) \sqrt{1 - \left| \frac{\mathbf{L}(\mathbf{x}_i)}{S} \right|^2} + \frac{\mathbf{L}(\mathbf{x}_i)}{S} \right] \cdot \left[\eta_j \hat{\mathbf{n}}(\mathbf{x}_j) \sqrt{1 - \left| \frac{\mathbf{L}(\mathbf{x}_j)}{S} \right|^2} + \frac{\mathbf{L}(\mathbf{x}_j)}{S} \right] \\ &\approx \left[\eta_i \hat{\mathbf{n}}(\mathbf{x}_i) \left(1 - \frac{1}{2} \left| \frac{\mathbf{L}(\mathbf{x}_i)}{S} \right|^2 \right) + \frac{\mathbf{L}(\mathbf{x}_i)}{S} \right] \cdot \left[\eta_j \hat{\mathbf{n}}(\mathbf{x}_j) \left(1 - \frac{1}{2} \left| \frac{\mathbf{L}(\mathbf{x}_j)}{S} \right|^2 \right) + \frac{\mathbf{L}(\mathbf{x}_j)}{S} \right] \\ &\approx \left[\eta_i \hat{\mathbf{n}}(\mathbf{x}_i) - \eta_i \hat{\mathbf{n}}(\mathbf{x}_i) \frac{1}{2} \left| \frac{\mathbf{L}(\mathbf{x}_i)}{S} \right|^2 + \frac{\mathbf{L}(\mathbf{x}_i)}{S} \right] \cdot \left[\eta_j \hat{\mathbf{n}}(\mathbf{x}_j) - \eta_j \hat{\mathbf{n}}(\mathbf{x}_j) \frac{1}{2} \left| \frac{\mathbf{L}(\mathbf{x}_j)}{S} \right|^2 + \frac{\mathbf{L}(\mathbf{x}_j)}{S} \right] \\ &\approx \eta_i \eta_j \hat{\mathbf{n}}(\mathbf{x}_i) \cdot \hat{\mathbf{n}}(\mathbf{x}_j) - \eta_i \eta_j \hat{\mathbf{n}}(\mathbf{x}_i) \cdot \hat{\mathbf{n}}(\mathbf{x}_j) \frac{1}{2} \left| \frac{\mathbf{L}(\mathbf{x}_j)}{S} \right|^2 + \eta_i \hat{\mathbf{n}}(\mathbf{x}_i) \cdot \frac{\mathbf{L}(\mathbf{x}_j)}{S} \\ &\quad - \eta_i \hat{\mathbf{n}}(\mathbf{x}_i) \frac{1}{2} \left| \frac{\mathbf{L}(\mathbf{x}_i)}{S} \right|^2 \cdot \eta_j \hat{\mathbf{n}}(\mathbf{x}_j) - \eta_i \hat{\mathbf{n}}(\mathbf{x}_i) \frac{1}{2} \left| \frac{\mathbf{L}(\mathbf{x}_i)}{S} \right|^2 \cdot \frac{\mathbf{L}(\mathbf{x}_j)}{S} \\ &\quad + \frac{\mathbf{L}(\mathbf{x}_i)}{S} \cdot \eta_j \hat{\mathbf{n}}(\mathbf{x}_j) - \frac{\mathbf{L}(\mathbf{x}_i)}{S} \cdot \eta_j \hat{\mathbf{n}}(\mathbf{x}_j) \frac{1}{2} \left| \frac{\mathbf{L}(\mathbf{x}_j)}{S} \right|^2 + \frac{\mathbf{L}(\mathbf{x}_i)}{S} \cdot \frac{\mathbf{L}(\mathbf{x}_j)}{S} + \mathcal{O}\left(\frac{\mathbf{L}}{S}\right)^4 \end{aligned}$$

$$\begin{aligned}
(4.2.3.5) \quad & \approx \frac{\eta_i \eta_j}{2} \left(\hat{\mathbf{n}}(\mathbf{x}_i)^2 + \hat{\mathbf{n}}(\mathbf{x}_j)^2 - (\hat{\mathbf{n}}(\mathbf{x}_i) - \hat{\mathbf{n}}(\mathbf{x}_j))^2 \right) + \frac{1}{S} \left(\eta_i \mathbf{L}_j \cdot \hat{\mathbf{n}}(\mathbf{x}_i) + \eta_j \mathbf{L}_i \cdot \hat{\mathbf{n}}(\mathbf{x}_j) \right) \\
& + \frac{1}{S^2} \left(\mathbf{L}_i \cdot \mathbf{L}_j - \frac{1}{2} \eta_i \eta_j (\hat{\mathbf{n}}(\mathbf{x}_i) - \hat{\mathbf{n}}(\mathbf{x}_j))^2 (\mathbf{L}_i^2 + \mathbf{L}_j^2) \right) \\
& \approx \eta_i \eta_j - \frac{\eta_i \eta_j}{2} (\hat{\mathbf{n}}(\mathbf{x}_i) - \hat{\mathbf{n}}(\mathbf{x}_j))^2 + \frac{1}{S^2} \left(\mathbf{L}_i \cdot \mathbf{L}_j - \frac{1}{2} \eta_i \eta_j (\hat{\mathbf{n}}(\mathbf{x}_i) - \hat{\mathbf{n}}(\mathbf{x}_j))^2 (\mathbf{L}_i^2 + \mathbf{L}_j^2) \right) \\
& + \frac{1}{S} \left(\eta_i \mathbf{L}_j \cdot \hat{\mathbf{n}}(\mathbf{x}_i) + \eta_j \mathbf{L}_i \cdot \hat{\mathbf{n}}(\mathbf{x}_j) \right) + \mathcal{O}(|\mathbf{L}|^2 |\hat{\mathbf{n}}(\mathbf{x}_i) - \hat{\mathbf{n}}(\mathbf{x}_j)|)
\end{aligned}$$

The differences between Néel fields can be approximated by derivatives

$$(4.2.3.6) \quad |\hat{\mathbf{n}}(\mathbf{x}_i) - \hat{\mathbf{n}}(\mathbf{x}_j)| \approx \partial_\ell \hat{\mathbf{n}}(\mathbf{x}) x_{ij}^\ell + \frac{1}{2} (\partial_{\ell k}^2 \hat{\mathbf{n}}(\mathbf{x})) x_{ij}^\ell x_{ij}^k, \quad \text{where } \mathbf{x}_{ij} = \mathbf{x}_i - \mathbf{x}_j, \text{ and where the summation is repeated over spatial indices, } \ell, k = 1, \dots, d.$$

From this and using [equation \(4.2.3.3\)](#), it is readily apparent that higher-order derivatives are higher order power in the small parameter $\Lambda R_J \ll 1$.

The first derivative cancels out in the cross terms by the symmetry of the Hamiltonian, which leads to the $\mathcal{O}(S^{-1})$ -terms being negligible. In effect,

$$\begin{aligned}
(4.2.3.7) \quad & \sum_{ij} \eta_i \mathbf{L}_j \cdot \hat{\mathbf{n}}(\mathbf{x}_i) \approx \sum_{ij} \eta_i J_{ij} \mathbf{L}_j \left[\hat{\mathbf{n}}(\mathbf{x}) + \partial_\ell \hat{\mathbf{n}}(\mathbf{x}) x_{ij}^\ell + \frac{1}{2} (\partial_{\ell k}^2 \hat{\mathbf{n}}(\mathbf{x})) x_{ij}^\ell x_{ij}^k \right] \\
& \approx \frac{1}{2} \sum_{ij} (\partial_{\ell k}^2 \hat{\mathbf{n}}(\mathbf{x})) x_{ij}^\ell x_{ij}^k \sim \mathcal{O}(\Lambda R_J)^2, \text{ which are negligible.}
\end{aligned}$$

The lattice sums can then be replaced by integrals, following the prescription

$$(4.2.3.8) \quad \sum_i F_i \rightarrow a^{-d} \int d^d x \sum_i \delta(\mathbf{x} - \mathbf{x}_i) F(\mathbf{x}),$$

leading to the continuum representation for the Hamiltonian, which reads

$$\begin{aligned}
(4.2.3.9) \quad & \mathbf{H} = \frac{S^2}{2} \sum_{ij} J_{ij} \left[\eta_i \eta_j - \frac{\eta_i \eta_j}{2} (\hat{\mathbf{n}}(\mathbf{x}_i) - \hat{\mathbf{n}}(\mathbf{x}_j))^2 \right] + \frac{1}{S^2} \left(\mathbf{L}_i \cdot \mathbf{L}_j - \frac{1}{2} \eta_i \eta_j (\hat{\mathbf{n}}(\mathbf{x}_i) - \hat{\mathbf{n}}(\mathbf{x}_j))^2 (\mathbf{L}_i^2 + \mathbf{L}_j^2) \right) \\
& = \frac{S^2}{2} \sum_{ij} J_{ij} \eta_i \eta_j - \frac{S^2}{2} \sum_{ij} J_{ij} \frac{\eta_i \eta_j}{2} \frac{1}{N a^d} \int d^d x \sum_\ell |\partial_\ell \hat{\mathbf{n}}(\mathbf{x})| |\mathbf{x}^2| + \frac{J_{ij}}{2} \frac{1}{a^d} \int d^d x \int d^d x' \sum_{ij} \delta(\mathbf{x} - \mathbf{x}_i) \delta(\mathbf{x} - \mathbf{x}_j) \mathbf{L}_{\mathbf{x}} \cdot \mathbf{L}_{\mathbf{x}'} \\
& + \frac{J_{ij}}{2} \frac{1}{a^d} \int d^d x \int d^d x' \sum_{ij} \eta_i \eta_j \delta(\mathbf{x} - \mathbf{x}') \delta(\mathbf{x} - \mathbf{x}_i) \\
& = \frac{S^2}{2} \sum_{ij} J_{ij} \eta_i \eta_j - \frac{S^2}{2 d N a^d} \sum_{ij} J_{ij} \eta_i \eta_j |\mathbf{x}|^2 \\
& + \int \frac{d^d x d^d x'}{2} \sum_{ij} J_{ij} \mathbf{L}_{\mathbf{x}} \left[\delta(\mathbf{x} - \mathbf{x}_i) \delta(\mathbf{x}' - \mathbf{x}_j) - \delta(\mathbf{x}' - \mathbf{x}) \delta(\mathbf{x} - \mathbf{x}_i) \eta_i \eta_j \right] \mathbf{L}_{\mathbf{x}'} \\
& = E_0^{\text{cl}} + \int \frac{d^d x}{2} \left[\begin{aligned} & \rho_s \sum_\ell |\partial_\ell \hat{\mathbf{n}}(\mathbf{x})|^2 \\ & + \int d^d x' (\mathbf{L}_{\mathbf{x}} \chi_{xx'}^{-1} \mathbf{L}_{\mathbf{x}'}) \end{aligned} \right], \quad E_0^{\text{cl}} = \frac{S^2}{2} \sum_{ij} J_{ij} \eta_i \eta_j, \quad \rho_s = -\frac{S^2}{2 d N a^d} \sum_{ij} J_{ij} \eta_i \eta_j |\mathbf{x}_i - \mathbf{x}_j|^2, \\
& \quad \chi_{\mathbf{x}, \mathbf{x}'}^{-1} = \frac{1}{N a^d} \sum_{ij} J_{ij} \left[\delta(\mathbf{x} - \mathbf{x}_i) \delta(\mathbf{x}' - \mathbf{x}_j) - \delta(\mathbf{x}' - \mathbf{x}) \delta(\mathbf{x} - \mathbf{x}_i) \eta_i \eta_j \right],
\end{aligned}$$

where E_0^{cl} is the classical energy, ρ_s is the stiffness constant and where $\chi_{\mathbf{x}, \mathbf{x}'}^{-1}$ is the inverse susceptibility. In Fourier space, the canting field term is given by

$$(4.2.3.10) \quad \int d^d x \int d^d x' \mathbf{L}_x \chi_{xx'}^{-1} \mathbf{L}_{x'} = \int_{q \leq \Lambda_{\mathfrak{B}3}} \frac{d^d q}{(2\pi)^d} \mathbf{L}_q \chi(\mathbf{q})^{-1} \mathbf{L}_{-\mathbf{q}}, \text{ where } \chi(\mathbf{q}) = \frac{1}{J(\mathbf{q}) - J(\vec{\pi})} \\ J(\mathbf{q}) = \sum_j e^{i\mathbf{q} \cdot \mathbf{x}_{ij}} J_{ij}.$$

4.2.3.3. The Kinetic Term. The Kinetic Berry phase term in the spin path integral is given by [equation \(4.1.1.6\)](#), reading

$$(4.2.3.11) \quad -iS \sum_i \omega_i = -iS \int_0^\beta d\tau \sum_i \mathbf{A}(\hat{\Omega}_i) \cdot \hat{\Omega}_i.$$

It is convenient to consider a inversion-symmetric vector potential, i.e. a vector potential such that $\mathbf{A}(\hat{\Omega}) = \mathbf{A}(-\hat{\Omega})$, such as

$$\mathbf{A} = -\frac{\cos \theta}{\sin \theta} \hat{\phi}.$$

Then, expanding the kinetic Berry phase in terms of the Néel and canting fields yields

$$(4.2.3.12) \quad \begin{aligned} -iS \sum_i \omega_i &= -iS \sum_i \eta_i \omega \left[\hat{\mathbf{n}}_i + \eta_i \frac{\mathbf{L}_i}{S} \right] = -iS \sum_i \eta_i \omega [\hat{\mathbf{n}}_i] + \frac{\delta \omega}{\delta \hat{\mathbf{n}}_i} \frac{\mathbf{L}_i}{S} \\ &= -iS \sum_i \left[\eta_i \omega [\hat{\mathbf{n}}_i] + \frac{\delta}{\delta \hat{\mathbf{n}}_i} \int_0^\beta d\tau \hat{\Omega} \cdot (\dot{\hat{\Omega}} \times \delta \hat{\Omega}) \right] \\ &= -iS \sum_i \left[\eta_i \omega [\hat{\mathbf{n}}_i] + \frac{\delta}{\delta \hat{\mathbf{n}}_i} \int_0^\beta \left(\hat{\mathbf{n}}_i + \eta_i \frac{\mathbf{L}_i}{S} \right) \cdot \left\{ \left(\partial_\tau \hat{\mathbf{n}}_i + \eta_i \frac{\partial_\tau \mathbf{L}_i}{S} \right) \times \left(\delta \hat{\mathbf{n}}_i + \eta_i \frac{\delta \mathbf{L}_i}{S} \right) \right\} \right] \\ &= -i\Gamma - i \int_0^\beta d\tau \sum_i \hat{\mathbf{n}}_i \times \partial_\tau \hat{\mathbf{n}}_i \cdot \mathbf{L}_i, \quad \text{where } \Gamma[\hat{\mathbf{n}}] = S \sum_i \eta_i \omega [\hat{\mathbf{n}}_i] \end{aligned}$$

Here, Γ is a topological Berry phase associated with the Néel field $\hat{\mathbf{n}}$. In the rotationally symmetric phases, interference between Berry phases can give rise to dramatic consequences on ground state degeneracy and low-energy spectrum.

4.2.3.4. Partition Function and Correlations. The second term in [equation \(4.2.3.12\)](#) couples the canting field to $\hat{\mathbf{n}} \times \partial_\tau \hat{\mathbf{n}}$, implying that said fields are canonical conjugates, akin to the $\frac{1}{2}(\mathbf{p} \cdot \dot{\mathbf{q}} - \dot{\mathbf{q}} \cdot \mathbf{p})$ in [equation \(4.2.1.4\)](#). The constraint $\delta(\mathbf{L} \cdot \hat{\mathbf{n}})$ is automatically satisfied in the second term of [equation \(4.2.3.12\)](#). From these, the partition function can be obtained from [equation \(4.1.1.5\)](#) via Gaussian integration

$$(4.2.3.13) \quad \begin{aligned} \mathcal{Z}[j] &= \int \mathcal{D}\hat{\Omega}(\tau) \exp \left[-iS \sum_i \omega[\hat{\Omega}_i] + \int_0^\beta \mathcal{H}[\hat{\Omega}(\tau)] \right] \\ &= \int \mathcal{D}\hat{\Omega}(\tau) e^{i\Gamma[\hat{n}]} \exp \left[-i \int_0^\beta d\tau \left(\sum_i \hat{\mathbf{n}}_i \times \partial_\tau \hat{\mathbf{n}}_i \cdot \mathbf{L}_i + E_0^{\text{cl}} + \int \frac{d^d x}{2} \rho_s \sum_\ell |\partial_\ell \hat{\mathbf{n}}(\mathbf{x})|^2 + \int d^d x' (\mathbf{L}_x \chi_{xx'}^{-1} \mathbf{L}_{x'}) \right) \right] \\ &= \int \mathcal{D}\hat{\mathbf{n}} e^{i\Gamma[\hat{n}]} \exp \left(- \int_0^\beta d\tau \int_\Lambda d^d x \frac{\rho_s}{2} \sum_{\ell=1}^d |\partial_\ell \hat{\mathbf{n}}(\mathbf{x})|^2 \right) \\ &\quad \times \underbrace{\int_{\Lambda_{\mathfrak{B}3}} \mathcal{D}\mathbf{L}^\perp \exp \left[- \int_0^\beta d\tau \int_\Lambda d^d x \frac{d^d q}{(2\pi)^d} (\hat{\mathbf{n}} \times \partial_\tau \hat{\mathbf{n}})_{-\mathbf{q}} \cdot \mathbf{L}_q - \frac{1}{2} \int_\Lambda \frac{d^d q}{(2\pi)^d} \chi^{-1}(\mathbf{q}) \mathbf{L}_q \cdot \mathbf{L}_{-\mathbf{q}} \right]}_{\propto \exp \left[-\frac{1}{2} \int d\tau \int^\Lambda \frac{d^d q}{(2\pi)^d} \chi(\mathbf{q}) (\hat{\mathbf{n}} \times \partial_\tau \hat{\mathbf{n}})_{\mathbf{q}} \cdot (\hat{\mathbf{n}} \times \partial_\tau \hat{\mathbf{n}})_{-\mathbf{q}} \right]} \end{aligned}$$

Inspecting [equation \(4.2.3.10\)](#), is it readily apparent that the variation of $\chi(\mathbf{q})$ is small on the scale of Λ , allowing itself to be replaced by its zero-momentum-limit. First, note that

$$\begin{aligned}
 |\hat{\mathbf{n}}(\mathbf{x}) \times \partial_\tau \hat{\mathbf{n}}(\mathbf{x})|^2 &= \delta^{ij} (\hat{\mathbf{n}}(\mathbf{x}) \times \partial_\tau \hat{\mathbf{n}}(\mathbf{x}))^i (\hat{\mathbf{n}}(\mathbf{x}) \times \partial_\tau \hat{\mathbf{n}}(\mathbf{x}))^j \\
 &= \delta^{ij} (\epsilon_{abi} \hat{\mathbf{n}}_a \partial_\tau \hat{\mathbf{n}}_b) (\epsilon^{cdj} \hat{\mathbf{n}}_c \partial_\tau \hat{\mathbf{n}}_d) \\
 &= \epsilon_{abi} \epsilon^{cdj} (\hat{\mathbf{n}}_a \hat{\mathbf{n}}_c \partial_\tau \hat{\mathbf{n}}_b \partial_\tau \hat{\mathbf{n}}_d) \\
 &= (\delta_a^c \delta_b^d - \delta_a^d \delta_b^c) (\hat{\mathbf{n}}_a \hat{\mathbf{n}}_c \partial_\tau \hat{\mathbf{n}}_b \partial_\tau \hat{\mathbf{n}}_d) \\
 &= \delta_a^c \delta_b^d \hat{\mathbf{n}}_a \hat{\mathbf{n}}_c \partial_\tau \hat{\mathbf{n}}_b \partial_\tau \hat{\mathbf{n}}_d - \delta_a^d \delta_b^c \hat{\mathbf{n}}_a \hat{\mathbf{n}}_c \partial_\tau \hat{\mathbf{n}}_b \partial_\tau \hat{\mathbf{n}}_d \\
 &= \hat{\mathbf{n}}_a \hat{\mathbf{n}}^a \partial_\tau \hat{\mathbf{n}}_b \partial_\tau \hat{\mathbf{n}}^b - \hat{\mathbf{n}}_a \hat{\mathbf{n}}_b \partial_\tau \hat{\mathbf{n}}^b \partial_\tau \hat{\mathbf{n}}^a = |\partial_\tau \hat{\mathbf{n}}|^2.
 \end{aligned}
 \tag{4.2.3.14}$$

Hence, the underbraced term in the previous equation for the partition function, allows it to be rewritten as a local interaction on the unimodular fields,

$$\begin{aligned}
 &\propto \exp \left[-\frac{1}{2} \int d\tau \int^\Lambda \frac{d^d q}{(2\pi)^d} \chi(\mathbf{q}) (\hat{\mathbf{n}} \times \partial_\tau \hat{\mathbf{n}})_\mathbf{q} \cdot (\hat{\mathbf{n}} \times \partial_\tau \hat{\mathbf{n}})_{-\mathbf{q}} \right] \\
 &\approx \exp \left[-\frac{1}{2} \int_0^\beta d\tau \int_\Lambda d^d x \chi_0 |\partial_\tau \hat{\mathbf{n}}|^2 \right], \quad \text{where } \chi_0 = a^{-d} \chi(0),
 \end{aligned}
 \tag{4.2.3.15}$$

and where $\int_\Lambda d^d x$ is understood to contain only the Fourier components of the integrand which obey $|\mathbf{q}| \leq \Lambda$. Hence, the partition function can be written as

$$\mathcal{Z} \propto \int_\Lambda \mathcal{D}\hat{\mathbf{n}} e^{i\Gamma[\hat{\mathbf{n}}]} \exp \left[-\frac{1}{2} \int_0^\beta d\tau \int_\Lambda d^d x \left(\chi_0 |\partial_\tau \hat{\mathbf{n}}|^2 + \rho_s \sum_{\ell=1}^d |\partial_\ell \hat{\mathbf{n}}|^2 \right) \right].
 \tag{4.2.3.16}$$

This can be expressed in a more compact form, by defining the spin wave velocity $c = \sqrt{\rho_s \chi_0^{-1}}$, and by defining a special coordinate system,

$$(x_1, \dots, x_d, c\tau) \rightarrow (x_1, \dots, x_{d+1}),
 \tag{4.2.3.17}$$

which allows the action of \mathcal{Z} of the Non-linear sigma model NLSM in $(d+1)$ -dimensions to be written in terms of Berry phases as,

$$\mathcal{Z}_{\text{NLSM}} \propto \int_\Lambda \mathcal{D}\hat{\mathbf{n}} e^{i\Gamma} \exp \left(-\int d^{d+1}x \mathcal{L}_{\text{NLSM}}^{d+1} \right), \quad \mathcal{L}_{\text{NLSM}}^D = \frac{\Lambda^{D-2}}{2f_D} \sum_{\mu=1}^D \partial_\mu \hat{\mathbf{n}} \cdot \partial_\mu \hat{\mathbf{n}}, \quad \begin{matrix} |\hat{\mathbf{n}}(\mathbf{x})| = 1, \\ \mathbf{x} \in \mathbb{R}^d. \end{matrix}
 \tag{4.2.3.18}$$

upto to normalization constants and where $f = (c/\rho_s)\Lambda^{D-2}$ is a dimensionless coupling constant. Using [equation \(4.2.3.18\)](#), several observables of physical relevance can be calculated, eg.

- the spin correlation function for the Heisenberg antiferromagnet at long-distances and at low temperature can be easily calculated as

$$\langle \mathbf{S}_i^\alpha \mathbf{S}_j^\alpha(\tau) \rangle_d \approx \frac{e^{i\vec{\pi}(\mathbf{x}_i - \mathbf{x}_j)} S^2}{N \mathcal{Z}} \int_\Lambda \mathcal{D}\hat{\mathbf{n}} e^{i\Gamma} \exp \left(-\int_0^{c\beta} dx^{d+1} \int d^d x \mathcal{L}_{\text{NLSM}}^{d+1} \right) \hat{\mathbf{n}}^\alpha(0, 0) \hat{\mathbf{n}}^\alpha(\mathbf{x}_j, c\tau),
 \tag{4.2.3.19}$$

- with the imaginary-time ground state correlations of the quantum Heisenberg antiferromagnet given by

$$\frac{1}{\mathcal{Z}} \sum_m |\langle \Psi_0 | \hat{\mathbf{n}}(0) | \Psi_0 \rangle|^2 e^{-(E_m - E_0)\tau} = \langle \hat{\mathbf{n}}^\alpha(0, 0) \hat{\mathbf{n}}^\alpha(0, c\tau) \rangle.
 \tag{4.2.3.20}$$

The NLSM is rotationally symmetric in spacetime. If the correlations exponentially-decay at large distances with ε -correlation length, the imaginary time decay period is given by ε/c . According to [equation \(4.2.3.20\)](#), the lowest excitation energy Δ determines the fastest decay rate, hence

$$\Delta = \min_m E_m - E_0 \approx \frac{c}{\varepsilon}.$$

If Δ does not vanish with the size of the system, it is a thermodynamically meaningful gap in the excitation spectrum called Haldane's gap.

In the absence of topological Γ -Berry phases, the ground state of the quantum antiferromagnet is described by the classical energy of the $(d+1)$ -dimensional NLSM, where f is a coupling constant of the classical problem. At finite, but low, β^{-1} -temperature, the quantum antiferromagnet is thus mapped onto the NLSM on a slab of finite width $c\beta$ in the imaginary-time dimension.

Caution must be taken, in that the Haldane mapping is valid under the restrictive conditions imposed by [equation \(4.2.3.3\)](#), which can be achieved in the $(S \gg 1)$ -limit. In said regime, the large- S semiclassical limit corresponds to the weak coupling ($f \ll 1$)-limit of the NLSM.

4.2.4. NonLinear Sigma Model: Weak Coupling. The NonLinear Sigma model in d -dimension emerges as the continuum approximation to the quantum Heisenberg antiferromagnet in $(d-1)$ -dimensions, with additional Berry phases, wherein the partition function $\mathcal{Z}_{\text{NLSM}}$ is given by [equation \(4.2.3.18\)](#) and where Λ is a momentum cutoff, i.e. the shortest wavelength included in the $\mathcal{D}\hat{\mathbf{n}}$ -field.

4.2.4.1. The Lattice Regularization. One can, in fact, rigorously prove the connection between the NLSM in d -dimensions and the quantum Heisenberg antiferromagnet in $(d-1)$ -dimensions via a regularization procedure. But note that [equation \(4.2.3.18\)](#)'s path integral is not well-defined until a regularization procedure is established.

Hence, first consider a d -dimensional classical Heisenberg model on a cubic-lattice, with N sites and lattice constant a . Its partition function reads

(4.2.4.1)

$$\mathcal{Z}_{\text{CMH}} = \int \prod_{i=1}^N d\hat{\Omega}_i \exp \left(\frac{\bar{J}}{T} \sum_{\langle ij \rangle} \hat{\Omega}_i \cdot \hat{\Omega}_j \right), \quad \begin{array}{l} \text{where the continuum} \\ \text{limit of } \mathcal{Z}_{\text{CMH}} \text{ is} \\ \text{given by the substitutions} \end{array} \quad \text{with } \eta_i = \begin{cases} \hat{\Omega}_i \rightarrow \eta_i \hat{\mathbf{n}}(\mathbf{x}_i), \\ 1 & \text{ferromagnet} \\ e^{i\mathbf{x}_i \cdot \vec{\pi}} & \text{antiferromagnet} \end{cases}$$

Moreover, sums and differences are substituted by lattice integrals and spatial derivatives, as follows

$$(4.2.4.2) \quad \sum_i F_i(\mathbf{x}_i) \rightarrow a^{-d} \int d^d x F(\mathbf{x}), \quad \hat{\Omega}(\mathbf{x}_i + a\hat{x}_\mu) - \hat{\Omega}(\mathbf{x}_i) \rightarrow a\partial_\mu \hat{\mathbf{n}},$$

and a dimensionless coupling constant $\frac{T}{J} \leftrightarrow f$. The measure is replaced by

$$(4.2.4.3) \quad \mathcal{D}\hat{\Omega} \rightarrow \prod_{\mathbf{q} \leq \Lambda} d\hat{\mathbf{n}}_{\mathbf{q}}, \quad \begin{array}{l} \text{where } \Lambda \text{ is the radius of a spherical Brillouin zone that has the same} \\ \text{number of degrees of freedom as the cubic zone of the lattice,} \end{array} \quad \frac{1}{N} \sum_{|k_\mu| \leq \frac{\pi}{a}}^{\text{cube}} = \frac{1}{N} \sum_{|k| \leq \Lambda}^{\text{sph}} = 1.$$

This determines the spherical Brillouin zone radius to be $\Lambda \sim a^{-1}$ for $d = 1, 2, 3$, with different constants each. A second-order expansion of the Hamiltonian in [equation \(4.2.4.1\)](#) in terms of gradients of the unimodular Néel field yields the following relation

$$(4.2.4.4) \quad \mathcal{Z}_{\text{CMH}} \left(\frac{T}{J}, a \right) \approx \mathcal{Z}_{\text{NLSM}}(f, \Lambda^{-1}) e^{\frac{aN}{f}},$$

which is valid at low-temperatures $f \ll 1$. The $e^{\frac{dN}{f}}$ -factor is the Boltzmann weight for the classical ground state. Note that the continuum approximation holds for the generating functional and for the long-wavelength spin correlations as well, as long as the spin correlation length ε is much larger than Λ^{-1} . Hence, in this regime, the NLSM properties are weakly-dependent on the regularization procedure, and therefore it serves as a model for diverse Heisenberg models. The key common feature of these Heisenberg models is the $O(3)$ -symmetry of the ground state manifold and the short-range nature of the interactions. Other details will primarily affect the choice of f and Λ .

The "nonlinearity" property of the NLSM is due to the Néel's field unimodular constraint, which may be expressed as

$$(4.2.4.5) \quad n^{(1)}(\mathbf{x}) = \sqrt{1 - [n^{(2)}(\mathbf{x})]^2 - [n^{(3)}(\mathbf{x})]^2}.$$

The measure is then replaced by

$$(4.2.4.6) \quad \int \mathcal{D}\hat{\mathbf{n}} = \int \mathcal{D}n^{(2)}(\mathbf{x}) \mathcal{D}n^{(3)}(\mathbf{x}) \frac{\Theta[1 - [n^{(2)}(\mathbf{x})]^2 - [n^{(3)}(\mathbf{x})]^2]}{\sqrt{1 - [n^{(2)}(\mathbf{x})]^2 - [n^{(3)}(\mathbf{x})]^2}},$$

wherein the first component has been written in terms of the other two components. This constraint induces anharmonic interactions between the field components. Thus, the correlation of the NLSM will differ, at least in principle, from those of a free Gaussian model, where the measure simply reads $\mathcal{D}n^{(2)}(\mathbf{x}) \mathcal{D}n^{(3)}(\mathbf{x})$ ²³. Moreover, given that the $O(3)$ -NLSM is the continuum limit of the classical Heisenberg model, it must obey the Mermin-Wagner theorem, which forbids spontaneously broken symmetries at any finite temperature (e.g. $f > 0$) for $(d = 1, 2)$ -dimensions. Now, it will be proven that the $O(n)$ -NLSM is disordered by fluctuations of the order parameter using a weak-coupling expansion in powers of f .

4.2.4.2. Naïve Weak Coupling Expansion. The $O(3)$ -NLSM can be easily generalized to a more general $O(n)$ -symmetry. Let the n orthonormal basis vectors be

$$\mathbf{e}^\alpha, \alpha = 0, \dots, n-1, \quad n^\alpha = \hat{\mathbf{n}} \cdot \mathbf{e}^\alpha.$$

Furthermore, the partition function and integration measure read

$$(4.2.4.7) \quad \mathcal{Z}_{O(n)} = \int \mathcal{D}\hat{\mathbf{n}} \exp \left[-\frac{\Lambda^{d-2}}{2f} \int d^d x \sum_{\mu\alpha} (\partial_\mu n^\alpha)^2 \right]$$

$$\mathcal{D}\hat{\mathbf{n}} = \prod_{\alpha} \mathcal{D}\hat{n}^\alpha \delta \left[1 - \sum_{\alpha=0}^{n-1} (n^\alpha)^2 \right].$$

At low temperatures, i.e. weak coupling $f \ll 1$, it is expected that the dominant configurations will be close to an ordered ground state. This requires for the breaking of the $O(n)$ -symmetry in the path integral and to choose a particular uniform ground state configuration, e.g. $\hat{\mathbf{n}} = \hat{\mathbf{e}}^0$. Then, this direction will be referred to as the longitudinal direction, with the other $n-1$ possible directions being the transverse fluctuations. The small fluctuations, e.g. those s.t. $|\hat{\mathbf{n}} - \hat{\mathbf{e}}^0| \ll 1$, are parameterized by

$$(4.2.4.8) \quad \hat{\mathbf{n}} = \hat{\mathbf{e}}^0 \sqrt{1 - \bar{\phi}^2} + \sum_{\alpha=1}^{n-1} \phi^a \hat{\mathbf{e}}^a, \quad \text{where } \bar{\phi}^2 = \sum_{\alpha=1}^{n-1} (\phi^a)^2.$$

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The expansion of the NLSM to leading order in $|\delta\hat{\mathbf{n}}|$ yields

(4.2.4.9)

$$\mathcal{Z} \approx \int \mathcal{D}\phi \exp \left(-\mathcal{S}^{(2)}[\phi] \right) [1 + \mathcal{O}(f)],$$

$$\mathcal{S}^{(2)} = \frac{\Lambda^{d-2}}{2f} \int d^d x \sum_{a,\mu} (\partial_\mu \phi^a)^2 + \mathcal{O}(\phi^3) = \frac{1}{2f\Lambda^2 N} \sum_{\mathbf{k},a} \mathbf{k}^2 \phi_{\mathbf{k}}^a \phi_{-\mathbf{k}}^a, \text{ where } \phi_{\mathbf{k}}^a = \Lambda^d \int d^d x e^{-i\mathbf{x}\cdot\mathbf{k}} \phi^a(\mathbf{x}).$$

In the previous expressions, $\mathcal{S}^{(2)}$ is the harmonic or "spin wave" energy functional. Using the weak-coupling parameterization of the small fluctuations, one can write the spin wave contribution to the local fluctuations

(4.2.4.10)

$$\langle |\delta\hat{\mathbf{n}}|^2 \rangle \approx \frac{1}{\mathcal{Z}N} \sum_{\mathbf{k},a} \int \mathcal{D}\hat{\mathbf{n}} \phi_{\mathbf{k}}^a \phi_{-\mathbf{k}}^a \exp \left(-\mathcal{S}^{(2)}[\phi] \right)$$

$$= \frac{(n-1)\Lambda^2 f}{(2\pi)^d} \lim_{\Lambda \rightarrow 0} \int_{\bar{\Lambda}}^{\Lambda} \frac{d^d x}{|\mathbf{k}|^2} = \infty \text{ in } d = 1, 2.$$

The spin fluctuations diverge for $n \geq 2$ in one and two dimensions for all $f > 0$. Thus, the naive weak-coupling expansion is inadequate for describing the ground state since the assumption of spontaneously broken symmetry is invalid. This is in accordance with Mermin-Wagner's theorem [theorem 5](#). Moreover, by Haldane's mapping, the ground state of the quantum Heisenberg antiferromagnet in one dimension should be disordered, a spin liquid, for all finite spin sizes $S < \infty$ (see for further details the "Physical Context" box in ?? 3.5.0.2).

More sophisticated methods can be used to treat infrared divergences.

4.2.4.3. "Simple" Renormalization Procedure. The infrared divergences, alluded to in the previous section, indicate the necessity to abandon the assumption of broken symmetry, and to preserve the exact $O(n)$ -symmetry of the path integral. However, it is possible to use a perturbative approach, in particular a weak-coupling approximation, to calculate the correlations in the disordered phase. In effect, this can be done by expanding the integrand around *slowly varying configurations*, $\hat{\mathbf{n}}^0(\mathbf{x})$, of the Néel field, which still need to be integrated over. The renormalization of the energy functional is carried out by sequentially eliminating the large-momenta modes, the "fast modes", leaving only the remaining "slow modes". This can be done in a procedure much similar to that of a quantum field theory, giving rise to a β -function, which will determine the correlation length.

Polyakov's renormalization technique goes as follows:

- First, the Néel field is separated into *slow* and *fast* degrees of freedom, labelled $\hat{\mathbf{n}}$, ϕ^a respectively, given by

(4.2.4.11)

$$\hat{\mathbf{n}}(\mathbf{x}) = \hat{\mathbf{n}}^0(\mathbf{x}) \sqrt{1 - \bar{\phi}^2} + \sum_{a=1}^{n-1} \phi^a \hat{\mathbf{e}}^a(\mathbf{x}), \quad \bar{\phi}^2 = \sum_{a=1}^{n-1} (\phi^a)^2,$$

with a comoving basis defined as

(4.2.4.12)

$$\hat{\mathbf{e}}^{(0)}(\mathbf{x}) = \hat{\mathbf{n}}^0(\mathbf{x}), \quad \hat{\mathbf{e}}^{(\alpha)}(\mathbf{x}) \cdot \hat{\mathbf{e}}^{(\beta)}(\mathbf{x}) = \delta^{\alpha\beta}, \alpha = 0, 1, \dots, n-1,$$

wherein Greek indices include the longitudinal $\hat{\mathbf{n}}^{(0)}$ -direction and where Latin indices are strictly restricted to the transverse directions $a, b = 1, \dots, n-1$. The relation between the coordinate systems at nearby points in space defines the **Gauge potentials**, defined as follows²⁴

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In effect, it's easy to see how the relationships between the gauge potentials and the moving orthonormal basis can be inverted,

(4.2.4.13)

$$\tilde{A}_\mu^{\alpha\beta} = \hat{\mathbf{e}}^{(\alpha)} \cdot \partial_\mu \hat{\mathbf{e}}^{(\beta)} \rightarrow \tilde{A}_\mu^{\alpha 0} = \mathbf{e}^\alpha \cdot \partial_\mu \mathbf{e}^0 \rightarrow \tilde{A}_\mu^{\alpha 0} = \mathbf{e}^\alpha \cdot \partial_\mu \hat{\mathbf{n}}^0 \Rightarrow \partial_\mu \hat{\mathbf{n}}^0 = \sum_a \tilde{A}_\mu^{\alpha 0} \mathbf{e}^\alpha$$

$$(4.2.4.14) \quad \tilde{A}_\mu^{\alpha\beta} = \hat{\mathbf{e}}^{(\alpha)} \cdot \partial_\mu \hat{\mathbf{e}}^{(\beta)}, \quad \text{which describe the variations of the moving orthonormal basis by } \partial_\mu \hat{\mathbf{n}}^0 = \sum_a \tilde{A}_\mu^{a0} \hat{\mathbf{e}}^a, \quad \partial_\mu \hat{\mathbf{e}}^a = \sum_b (\tilde{A}_\mu^{ba} \hat{\mathbf{e}}^b) - \tilde{A}_\mu^{a0} \hat{\mathbf{n}}^0$$

In the context of differential geometry, these fields are connections between principal bundles. Since the $\hat{\mathbf{e}}^a$ -vectors are orthonormal, the gauge potentials $\tilde{A}_\mu^{\alpha\beta}$ must be antisymmetric, namely $\tilde{A}_\mu^{\alpha\beta} = -\tilde{A}_\mu^{\beta\alpha}$.

- Next, an intermediate momentum scale $\tilde{\Lambda}$ is chosen, s.t.

$$0 < \tilde{\Lambda} \ll \Lambda,$$

which separates the fast and slow fields, which are defined as follows:

$$(4.2.4.15) \quad \begin{aligned} \text{Fast fields: } \phi^a(\mathbf{x}) &= \sum_{\tilde{\Lambda} \leq |\mathbf{k}| \leq \Lambda} e^{i\mathbf{k} \cdot \mathbf{x}} \phi_{\mathbf{k}}^a, \\ \text{Slow fields: } \mathbf{X}(\mathbf{x}) &= \sum_{|\mathbf{k}| \leq \tilde{\Lambda}} e^{i\mathbf{k} \cdot \mathbf{x}} \mathbf{X}_{\mathbf{k}}, \quad \mathbf{X} = \hat{\mathbf{e}}^a, \tilde{A}_\mu^{\alpha\beta}. \end{aligned}$$

- With these reparametrizations in mind, specifically those defined in [equation \(4.2.4.11\)](#), the NLSM Lagrangian now reads

$$(4.2.4.16) \quad \begin{aligned} \mathcal{L} &= \frac{\Lambda^{d-2}}{2f} \int d^d x \, \mathcal{J}[\hat{\mathbf{n}}; \hat{\mathbf{n}}^0, \phi^a] \sum_{\mu=1,d} \partial_\mu \hat{\mathbf{n}} \cdot \partial_\mu \hat{\mathbf{n}} \\ &= \frac{\Lambda^{d-2}}{2f} \int d^d x \sum_{\mu=1,d} \left[(\partial_\mu \hat{\mathbf{n}}^0) \sqrt{1 - \bar{\phi}^2} + \hat{\mathbf{n}}^0 \partial_\mu \sqrt{1 - \bar{\phi}^2} + \sum_a (\partial_\mu \phi^a) \hat{\mathbf{e}}^a + \sum_a \phi^a \partial_\mu \hat{\mathbf{e}}^a \right]^2 \\ &= \frac{\Lambda^{d-2}}{2f} \int d^d x \sum_{\mu=1,d} \left[\sqrt{1 - \bar{\phi}^2} \sum_a \tilde{A}_\mu^{\alpha\beta} \hat{\mathbf{e}}^a + (\partial_\mu \sqrt{1 - \bar{\phi}^2}) \hat{\mathbf{n}}^0 + \sum_a (\partial_\mu \phi^a) \hat{\mathbf{e}}^a + \sum_a \phi^a \sum_b (\tilde{A}_\mu^{ba} \hat{\mathbf{e}}^b - \tilde{A}_\mu^{a0} \hat{\mathbf{n}}^0) \right]^2 \\ &= \frac{\Lambda^{d-2}}{2f} \sum_{\mu ab} \left[\left(\partial_\mu \sqrt{1 - \bar{\phi}^2} - \tilde{A}_\mu^{a0} \phi^a \right)^2 + \left(\partial_\mu \phi^a + \tilde{A}_\mu^{ab} \phi^b + \tilde{A}_\mu^{a0} \sqrt{1 - \bar{\phi}^2} \right)^2 \right]^2 \\ &= \frac{\Lambda^{d-2}}{2f} \sum_{\mu ab} \left[\begin{aligned} &\left\{ \left(\partial_\mu \sqrt{1 - \bar{\phi}^2} \right)^2 - 2 \partial_\mu \sqrt{1 - \bar{\phi}^2} \times \tilde{A}_\mu^{a0} \phi^a \right\} + \tilde{A}_\mu^{a0} \tilde{A}_\mu^{b0} \phi^a \phi^b \\ &+ \left(\partial_\mu \phi^a + \tilde{A}_\mu^{ab} \phi^b \right)^2 + 2 \left\{ \left(\partial_\mu \phi^a + \tilde{A}_\mu^{ab} \phi^b \right)^2 \times \tilde{A}_\mu^{a0} \sqrt{1 - \bar{\phi}^2} \right\} \\ &+ \left(\tilde{A}_\mu^{a0} \sqrt{1 - \bar{\phi}^2} \right)^2 \end{aligned} \right] \\ &= \frac{\Lambda^{d-2}}{2f} \sum_{\mu ab} \left[\left(\partial_\mu \phi^a + \tilde{A}_\mu^{ab} \phi^b \right)^2 + \left(\tilde{A}_\mu^{a0} \phi^a + \mathcal{O}(\phi^2, f^{-1} \phi^3) \right)^2 \right] \\ &\Rightarrow \mathcal{Z} \approx \int_{\tilde{\Lambda}} \mathcal{D}\hat{\mathbf{n}} \exp \left[- \int_{\tilde{\Lambda}} d^d x \, \mathcal{L}[\hat{\mathbf{n}}^0] \mathcal{Z}^{(2)}[\hat{\mathbf{n}}^0] \left\{ 1 + \mathcal{O}(\phi^2, f^{-1} \phi^3) \right\} \right] \end{aligned}$$

where, in the ante-penultimate line in the last derivations ²⁵, the terms in brackets $\{\}$ are of cubic and higher powers of ϕ and thus, can be neglected for the time being. Here, the $\mathcal{Z}^{(2)}[\hat{\mathbf{n}}^0]$ -factor is given in terms of the fast-spin wave Lagrangian $\mathcal{L}^{(2)}[\hat{\mathbf{n}}^0, \phi]$, as follows

²⁵Note that, in the ante-penultimate line,

$$(4.2.4.17) \quad \left(\tilde{A}_\mu^{a0} \sqrt{1 - \bar{\phi}^2} \right)^2 = \tilde{A}_\mu^{a0} \tilde{A}_\mu^{b0} \phi^a \phi^b (1 - \bar{\phi}^2).$$

(4.2.4.18)

$$\mathcal{Z}^{(2)} = \int_{\Lambda} \mathcal{D}\hat{\mathbf{n}} \exp \left[- \int_{\Lambda} d^d x \mathcal{L}^{(2)}[\hat{\mathbf{n}}^0, \phi] \right], \quad \mathcal{L}^{(2)} = \frac{\Lambda^{d-2}}{2f} \sum_{\mu ab} \left[\left(\partial_{\mu} \phi^a + \tilde{A}_{\mu}^{ab} \phi^b \right)^2 + \tilde{A}_{\mu}^{a0} \tilde{A}_{\mu}^{b0} (\phi^a \phi^b - \delta_{ab} \bar{\phi}^2)^2 \right].$$

In the context of the renormalization group theory, a one-loop-level calculation is being performed, since $\mathcal{O}[f^{-1}\phi^3]$ terms are neglected entirely. Moreover, cross terms of the sort $(\tilde{A}_{\mathbf{q}}^{ab} \phi_{\mathbf{k}-\mathbf{q}}^b \tilde{A}_{-\mathbf{k}}^{a0})$ were neglected as well, since these couple slow and fast fields. By letting $\tilde{\Lambda} \ll \Lambda$, the contribution of such terms is small since they vanish for all $\mathbf{k} > 2\tilde{\Lambda}$, i.e. for most of the Brillouin zone.

The fast spin wave Lagrangian $\mathcal{L}^{(2)}$ describes the high-momenta fluctuations around the classical Néel field $\hat{\mathbf{n}}^0$. There is a certain degrees of arbitrariness in the choice of the transverse unit vectors $\hat{\mathbf{e}}^a$, given in [equation \(4.2.4.11\)](#) in the $(n > 2)$ -case. In effect, in the $(n > 2)$ -case, one can continuously rotate the transverse coordinate system in the $(n - 1)$ -dimensional case at each point in space using any $(n - 1) \times (n - 1)$ -dimensional matrix R s.t.

(4.2.4.19)

$$\begin{aligned} \phi^a &\rightarrow R^{ab} \phi^b && \text{which transforms the covariant derivative in} \\ R^T R &= \mathbf{id}_{\mathbb{R}^{(n-1) \times (n-1)}}, && \text{the first term in [equation \(4.2.4.18\)](#) as} \end{aligned} \quad \partial_{\mu} \delta_{ab} + \tilde{A}_{\mu}^{ab} \rightarrow \partial_{\mu} \delta_{ab} + \tilde{A}_{\mu}^{ab} + [(\partial_{\mu} R) R^{-1}]^{ab}$$

The manifest invariance of the fast spin wave Lagrangian $\mathcal{L}^{(2)}$, under the [equation \(4.2.4.19\)](#), is an example of a *gauge symmetry* of the fast modes Lagrangian. This Lagrangian can only depend on the derivatives of the gauge potentials \tilde{A}_{μ}^{ab} , which are slow fields. Hence, these terms must be higher-order derivatives of $\hat{\mathbf{n}}^0$ than the leading term $\mathcal{L}[\hat{\mathbf{n}}^0]$. This allows for the neglecting of the gauge potentials in $\mathcal{L}^{(2)}$. It turns out that these gauge potentials are important when expanding a slow field with finite topological density.

4.2.4.4. .

5. ADDITIONAL TOPICS

Additional Topics Additional Topics

6. VARIATIONAL WAVE FUNCTIONS AND PARENT HAMILTONIANS

Physical Context

The ground state for most non-ferromagnetic Heisenberg models are not known. Even those that are known in analytical form, such as the Bethe solution in one dimension, require numerical simulation computations for their spin correlations. Variational wave functions provide an educated guess for the ground state. Physical insight can be gained by a variational calculation, where the energy is minimized with respect to a chosen set of parameters. The variational approach go further beyond the regimes of any particular expansion scheme -semiclassical, large- N , or others.

A fundamental idea behind the Variational approach is that of the parent Hamiltonian. This is the Hamiltonian for which a particular wave function happens to be the exact ground state. Although the Hamiltonian may differ from the physical model by extra interactions, it serves to enhance the understanding of the system by bringing into light the relation between interactions and ground state correlations.

6.0.0.1. Valence Bond States. The valence bond states are variational wave functions for antiferromagnetic Heisenberg models, having been studied extensively in the context of quantum magnetism and high-temperature superconductivity. Their general form is

$$(6.0.0.1) \quad |\{c_\alpha\}, S\rangle = \sum_{\alpha} c_{\alpha} |\alpha\rangle, \quad \text{where } c_{\alpha} \text{ are the variational parameters and } |\alpha\rangle = \prod_{i,j \in \Lambda_{\alpha}} (\mathbf{a}_i^{\dagger} \mathbf{b}_j^{\dagger} - \mathbf{b}_i^{\dagger} \mathbf{a}_j^{\dagger}) |0\rangle,$$

where $\mathbf{a}_i, \mathbf{b}_i$ are Schwinger bosons on the i -th site, and where Λ_{α} is a particular configuration of the (ij) -bonds on the lattice. The condition on Λ_{α} is that precisely $2S$ bonds will emanate from each site. In certain cases, the sum on [equation \(6.0.0.1\)](#) is dominated by a finite number of configurations in the large lattice limit²⁶.

From inspection of [equation \(6.0.0.1\)](#), it is clear that all bond operators $\mathbf{a}_i^{\dagger} \mathbf{b}_j^{\dagger} - \mathbf{b}_i^{\dagger} \mathbf{a}_j^{\dagger}$ are invariant under global spin rotations, making $|\{c_{\alpha}\}, S\rangle$ a singlet of total spin. A special class of valence bond states is given by

$$(6.0.0.2) \quad |\hat{u}, S\rangle = \sum_{\alpha} \left(\prod_{(ij) \in \alpha} u_{ij} \right) |\alpha\rangle,$$

wherein there are upto $\frac{N(N-1)}{2}$ independent variational parameters u_{ij} , with N being the number of lattice sites. Furthermore, if the bond parameters are both bipartite and positive, this is

$$(6.0.0.3) \quad u_{ij} = \begin{cases} u_{ij} > 0 & \text{if } \begin{matrix} i \in A \\ j \in B \end{matrix} \\ 0 & \text{if } i, j \text{ are on the same sublattice} \end{cases},$$

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The cases where there are macroscopically many configurations were denoted resonating valence bonds states (RVB) by Anderson.

then the $|\hat{u}, S\rangle$ -states obey the Marshall sign criterion **theorem 2**, as required for the ground state of a bipartite Heisenberg antiferromagnet.

A Schwinger boson mean field state is defined as

$$(6.0.0.4) \quad |\hat{u}\rangle = \exp \left[\frac{1}{2} \sum_{ij} u_{ij} \left(\mathbf{a}_i^\dagger \mathbf{b}_i^\dagger - \mathbf{b}_i^\dagger \mathbf{a}_i^\dagger \right) \right], \text{ with } u_{ij} = -u_{ji}.$$

Such states are generated by the Schwinger boson mean field theory of the Heisenberg antiferromagnet, which will be reviewed later on. The $|\hat{u}\rangle$ -state contains contributions of different spin sizes at all sites, and therefore is not a bona fide spin state. It can be transformed, by a Bogoliubov transformation on the Schwinger bosons, into a factorizable Fock state, which is in fact a vacuum of the transformed bosons. Moreover, the correlations of the mean field states can be evaluated analytically. The valence bond state, given by [equation \(6.0.0.2\)](#), can be constructed from [equation \(6.0.0.4\)](#) via a projection. Said projection is given by the operator defined in [equation \(3.6.0.4\)](#),

$$|\hat{u}, S\rangle = \mathbf{P}_S |\hat{u}\rangle,$$

where the projected state is non-factorizable, due to correlations introduced by \mathbf{P}_S , said correlation being much easier to evaluate in $|\hat{u}\rangle$ -the Schwinger boson mean field state- than in the special class of valence bond states, $|\hat{u}, S\rangle$.

6.0.0.2. *Spin one-half States.* Consider the spinor states for $S = \frac{1}{2}$ -states, given by

$$(6.0.0.5) \quad |\uparrow_i\rangle = \mathbf{a}_i^\dagger |0\rangle, \quad |\downarrow_i\rangle = \mathbf{b}_i^\dagger |0\rangle.$$

Any bipartite valence bond configuration, those given by [equation \(6.0.0.1\)](#), that obeys the Marshall sign criterion can be written in a normalized form as a product of singlet bonds,

$$(6.0.0.6) \quad |\alpha\rangle_{S=\frac{1}{2}} = \prod_{\substack{i \in A \\ i \in B \\ (ij) \in \Lambda_\alpha}} \frac{|\uparrow_i\rangle |\downarrow_j\rangle - |\downarrow_i\rangle |\uparrow_j\rangle}{\sqrt{2}}, \quad \text{where the spin correlations are simply } \langle \mathbf{S}_i \mathbf{S}_j \rangle = \begin{cases} \frac{3}{4} & i = j \\ -\frac{3}{4} & (ij) \in \Lambda_\alpha \\ 0 & (ij) \notin \Lambda_\alpha \end{cases}.$$

Thus, if the bonds in Λ_α are short range, $|\alpha\rangle_{S=\frac{1}{2}}$ is a disordered spin liquid state. Different valence bond configurations are not orthogonal, for their finite overlap is given by

$$(6.0.0.7) \quad \langle \alpha | \beta \rangle = \prod_{\ell \in \langle \alpha | \beta \rangle} 2^{1-\frac{L_\ell}{2}} = 2^{N_L^{(\alpha\beta)} - N}, \quad \text{where, in the first equality, the product is made over all loops } \ell \text{ of } L_\ell\text{-length found in the overlap in the two configurations.}$$

For two identical bonds in α, β , a length $L = 2$ loop is produced. In the second identity of the preceding equation, N_L is the total number of loops and N is the number of sites on the lattice. Since $|\alpha\rangle$ contains one bond per site, and the coordination number of all simple lattices is equal to or larger than two, the $|\alpha\rangle$ -states break lattice translational symmetry. This symmetry can be restored in [equation \(6.0.0.1\)](#) by summing over α .

From now, the discussion will be restricted to nearest-neighbour bonds or dimers, for the one and two-dimensional cases separately.

6.0.0.3. *The Majumdar-Ghosh Hamiltonian.* Majumdar and Ghosh introduced the Hamiltonian,

$$(6.0.0.8) \quad \mathbf{H}^{\text{MG}} = \frac{4|K|}{3} \sum_{i=1}^N \left(\mathbf{S}_i \cdot \mathbf{S}_{i+1} + \frac{1}{2} \mathbf{S}_i \cdot \mathbf{S}_{i+2} \right) + \frac{1}{N}, \text{ where } \mathbf{S}_{N+1} = \mathbf{S}_1,$$

and where i labels the sites of one-dimensional chain with even number of sites. It can be proved that this Hamiltonian is the parent Hamiltonian of the two-dimer states,

$$(6.0.0.9) \quad |d\rangle_{\pm} = \prod_{n=1}^{\frac{N}{2}} \frac{|\uparrow_{2n}\rangle |\downarrow_{2n\pm 1}\rangle - |\downarrow_{2n}\rangle |\uparrow_{2n\pm 1}\rangle}{\sqrt{2}}, \text{ where } \mathbf{H}^{\text{MG}} |d\rangle_{\pm} = 0,$$

and that all other eigenenergies are positive. Thus, should this statement be true, $|d\rangle_{\pm}$ span the ground state manifold of the Majumdar-Ghosh Hamiltonian.

This Hamiltonian includes antiferromagnetic interactions between next-nearest neighbours, which partially frustrate the nearest neighbour correlations. Thus, the ground state can be expected to be more disordered than the pure nearest-neighbour model. Indeed, the correlations of the Bethe wave function of the nearest-neighbour model decay as an inverse power of distance, while [equation \(6.0.0.6\)](#) states that the dimer state correlations vanish beyond a single lattice constant.

Proof. Consider a total spin triad of spins at sites $(i-1, i, i+1)$ is given by

$$(6.0.0.10) \quad \mathbf{J}_i = \mathbf{S}_{i-1} + \mathbf{S}_i + \mathbf{S}_{i+1},$$

whose squares has eigenvalues $J(J+1)$, where $J = \frac{1}{2}, \frac{3}{2}$. The main pinning point of this proof is that the Majumdar-Ghosh Hamiltonian can be written as a sum of projection operators,

$$(6.0.0.11) \quad \mathbf{H}^{\text{HG}} = |K| \sum_i \mathcal{P}_{\frac{3}{2}}(i-1, i, i+1), \text{ where } \mathcal{P}_{\frac{3}{2}}(i-1, i, i+1) = \frac{1}{3} \left(\mathbf{J}_i^2 - \frac{3}{4} \right) \\ = \frac{1}{2} + \frac{2}{3} (\mathbf{S}_{i-1} \cdot \mathbf{S}_i + \mathbf{S}_{i-1} \cdot \mathbf{S}_{i+1} + \mathbf{S}_i \cdot \mathbf{S}_{i+1}).$$

Since \mathbf{J}_i^2 's eigenvalues are $J(J+1)$, this $\mathcal{P}_{\frac{3}{2}}(i-1, i, i+1)$ -operator clearly annihilates any state with total spin $J = \frac{1}{2}$ of the triad $(i-1, i, i+1)$. Dimer states, given by [equation \(6.0.0.9\)](#), do not contain states with total $\mathbf{J}^z > \frac{1}{2}$ of any three sites, since two of the three spins have cancelling \mathbf{S}^z quantum numbers.

Consider a triad with $\mathbf{J}^z = \frac{1}{2}$ but $J = \frac{3}{2}$ and a global $\mathfrak{su}(2)$ -rotation on $|d\rangle$. This will admix, by application of \mathbf{J}_i^+ , a component of $\mathbf{J}_i^z = \frac{3}{2}$ into wave function, which contradicts the rotational invariance of $|d\rangle$ ²⁷. Then, there are no triads with total spin $J > \frac{1}{2}$ and therefore each $\mathcal{P}_{\frac{3}{2}}(i-1, i, i+1)$ operator annihilates the dimer states, $|d\rangle_{\pm}$. And given that the $|K| \mathcal{P}_{\frac{3}{2}}(i-1, i, i+1)$ -operators are non-negative operators, for these are projection operators, $|d\rangle_{\pm}$ span the ground state manifold of the Majumdar-Ghosh Hamiltonian.

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6.0.0.4. Square Lattice RVB States. On the square $S = \frac{1}{2}$ -lattice, the resonating valence bond states, given by [equation \(6.0.0.1\)](#), fulfill all the requirements of Marshall's theorems, [theorem 2](#) and [theorem 3](#), for the quantum antiferromagnet²⁸. For RVB states, it is not easy at all to evaluate correlation for the $|\alpha\rangle$ -components are not orthogonal. The number of different bond coverings on the square lattice increases exponentially with the number of sites. Moreover, M.E. Fisher has calculated that the number of dimer configurations on the square lattice grows

²⁷Remember that dimer states are given in terms of Schwinger boson creation operators acting on the system's ground state. These Schwinger boson creation operators, as stated in [equation \(3.6.0.7\)](#), transform as $SU(2)$ -vectors.

²⁸

Historical Context

These states were proposed as candidates for spin liquid ground states by P.W. Anderson. Although the nearest-neighbour antiferromagnets, on the square lattice is ordered at zero temperature, these states are very popular in the context of doped antiferromagnets and high-temperature superconductivity.

exponentially with the lattice size.

Numerical Monte Carlo simulations by Liang, Doucot, and Anderson on finite lattices suggest that the RVB state has no long-range order for bonds that decay at least as rapidly as

$$u_{ij} \approx |\mathbf{x}_i - \mathbf{x}_j|^{-p}, \text{ with } p \geq 5.$$

The RVB states can thus be used as variational ground states for both ordered and disordered phases, making these appealing candidates for studying the transition from the Néel antiferromagnet to a paramagnetic phase.

6.0.0.5. Valence Bond Solids and AKLT Models. The valence bond solids are defined as follows

$$(6.0.0.12) \quad |\Psi^{\text{VBS}}\rangle = \prod_{\langle ij \rangle} \left(\mathbf{a}^\dagger \mathbf{b}^\dagger - \mathbf{b}^\dagger \mathbf{a}^\dagger \right)^M |0\rangle, \quad \text{where } \langle ij \rangle \text{ are all the nearest neighbour bonds of the lattice} \\ \text{and where } M \in \mathbb{Z} \text{ such that } M = \frac{2S}{z} \text{ where } z \text{ is the lattice coordination number}.$$

From the definition it is clear that S is restricted to values depending on the lattice structure. For instance for $(d = 1)$ -lattices, $S = 1, 2, \dots$ and for $(d = 2)$ -lattices, $S = 2, 4, \dots$.

Affleck, Kennedy, Lieb and Tasaki (AKLT) constructed the parent Hamiltonians for the VBS states, given by

$$(6.0.0.13) \quad \mathbf{H}^{\text{AKLT}} = \sum_{\langle ij \rangle} \sum_{J=2S-M+1}^{2S} K_J \mathcal{P}_J(ij), \quad K_J \geq 0.$$

The bond-projector $\mathcal{P}_J(ij)$ projects the bond spin $\mathbf{J}_{ij} = \mathbf{S}_i + \mathbf{S}_j$ onto the subspace of J -magnitude. Any power of $m = 0, 1, 2, \dots$ of $\mathbf{S}_i \cdot \mathbf{S}_j$ can be written in terms of powers of \mathbf{J}_{ij} and expanded as a linear combination of bond projector operators

$$(6.0.0.14) \quad (\mathbf{S}_i \cdot \mathbf{S}_j)^m = \sum_{J=0}^{2S} \left[\frac{1}{2} J(J+1) - S(S+1) \right]^m \mathcal{P}_J(ij),$$

this relation being useful for inverting the bond-projection operators as polynomials of $\mathbf{S}_i \cdot \mathbf{S}_j$.

Lemma 5. *The valence bond solids, defined in [equation \(6.0.0.12\)](#), are the eigenstates of the AKLT-Hamiltonian,*

$$(6.0.0.15) \quad \mathbf{H}^{\text{AKLT}} |\Psi^{\text{VBS}}\rangle = 0,$$

which holds if $|\Psi^{\text{VBS}}\rangle$ has no component with bond spin $J_{ij} > 2S - M$, for any (ij) .

Proof. In effect, consider the contribution to $|\Psi^{\text{VBS}}\rangle$ made by the term with maximally possible number of creation Schwinger operators, \mathbf{a}^\dagger in particular, on a particular bond (ij) ,

$$(6.0.0.16) \quad \dots (\mathbf{a}_i^\dagger)^{2S-M} \left(\mathbf{a}^\dagger \mathbf{b}^\dagger - \mathbf{b}^\dagger \mathbf{a}^\dagger \right)^M (\mathbf{a}_j^\dagger)^{2S-M} \dots$$

By counting the power of the \mathbf{a}^\dagger -operators minus the number of \mathbf{b}^\dagger -operators, the maximum eigenvalues of \mathbf{J}_{ij}^z acting on [equation \(6.0.0.12\)](#) is

$$(6.0.0.17) \quad \mathbf{J}_{\text{max}}^z = 2S - M$$

TODO: continue it.

7. GROUND STATES AND EXCITATIONS

Physical Context

The focus of previous sections has been primarily on the ground state of the Heisenberg model. Knowledge of the exact ground state (or in lack of the exact ground state, a good variational state) allows for the calculation of the equal time spin correlations at zero temperature. Experiments, however, can only measure dynamical responses at finite frequencies and finite temperatures. As it turns out, these dynamical correlations depend on the excited states and energies.

The ground state correlations can be used to construct certain approximate low-lying excitations. This approach is called the **single mode approximation** (SMA). Furthermore, the Heisenberg antiferromagnet provides an opportunity to demonstrate the SMA for a system with a highly-correlated ground state.

Historical Context

The SMA was originally devised by Bijl and Feynmann to determine the photon-roton dispersion curve in superfluid He4.

In this section, the information about spin excitations and correlations is embodied primarily by the dynamical structure factor and some derived quantities,

(7.0.0.1)

$$\begin{aligned}\mathcal{S}(\mathbf{q}, \omega) &= \frac{1}{N} \int_{\mathbb{R}} dt e^{i\omega t} \sum_{ij} e^{i\mathbf{q} \cdot (\mathbf{x} - \mathbf{x}_j)} \langle \mathbf{S}_i^z(t) \mathbf{S}_j^z(0) \rangle = \frac{2\pi}{NZ} \sum_{\alpha, \beta} e^{-\beta E_\alpha} \langle \alpha | \mathbf{S}_{\mathbf{q}}^z | \beta \rangle \langle \beta | \mathbf{S}_{-\mathbf{q}}^z | \alpha \rangle \delta(\omega + E_\alpha - E_\beta), \\ \mathcal{S}(\mathbf{q}) &= \int_{\mathbb{R}} \frac{d\omega}{2\pi} \mathcal{S}(\mathbf{q}, \omega) = \frac{1}{N} \langle \mathbf{S}_{\mathbf{q}}^z \mathbf{S}_{-\mathbf{q}}^z \rangle, \quad \text{which is the equal-time correlation function} \\ \mathcal{F}(\mathbf{q}) &= \frac{1}{N} \left\langle \left[\mathbf{S}_{\mathbf{q}}^z, [\mathbf{S}_{-\mathbf{q}}^z, \mathbf{H}] \right] \right\rangle = \int_{\mathbb{R}} \frac{d\omega}{2\pi} \omega \mathcal{S}(\mathbf{q}, \omega), \quad \begin{array}{l} \text{which is the double commutator function, previously} \\ \text{encountered in the proof of the Mervin-Wagner} \\ \text{theorem, \textcolor{red}{theorem 5}.} \end{array}\end{aligned}$$

Note that, even though $\mathcal{F}(\mathbf{q})$ is a static ground state expectation value, it contains dynamical information through the Hamiltonian. These relations thus show that the $\mathcal{S}(\mathbf{q})$ and $\mathcal{F}(\mathbf{q})$ describe the average and the first frequency moment of $\mathcal{S}(\mathbf{q}, \omega)$, respectively. Hence, a characteristic frequency for spin excitations at \mathbf{q} -momentum is defined as

$$(7.0.0.2) \quad \bar{\omega}_{\mathbf{q}} = \frac{\mathcal{F}(\mathbf{q})}{\mathcal{S}(\mathbf{q})}.$$

7.0.0.1. The Single Mode Approximation. Consider a quantum system with ground state $|0\rangle$. By translational invariance, all excitations may be labelled by the lattice \mathbf{q} -momentum. A "single mode" is constructed then as

$$(7.0.0.3) \quad |\mathbf{q}\rangle = \mathbf{S}_{\mathbf{q}}^z |0\rangle.$$

The single mode state approximates a true excitation of the system if $\mathcal{S}(\mathbf{q}, \omega)$ at $T = 0$ is sharply peaked about $\omega = \bar{\omega}_{\mathbf{q}}$, i.e.

$$(7.0.0.4) \quad \mathcal{S}(\mathbf{q}, \omega) \approx 2\pi \mathcal{S}(\mathbf{q}) \delta(\omega - \omega_{\mathbf{q}}),$$

which will be precise if

$$(7.0.0.5) \quad (\mathbf{H} - E_0) |\mathbf{q}\rangle = \omega_{\mathbf{q}} |\mathbf{q}\rangle.$$

In general, even if this holds exactly, the single mode states are only a small subset of all excitations -i.e. those connected to the ground state via $\mathbf{S}_{\mathbf{q}}^z$. If their products are also elementary excitations, they can be considered as weakly interacting elementary excitations. Thus, their energies approximately determine the system's low-frequency response and low-temperature thermodynamics. In ?? 3.5.0.2, it was mentioned that a non-zero gap in the excitation spectrum in the Heisenberg model implies the absence of long-range order at zero temperature (ie. in said conditions, the ground state of the Heisenberg model must be disordered). In this case, the dynamical structure factor must vanish below the gap frequency

$$(7.0.0.6) \quad \mathcal{S}(\mathbf{q}, \omega) = 0, \text{ for } \omega \in (0, \Delta).$$

Then, the single-mode energy $\bar{\omega}_{\mathbf{q}}$ is an upper bound on the lowest-excitation energy at \mathbf{q} -momentum,

$$(7.0.0.7) \quad \min_{\alpha} E_{\alpha}(\mathbf{q}) \leq \bar{\omega}_{\mathbf{q}}.$$

However, the existence of an energy gap Δ^{SMA} in the single-mode spectrum,

$$(7.0.0.8) \quad \Delta^{\text{SMA}} \leq \bar{\omega}_{\mathbf{q}}, \forall \mathbf{q}.$$

does not imply the existence of a true gap for the exact eigenenergies E_{α} .

7.0.0.2. Goldstone's Theorem in Condensed Matter Physics. The [equation \(7.0.0.7\)](#)'s inequality serves not to prove the existence of a gap in the excitation spectrum, but rather to disprove its existence ie. it can be used to prove gaplessness. Thus, the lattice-version of the Goldstone theroem can be used on the Heisenberg model.

Physical Context

The essence of Goldstone's theorem is that for a Hamiltonian with short-range interactions and no gauge fields, spontaneously broken continuous symmetry implies the existence of low-energy excitations called Goldstone modes. If the ground state has $\bar{\mathbf{q}}$ -momentum, the energy of the Goldstone mode vanishes as $\mathbf{q} \rightarrow \bar{\mathbf{q}}$.

Example

For example,

- $\bar{\mathbf{q}} = 0$ for the ferromagnet,
- $\bar{\mathbf{q}} = \pi$ for the Néel antiferromagnet in the cubic lattice.

In relativistic field theories, Goldstone modes represent massless particles. In condensed matter physics on the other hand, Goldstone modes appear in many systems, eg.

- acoustic phonons in solids that break translational symmetry,
- spin waves in $O(n)$ Heisenberg models, or $n > 1$ with n the dimension of the spins.

Consider the short-range Heisenberg Hamiltonian, given in [equation \(3.5.0.8\)](#), which satisfies the conditions of the Mervin-Wagner theorem, [theorem 5](#). Then, the following theorem holds

Theorem 6 *Goldstone's Theorem*

If the spin correlation diverges at some $\bar{\mathbf{q}}$ -wave vector, ie.

$$\lim_{\mathbf{q} \rightarrow \bar{\mathbf{q}}} \mathcal{S}(\mathbf{q}) \rightarrow \infty,$$

then there exists a Goldstone mode labelled by the \mathbf{q} -momentum, whose energy $E(\mathbf{q})$ vanishes at $\bar{\mathbf{q}}$, ie.

$$\lim_{\mathbf{q} \rightarrow \bar{\mathbf{q}}} E(\mathbf{q}) = 0.$$

Proof. In effect, consider the characteristic frequency for spin excitations at \mathbf{q} -momentum, defined in [equation \(7.0.0.2\)](#) and [equation \(7.0.0.1\)](#). The double commutator, which appears in both relations, has an upper bound given by [equation \(3.5.0.13\)](#), which in turn induces a lower bound on the characteristic frequency. Thus, for $h = 0$,

$$(7.0.0.9) \quad \mathcal{F}(\mathbf{q}) \leq \cancel{h\mathcal{M}_{\mathbf{q}}} + S(S+1)\bar{J}|\mathbf{q}|^2 \implies \bar{\omega}_{\mathbf{q}} \leq \frac{2S(S+1)\bar{J}|\mathbf{q}|^2}{\mathcal{S}(\mathbf{q})}.$$

Then, consider the case in which $\lim_{\mathbf{q} \rightarrow \bar{\mathbf{q}}} \mathcal{S}(\mathbf{q}) \rightarrow \infty$. Furthermore, [equation \(7.0.0.7\)](#) induces an upper bound on the lowest-level excitation energy at a fixed $\bar{\mathbf{q}}$ -momentum

$$\min_{\alpha} E_{\alpha}(\mathbf{q}) \leq \bar{\omega}_{\mathbf{q}} \rightarrow \infty \text{ which implies that } \lim_{\mathbf{q} \rightarrow \bar{\mathbf{q}}} E(\mathbf{q}) = 0.$$

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Corollary 3. *If the ground state has broken symmetry, there exists Goldstone modes, as in [theorem 6](#).*

Proof. This statement follows directly from [section 3.5](#)'s red box titled "Spontaneously broken symmetry and true long-range order", in which it is proven that an spontaneously broken symmetry implies true long-range order and the divergence of $\mathcal{S}(\mathbf{q})$ as $N \rightarrow \infty$.

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Physical Context

It is important to note that the converse of the Goldstone theorem is false, ie. the existence of gapless excitations does not imply true long-range order. Notable counterexamples are,

- the Fermi gas, which has gapless particle-hole excitations but no long-range order,
- and the $S = \frac{1}{2}$ -one-dimensional Heisenberg model has gapless excitations which vanish at $q = \{0, \pi\}$ with $\omega_q \propto |\sin q|$ but its correlation decay algebraically at long distances.
- The Haldane gap in the $S = 1$ -AKLT models.

The ground states of the AKLT models are valence bond solids, with single mode energies given by

$$\bar{\omega}_q \geq 0.370K.$$

As discussed earlier, although the single-mode spectrum has a gap of magnitude $0.370K$, this doesn't imply that a gap actually exists in the exact spectrum. However, numerical simulations for the $S = 1$ -AKLT model support the existence of a gap of magnitude $\Delta = 0.350K$. The same simulation for the standard Heisenberg model shows the existence of a gap of size $0.325K$. Since the AKLT model has a gap in the SMA-spectrum, this numerical evidence suggests that the additional biquadratic terms $\frac{K}{3}(\mathbf{S}_i \cdot \mathbf{S}_j)^2$ do not drastically alter neither the ground state correlations nor the low lying excitation of the $S = 1$ -Heisenberg model in one dimension. Furthermore, the Haldane gap survives in the thermodynamical limit, this being a general feature of integer spin quantum antiferromagnets in low dimensions.

8. A REVIEW OF QUANTUM FIELD THEORY

A Review of Quantum Field Theory

8.1. Motivation for the derivation of a Quantum Field Theoretic Approach to Condensed Matter Theory.

Consider a quantum many-body system governed by a time-dependent Hamiltonian $\mathbf{H}(t)$, s.t. in the distant past the system was in a state specified by a many-body density matrix $\rho(-\infty)$, whose identity is non-important. It may be thought of as the equilibrium density matrix associated with the Hamiltonian $\mathbf{H}(-\infty)$. Without any loss in generality, one may consider that $\mathbf{H}(-\infty)$ describes a free (Gaussian) theory, i.e. a theory of non-interacting particles. The interactions may be **adiabatically** switched on to reach their actual physical strength sometime prior to the observation time. This idea of adiabaticity is paramount to derive useful analytical methods in Quantum Field Theory.

8.1.0.1. Closed Time Contour Evolution. The system's density matrix evolves according to the von Neumann equation

$$\begin{aligned} \frac{\partial}{\partial t} \rho(t) &= -i[\mathbf{H}(t), \rho(t)], \text{ with formal solution } \rho(t) = \mathcal{U}(t, \infty) \rho(-\infty) \mathcal{U}^\dagger(t, \infty) \\ &= \mathcal{U}(t, \infty) \rho(-\infty) \mathcal{U}(-\infty, t), \\ (8.1.0.1) \quad &\partial_t \mathcal{U}(t, t') = -i\mathbf{H}(t) \mathcal{U}(t, t') \\ &\text{with the evolution operators obeying} \\ &\partial_{t'} \mathcal{U}(t, t') = i\mathcal{U}(t, t') \mathbf{H}(t') \end{aligned}$$

Given that the Hamiltonian operators at different times need not to commute, the \mathcal{U} -evolution operator must be understood as an infinite product of incremental evolution operators with instantaneous locally constant Hamiltonians, i.e.

$$\begin{aligned} \mathcal{U}(t, t') &= \lim_{N \rightarrow \infty} e^{i\mathbf{H}(t-\delta t) \delta t} e^{i\mathbf{H}(t-2\delta t) \delta t} \dots e^{i\mathbf{H}(t-N\delta t) \delta t} e^{i\mathbf{H}(t') \delta t}, \quad \delta t = \frac{t-t'}{N}, \\ (8.1.0.2) \quad &= \mathcal{T} \exp \left(-i \int_{t'}^t dt \mathbf{H}(t) \right). \end{aligned}$$

Physical observables at t -time can be computed as

$$(8.1.0.3) \quad \langle \mathbf{O}(t) \rangle \equiv \frac{\text{Tr} \rho(t) \mathbf{O}(t)}{\text{Tr} \rho(t)} \text{Tr} \left[\mathcal{U}(-\infty, t) \mathbf{O} \mathcal{U}(t, -\infty) \rho(-\infty) \right],$$

where the trace is performed over the many-body Hilbert space. The previous expression, in its last form, reads as an evolution from $t = -\infty$ (wherein the initial state is known), towards t (wherein the observable is calculated), and then back to $t = -\infty$. Hence, calculation of an observable implies evolving the initial state both forward and backward.

Such forward-backward evolution can be avoided, in the equilibrium quantum field theory, with a special trick. In the zero-temperature equilibrium formalism, the expectation values are taken w.r.t. the ground state of the interacting many-body system, $|\text{GS}\rangle$. This state is obtained from the known, and simple ground state of the corresponding non-interacting system, $|0\rangle$, by action with the evolution operator of the fully-interacting theory, e.g.

$$|\text{GS}\rangle = \mathcal{U}(t, -\infty) |0\rangle.$$

This mapping holds if and only if the only allowed time-dependence for the Hamiltonian is an **adiabatic switching** of the interactions on and off the distant past and distant future, respectively, since the idea is to only deal with systems in equilibrium. The notion of adiabatic-switching will be discussed in detail in **The concept of Adiabaticity in Quantum Mechanics**.

If such a mapping holds, then it is clear that

$$\langle \text{GS} | \mathbf{O} | \text{GS} \rangle = \langle 0 | \mathcal{U}(-\infty, t) \mathbf{O} \mathcal{U}(t, -\infty) | 0 \rangle.$$

The trick lies in arguing that

$$\mathcal{U}(+\infty, -\infty) | 0 \rangle = e^{iL} | 0 \rangle, \quad \langle 0 | \mathcal{U}(+\infty, -\infty) = \langle 0 | e^{iL}.$$

This is, the evolution of the non-interacting ground state -upon adiabatic switching-on of the interactions and subsequent adiabatic switching-off- brings the system back into the state, upto a phase factor e^{iL} . This statement is based on the hypothesis that the adiabatic perturbation keeps the system in its evolving ground state at all times. If so, the normalization condition $\langle 0 | 0 \rangle = 1$, implies that the only possible change is the phase of the non-interacting ground state $e^{iL} = \langle 0 | \mathcal{U}(+\infty, -\infty) | 0 \rangle$. Similarly, it follows that

$$\begin{aligned} \langle \text{GS} | \mathbf{O} | \text{GS} \rangle &= \langle 0 | \mathcal{U}(-\infty, t) \mathbf{O} \mathcal{U}(t, -\infty) | 0 \rangle \\ &= e^{-iL} \langle 0 | e^{iL} \mathcal{U}(-\infty, t) \mathbf{O} \mathcal{U}(t, -\infty) | 0 \rangle \\ (8.1.0.4) \quad &= e^{-iL} \langle 0 | \mathcal{U}(+\infty, -\infty) \mathcal{U}(-\infty, t) \mathbf{O} \mathcal{U}(t, -\infty) | 0 \rangle | 0 \rangle \\ &= \frac{\langle 0 | \mathcal{U}(+\infty, t) \mathbf{O} \mathcal{U}(t, -\infty) | 0 \rangle}{\langle 0 | \mathcal{U}(+\infty, -\infty) | 0 \rangle}, \quad \text{where } \mathcal{U}(+\infty, -\infty) \mathcal{U}(-\infty, t) = \mathcal{U}(+\infty, t). \end{aligned}$$

The result of this procedure is that one needs to consider only the forward evolution. Indeed, the numerator in the last expression calls for evolving the non-interacting ground state $|0\rangle$ from the distant past to the observation time, wherein the observable acts, and then proceeding towards the distant futures, where the overlap with the same known state $|0\rangle$ is evaluated. This formalism will work in the finite-temperature equilibrium formalism, as will be seen in **Finite Temperature Green Functions**, wherein the equilibrium density matrix is $\rho \propto e^{-\beta \mathbf{H}}$ is treated as an evolution operator in imaginary time $\tau \in [0, \beta]$. The observables are then evaluated at imaginary times τ_1, τ_2, \dots and the result must be analytically continued back to the real-time axis. Since the adiabatic switching of the interaction does not drive the system out of equilibrium, a statement similar to **equation (8.1.0.4)** holds as well, in which one needs only to describe the forward evolution, albeit along the finite time interval in the imaginary direction.

Denominator

8.2. Green Functions for Many-Body systems. Green function techniques play a fundamental role in the treatment of many-body systems, for many properties of a many body system -eg. ground state energy, excitation spectra, response functions, electric and magnetic susceptibilities and thermodynamic quantities- can be derived from many-body Green functions.

Consider the ground state $|\Omega\rangle$ of a system. At some time t' , a perturbation is enacted upon the system, said perturbation represented by a local operator $\mathbf{O}^\dagger(\mathbf{x}, t)$, creating an initial state $\mathbf{O}^\dagger(\mathbf{x}', t') |\Omega\rangle$. Thus, the system has been deviated from the initial some by some local excitations, a disturbance, and then this disturbance will evolve in space and time. Such a scenario describes typical local measurements done on a quantum many-particle system. Since an explicit time dependence is necessary for this formalism, the Heisenberg picture is useful. To determine physically interesting quantities, the quantum-mechanical amplitude to find the system in state $\mathbf{O}^\dagger(\mathbf{x}', t') |\Omega\rangle$ at time $t > t'$ namely,

$$\mathcal{A} = \langle \Omega | \mathbf{O}(\mathbf{x}, t) \mathbf{O}^\dagger(\mathbf{x}', t') | \Omega \rangle.$$

Many quantum-mechanical operators may be written in bilinear forms of the field operators themselves, $\hat{\Psi}(\mathbf{x}, t)$. However, it turns out that correlation functions of the field operators are related to experiments such as photoemission and tunnelling spectroscopies, which involve extraction or injection of a single electron. These field correlation functions can be identified as propagators which satisfy the Green function equation.

8.2.1. The One-Particle Green Function of Many-Body Systems.

8.2.1.1. Non-interacting Particle Propagator. Consider a particle in free space described by a single-particle time-independent Hamiltonian \mathbf{H}_1 . Its eigenstates and eigenenergies are

$$\mathbf{H}_1 |\phi_n\rangle = \varepsilon_n |\phi_n\rangle.$$

In general, if a particle is in the $|\phi\rangle_i$ -state, it will remain in that state forever. Instead, consider a system in a generic state $|\psi_{\text{trial}}\rangle$. If the trial state is created at time t' , the wave-function at a later time t is given by

$$|\psi(t)\rangle = e^{-\frac{i\mathbf{H}(t-t')}{\hbar}} |\psi_{\text{trial}}\rangle = \sum_n |\psi_n\rangle e^{-\frac{i\varepsilon_n(t-t')}{\hbar}} \langle\phi_n|\psi_{\text{trial}}\rangle.$$

Then, at time t , the probability amplitude that a measurement would find the particle at position \mathbf{x} can be found as

(8.2.1.1)

$$\begin{aligned} \langle\mathbf{x}|\psi(t)\rangle \Theta(t-t') &= \int_{\mathbb{R}^3} d\mathbf{x}' \langle\mathbf{x}| e^{-\frac{i\mathbf{H}(t-t')}{\hbar}} |\mathbf{x}'\rangle \langle\mathbf{x}'|\psi_{\text{trial}}\rangle \Theta(t-t') \\ &= \int_{\mathbb{R}^3} d\mathbf{x}' \sum_{n \in \mathbb{N}} \langle\mathbf{x}|\phi_n\rangle e^{-\frac{i\varepsilon_n(t-t')}{\hbar}} \langle\phi_n|\mathbf{x}'\rangle \langle\mathbf{x}'|\psi_{\text{trial}}\rangle \Theta(t-t') \\ &= i \int_{\mathbb{R}^3} d\mathbf{x}' \mathfrak{G}^{\text{R}}(\mathbf{x}t, \mathbf{x}'t') \psi_{\text{trial}}(\mathbf{x}', t'), \end{aligned}$$

where $\mathfrak{G}^{\text{R}}(\mathbf{x}t, \mathbf{x}'t') = -i\Theta(t-t') \langle\mathbf{x}| e^{-\frac{i\mathbf{H}(t-t')}{\hbar}} |\mathbf{x}'\rangle$
is the retarded Green function, or propagator,
in coordinate representation.

Note that the time translational invariance is apparent due to the Hamiltonian's dependence on $\tau = t - t'$. Therefore, if the Green function $\mathfrak{G}^{\text{R}}(\mathbf{x}t, \mathbf{x}'t')$ is known, then it can be used to calculate the evolution of any initial state. Interesting information can be extracted from the propagator.

- The bracket $\langle\phi_n|\mathbf{x}\rangle = \langle\phi_n|\psi_n^\dagger(\mathbf{x})|0\rangle$ is just the probability amplitude of forcing the system to collapse into the eigenstate $|\psi_n\rangle$ when a particle is inserted at position \mathbf{x} and its energy is measured immediately afterwards.
- Its resolvent can also be found. From [equation \(8.2.1.1\)](#), note that the time evolution is a superposition of waves propagating with different energies. In principle, it could be used to determine the eigenspectrum. This can be done via a Fourier transform to the frequency domain²⁹,

$$\begin{aligned} \mathfrak{G}^{\text{R}}(\mathbf{x}, \mathbf{x}'; \omega) &= \mathcal{F}[\mathfrak{G}^{\text{R}}(\mathbf{x}t, \mathbf{x}'t')](\mathbf{x}, \mathbf{x}'; \omega) \\ (8.2.1.2) \quad &= \int dt e^{i(\omega+i\eta)\tau} \langle\mathbf{x}| e^{-\frac{i\mathbf{H}\tau}{\hbar}} |\mathbf{x}'\rangle = \left\langle \mathbf{x} \left| \frac{1}{\omega - \mathbf{H} + i\eta} \right| \mathbf{x}' \right\rangle, \end{aligned}$$

where the Sokhotski-Plemelj formula has been used. This is just the resolvent whose poles constitute the quantum system's spectrum.

Operationally, we can devise a gedankenexperiment in which the particle is initially inserted at \mathbf{x}' -position and picked up at \mathbf{x} -position after some time t . When performing this recipe with adequate resolution and for a sufficiently large number of different positions and elapsed times, the Fourier transform of the record data would provide the full eigenvalue spectrum. Such an experiment would give complete information on our particle.

This scenario thus reveals that all observables can be expressed in terms of the Green function and thus avoid the explicit use of the wavefunctions.

²⁹Note that the Fourier transform is ill-defined, given the oscillatory behaviour for $t \rightarrow \infty$. Therefore, according to the Sokhotski-Plemelj formula, the new function

$$e^{-\eta|t|} \mathfrak{G}^{\text{R}}(\mathbf{x}, \mathbf{x}'; \omega), \text{ with } \eta > 0,$$

is considered.

Differential Equation for the Green Function

Consider [equation \(8.2.1.1\)](#). The LHS's time derivative will yield

(8.2.1.3)

$$\begin{aligned}
 \langle \mathbf{x} | \psi(t) \rangle \Theta(t - t') &= i \int_{\mathbb{R}^3} d\mathbf{x}' \mathfrak{G}^R(\mathbf{x}t, \mathbf{x}'t') \psi_{\text{trial}}(\mathbf{x}', t') \\
 i\delta(t - t') \psi(\mathbf{x}, t) + i\Theta(t - t') \frac{\partial \psi(\mathbf{x}, t)}{\partial t} &= i \frac{\partial}{\partial t} \left(i \int_{\mathbb{R}^3} d\mathbf{x}' \mathfrak{G}^R(\mathbf{x}t, \mathbf{x}'t') \psi(\mathbf{x}', t') \right) \\
 i\delta(t - t') \int_{\mathbb{R}^3} d\mathbf{x}' \delta(\mathbf{x} - \mathbf{x}') \psi(\mathbf{x}, t) + \mathbf{H} \Theta(t - t')(\mathbf{x}) \frac{\partial \psi(\mathbf{x}, t)}{\partial t} &= i \frac{\partial}{\partial t} \left(i \int_{\mathbb{R}^3} d\mathbf{x}' \mathfrak{G}^R(\mathbf{x}t, \mathbf{x}'t') \psi(\mathbf{x}', t') \right) \\
 i\delta(t - t') \int_{\mathbb{R}^3} d\mathbf{x}' \delta(\mathbf{x} - \mathbf{x}') \psi(\mathbf{x}, t) + \mathbf{H} \left(i \int_{\mathbb{R}^3} d\mathbf{x}' \mathfrak{G}^R(\mathbf{x}t, \mathbf{x}'t') \psi_{\text{trial}}(\mathbf{x}', t') \right) &= i \frac{\partial}{\partial t} \left(i \int_{\mathbb{R}^3} d\mathbf{x}' \mathfrak{G}^R(\mathbf{x}t, \mathbf{x}'t') \psi(\mathbf{x}', t') \right) \\
 \int_{\mathbb{R}^3} d\mathbf{x}' \left(i \frac{\partial}{\partial t} - \mathbf{H} \right) \mathfrak{G}^R(\mathbf{x}t, \mathbf{x}'t') \psi(\mathbf{x}', t') &= i\delta(t - t') \int_{\mathbb{R}^3} d\mathbf{x}' \delta(\mathbf{x} - \mathbf{x}') \psi(\mathbf{x}, t) \\
 \Rightarrow \left(i \frac{\partial}{\partial t} - \mathbf{H} \right) \mathfrak{G}^R(\mathbf{x}t, \mathbf{x}'t') &= \delta(t - t') \delta(\mathbf{x} - \mathbf{x}'),
 \end{aligned}$$

which is just the definition of the Green function for the Schrödinger equation in differential equation form

8.2.2. One-particle Green Function of Many-Body Systems. The notion of a Green function can be applied to situations in which a particle is injected into a system containing many similar interacting particles and subsequently remove it from the system. Conversely, a particle can be removed from the system and then inject it back. This process may induce non-trivial collective excitations on the system. Physically, an interesting quantity would be the correlation function.

Consider a zero-temperature system. The one-particle Green function is defined, in the position representation, as

$$(8.2.2.1) \quad i\mathfrak{G}_{\alpha\beta}(\mathbf{x}t, \mathbf{x}'t') = \frac{\langle \Omega | \mathcal{T}[\hat{\psi}_{H\alpha}(\mathbf{x}, t) \hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t')] | \Omega \rangle}{\langle \Omega | \Omega \rangle},$$

where

- $|\Omega\rangle$ is the ground state of the interacting system in the Heisenberg picture,
- and $\hat{\psi}_{H\alpha}(\mathbf{x}, t)$ is a Heisenberg field operator, which can be written in terms of the corresponding Schrödinger field operator as $\hat{\psi}_{H\alpha}(\mathbf{x}, t) = e^{\frac{i\mathbf{H}t}{\hbar}} \hat{\psi}_{H\alpha}(\mathbf{x}, t) e^{-\frac{i\mathbf{H}t}{\hbar}}$,
- and where α represents a set of quantum numbers associated with the particle, eg. spin, momentum etc.
- \mathcal{T} is the time-ordering operator, whose definition is as follows

$$\begin{aligned}
 (8.2.2.2) \quad \mathcal{T}[\hat{\psi}_{H\alpha}(\mathbf{x}, t) \hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t')] &= \begin{cases} \hat{\psi}_{H\alpha}(\mathbf{x}, t) \hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t') & t > t' \\ -\hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t') \hat{\psi}_{H\alpha}(\mathbf{x}, t) & t' > t \end{cases} \\
 &= \Theta(t - t') \hat{\psi}_{H\alpha}(\mathbf{x}, t) \hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t') - \Theta(t' - t) \hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t') \hat{\psi}_{H\alpha}(\mathbf{x}, t),
 \end{aligned}$$

where the minus sign arises from the anticommutation of the fermionic field operators.

Now, [equation \(8.2.2.1\)](#)'s one-particle nature is due to the fact that it represents a propagator of a particle created at \mathbf{x}', t' with quantum numbers β , and annihilated at \mathbf{x}, t with quantum numbers α . Note, however, that these so-called single-particle Green functions are truly many-body objects since they describe the propagation of a single particle obeying the full many-body Hamiltonian and therefore contain all effects of interactions with the other particles in the system.

Depending on the problem at hand, there are a number of useful Green functions. The greater and lesser Green functions are defined as

$$(8.2.2.3) \quad \begin{aligned} \mathfrak{G}^>(\mathbf{x}t, \mathbf{x}'t') &= -i \langle \Omega | \hat{\psi}_{H\alpha}(\mathbf{x}, t) \hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t') | \Omega \rangle \\ \mathfrak{G}^<(\mathbf{x}t, \mathbf{x}'t') &= \pm i \langle \Omega | \hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t') \hat{\psi}_{H\alpha}(\mathbf{x}, t) | \Omega \rangle \end{aligned} \quad \begin{array}{l} \text{where the + and - signs are for} \\ \text{fermions and bosons, respectively.} \end{array}$$

The retarded and advanced Green functions are

$$(8.2.2.4) \quad \begin{aligned} \mathfrak{G}^R(\mathbf{x}t, \mathbf{x}'t') &= \begin{cases} -i\Theta(t-t') \langle \Omega | [\hat{\psi}_{H\alpha}(\mathbf{x}, t), \hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t')]_{\pm} | \Omega \rangle \\ \Theta(t-t') [\mathfrak{G}^>(\mathbf{x}t, \mathbf{x}'t') - \mathfrak{G}^<(\mathbf{x}t, \mathbf{x}'t')] \end{cases} \\ \mathfrak{G}^A(\mathbf{x}t, \mathbf{x}'t') &= \begin{cases} i\Theta(t-t') \langle \Omega | [\hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t'), \hat{\psi}_{H\alpha}(\mathbf{x}, t)]_{\pm} | \Omega \rangle \\ \Theta(t-t') [\mathfrak{G}^<(\mathbf{x}t, \mathbf{x}'t') - \mathfrak{G}^>(\mathbf{x}t, \mathbf{x}'t')] \end{cases} \end{aligned}$$

For compactness, $\langle \Omega | \Omega \rangle = 1$. Note that, from these definitions, $\mathfrak{G}^R = 0$, for $t < t'$ and $\mathfrak{G}^A = 0$, for $t > t'$. So far these Green functions are the space-time Green functions, since they involve creation and annihilation of particles at definite locations in space and time. Analogous Green functions may be defined in other basis. For example, in translationally invariant systems in space-time, the momentum-frequency domain may be useful.

As a many-body correlation function, a Green function conveys only part of the full information available in the many-body wavefunctions of the system, but it provides the relevant information for the problem at hand.

8.2.2.1. *Equations of Motion for the One-Particle Green function.* Consider a system with Hamiltonian,

$$(8.2.2.5) \quad \begin{aligned} \mathbf{H}^0 &= \sum_{\alpha \in \Lambda} \varepsilon_{\alpha} \mathbf{c}_{\alpha}^{\dagger} \mathbf{c}_{\alpha}, \\ \mathbf{H} &= \mathbf{H}^0 + \mathbf{H}', \\ \mathbf{H}' &\simeq \mathcal{O}\left((\mathbf{c}^{\dagger} \mathbf{c})^2\right) \end{aligned}$$

In the α -basis, the retarded Green function reads

$$\begin{aligned} \mathfrak{G}^R(\alpha t, \beta t') &= -i\Theta(t-t') \left\langle [\mathbf{c}_{\alpha}(t), \mathbf{c}_{\beta}^{\dagger}(t')]_{\pm} \right\rangle_{\Omega} \\ &\Rightarrow \left[i \frac{\partial}{\partial t} - \epsilon_{\alpha} \right] \mathfrak{G}^R(\alpha t, \beta t') = i \frac{\partial}{\partial t} \mathfrak{G}^R(\alpha t, \beta t') - \epsilon_{\alpha} \mathfrak{G}^R(\alpha t, \beta t') \\ &= i(-i) \left[\left(\frac{\partial}{\partial t} \Theta(t-t') \right) \left\langle [\mathbf{c}_{\alpha}(t), \mathbf{c}_{\beta}^{\dagger}(t')]_{\pm} \right\rangle_{\Omega} \right. \\ &\quad \left. - i\Theta(t-t') \left\langle \left[\frac{\partial \mathbf{c}_{\alpha}(t)}{\partial t}, \mathbf{c}_{\beta}^{\dagger}(t') \right]_{\pm} \right\rangle_{\Omega} \right] - \epsilon_{\alpha} \mathfrak{G}^R(\alpha t, \beta t') \\ &= \left[\delta(t-t') \left\langle [\mathbf{c}_{\alpha}(t), \mathbf{c}_{\beta}^{\dagger}(t')]_{\pm} \right\rangle_{\Omega} \right. \\ &\quad \left. - i\Theta(t-t') \left\langle [\mathbf{c}_{\alpha}(t), \mathbf{H}_0 + \mathbf{H}', \mathbf{c}_{\beta}^{\dagger}(t')]_{\pm} \right\rangle_{\Omega} \right] - \epsilon_{\alpha} \mathfrak{G}^R(\alpha t, \beta t') \\ &= \left[\delta(t-t') \delta_{\alpha\beta} - i\Theta(t-t') \left\langle [\mathbf{c}_{\alpha}(t), \mathbf{H}', \mathbf{c}_{\beta}^{\dagger}(t')]_{\pm} \right\rangle_{\Omega} \right. \\ &\quad \left. - \epsilon_{\alpha} i\Theta(t-t') \left\langle [\mathbf{c}_{\alpha}(t), \mathbf{c}_{\beta}^{\dagger}(t')]_{\pm} \right\rangle_{\Omega} \right] - \epsilon_{\alpha} \mathfrak{G}^R(\alpha t, \beta t') \\ &= \left[\delta(t-t') \delta_{\alpha\beta} - i\Theta(t-t') \left\langle [\mathbf{c}_{\alpha}(t), \mathbf{H}', \mathbf{c}_{\beta}^{\dagger}(t')]_{\pm} \right\rangle_{\Omega} \right], \end{aligned}$$

where $[\mathbf{c}_\alpha, \mathbf{H}_0] = -\varepsilon_\alpha \mathbf{c}_\alpha$ has been used. Now, if $\mathcal{O}\left((\mathbf{c}^\dagger \mathbf{c})^2\right)$, then $[\mathbf{c}_\alpha(t), \mathbf{H}']$ will be cubic in the \mathbf{c} -operators. Then,

$$\mathfrak{G}_2^R(\alpha t, \beta t') = i\Theta(t - t') \left\langle \left[[\mathbf{c}_\alpha(t), \mathbf{H}'], \mathbf{c}_\beta^\dagger(t') \right]_{\pm} \right\rangle_{\Omega}$$

is a two-particle Green function. Then, the dynamics of a one-particle Green function becomes coupled to that of a two-particle Green function. This two-particle Green function's dynamic will become coupled to a three-particle Green function and so on. Unless the couplings terminate at a reasonable order, there is an infinite hierarchy of equations and corresponding Green functions.

8.2.2.2. Equations of Motion and Microscopic Self-Interaction. Consider the many-electron Hamiltonian

$$(8.2.2.6) \quad \begin{aligned} \mathbf{H} = & \int_{\mathbb{R}^3} d\mathbf{x} \sum_{\sigma \in \mathbb{C}^2} \psi_\sigma^\dagger(\mathbf{x}) \left[-\frac{\hbar^2}{2m} \nabla^2 + V(\mathbf{x}) \right] \psi_\sigma(\mathbf{x}) \\ & + \frac{1}{2} \int_{\mathbb{R}^6} d\mathbf{x} d\mathbf{x}' \sum_{\sigma \sigma' \in \mathbb{C}^2} \psi_\sigma^\dagger(\mathbf{x}) \psi_{\sigma'}^\dagger(\mathbf{x}') V(|\mathbf{x} - \mathbf{x}'|) \psi_{\sigma'}(\mathbf{x}) \psi_\sigma(\mathbf{x}), \end{aligned}$$

where V represents the pair interaction. This Hamiltonian includes a one-electron potential U . The Green function's equation of motion may be found from

$$(8.2.2.7) \quad \left[i \frac{\partial}{\partial t} - \frac{1}{\hbar} \left(-\frac{\hbar^2}{2m} \nabla^2 + U(\mathbf{x}) \right) \right] \mathfrak{G}_{\sigma\sigma'}(\mathbf{x}t; \mathbf{x}'t') = i \frac{\partial}{\partial t} \mathfrak{G}_{\sigma\sigma'}(\mathbf{x}t; \mathbf{x}'t') - \left(-\frac{\hbar^2}{2m} \nabla^2 + U(\mathbf{x}) \right) \mathfrak{G}_{\sigma\sigma'}(\mathbf{x}t; \mathbf{x}'t').$$

Then, the first term may be rewritten as

$$(8.2.2.8) \quad \begin{aligned} i \frac{\partial}{\partial t} \mathfrak{G}_{\sigma\sigma'}(\mathbf{x}t; \mathbf{x}'t') = & (-i) i \frac{\partial}{\partial t} \left\langle \mathcal{T} \hat{\psi}_{H\sigma}(\mathbf{x}, t) \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \right\rangle_{\Omega} \\ = & \frac{\partial}{\partial t} \left(\Theta(t - t') \left\langle \hat{\psi}_{H\sigma}(\mathbf{x}, t) \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \right\rangle_{\Omega} - \Theta(t' - t) \left\langle \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \hat{\psi}_{H\sigma}(\mathbf{x}, t) \right\rangle_{\Omega} \right) \\ = & \delta(t - t') \left\langle \hat{\psi}_{H\sigma}(\mathbf{x}, t) \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') + \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \hat{\psi}_{H\sigma}(\mathbf{x}, t) \right\rangle_{\Omega} \\ & + \Theta(t - t') \left\langle \frac{\partial \hat{\psi}_{H\sigma}(\mathbf{x}, t)}{\partial t} \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \right\rangle_{\Omega} - \Theta(t' - t) \left\langle \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \frac{\partial \hat{\psi}_{H\sigma}(\mathbf{x}, t)}{\partial t} \right\rangle_{\Omega} \\ = & \delta(t - t') \left\langle \hat{\psi}_{H\sigma}(\mathbf{x}, t) \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') + \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \hat{\psi}_{H\sigma}(\mathbf{x}, t) \right\rangle_{\Omega} \\ & - \Theta(t - t') \left\langle [\hat{\psi}_{H\sigma}(\mathbf{x}, t), \mathbf{H}] \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \right\rangle_{\Omega} + \Theta(t' - t) \left\langle \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') [\hat{\psi}_{H\sigma}(\mathbf{x}, t), \mathbf{H}] \right\rangle_{\Omega} \\ = & \delta(t - t') \left\langle e^{i\mathbf{H}t} \{ \hat{\psi}_{S\sigma}(\mathbf{x}), \hat{\psi}_{S\sigma'}^\dagger(\mathbf{x}') \} e^{-i\mathbf{H}t} \right\rangle_{\Omega} + \mathcal{T} \left\langle \left[-i [\hat{\psi}_{H\sigma}(\mathbf{x}, t), \mathbf{H}] \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \right] \right\rangle_{\Omega} \\ = & \delta(\mathbf{x} - \mathbf{x}') \delta(t - t') \delta_{\sigma\sigma'} + \mathcal{T} \left\langle \left[-i [\hat{\psi}_{H\sigma}(\mathbf{x}, t), \mathbf{H}] \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \right] \right\rangle_{\Omega} \\ = & \text{id}_{\mathcal{F}} + \mathcal{T} \left\langle \left[-i [\hat{\psi}_{H\sigma}(\mathbf{x}, t), \mathbf{H}] \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \right] \right\rangle_{\Omega}, \end{aligned}$$

where the commutator in the Heisenberg picture has been transformed to one on the Schrödinger picture

(8.2.2.9)

$$\begin{aligned}
[\hat{\psi}_{H\sigma}^\dagger(\mathbf{x}, t), \mathbf{H}] &= e^{i\mathbf{H}t} [\hat{\psi}_{S\sigma}^\dagger(\mathbf{x}), \mathbf{H}] e^{-i\mathbf{H}t} \\
\Rightarrow [\hat{\psi}_{S\sigma}^\dagger(\mathbf{x}), \mathbf{H}] &= \left[-\frac{\hbar^2}{2m} \nabla^2 + U(\mathbf{x}) \right] \hat{\psi}_{S\sigma}(\mathbf{x}) + \int_{\mathbb{R}^3} d\mathbf{x}'' \sum_{\sigma'' \in \mathbb{C}^2} \psi_{S\sigma''}^\dagger(\mathbf{x}'') V(|\mathbf{x} - \mathbf{x}''|) \psi_{S\sigma''}(\mathbf{x}'') \hat{\psi}_{S\sigma}(\mathbf{x}) \\
\Rightarrow [\hat{\psi}_{h\sigma}^\dagger(\mathbf{x}, T), \mathbf{H}] &= \left[-\frac{\hbar^2}{2m} \nabla^2 + U(\mathbf{x}) \right] \hat{\psi}_{H\sigma}(\mathbf{x}, t) + \int_{\mathbb{R}^3} d\mathbf{x}'' \sum_{\sigma'' \in \mathbb{C}^2} \psi_{H\sigma''}^\dagger(\mathbf{x}'', t) V(\mathbf{x}t; \mathbf{x}'t') \psi_{H\sigma''}(\mathbf{x}'', t) \hat{\psi}_{H\sigma}(\mathbf{x}, t)
\end{aligned}$$

Plugging this result into the previous equation yields

$$\begin{aligned}
\text{i}\partial_{\mathcal{F}} &= i \frac{\partial}{\partial t} \mathfrak{G}_{\sigma\sigma'}(\mathbf{x}t; \mathbf{x}'t') - \mathcal{T} \left\langle \left[-i \left[\hat{\psi}_{H\sigma}(\mathbf{x}, t), \mathbf{H} \right] \hat{\psi}_{H\sigma'}^\dagger(\mathbf{x}', t') \right] \right\rangle_{\Omega} \\
(8.2.2.10) \quad \text{i}\partial_{\mathcal{F}} &= \left\{ i \frac{\partial}{\partial t} - \left[-\frac{\hbar^2}{2m} \nabla^2 + U(\mathbf{x}) \right] \right\} \mathfrak{G}_{\sigma\sigma'}(\mathbf{x}t; \mathbf{x}'t') \\
&\quad + i \int_{\mathbb{R}^4} d\mathbf{x}'' dt'' V(\mathbf{x}t; \mathbf{x}''t'') \left\langle \mathcal{T} \left[\psi_{H\sigma''}^\dagger(\mathbf{x}'', t'') \psi_{H\sigma''}(\mathbf{x}'', t'') \psi_{H\sigma}(\mathbf{x}, t) \psi_{H\sigma'}^\dagger(\mathbf{x}', t') \right] \right\rangle_{\Omega}.
\end{aligned}$$

Then, in the absence of interactions,

$$(8.2.2.11) \quad \left\{ i \frac{\partial}{\partial t} - \left[-\frac{\hbar^2}{2m} \nabla^2 + U(\mathbf{x}) \right] \right\} \mathfrak{G}_{\sigma\sigma'}^{(0)}(\mathbf{x}t; \mathbf{x}'t') = \text{i}\partial_{\mathcal{F}}$$

$$(8.2.2.12) \quad \mathfrak{G}_{\sigma\sigma'}^{(0)}(\mathbf{x}t; \mathbf{x}'t') = \left\{ i \frac{\partial}{\partial t} - \left[-\frac{\hbar^2}{2m} \nabla^2 + U(\mathbf{x}) \right] \right\}^{-1}$$

for a given periodic potential $U(\mathbf{x})$ an analytic expression may be found. Note that the integro-differential equation for the Green functions, written in [equation \(8.2.2.10\)](#), has no general explicit solution. But it can be seen, from a matrix perspective, as having an infinite number of continuous and discrete indices. Let,

$$\begin{aligned}
i \int_{\mathbb{R}^4} d\mathbf{x}'' dt'' V(\mathbf{x}t; \mathbf{x}''t'') \left\langle \mathcal{T} \left[\psi_{H\sigma''}^\dagger(\mathbf{x}'', t'') \psi_{H\sigma''}(\mathbf{x}'', t'') \psi_{H\sigma}(\mathbf{x}, t) \psi_{H\sigma'}^\dagger(\mathbf{x}', t') \right] \right\rangle_{\Omega} &= \\
&= - \int_{\mathbb{R}^4} d\mathbf{x}'' dt'' \Sigma_{\sigma\sigma''}^*(\mathbf{x}t; \mathbf{x}''t'') \mathfrak{G}_{\sigma''\sigma'}^{(0)}(\mathbf{x}''t''; \mathbf{x}'t').
\end{aligned}$$

Then, ?? may be rewritten in matrix notation as

$$(8.2.2.13) \quad \left[(\mathfrak{G}^{(0)})^{-1} - \Sigma^* \right] \mathfrak{G} = \text{i}\partial_{\mathcal{F}} \Rightarrow \mathfrak{G} = \left[(\mathfrak{G}^{(0)})^{-1} - \Sigma^* \right]^{-1} = \mathfrak{G}^{(0)} + \mathfrak{G}^{(0)} \Sigma^* \mathfrak{G},$$

which is the matrix form of the Dyson equations for the one-particle Green function, or the Dyson series. It will be shown that Σ^* is the one-particle irreducible self-energy.

8.2.2.3. Physical Interpretation of the One-particle Green function and the Self-Energy. Let $|\Omega\rangle$ be the (normalized) system's ground state, such that the system's Green function can be written as

$$(8.2.2.14) \quad i\mathfrak{G}_{\alpha\beta}(\mathbf{x}t, \mathbf{x}'t') = \langle \Omega | \mathcal{T} [\hat{\psi}_{H\alpha}(\mathbf{x}, t) \hat{\psi}_{H\beta}^\dagger(\mathbf{x}', t')] | \Omega \rangle.$$

The complete set of eigenstates of the Hamiltonian, defined on the Fock space and thus containing any number of particles, is given by

$$\sum_n \langle \Psi_n | \Psi_n \rangle = \text{id}_{\mathcal{F}},$$

$e^{-i\mathbf{H}t} |\Psi_n^N\rangle = e^{-iE_n(N)t} |\Psi_n^N\rangle$, where $|\Psi_n^N\rangle$ is the n -th eigenstate in the sub-Hilbert space for N particles with $E_n(N)$, its corresponding eigenvalue

$$\Rightarrow i\mathfrak{G}_{\sigma\sigma'}(\mathbf{x}t, \mathbf{x}'t') = \sum_n \begin{bmatrix} \Theta(t-t')e^{-i(E_n-E_0)(t-t')} \langle \Psi_0 | \psi_{\sigma}(\mathbf{x}) | \Psi_0 \rangle \langle \Psi_0 | \psi_{\sigma'}^{\dagger}(\mathbf{x}') | \Psi_0 \rangle \\ -\Theta(t-t')e^{-i(E_n-E_0)(t-t')} \langle \Psi_0 | \psi_{\sigma'}^{\dagger}(\mathbf{x}') | \Psi_0 \rangle \langle \Psi_0 | \psi_{\sigma}(\mathbf{x}) | \Psi_0 \rangle \end{bmatrix}$$

$$\begin{aligned} \Rightarrow \mathcal{F}[i\mathfrak{G}_{\sigma\sigma'}(\mathbf{x}t, \mathbf{x}'t')](\mathbf{x}; \mathbf{x}', \omega) &= \int_{\mathbb{R}} d\tau e^{i\omega\tau} \mathfrak{G}_{\sigma\sigma'}(\mathbf{x}t, \mathbf{x}'t') \rightarrow (-i) \int_{\mathbb{R}} d\tau e^{i\omega\tau} \Theta(\tau) e^{-i(E_n-E_0)\tau} \\ &= i \int_{\mathbb{R}} \frac{d\omega'}{2\pi i} \frac{1}{\omega' + i\eta} \int_{\mathbb{R}} d\tau e^{i\omega\tau} e^{-i\omega'\tau} e^{-i(E_n-E_0)\tau} \\ &= \frac{1}{\omega - (E_n - E_0) + i\eta}. \end{aligned}$$

Consider the first term of the matrix elements of the Green function,

$$\Theta(t-t')e^{-i(E_n-E_0)(t-t')} \langle \Psi_0 | \psi_{\sigma}(\mathbf{x}) | \Psi_0 \rangle \langle \Psi_0 | \psi_{\sigma'}^{\dagger}(\mathbf{x}') | \Psi_0 \rangle.$$

Examination of this matrix element shows that the $|\Psi_n\rangle$ -state must correspond to an excited state with $(N+1)$ particles, since $|\Psi_0\rangle$ is the ground state for the N -body system. Consequently, its eigenvalue E_n is the energy of an excited state of the $(N+1)$ -particle system. It is denoted explicitly as $E_n \rightarrow E_n(N+1)$. Thus, the denominator in the Green function's Fourier transform can be written as

$$E_n - E_0 \rightarrow E_n(N+1) - E_0(N).$$

Similarly for the second term, the energies of the denominator read $E_0 - E_n \rightarrow E_0(N) - E_n(N-1)$, thus yielding the final result

$$(8.2.2.15) \quad \Rightarrow i\mathfrak{G}_{\sigma\sigma'}(\mathbf{x}t, \mathbf{x}'t') = \sum_n \left[\frac{\langle \Psi_0 | \psi_{\sigma}(\mathbf{x}) | \Psi_0 \rangle \langle \Psi_0 | \psi_{\sigma'}^{\dagger}(\mathbf{x}') | \Psi_0 \rangle}{\omega - (E_n(N+1) - E_0(N)) + i\eta} + \frac{\langle \Psi_0 | \psi_{\sigma'}^{\dagger}(\mathbf{x}') | \Psi_0 \rangle \langle \Psi_0 | \psi_{\sigma}(\mathbf{x}) | \Psi_0 \rangle}{\omega - (E_0(N) - E_n(N-1)) - i\eta} \right].$$

The one-particle Green function's poles are precisely the one-particle excitations of the interacting system. The first term describes the excitations produced by adding one particle to the system, therefore defining particle excitations with energies above the chemical potential, namely

$$E_n(N+1) - E_0 = E_n(N+1) + E_0(N+1) - E_0(N) = \epsilon_n(N+1) + \mu,$$

where $E_0(N+1)$ is the ground state of the $(N+1)$ -system. The quantity $E_0(N+1) - E_0(N)$ is identified with the chemical potential μ , since the volume of the system is kept constant. Consequently, the energy difference $E_n(N+1) - E_0(N+1) = \epsilon_n(N+1)$ must be just the excitation energy of the $(N+1)$ -system, provided $\epsilon_n(N+1)$ is non-negative. The second term describes the transition to a system with one particle less, with manifest hole excitations having energies below the chemical potential. The merger of hole and electron excitations in a single function is one of the notable advantages of the time-ordered Green function³⁰.

8.2.2.4. Green Function for the Fermi Gas. Consider the Sommerfeld gas, for which all excited states are known. The corresponding field operators are given by

$$(8.2.2.16) \quad \Psi(\mathbf{x}, t) = \frac{1}{\sqrt{\Omega}} \sum_{\substack{\mathbf{k} \in \Lambda \\ \sigma \in \mathbb{C}^2}} e^{i\mathbf{k} \cdot \mathbf{x}} \mathbf{c}_{\mathbf{k}\sigma}, \quad \Psi^{\dagger}(\mathbf{x}, t) = \frac{1}{\sqrt{\Omega}} \sum_{\substack{\mathbf{k} \in \Lambda \\ \sigma \in \mathbb{C}^2}} e^{-i\mathbf{k} \cdot \mathbf{x}} \mathbf{c}_{\mathbf{k}\sigma}^{\dagger}$$

³⁰An alternative identification may be given, in terms of the electron affinity and the ionization energy.

$$\begin{aligned}
(8.2.2.17) \quad \Rightarrow \mathcal{M} &= \langle \Psi_0 | \psi_\sigma(\mathbf{x}) | \Psi_n \rangle \langle \Psi_n | \psi_{\sigma'}^\dagger(\mathbf{x}') | \Psi_0 \rangle = \frac{1}{\Omega} \sum_{\mathbf{k}, \mathbf{k}' \in \Lambda} e^{i\mathbf{k} \cdot \mathbf{x}} e^{-i\mathbf{k}' \cdot \mathbf{x}'} \langle \Psi_0 | \mathbf{c}_{\mathbf{k}\sigma} | \Psi_n \rangle \langle \Psi_n | \mathbf{c}_{\mathbf{k}'\sigma'}^\dagger | \Psi_0 \rangle \\
&= \frac{1}{\Omega} \sum_{\mathbf{k}, \mathbf{k}' \in \Lambda} e^{i\mathbf{k} \cdot (\mathbf{x} - \mathbf{x}')} \langle \Psi_0 | \mathbf{c}_{\mathbf{k}\sigma} | \mathbf{k}\sigma \rangle \langle \mathbf{k}\sigma | \mathbf{c}_{\mathbf{k}'\sigma'}^\dagger | \Psi_0 \rangle.
\end{aligned}$$

The translational invariance of the Sommerfeld gas leads to an $(\mathbf{x} - \mathbf{x}')$ -dependence, thus the spatiotemporal Fourier transformation of the Green function is

$$(8.2.2.18) \quad \mathfrak{G}_{\sigma\sigma'}(\mathbf{k}, \omega) = \delta_{\sigma\sigma'} \left[\frac{|\langle \Psi_0 | \mathbf{c}_{\mathbf{k}\sigma} | \mathbf{k}\sigma \rangle|^2}{\omega - \varepsilon_{\mathbf{k}} + i\eta} + \frac{|\langle \Psi_0 | \mathbf{c}_{\mathbf{k}\sigma}^\dagger | \bar{\mathbf{k}}\sigma \rangle|^2}{\omega - \varepsilon_{\mathbf{k}} - i\eta} \right],$$

where $|\bar{\mathbf{k}}\sigma\rangle = \mathbf{c}_{\mathbf{k}\sigma} |\Psi_0\rangle$ represents the ground state with one particle at $\varepsilon_{\mathbf{k}}$ removed. The poles of the one-particle Green function are, thus, just the one-particle excitation energies for particles and holes, with the energies of particle states above the chemical potential, and the states for holes energies below the chemical potential, satisfying $\varepsilon_{\mathbf{k}} > 0$. Since the Sommerfeld's gas ground state is just the Fermi sea, $|\Psi_0\rangle = \prod_{|\mathbf{k}| < k_F, \sigma} \mathbf{c}_{\mathbf{k}\sigma}^\dagger |0\rangle$, then the following identities hold

$$(8.2.2.19) \quad |\langle \Psi_0 | \mathbf{c}_{\mathbf{k}\sigma} | \mathbf{k}\sigma \rangle|^2 = \Theta(|\mathbf{k}| - k_F), \quad |\langle \Psi_0 | \mathbf{c}_{\mathbf{k}\sigma}^\dagger | \bar{\mathbf{k}}\sigma \rangle|^2 = \Theta(k_F - |\mathbf{k}|).$$

8.2.2.5. Spectral functions. The spectral functions or spectral weights, $\mathcal{A}(\nu, \omega)$, are the quantum state resolution of a particle with a given energy ω or as the energy resolution for a particle in a given quantum number ν . It quantifies how suitable the excitation created by adding or removing a particle in $|\nu\rangle$ -state can be described by a free non-interacting particle or hole, respectively.

In particular, let $\mathcal{A}^+(\mathbf{k}, \omega)$ and $\mathcal{A}^-(\mathbf{k}, \omega)$ be the spectral functions which quantify the probability of a particle (+) or a hole (-) with \mathbf{k} -momentum and ω -energy to be in an exact eigenstate of the system with either $N + 1$ or $N - 1$ particles, respectively, and are defined as follows

$$\begin{aligned}
(8.2.2.20) \quad \mathcal{A}^+(\mathbf{k}, \omega) &= \frac{1}{\Omega} \sum_n |\langle \Psi_0 | \mathbf{c}_{\mathbf{k}\sigma} | \Psi_n \rangle|^2 \delta(\omega - [E_N(N + 1) - E_0(N)]) \\
\mathcal{A}^-(\mathbf{k}, \omega) &= \frac{1}{\Omega} \sum_n |\langle \Psi_0 | \mathbf{c}_{\mathbf{k}\sigma}^\dagger | \Psi_n \rangle|^2 \delta(\omega - [E_0(N) - E_N(N + 1)]),
\end{aligned}$$

from which it is clear that

$$\mathcal{A}^+(\mathbf{k}, \omega) + \mathcal{A}^-(\mathbf{k}, \omega) = \Im \mathfrak{m}[\mathfrak{G}(\mathbf{k}, \omega)],$$

where the spectral functions are manifestly both real and positive definite, which allows their interpretation as a probability to find a single particle excitation with energy ω and \mathbf{k} -momentum. Note as well that

$$\int_{\mathbb{R}} d\omega \mathcal{A}(\mathbf{k}, \omega) = \langle \Psi_0 | \Psi_0 \rangle = 1.$$

Thus, the spectral functions satisfy the probability rule and is thus a probability, irrespective from the Hamiltonian. These spectral functions include only the states involving one particle more or less in the system. Consequently, these specify the one-particle density of states resolved in momentum. It turns out that spectral functions are experimentally accessible with the techniques of angle-resolved photoemission spectroscopy (ARPES) and inverse photoemission spectroscopy.

The density of states of one-particle excitations is obtained by summing over momenta,

$$(8.2.2.21) \quad \mathcal{D}(\omega) = \sum_{\mathbf{k} \in \Lambda} \left[\mathcal{A}^+(\mathbf{k}, \omega) + \mathcal{A}^-(\mathbf{k}, \omega) \right],$$

and covers states with energies both above and below the Fermi energy. The spectral functions are, in fact, related to the Green function as

$$(8.2.2.22) \quad \begin{aligned} \mathfrak{G}(\mathbf{k}, \omega) &= \int_{\mathbb{R}} d\omega' \left[\frac{\mathcal{A}^+(\mathbf{k}, \omega')}{\omega - \omega' + i\eta} + \frac{\mathcal{A}^-(\mathbf{k}, \omega')}{\omega - \omega' - i\eta} \right] \\ &= \mathcal{P} \int_{\mathbb{R}} d\omega' \frac{\mathcal{A}^+(\mathbf{k}, \omega')}{\omega - \omega' + i\eta} - i\pi \mathcal{A}^+(\mathbf{k}, \omega'), \text{ for } \omega > \mu = E_F, \end{aligned}$$

wherein Plemelj-Sokhotolksi's formula was used, $\frac{1}{\omega \pm i\eta} = \mathcal{P} \frac{1}{\omega} \mp i\pi \delta(\omega)$. The spectral functions' reality implies that

$$(8.2.2.23) \quad \begin{aligned} \mathcal{A}^+(\mathbf{k}, \omega') &= -\frac{1}{\pi} \Im[\mathfrak{G}(\mathbf{k}, \omega)] \Rightarrow \mathcal{D}(\omega) = -\frac{1}{\pi} \Im[\mathfrak{G}(\mathbf{x}, \mathbf{x}', \omega)] \text{ for } \omega > \mu \\ \mathcal{A}^-(\mathbf{k}; \omega') &= \frac{1}{\pi} \Im[\mathfrak{G}(\mathbf{k}, \omega)] \Rightarrow \mathcal{D}(\omega) = -\frac{1}{\pi} \Im[\mathfrak{G}(\mathbf{x}, \mathbf{x}'; \omega)] \text{ for } \omega < \mu. \end{aligned}$$

Spectral signatures in the one-particle Green Function.

Consider a non-interacting fermionic system. Its Green function may be written as

$$(8.2.2.24) \quad \mathfrak{G}_0(\mathbf{k}, \omega) = \frac{1}{\Omega} \left[\frac{\Theta(|\mathbf{k}| - k_F)}{\omega - \varepsilon_0(\mathbf{k}) + i\eta} + \frac{\Theta(k_F - |\mathbf{k}|)}{\omega - \varepsilon_0(\mathbf{k}) - i\eta} \right] \Rightarrow \begin{aligned} A_0^+(\mathbf{k}, \omega) &= \frac{1}{\Omega} \delta(\omega - \varepsilon_0(\mathbf{k})) & \text{for } |\mathbf{k}| > k_F \\ A_0^-(\mathbf{k}, \omega) &= \frac{1}{\Omega} \delta(\omega - \varepsilon_0(\mathbf{k})) & \text{for } |\mathbf{k}| < k_F, \end{aligned}$$

which shows that, for a non-interacting system, every single-particle state has unity weight, aside from a normalization factor. It confirms that all the weight is on the one-particle, or hole, state. The excitation with energy ω can only happen by adding an electron, or a hole, to the $|\mathbf{k}\rangle$ -state such that $\varepsilon_0(\mathbf{k}) = \omega$.

In an interacting system, the Green function is given by

$$(8.2.2.25) \quad \mathfrak{G}(\mathbf{k}, \mathbf{k}'; \omega) = \frac{1}{[\mathfrak{G}(\mathbf{k}, \omega)]^{-1} - \Sigma^*(\mathbf{k}, (\mathbf{k}'; \omega))}, \text{ where } [\mathfrak{G}(\mathbf{k}, \omega)]^{-1} = \omega - \varepsilon_0(\mathbf{k}).$$

Both \mathfrak{G} and the self-energy Σ^* are, in general, non-local function that depend on two spatial or momentum points. However, in the presence of systems with translational or periodic invariance, these become diagonal in momentum, thus the following equality holds

$$(8.2.2.26) \quad \mathfrak{G}(\mathbf{k}, \omega) = \frac{1}{\omega - \varepsilon_0(\mathbf{k}) - \Sigma^*(\mathbf{k}, \omega)} = \frac{1}{\omega - \varepsilon_0(\mathbf{k}) - \Re[\Sigma^*(\mathbf{k}, \omega)] - \Im[\Sigma^*(\mathbf{k}, \omega)]},$$

where $\Sigma^*(\mathbf{k}, \omega)$ could be a messy complex function, but its real part modifies, or renormalizes, the non-interacting dispersion $\varepsilon_0(\mathbf{k})$, while its imaginary part replaces the delta function of the pole with a finite width profile, a Lorentzian line shape.

It turns out that the imaginary part accounts for a finite lifetime in the (\mathbf{k}, ω) -state, and arises from electron-electron interactions, electron-phonon interactions or other scattering events. Furthermore, only states close to the Fermi surface are relevant, and for such states, scattering events due to electron-electron interactions lead to quadratic contribution to the line broadening, namely $-\Im \Sigma^*(\mathbf{k}, \omega) \propto (\omega - E_f)^2$. In effect, consider for simplicity $E_F = 0$. Since here the interest lies only in the low-energy excitations close to the Fermi level, a Taylor-expansion of $\varepsilon(\mathbf{k}) = \varepsilon_0(\mathbf{k}) + \Re \Sigma^*(\mathbf{k}, \omega)$ around $k = k_F$ and $\omega = E_F = 0$. In particular, if only linear terms $\mathcal{O}(\Delta \mathbf{k})$ are retained,

$$(8.2.2.27) \quad \varepsilon(\mathbf{k}) = \varepsilon_0(\mathbf{k}) + \Re \Sigma^*(\mathbf{k}, \omega) = \left[\frac{k_F}{m} + \frac{\partial \Re \Sigma^*}{\partial k} \Big|_{k=k_F} \right] \Delta \mathbf{k} + \frac{\partial \Re \Sigma^*}{\partial \omega} \Big|_{\omega=0},$$

where the derivatives are taken along a direction perpendicular to the Fermi surface. Substituting this result in the previous equation for the interacting Green function, yields an approximation valid for small ω and $\Delta \mathbf{k}$ of the form,

$$(8.2.2.28) \quad \mathfrak{G}(\mathbf{k}, \omega) = \frac{\mathcal{Z}}{\omega - \varepsilon(\mathbf{k})},$$

which formally resembles the expression for \mathfrak{G}^0 , with a few modifications,

- Here, the spectral weight factor is

$$\mathcal{Z} = \left(1 - \frac{\partial \Re \Sigma^*}{\partial \omega} \Big|_{\omega=0} \right)^{-1},$$

- with a mass renormalization given by

$$\varepsilon_0(\mathbf{k}) = (k - k_F) \frac{k_F}{m} \Rightarrow \varepsilon(\mathbf{k}) = (k - k_F) \frac{k_F}{m^*},$$

such that

$$\frac{m^*}{m} = \frac{1}{\mathcal{Z}} \left(1 + \nabla_{\mathbf{k}} \Re \Sigma^* \Big|_{k=k_F} \right)^{-1},$$

which indicates that the dispersion relation has been changed by the interactions. Therefore, a different mass from one of the independent particles emerges. This renormalization is consistent with experimental results for fermionic systems, as for example in the measurements of electronic specific heat of Fermi liquids.

A particle whose lifetime becomes finite and/or its energy gets renormalized is referred to as a quasiparticle, in order to differentiate from its kind in a non-interacting environment. The quasiparticle consists of the original real, individual particle, plus a cloud of disturbed neighbours interacting with it. It behaves very much like a real particle but with an effective mass and a lifetime.

Particle lifetime in the vicinity of a Fermi surface

Consider the consequences of injecting a quasiparticle into a state above, but close to the Fermi surface,

$$\varepsilon = E_F + \delta\varepsilon, \mathbf{k}, \text{ where } k > k_F, \frac{\delta\varepsilon}{E_F} \ll 1.$$

The quasiparticle is expected to be scattered by particles in the Fermi sea, with $(\varepsilon' < E_F, k' < k_F)$. Since all states in the Fermi sea are occupied, both particles should be scattered to states outside the Fermi sea. Momentum conservation implies that $p^\mu + p'^\mu = p_f^\mu + p'_f{}^\mu$. The scattering rate $\gamma_{\mathbf{k}}$ is expected to be proportional to the accessible phase space, namely

$$(8.2.2.29) \quad \begin{aligned} \gamma_{\mathbf{k}} &\propto \int_{\mathbb{R}^3} d\mathbf{k}' \int_{\mathbb{R}^3} d\mathbf{k}_f \int_{\mathbb{R}^3} d\mathbf{k}'_f \delta(\mathbf{k} + \mathbf{k}' - \mathbf{k}_f - \mathbf{k}'_f) \\ &\propto \int_{-\infty}^{E_F} d\varepsilon' \mathcal{D}(\varepsilon') \int_{E_F}^{\infty} d\varepsilon_f \mathcal{D}(\varepsilon_f) \int_{E_F}^{\infty} d\varepsilon'_f \mathcal{D}(\varepsilon'_f) \delta(\varepsilon + \varepsilon' - \varepsilon_f - \varepsilon'_f) \\ &\sim \mathcal{D}^3(E_F) (\delta\varepsilon)^2. \end{aligned}$$

Since $\gamma \approx (\varepsilon - E_F)^2$ in the vicinity of the Fermi energy implies that the existence of the Fermi surface allows quasiparticles to become more clearly defined as the Fermi surface is approached. Furthermore, the lifetime becomes infinite on the Fermi surface itself, since γ has a sign change as it crosses the Fermi surface. The same phase space argument reveals that the particle lifetime approaches zero as the quasiparticle's energy distances from the

vicinity of the Fermi surface. The immediate implication is that particles injected into a state with $\varepsilon \gg E_F$ are unstable and will immediately scatter into other states, leaving no trace of coherent information about its propagation, ie. its propagation is incoherent. There are no poles in the Green function for such scattering events. Note that fermions close to the Fermi surface will scatter very little, allowing for a qualitatively correct non-interacting treatment, which is a truly remarkable consequence. Many experimentally measurable properties are recorded at low temperatures, and thus essentially involve fermions at or near the Fermi surface. These properties are not qualitatively changed by electron-electron interactions.

Spectral manifestations of Particle Lifetime

So far, interparticle interactions are responsible for replacing the ideal delta function in the spectral function with a broadened profile. Said interaction are expected to redistribute the spectral weight, due to the energy exchanges between the injected particle and the quasiparticles in the systems. Traditionally, the interactions of interest are electron-phonon and electron-electron interactions. Consider a system which, for simplicity, has only a single electronic state with energy $\varepsilon(\mathbf{k})$ close to the Fermi energy E_F , with an inverse lifetime or line-broadening given by

$$\Gamma_{\mathbf{k}} = \mathcal{Z} C_{\mathbf{k}} (\varepsilon(\mathbf{k}) - E_F)^2,$$

where $C_{\mathbf{k}}$ is a proportionality constant. Then,

$$(8.2.2.30) \quad \mathfrak{G}(\mathbf{k}, \omega) = \frac{\mathcal{Z}}{\omega - \varepsilon(\mathbf{k}) + i\Gamma_{\mathbf{k}}} \underbrace{\mathfrak{G}_{\text{coh}}(\mathbf{k}, \omega)} + \mathfrak{G}_{\text{incoh}}(\mathbf{k}, \omega).$$

Here $\mathfrak{G}_{\text{incoh}}$ contains the rest of the spectrum, describing incoherent scattering events that have no corresponding poles. In contrast, $\mathfrak{G}_{\text{coh}}$, the coherent term corresponds to a coherently propagating particle at least for times $t < \frac{1}{\Gamma_{\mathbf{k}}}$. To confirm this fact, consider the Fourier transform $\mathfrak{G}_{\text{coh}}(\mathbf{k}, \omega)$ to study its time evolution

$$(8.2.2.31) \quad \mathfrak{G}_{\text{coh}}(\mathbf{k}, t) = \int_{\mathbb{R}} d\omega e^{-i\omega t} \frac{\mathcal{Z}}{\omega - \varepsilon(\mathbf{k}) + i\Gamma_{\mathbf{k}}} \sim \mathcal{Z} \exp\left(-i[\varepsilon(\mathbf{k}) - i\Gamma_{\mathbf{k}}]t\right),$$

where the integral has been performed using Jordan's theorem, considering $t > 0$ and $\Gamma_{\mathbf{k}} > 0$, so that the contour integral can be performed in the lower half-plane. Thus, the coherent part evolves as a decaying plane wave with $\frac{1}{\Gamma}$ -lifetime, with a particle weight given by the pole residue \mathcal{Z} , the quasiparticle weight. When \mathcal{Z} vanishes, the electron quasiparticle disappears.

$\mathfrak{G}_{\text{coh}}$ can be rewritten as

$$(8.2.2.32) \quad \mathfrak{G}_{\text{coh}} = \frac{\mathcal{Z}(\omega - \varepsilon(\mathbf{k}) - i\Gamma_{\mathbf{k}})}{(\omega - \varepsilon(\mathbf{k})^2 + \Gamma_{\mathbf{k}}^2)} \Rightarrow \mathcal{A}_{\text{coh}}^+(\mathbf{k}, \omega) = -\frac{1}{\pi} \Im \mathfrak{G}_{\text{coh}}(\mathbf{k}, \omega) \rightarrow \frac{\mathcal{Z} \frac{\Gamma_{\mathbf{k}}}{\pi}}{(\omega - \varepsilon(\mathbf{k})^2 + \Gamma_{\mathbf{k}}^2)}$$

Naturally, the full spectral function may be written as $\mathcal{A}^+(\mathbf{k}, \omega) = \mathcal{A}_{\text{coh}}^+(\mathbf{k}, \omega) + \mathcal{A}_{\text{incoh}}^+(\mathbf{k}, \omega)$, which is schematically represented in [figure 3](#). The contributions from the coherent terms appear as a Lorentzian line shape, while the background arises from the incoherent part. If the quasiparticle weight vanishes, then the background only manifests itself in the spectral function.

The coherent terms of the spectral function have the following features:

- 1) It has a Lorentzian profile peaked at $\varepsilon(\mathbf{k})$.
- 2) The single-particle energies are renormalized, $\varepsilon_0(\mathbf{k}) \rightarrow \varepsilon(\mathbf{k})$.
- 3) Its width is given by the inverse of the lifetime.
- 4) Its magnitude is proportional to the quasiparticle weight \mathcal{Z} , a positive number between 0 and 1.
- 5) The area under the peak has decreased, from 1 and \mathcal{Z} .

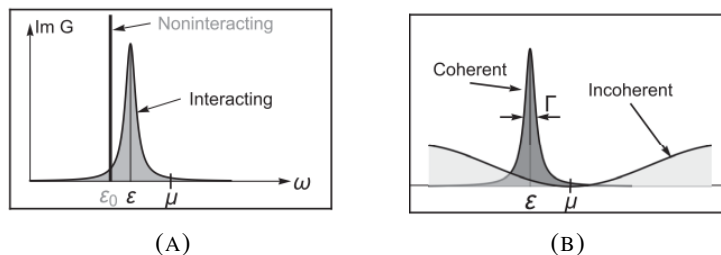


FIGURE 3. Fig. A: Spectral function for a particle in a non-interacting system (black delta function) and of the corresponding quasiparticle after turning on interactions (gray filled curve). Notice the shift in energy and the linewidth broadening.

Fig. B: Typical spectral function for a quasiparticle displaying its coherent component and the coherent background.

The incoherent terms, in contrast, presents a continuum, not a peak. It must be present if $Z \neq 1$, in order to satisfy the sum rule/normalization of the spectral function. A similar result is obtained for the hole spectral function, wherein $\Gamma < 0$.

Physical Context

Fermionic systems for which the previously laid picture holds are called Fermi liquids, its theory phenomenologically developed by Landau (1957), and in the subsequent years a microscopic formulation with the help of many-body perturbation theory was developed as well.

The identification and investigation of interacting fermionic systems which do not obey Fermi liquid theory is an important research area in modern many-body physics. Such non-Fermi liquids have, by definition, spectral functions which cannot be approximated by [equation \(8.2.2.32\)](#) and, therefore, cannot be understood in terms of a picture of non-interacting fermions. A prominent example of a non-Fermi liquid is the so-called Luttinger fermions, which occur in one-dimension.

8.3. Quantum Field Theory and Feynmann Rules. In the definition of the Green function, the average values are taken with respect to the Heisenberg ground-state vector, $|\psi_H\rangle$, which is a time-independent and generally intractable object for interacting systems. The strategy is to set

$$|\psi_H\rangle = |\psi_I(0)\rangle = \mathcal{S}^\dagger(t, 0) |\psi_I(t)\rangle,$$

and use the time $t \rightarrow -\infty$ to interpolate the non-interacting ground state with the aid of the principle of adiabatic switching on.

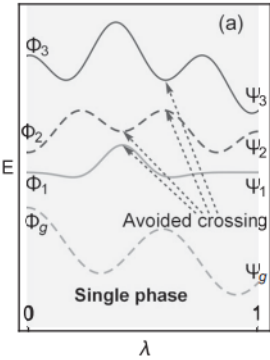
8.3.0.1. The concept of Adiabaticity in Quantum Mechanics. The construct of adiabaticity plays an important role in quantum many-body theory, by paving a path to understand a many-body problem, even when there is only an approximate solution.

In order to examine the usefulness of adiabaticity, consider a many-body quantum system with Hamiltonian $\mathbf{H} = \mathbf{H}_0 + \mathbf{H}'$, for which no exact solution of its ground state $|\Psi_g\rangle$ can be obtained but the exact ground state for \mathbf{H}_0 , labelled as $|\Phi_g\rangle$, is known. To connect the two systems, the \mathbf{H}_0 -Hamiltonian is ever present, with ground state $|\Phi_g\rangle$, and adiabatically slowly turn on \mathbf{H}' until the total Hamiltonian is reached. This can be done by setting

$$(8.3.0.1) \quad \mathbf{H}(t) = \mathbf{H} + \lambda(t)\mathbf{H}', \text{ where } \lambda(t) = e^{-\eta|t|} \text{ and } \eta > 0, \text{ is arbitrary small,}$$

and varying $-\infty \leq t \rightarrow 0$. The efficacy of this approach hinges on whether $|\Phi_g\rangle$ has the same basic symmetries as $|\Psi_g\rangle$, ie. if these share the same quantum number labels. There are two possible paths for an adiabatic switching of a many-body system, which are explored as follows.

- *The ground-state symmetry remains invariant throughout the adiabatic evolution,*



as λ increases from 0 to 1. In such cases, the corresponding evolution of the energy spectrum, within the symmetry subspace of the ground state. In **figure 4**, subfigure a), a discrete spectrum with a series of avoided crossing, from $\lambda = 0$ to $\lambda = 1$, is shown. These crossings are avoided since the states have the same symmetries. The avoided crossing arise since an intersection of two levels having the same symmetry is removed because, in general, the off-diagonal matrix elements between the two states will not vanish, and thus will give rise to mutual repulsion.

FIGURE 4. Fig a): Adiabatic evolution of a discrete spectrum within the symmetry subspace of the ground state. The ground state can be adiabatically evolved all the way to $\lambda = 1$, from $|\Phi_g\rangle$ to $|\Psi_g\rangle$.

- The ground-state symmetry changes during the evolution due to a crossing of an energy level with different symmetry .

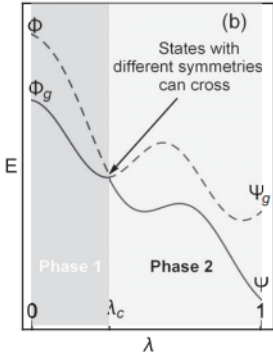


FIGURE 5. Fig b): Adiabatic evolution of the Hilbert space with λ . A phase transition occurs at $\lambda = \lambda_c$, where an excited state of the unperturbed Hamiltonian, with a symmetry different from that of the ground state, crosses below the ground state.

In this case, it is possible for the energy levels of states with different symmetries to cross, since selection rules give vanishing matrix elements between the states, and thus prevent them from mixing. This scenario leads to an adiabatic evolution, where a level crossing actually occurs at some $\lambda = \lambda_c$. Within such scenario, an excited state $|\Phi\rangle$ of the unperturbed Hamiltonian, with a symmetry different from that of the ground state $|\Psi_g\rangle$, may cross at λ_c to a lower energy than the ground state, resulting in a symmetry change of the ground state. A simple example is when a ferromagnetic ground state becomes stabilized by interactions, however, in this case a continuous rotational symmetry of the spin is broken (ie. spontaneous symmetry breaking).

Thus, the switching-on procedure may become troublesome. Specially if it leads to instabilities at some critical value $\lambda = \lambda_c$. The presence of a λ_c ushers a phase transition whereby the ground state changes its symmetry. If the transition is non-analytic, it signals a quantum critical transition initiated at a quantum critical point.

In many cases of interest, there is no symmetry-changing phase transition as the perturbation Hamiltonian is turned on, and the procedure of adiabatic evolution can be used to turn on the interaction and to analytically evolve the ground state from $|\Phi_g\rangle$ to $|\Psi_g\rangle$. The adiabatic concept, together with the Gell-Mann-Low theorem, are key elements to perturbation theory.

8.3.0.2. Gell-Mann-Low Theorem.

Historical Context

In a 1951 seminal paper, Murray Gell-Mann and Francis Low used the adiabatic construct to establish the connection between the interacting and non-interacting Green functions, in the famous Gell-Mann-Low theorem.

The ground state $|\Psi_0\rangle$, which appears in the definition of the Green function of the interacting system, must be a time independent state in the Heisenberg picture. So far, only expectation values about the non-interacting ground state $|\Phi_0\rangle$ can be calculated. Thus, a path is needed to connect the states of the exactly solvable unperturbed Hamiltonian \mathbf{H}_0 with the ones of the full Hamiltonian. For this purpose, consider the time-dependent Hamiltonian given by [equation \(8.3.0.1\)](#), which allows for the adiabatic switching on:

- For $|t| \rightarrow \infty$, the Hamiltonian reduces to \mathbf{H}_0 .
- At $t = 0$, it becomes the interacting problem of interest.
- In the $\eta \rightarrow 0$ -limit, the interaction is switched on adiabatically.

The scattering operator, in this context, is thus defined as

$$(8.3.0.2) \quad \mathcal{S}_\eta(t, t_0) = \sum_{n=0}^{\infty} \frac{1}{n!} (-i)^n \int_{t_0}^t dt_1 \cdots \int_{t_0}^t dt_n e^{-\eta(|t_1| + \cdots + |t_n|)} \mathcal{T}[\mathbf{H}_I(t_1) \cdots \mathbf{H}_I(t_n)].$$

As $t_0 \rightarrow -\infty$, the effect of the interaction vanishes and the problem reduces to that of the unperturbed Hamiltonian,

$$t_0 \rightarrow -\infty \iff \begin{cases} \mathbf{H}_0 |\Psi_0\rangle = E^{(0)} |\Psi_0\rangle \\ |\Psi_I(t_0)\rangle \rightarrow |\Psi_0\rangle. \end{cases}$$

Now, recall the relation among state vector in different pictures,

$$|\Psi_H\rangle = |\Psi_S(t=0)\rangle = |\Psi_I(t=0)\rangle,$$

where $t = 0$ corresponds to the fully interacting Hamiltonian. In particular, the following relation between the Heisenberg and interaction state vectors is found

$$|\Psi_H\rangle = |\Psi_I(0)\rangle = \mathcal{S}_\eta(0, -\infty) |\Psi_0\rangle.$$

The previous equation gives the relationship between the eigenstates of the non-interacting system with those of the fully interacting one, albeit via the adiabatic switching process. Then, the following theorem arises

Theorem 7 Gell-Mann-Low theorem

In the absence of any symmetry breaking transition, and in the $(\eta \rightarrow 0)$ -limit, $\mathcal{S}_\eta \rightarrow \mathcal{S}_I$, and provides the right solution. To be more precise, let $|\Psi_0\rangle$ be an eigenstate of the unperturbed Hamiltonian \mathbf{H}_0 with energy $E^{(0)}$, and consider the interacting Hamiltonian given by

$$\mathbf{H} = \mathbf{H}_0 + g e^{-\eta|t|} \mathbf{H}',$$

which effectively interpolates between \mathbf{H} and \mathbf{H}_0 in the limit $\eta \rightarrow 0$ and $|t| \rightarrow 0$. Let \mathcal{S}_η denote the evolution operator in the interaction limit. Then, in the $(\eta \rightarrow 0)$ -limit,

$$|\Psi_\eta^{(\pm)}\rangle = \frac{\mathcal{S}(0, \pm\infty) |\Psi_0\rangle}{\langle \Psi_0 | \mathcal{S}(0, \pm\infty) | \Psi_0 \rangle},$$

exists, is well defined with the $|\Psi_\eta^{(\pm)}\rangle$ -states being the eigenstates of the full Hamiltonian.

Proof. In effect, let the coupling constant be $g = e^{\eta\theta}$ and consider the Dyson expansion of the evolution operator. From Schrödinger's equation for the time evolution operator, it is known that

$$i\partial_{t_1} \mathcal{S}_\eta(t_1, t_2) = \mathbf{H} \mathcal{S}_\eta(t_1, t_2), \text{ with } \mathcal{S}_\eta(t_2, t_2) = \text{id},$$

and the adjoint Schrödinger equation establishes that

$$-i\partial_{t_1} \mathcal{S}_\eta(t_1, t_2) = \mathcal{S}_\eta(t_1, t_2) \mathbf{H}(t_1).$$

Then, formally

$$\mathcal{S}_\eta(t_1, t_2) = \text{id} - i \int_{t_2}^{t_1} dt' (\mathbf{H}_0 + e^{-\eta(|t'|-\theta)} \mathbf{H}') \mathcal{S}_\eta(t', t_2).$$

Then, for $t_2 \leq t_1 \leq 0$ and with $\tau = t' + \theta$, it follows that

$$\begin{aligned} \mathcal{S}_\eta(t_1, t_2) &= \text{id} - i \int_{\theta+t_2}^{\theta+t_1} dt' (\mathbf{H}_0 + e^{\eta\tau} \mathbf{H}') \mathcal{S}_\eta(\tau - \theta, t_2) \\ (8.3.0.3) \quad &\Rightarrow \partial_\theta \mathcal{S}_\eta(t_1, t_2) \equiv \eta \partial_g \mathcal{S}_\eta(t_1, t_2) = \partial_{t_1} \mathcal{S}_\eta(t_1, t_2) + \partial_{t_2} \mathcal{S}_\eta(t_1, t_2) \\ &= \mathbf{H}(t_1) \mathcal{S}_\eta(t_1, t_2) - \mathcal{S}_\eta(t_1, t_2) \mathbf{H}(t_2). \end{aligned}$$

The corresponding equation in the interaction picture can be obtained from the previous equation by pre-multiplying both sides with $e^{i\mathbf{H}t_1}$ and post-multiplying with $e^{i\mathbf{H}t_2}$ and using the definition of the interaction picture, namely that

$$\mathcal{S}_{\eta,I}(t_1, t_2) = e^{i\mathbf{H}_0 t_1} \mathcal{S}_\eta(t_1, t_2) e^{i\mathbf{H}_0 t_2}.$$

The other case, in which $t_2 \geq t_1 \geq 0$, can be obtained in a similar fashion, yielding an additional global minus sign in front of the commutator. The mixed-sign case is not of interest here. Then, using the definition of the $|\Psi_\eta^{(\pm)}\rangle$ and employing [equation \(8.3.0.3\)](#)

(8.3.0.4)

$$\begin{aligned}
& \left(\mathbf{H}_\eta(t=0) - E_0 \pm i\eta g \partial_g \right) \mathcal{S}_{\eta,I}(0, \pm\infty) |\Psi_0\rangle = 0. \\
& \Rightarrow \mp i\eta g \partial_g \mathcal{S}_{\eta,I}(0, \pm\infty) |\Psi_\eta^{(\pm)}\rangle = (\mathbf{H}_\eta - E_0) \mathcal{S}_{\eta,I}(0, \pm\infty) |\Psi_\eta^{(\pm)}\rangle \\
& \Rightarrow \mp i\eta g \partial_g |\Psi_0\rangle = \frac{(\mathbf{H} - E_0) \mathcal{S}_\eta(t_1, t_2) |\Psi_0\rangle}{\langle \Psi_0 | \mathcal{S}_\eta(t_1, t_2) | \Psi_0 \rangle} - \frac{\mathcal{S}_\eta(t_1, t_2) |\Psi_0\rangle}{\langle \Psi_0 | \mathcal{S}_\eta(t_1, t_2) | \Psi_0 \rangle} \langle \Psi_0 | (\mathbf{H} - E_0) \mathcal{S}_\eta(t_1, t_2) | \Psi_0 \rangle
\end{aligned}$$

$$\begin{aligned}
\Rightarrow \mp i\eta g \partial_g |\Psi_0\rangle &= (\mathbf{H} - E_0) |\Psi_\eta^{(\pm)}\rangle - \langle \Psi_\eta^{(\pm)} | \langle \Psi_0 | \mathbf{H} - E_0 | \Psi_\eta^{(\pm)} \rangle \\
&= [\mathbf{H} - \mathcal{E}^{(\pm)}] |\Psi_\eta^{(\pm)}\rangle,
\end{aligned}$$

where $\mathcal{E}^{(\pm)} = E_0 + \langle \Psi_0 | \mathbf{H} - \mathbf{H}_0 | \Psi_\eta^{(\pm)} \rangle$. In the $\eta \rightarrow 0^+$ -limit, by assumption, $g \partial_g |\Psi_\eta^{(\pm)}\rangle$ is finite. Then, it is clear that $|\Psi_\eta^{(\pm)}\rangle$ is an eigenstate of the full Hamiltonian, completing the proof.

*quod
erat
demonstrandum* ■

8.3.0.3. One-particle Green functions as a power series in the interaction. The Gell-Mann-Low theorem connects the eigenstates of the non-interacting system with those of the fully interacting one, with the aid of the interaction picture. The Heisenberg operators, appearing in the Green function, remain to be related to those in the corresponding interaction picture.

To that end, consider the relationship between operators in the Heisenberg and interaction pictures,

$$(8.3.0.5) \quad \mathbf{O}_H(t) = e^{i\mathbf{H}t} \mathbf{O}_S e^{-i\mathbf{H}t} = e^{i\mathbf{H}t} e^{-i\mathbf{H}_0 t} \mathbf{O}_I e^{i\mathbf{H}_0 t} e^{-i\mathbf{H}t}.$$

Furthermore, the states in the interaction picture are given by

$$(8.3.0.6) \quad |\Psi_I(t)\rangle = e^{i\mathbf{H}_0 t} e^{-i\mathbf{H}t} |\Psi_I(0)\rangle = \mathcal{S}_I(t, 0) |\Psi_I(0)\rangle.$$

Then,

$$(8.3.0.7) \quad \mathbf{O}_H(t) = \mathcal{S}_I(0, t) \mathbf{O}_I(t) \mathcal{S}_I(t, 0).$$

With these ingredients, the product of two Heisenberg operators, in the interaction picture, reads

$$\begin{aligned}
(8.3.0.8) \quad \mathbf{O}_H(t) \mathbf{O}_H(t_1) &= \mathcal{S}_I(0, t) \mathbf{O}_I(t) \mathcal{S}_I(t, 0) \mathcal{S}_I(0, t_1) \mathbf{O}_I(t_1) \mathcal{S}_I(t_1, 0) \\
&= \mathcal{S}_I(0, t) \mathbf{O}_I(t) \mathcal{S}_I(t, t_1) \mathbf{O}_I(t_1) \mathcal{S}_I(t_1, 0).
\end{aligned}$$

Then, the Green function in terms of interaction-picture operators reads

$$\begin{aligned}
(8.3.0.9) \quad i\mathfrak{G}(x; x') &= \langle \Psi_H | \mathcal{T}[\psi_H(x) \psi_H^\dagger(x')] | \Psi_H \rangle \\
&= \frac{\langle \Phi_0 | \mathcal{S}_I(\infty, 0) \mathcal{T}[\mathcal{S}_I(0, t) \psi_I(x) \mathcal{S}_I(t, t') \psi_I(x') \mathcal{S}_I(t', 0)] \mathcal{S}_I(0, -\infty) | \Phi_0 \rangle}{\langle \Phi_0 | \mathcal{S}_I(\infty, 0) \mathcal{S}_I(0, -\infty) | \Phi_0 \rangle},
\end{aligned}$$

where $x \equiv (\mathbf{x}, t, \sigma)$ for compactness. The (ψ_H, ψ_H^\dagger) -Heisenberg operators have been replaced by their interaction-picture counterparts, which simply evolve under the action of the unperturbed Hamiltonian \mathbf{H}_0 . The ground state of the interaction problem $|\Psi_0\rangle$ has been replaced by the non-interacting one, $|\Phi_0\rangle$. Furthermore, since the evolution operations \mathcal{S} form a semigroup, all evolution operators in the preceding equation can be combined into a single \mathcal{S} , for the time-ordering operator takes care of the proper time ordering anyhow, thus yielding

(8.3.0.10)

$$i\mathcal{G}^{\sigma\sigma'}(\mathbf{x}, t; \mathbf{x}', t') = \frac{\langle \Phi_0 | \mathcal{T}[\mathcal{S}_I(-\infty, \infty) \psi_{I\sigma}(\mathbf{x}, t) \psi_{I\sigma'}(\mathbf{x}', t')] | \Phi_0 \rangle}{\langle \Phi_0 | \mathcal{S}_I(\infty, -\infty) | \Phi_0 \rangle}$$

$$= \sum_{n=0}^{\infty} \frac{1}{n!} (-i)^n \int_{-\infty}^{\infty} dt_1 \cdots \int_{-\infty}^{\infty} dt_n \frac{\langle \Phi_0 | \mathcal{T}[\mathbf{H}_I(t_1) \cdots \mathbf{H}_I(t_n) \psi_{I\sigma}(\mathbf{x}, t) \psi_{I\sigma'}(\mathbf{x}', t')] | \Phi_0 \rangle}{\langle \Phi_0 | \mathcal{S}_I(\infty, -\infty) | \Phi_0 \rangle},$$

where the one-particle Green function is now explicitly written as an expansion in powers of the perturbation \mathbf{H}_I . For this reformulation to be useful, a method is needed to mitigate taking expectation values of large numbers of operators in a non-interacting system. This procedure is systematized with the application of Wick's theorem and simplified with the aid of Feynmann diagrams.

8.3.0.4. Wick's Theorem: Normal ordering and contractions. In the preceding identity for the one-particle Green function, the arraying of interaction Hamiltonians is fixed by time ordering. However, as the interaction Hamiltonian contains creation and annihilation operators, it is convenient to move all annihilation operators to the right, where these can eliminate particles. This ordering is called the *normal ordering*. Furthermore, observables have zero-expectation values in the vacuum state, it is convenient to construct them in normal order, with creation operators on the left of destruction operators. Wick's theorem provides a way to go from time-ordered to normal-ordered products, which is valid when the expectation values are with respect to a Hamiltonian that is quadratic in fermionic (bosonic) operators.

Recall from the study of superfluidity phenomenon that the Bogoliubov canonical transformation generated new non-interacting quasiparticles from the original interacting bosons, and the emerging ground state acted as a vacuum state for the new quasiparticles. In many physical situations, the ground state $|\Psi_g\rangle$, of some incipient effective Hamiltonian, will be taken to be the reference vacuum state. Such a state is filled with both particles and quasiparticles, and the effective Hamiltonian is expressed in a basis of canonical operators $\left\{(\alpha_a, \alpha_a^\dagger)\right\}_{a=1}$, where

$$(8.3.0.11) \quad \alpha_a |\Psi_g\rangle = 0, \quad \langle \Psi_g | \alpha_a^\dagger = 0,$$

where this action only holds for such systems, but not for the other creation and annihilation operators associated with the system.

Physical Context

In order to bring light upon these ideas, consider the non-interacting ground state $|\Phi_0\rangle$, describing the Fermi sea. Then, the field operators are given in terms of the canonical particle-hole transformation,

$$(8.3.0.12) \quad \psi(\mathbf{x}) = \sum_{\mathbf{k}\sigma} \phi_{\mathbf{k}\sigma} \mathbf{c}_{\mathbf{k}\sigma} = \begin{cases} \sum_{\mathbf{k}\sigma} \phi_{\mathbf{k}\sigma}(\mathbf{x}) \mathbf{d}_{\mathbf{k}\sigma} = \phi^{(-)}(\mathbf{x}) & \text{if } k > k_F \text{ annihilates particles} \\ \sum_{\mathbf{k}\sigma} \phi_{\mathbf{k}\sigma}(\mathbf{x}) \mathbf{h}_{-\mathbf{k}\sigma} = \phi^{(+)}(\mathbf{x}) & \text{if } k < k_F \text{ create holes} \end{cases},$$

then

$$(8.3.0.13) \quad \phi^{(-)}(\mathbf{x}) |\Phi_0\rangle = 0, \quad \phi^{(+)\dagger}(\mathbf{x}) |\Phi_0\rangle = 0.$$

The preceding relations indicate that $(\phi^{(-)}, \phi^{(+)\dagger})$ are annihilation parts, while $(\phi^{(+)}, \phi^{(-)\dagger})$ are the creation operators.

The following discussion will be restricted to fermionic operators in a non-interacting system.

Definition 3.

Consider a product of operators \mathbf{A}_i^\pm is normally ordered if all the \mathbf{A}_i^- -operators are on the right of the \mathbf{A}_i^+ -operators, i.e.

$$\mathbf{A}_1^+ \cdots \mathbf{A}_k^+ \mathbf{A}_1^- \cdots \mathbf{A}_n^-.$$

The usefulness of this definition is that the expectation value on the ground state of a normally ordered operator product is always zero.

It is clear that any product of operators $\mathbf{A}_1 \cdots \mathbf{A}_n$ can be written as a sum of normally ordered terms. In effect, every local operator can be written as $\mathbf{A}_i^+ + \mathbf{A}_i^-$, thus yielding 2^n terms. In each term, the components \mathbf{A}_i^- are brought to the right by successive anticommutation, yielding the desired result. Wick's theorem is an efficient answer to this particular problem: to write a product $\prod_i \mathbf{A}_i$ as a sum of normally ordered terms. The theorem is an extremely useful operator identity, with important corollaries.

Definition 4 - Normal Ordering.

The normal-ordering operator, \mathcal{N} brings a generic product into a normal form. If a product of operators contains k \mathbf{A}_i^+ -factors mixed with $n - k$ \mathbf{A}_i^- -factors, its action is given by

$$(8.3.0.14) \quad \mathcal{N}[\mathbf{A}_i^{(\pm)} \cdots \mathbf{A}_n^{(\pm)}] = (\pm 1)^P \mathbf{A}_{i_1}^+ \cdots \mathbf{A}_{i_k}^+ \cdots \mathbf{A}_{i_n}^-,$$

where P is the permutation which brings the $(1, \dots, n)$ -sequence to its normal ordered (i_1, \dots, i_n) -sequence.

At first glance, it may seem that the normal ordering is ambiguous, since the $(+)$ -operators and $(-)$ -operators can be given in different orders. However, the different expressions are actually the same operator, since \mathbf{A}^+ -operators anticommute among themselves, with the same being true for the \mathbf{A}^- -operators. For example, $\mathcal{N}[\mathbf{A}_1^+ \mathbf{A}_2^+] = \mathbf{A}_1^+ \mathbf{A}_2^+ = -\mathbf{A}_2^+ \mathbf{A}_1^+$. Now, any product $\prod_i \mathbf{A}_i$ can be written as a sum of normally ordered terms.

Example

Consider for example a product of two operators,

$$\begin{aligned} \mathbf{A}_1 \mathbf{A}_2 &= (\mathbf{A}_1^+ + \mathbf{A}_1^-)(\mathbf{A}_2^+ + \mathbf{A}_2^-) \\ &= \mathbf{A}_1^+ \mathbf{A}_2^+ + \mathbf{A}_1^- \mathbf{A}_2^+ + \mathbf{A}_1^+ \mathbf{A}_2^- + \mathbf{A}_1^- \mathbf{A}_2^- \\ &= \mathcal{N}[\mathbf{A}_1 \mathbf{A}_2] + \{\mathbf{A}_1^-, \mathbf{A}_2^+\}. \end{aligned}$$

The vanishing of the expectation value of a normally ordered product implies

$$\{\mathbf{A}_1^-, \mathbf{A}_2^+\} = \langle \Psi_g | \mathbf{A}_1 \mathbf{A}_2 | \Psi_g \rangle,$$

which is a c-number, the contraction of the two operators, denoted $\mathbf{A}_1^c \mathbf{A}_2^c$.

Using the previous example:

Definition 5 - Contraction.

The contraction of two operators, where a product of n' operators is inserted between them, may be defined by

$$(8.3.0.15) \quad \mathbf{A}_1^c (\mathbf{A}_{1'} \cdots \mathbf{A}_{n'}) \mathbf{A}_2^c = (-1)^{n'} \mathbf{A}_1^c \mathbf{A}_2^c (\mathbf{A}_{1'} \cdots \mathbf{A}_{n'}).$$

Wick's theorem provides a computationally efficient way for calculating operator contractions, by mapping it to a purely combinatorial problem.

Theorem 8 Wick's theorem

Consider a product of n operators, then the following equality holds:

$$(8.3.0.16) \quad \mathbf{A}_1 \cdots \mathbf{A}_n = \mathcal{N}[1 \cdots n] + \sum_c \mathcal{N}[\cdots i^c \cdots j^c \cdots n] + \sum_{c,d} \mathcal{N}[\cdots i^c \cdots r^d \cdots j^c \cdots s^d \cdots n] + \cdots,$$

where the first sum is performed over single contractions of pairs, the second sum is performed over double contractions and so on. If n is even, the result consists only of products of contractions (c -numbers).

Some important consequences of Wick's theorem readily follow,

Remark. As a general rule, the expectation value of a product of destruction and creation operators. Particle number conservation thus requires that the number of destruction operators equals that of creators.

Corollary 4. The following identity can be proved from Wick's theorem³¹,

$$(8.3.0.17)$$

$$\langle \Psi_g | \mathbf{A}_1 \cdots \mathbf{A}_{2n} | \Psi_g \rangle = \sum_{\mathcal{P}_2} (-1)^P \langle \mathbf{A}_{i_1} \mathbf{A}_{j_1} \rangle \cdots \langle \mathbf{A}_{i_n} \mathbf{A}_{j_n} \rangle,$$

where the sum is performed over the set \mathcal{P}_2 of partitions of $1, \dots, 2n$ into sets of pairs $\{(i_1, j_1), \dots, (i_n, j_n)\}$ with $(i, j) \sim (j, i)$ belonging to the same equivalency class.

An important variation of Wick's theorem deals with the normal-ordering of a time-ordered product, via a \mathcal{T} -contraction operator.

Theorem 9

Consider a product of two-operators. Then, its time-ordered product can be written as

$$(8.3.0.18) \quad \mathcal{T} \mathbf{A}_i(t_i) \mathbf{A}_j(t_j) = \mathcal{N}[\mathbf{A}_i(t_j) \mathbf{A}_j(t_i)] + \overbrace{\mathbf{A}_i(t_i) \mathbf{A}_j(t_j)},$$

then a \mathcal{T} (time ordered)-contraction is defined as

$$(8.3.0.19) \quad \overbrace{\mathbf{A}_i(t_i) \mathbf{A}_j(t_j)} = \langle \Psi_g | \mathcal{T} \mathbf{A}_i(t_i) \mathbf{A}_j(t_j) | \Psi \rangle,$$

Remark. The \mathcal{T} -contraction of two-operators is such that

$$(8.3.0.20) \quad \overbrace{\mathbf{A}_i(t_i) (\cdots) \mathbf{A}_j(t_j)} = (-1)^k \overbrace{\mathbf{A}_i(t_i) \mathbf{A}_j(t_j)} (\cdots).$$

Remark. The \mathcal{T} -contractions have a new property, not shared by ordinary contractions,

$$(8.3.0.21) \quad \overbrace{\mathbf{A}_i(t_i) \mathbf{A}_j(t_j)} = - \overbrace{\mathbf{A}_j(t_j) \mathbf{A}_i(t_i)}.$$

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Example

Using the previous corollary, it is trivial that

$$\langle \Psi_g | 1234 | \Psi_g \rangle = \langle 12 \rangle \langle 34 \rangle - \langle 1324 \rangle + \langle 14 \rangle \langle 23 \rangle.$$

Remark. From the previous definitions and remarks, the following explicit expressions can be given,

(8.3.0.22)

$$\overbrace{\psi(t_1)\psi^\dagger(t_2)} = \left\langle \mathcal{T} \psi(t_1)\psi^\dagger(t_2) \right\rangle = i\mathfrak{G}^{(0)}(1, 2), \quad \overbrace{\psi(t_1)\psi(t_2)} = \left\langle \mathcal{T} \psi(t_1)\psi(t_2) \right\rangle = i\mathfrak{F}^{(0)}(1, 2), \quad \overbrace{\psi^\dagger(t_1)\psi^\dagger(t_2)} = \left\langle \mathcal{T} \psi^\dagger(t_1)\psi^\dagger(t_2) \right\rangle = i\mathfrak{F}^{(0)\dagger}(1, 2).$$

For systems where $|\Psi_g\rangle$ has a definite number of particles, the anomalous correlators $\mathfrak{F}^{(0)}$, $\mathfrak{F}^{(0)\dagger}$ are zero. They are non-zero, for example, in the Bardeen-Cooper-Schrieffer theory.

8.3.1. Feynmann Diagrams. Consider the perturbation expansion of the Green function, given by [equation \(8.3.0.10\)](#) and consider only terms linear on the interaction Hamiltonian \mathbf{H}_I , ie.

$$(8.3.1.1) \quad i\mathfrak{G}_{\sigma\sigma'} \rightarrow i\mathfrak{G}_{\sigma\sigma'}^{(0)} + i\mathfrak{G}_{\sigma\sigma'}^{(1)}, \text{ with } t \quad \mathfrak{G}_{\sigma\sigma'}^{(1)} \propto -i \int_{\mathbb{R}} dt \langle \Psi_0 | \mathcal{T} [\mathbf{H}_I(t) \psi_{I\sigma}(\mathbf{x}, t) \psi_{I\sigma'}^\dagger(\mathbf{x}', t')] | \Psi_0 \rangle$$

$$\mathbf{H}_I = \frac{1}{2} \int_{\mathbb{R}^6} d\mathbf{x}_1 d\mathbf{x}_2 \psi_{\sigma_1}^\dagger(\mathbf{x}_1) \psi_{\sigma_2}^\dagger(\mathbf{x}_2) V(|\mathbf{x}_1 - \mathbf{x}_2|) \psi_{\sigma_2}(\mathbf{x}_2) \psi_{\sigma_1}(\mathbf{x}_1).$$

Consider the interaction of the form

$$V_{\sigma_1\sigma_1'}^{\sigma_2\sigma_2'}(\mathbf{x}_1, t_1; \mathbf{x}_2, t_2) = V(|\mathbf{x}_1 - \mathbf{x}_2|) \psi_{\sigma_2}(\mathbf{x}_2) \delta(t_1 - t_2) \delta_{\sigma_1\sigma_1'}^{\sigma_2\sigma_2'},$$

then

(8.3.1.2)

$$i\mathfrak{G}_{\sigma\sigma'}^{(1)} \propto -\frac{i}{2} \sum_{\substack{\sigma_1, \sigma_1' \\ \sigma_2, \sigma_2'}} \int_{\mathbb{R}^8} d^4x_1 d^4x_2 V_{\sigma_1\sigma_1'}^{\sigma_2\sigma_2'}(x_1, x_2) \langle \Psi_0 | \mathcal{T} \left[\psi_{\sigma_1}^\dagger(x_1) \psi_{\sigma_2}^\dagger(x_2) \psi_{\sigma_2'}(x_2) \psi_{\sigma_1'}(x_1) \psi_{\sigma}(x) \psi_{\sigma'}^\dagger(x') \right] | \Psi_0 \rangle$$

The problem of interacting electrons has been reduced to determining expectation values of its non-interacting counterparts. Now, Wick's theorem must be applied by considering all possible pair contractions. Each particle emerging from a creation operator in the previous equation, must be destroyed by an annihilation operator. Thus, every creation operator can be paired to an annihilation operator, keeping in mind the following simplifying points to obtain the final expression of the first-order Green function in terms of the free Green functions:

- the ground state is that of \mathbf{H}_0 ,
- the operators now have the form

$$\psi_{\sigma}^\dagger(\mathbf{x}, t) = e^{i\mathbf{H}_0 t} \psi_{S\sigma}^\dagger(\mathbf{x}) e^{-i\mathbf{H}_0 t}.$$

For eigenstates of the unperturbed Hamiltonian representing a crystalline solid, these can be written as

$$e^{i\mathbf{H}_0 t} \mathbf{c}_{n\mathbf{k}\sigma}^\dagger e^{-i\mathbf{H}_0 t} \rightarrow e^{i\varepsilon_{n\mathbf{k}} t} \mathbf{c}_{n\mathbf{k}\sigma}^\dagger.$$

- An injected electron in the $|n\mathbf{k}\sigma\rangle$ of a non-interacting system must be removed from the same state, since it evolves in time unperturbed. This allows for the insertion of $|\Psi_0\rangle \langle\Psi_0|$ between pairs of creation and annihilation operators to obtain

$$\langle \Psi_0 | \mathbf{c}_{n\mathbf{k}\sigma}(t) \mathbf{c}_{n\mathbf{k}\sigma}^\dagger(t_1) | \Psi_0 \rangle,$$

which is just the non-interacting single-particle Green function. This procedure can be generalized to all creation and annihilation operators a Green function contains, for they represent non-interacting systems. This is a direct consequence of Wick's theorem.

All in all, this procedures yield a total of six terms in the first-order Green function, given in ??, which can be obtained from all possible pair-wise contractions of a creation operator with an annihilation operator, yielding

(8.3.1.3)

$$\begin{aligned}
\langle \Psi_0 | \mathcal{T} \left[\psi_{\sigma_1}^\dagger(x_1) \psi_{\sigma_2}^\dagger(x_2) \psi_{\sigma'_2}(x_2) \psi_{\sigma'_1}(x_1) \psi_\sigma(x) \psi_{\sigma'}^\dagger(x') \right] | \Psi_0 \rangle &= \\
&= \left(\overbrace{\psi_{\sigma_1}^\dagger(x_1) \psi_{\sigma_2}^\dagger(x_2) \psi_{\sigma'_2}(x_2) \psi_{\sigma'_1}(x_1)} - \overbrace{\psi_{\sigma_1}^\dagger(x_1) \psi_{\sigma_2}^\dagger(x_2) \psi_{\sigma'_2}(x_2) \psi_{\sigma'_1}(x_1)} \right) \psi_\sigma(x) \psi_{\sigma'}^\dagger(x') \\
&+ \overbrace{\psi_{\sigma_1}^\dagger(x_1) \psi_{\sigma_2}^\dagger(x_2) \psi_{\sigma'_2}(x_2) \psi_{\sigma'_1}(x_1) \psi_\sigma(x) \psi_{\sigma'}^\dagger(x')} - \overbrace{\psi_{\sigma_1}^\dagger(x_1) \psi_{\sigma_2}^\dagger(x_2) \psi_{\sigma'_2}(x_2) \psi_{\sigma'_1}(x_1) \psi_\sigma(x) \psi_{\sigma'}^\dagger(x')} \\
&+ \overbrace{\psi_{\sigma_1}^\dagger(x_1) \psi_{\sigma_2}^\dagger(x_2) \psi_{\sigma'_2}(x_2) \psi_{\sigma'_1}(x_1) \psi_\sigma(x) \psi_{\sigma'}^\dagger(x')} - \overbrace{\psi_{\sigma_1}^\dagger(x_1) \psi_{\sigma_2}^\dagger(x_2) \psi_{\sigma'_2}(x_2) \psi_{\sigma'_1}(x_1) \psi_\sigma(x) \psi_{\sigma'}^\dagger(x')} \\
&= i\mathfrak{G}_{\sigma\sigma'}^{(x,x');(0)} \left(i\mathfrak{G}_{\sigma_1\sigma'_1}^{(x_1,x_1);(0)} i\mathfrak{G}_{\sigma_2\sigma'_2}^{(x_2,x_2);(0)} - i\mathfrak{G}_{\sigma'_1\sigma_2}^{(x_1,x_2);(0)} i\mathfrak{G}_{\sigma'_2\sigma_1}^{(x_2,x_1);(0)} \right) \\
&+ i\mathfrak{G}_{\sigma\sigma_1}^{(x,x_1);(0)} \left(i\mathfrak{G}_{\sigma'_1\sigma_2}^{(x_1,x_2);(0)} i\mathfrak{G}_{\sigma'_2\sigma_2}^{(x_2,x');(0)} - i\mathfrak{G}_{\sigma'_1\sigma'}^{(x_1,x');(0)} i\mathfrak{G}_{\sigma'_2\sigma_2}^{(x_2,x_2);(0)} \right) \\
&+ i\mathfrak{G}_{\sigma\sigma_2}^{(x,x_2);(0)} \left(i\mathfrak{G}_{\sigma'_2\sigma_1}^{(x_2,x_1);(0)} i\mathfrak{G}_{\sigma'_1\sigma'}^{(x_1,x');(0)} - i\mathfrak{G}_{\sigma'_2\sigma'}^{(x_2,x');(0)} i\mathfrak{G}_{\sigma'_1\sigma_2}^{(x_1,x_1);(0)} \right)
\end{aligned}$$

The, using Wick's theorem, each term in the perturbation expansion can be expressed as a product of non-interacting Green functions, which is the general methodology behind perturbation theory.

8.3.1.1. Equal-Time operators. In the previous calculation, the Green functions are ambiguous, since these represent contractions of creation and annihilation operators at equal times with the time-ordered product being undefined at equal times. Moreover, [equation \(8.2.2.10\)](#) has a discontinuity at equal times $t = t'$. Consequently, taking this limit from either side will yield different results. However, [equation \(8.3.1.2\)](#) shows that the operators of \mathbf{H}_I come in sets of four having identical times and a term such as $\mathfrak{G}^{(x_1,x_1);(0)}$ arising from a contraction of two fields within the perturbation Hamiltonian, where they appear in the form of $\psi^\dagger(x)\psi(x)$ with the adjoint field always occurring to the left of the field,

$$\mathbf{H}_I \propto \psi_{\sigma_1}^\dagger(\mathbf{x}_1) \psi_{\sigma_2}^\dagger(\mathbf{x}_2) \psi_{\sigma'_2}(\mathbf{x}_2) \psi_{\sigma'_1}(\mathbf{x}_1) \rightarrow \psi_{\sigma_1}^\dagger(\mathbf{x}_1) \psi_{\sigma'_1}(\mathbf{x}_1) \psi_{\sigma_2}^\dagger(\mathbf{x}_2) \psi_{\sigma'_2}(\mathbf{x}_2), \quad \mathbf{x}_1 \neq \mathbf{x}_2.$$

When this term is replaced in the propagators, the backward time propagation introduces a (-1)-factor for fermions. Therefore, Green functions at equal times for fermionic systems are to interpreted as follows

$$\begin{aligned}
\mathfrak{G}_{\sigma\sigma'}^{(0)}(x, x) &= \lim_{t' \rightarrow t} \langle \Psi_0 | \psi_\sigma(\mathbf{x}, t) \psi_{\sigma'}^\dagger(\mathbf{x}, t') | \Psi_0 \rangle \\
(8.3.1.4) \quad &= - \langle \Psi_0 | \psi_{\sigma'}^\dagger(\mathbf{x}) \psi_\sigma(\mathbf{x}) | \Psi_0 \rangle = -\delta_{\sigma\sigma'} \frac{n^{(0)}(\mathbf{x})}{2s+1},
\end{aligned}$$

for fermionic systems of s -spin, where $n^{(0)}(\mathbf{x})$ is the electron density in the non-interacting system, which may be different from $n(\mathbf{x})$ of the interacting system. For a uniform system, like the jellium, they coincide.

As a general rule, if two operators occur at equal time, they should be normal ordered. Operationally, this means that, when encountering Green functions of the form $\mathfrak{G}^{(0)}(\mathbf{x}, \mathbf{x}'; t = 0)$, these should be interpreted as

$$(8.3.1.5) \quad \mathfrak{G}^{(0)}(\mathbf{x}, \mathbf{x}'; t = 0) \equiv \lim_{\epsilon^-} \mathfrak{G}_{\sigma\sigma'}^{(0)}(\mathbf{x}, \mathbf{x}'; t + \epsilon)$$

8.3.2. The Feynmann Rules. The Feynmann rules allows for the calculation of the n -th-order term in the perturbation expansion of the Green function, which can be formulated in both direct space and in momentum space. The rules read

- Draw all topologically distinct connected diagrams with n interaction (wavy) lines V and $2n + 1$ Green function $\mathfrak{G}^{(0)}$ solid lines.

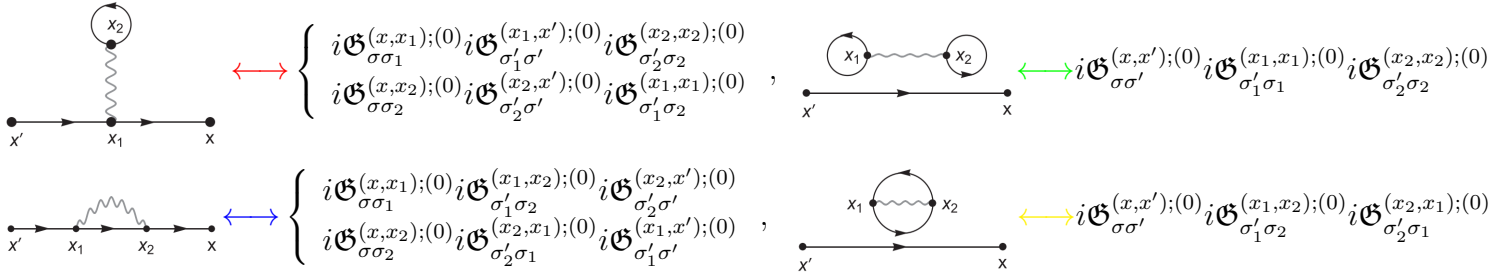
- Each wavy line corresponds to the interaction written in [equation \(8.3.1.2\)](#) (its Fourier transform in $(\mathbf{k}-\omega)$ -basis).
- Vertex and line labelling:
 - $\mathbf{x}-\mathbf{t}$ -basis: Each vertex is labelled by a point in spacetime x_i ; each solid line represents a Green function $\mathfrak{G}^{(0)}(x, x')$, running from x' to x .
 - $\mathbf{k}-\omega$ -basis: a direction is assigned to each interaction line; associated directed momentum and frequency are added to each line (both solid and wavy) and conservation of energy and momentum is imposed at each vertex.
- All internal variables are integrated over spacetime or frequency space.
- Each diagram carries a $(-1)^L$ -factor, where L is the number of closed fermion loops in the diagram.
- Each diagram contributes to \mathfrak{G} a factor

$$-i\left(\frac{-i}{\hbar}\right)^n i^{2n+1} = \left(\frac{i}{\hbar}\right)^n$$

- Feynmann rules for the first-order expansion in space-time representation

From this rules, the resulting diagrams for the first-order expansion of the Green function, shown in [equation \(8.3.1.2\)](#)-[equation \(8.3.1.3\)](#), read as follows

(8.3.2.1)



The **first** and **third** diagrams, the ones with red and blue arrows, depict a pair of Green function products, which have identical graphs but with x_1 and x_2 interchanged and said graphs are connected. In the **third** diagram, the one with the blue arrow, an exchange occurs in the contraction with x , whereas the x' -point has the same contraction. First, the injected electron travels to point x_1, t_1 where it interacts with another electron at point x_2, t_2 . This electron at x_2, t_2 then travels to the finish point, where it is taken out, with another electron at x_1, t_1 travelling to x_2, t_2 in order to leave the system in the same state it started. Consequently, such diagrams are referred to as exchange diagrams. In the **first** and **third** diagrams, there are lines starting and ending at the same vertex, thus closing a fermion loop. These describe scattering events at (x_i, t_i) , represented by a Green function $\mathfrak{G}^{(x_i, x_i);(0)}$. These can be thought of as spontaneous creation and re-annihilation of a particle. Recall that when these arise from an interaction Hamiltonian \mathbf{H}_I , the prescription is to interpret them as

$$\lim_{\epsilon \rightarrow 0} \mathfrak{G}_{\sigma'_1\sigma_1}^{(x_1, t_1; x_1, t_1 + \epsilon)},$$

and represent electronic densities.

In this representation, the diagrams have a nice physical interpretation. The Green function $\mathfrak{G}(\mathbf{x}, t; \mathbf{x}', t')$ is the probability amplitude for a particle to get from the \mathbf{x}' -point, at t' -time, to \mathbf{x} -point at t -time. The Feynmann expansion thus states that the electron can get between these two points in all possible ways, weighted by the appropriate amplitude, with the zero-th order Feynmann graph being the direct graph.

- Feynmann rules for the first-order expansion in momentum representation.

The interaction Hamiltonian can be written as

$$(8.3.2.2) \quad \mathbf{H}_I = \frac{1}{\Omega} \sum_{\mathbf{k}_1, \mathbf{k}'_1, \mathbf{q}} \frac{V(\mathbf{q})}{2} \mathbf{c}_{\mathbf{k}_1 - \mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}_2 + \mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}_2} \mathbf{c}_{\mathbf{k}_1}.$$

The only difference between momentum space and realspace is that integrations are carried out over internal momenta, rather than internal position, with momentum being conserved at each vertex, where an interaction links two points at equal time.

With the aid of such basic diagrammatic representation, a depiction can be given for the interacting Green function, which can be diagrammatically represented by a double line with an arrow, ie.

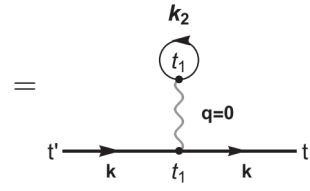
$$(8.3.2.3) \quad \mathfrak{G}(\mathbf{k}, t - t') = \text{t}' \xrightarrow[\mathbf{k}]{} \text{t} = \mathfrak{G}^{(0)} + \mathfrak{G}^{(1)} + \mathfrak{G}^{(2)} + \dots,$$

considering perturbation terms upto quadractic order in the interaction Hamiltonian, $\mathcal{O}(\mathbf{H}_I^2)$:

- Zero-th order The zero-th order expansion yields the following diagrammatic representation

$$(8.3.2.4) \quad \text{t}' \xrightarrow[\mathbf{k}]{G(\mathbf{k}, t - t')} \text{t} = \text{t}' \xrightarrow[\mathbf{k}]{G^{(0)}(\mathbf{k}, t - t')} \text{t}, \quad \text{namely, that the particle with } \mathbf{k}\text{-momentum injected at } t_1\text{-time travels unperturbed until it is extracted at } t\text{-time, without suffering any scattering effects.}$$

- First-order terms: the first order terms can then be interpreted as suffering a single scattering event. From [equation \(8.3.1.3\)](#) there are a total of six first-order terms contributing to $\mathfrak{G}^{(1)}$, which can be represented by six diagrams. Consider, eg., the first pairing,

$$(8.3.2.5) \quad \begin{aligned} & \left\langle \mathcal{T} \left[\overbrace{\mathbf{c}_{\mathbf{k}}(t) \mathbf{c}_{\mathbf{k}_1 - \mathbf{q}}^\dagger(t_1)}^{\mathbf{k}_1 - \mathbf{q} = \mathbf{k}} \overbrace{\mathbf{c}_{\mathbf{k}_2 + \mathbf{q}}(t_1)^\dagger \mathbf{c}_{\mathbf{k}_2}(t_1)}^{\mathbf{k}_2 = \mathbf{k}_2 + \mathbf{q}} \overbrace{\mathbf{c}_{\mathbf{k}_1}(t_1) \mathbf{c}_{\mathbf{k}}^\dagger(t')}^{\mathbf{k}_1 = \mathbf{k}} \right] \right\rangle_{\Phi_0} \\ & \Rightarrow \mathfrak{G}_a^{(1)}(\mathbf{k}, t - t') = - \frac{1}{\Omega} \int_{\mathbb{R}} dt \sum_{\mathbf{k}_1, \mathbf{k}'_1, \mathbf{q}} \frac{V(\mathbf{q})}{2} \mathcal{T} \left[\overbrace{\mathbf{c}_{\mathbf{k}}(t) \mathbf{c}_{\mathbf{k}_1 - \mathbf{q}}^\dagger(t_1)}^{\mathbf{k}_1 - \mathbf{q} = \mathbf{k}} \overbrace{\mathbf{c}_{\mathbf{k}_2 + \mathbf{q}}(t_1)^\dagger \mathbf{c}_{\mathbf{k}_2}(t_1)}^{\mathbf{k}_2 = \mathbf{k}_2 + \mathbf{q}} \overbrace{\mathbf{c}_{\mathbf{k}_1}(t_1) \mathbf{c}_{\mathbf{k}}^\dagger(t')}^{\mathbf{k}_1 = \mathbf{k}} \right] \\ & = \frac{i}{\Omega} \int_{\mathbb{R}} dt \sum_{\mathbf{k}_2} \mathfrak{G}^{(0)}(\mathbf{k}, t - t_1) \mathfrak{G}^{(0)}(\mathbf{k}_2, 0^-) \mathfrak{G}^{(0)}(\mathbf{k}, t_1 - t') \\ & = \end{aligned}$$


The six first-order diagrams then are isomorphic to those shown in [equation \(8.3.2.1\)](#). In fact, only the labels change yielding

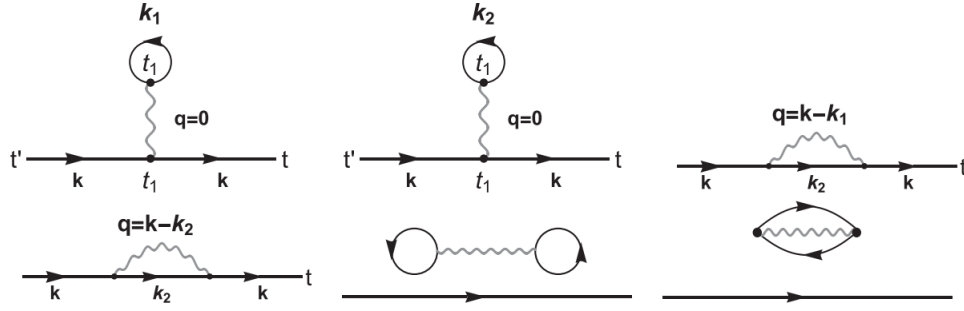


FIGURE 6. All six first-order Feynmann diagrams in momentum representation, derived from the six Wick contraction shown in [equation \(8.3.1.3\)](#).

An examination of the diagrams reveals that there are two pairs of identical diagrams and two disconnected diagrams.

- Second-order terms

The second order terms involve

$$(8.3.2.6) \quad \left\langle \mathcal{T} \left[\mathbf{c}_{\mathbf{k}}(t) \mathbf{H}_I(t_1) \mathbf{H}_I(t_2) \mathbf{c}_{\mathbf{k}}(t_1) \right] \right\rangle_{\Phi_0},$$

whose diagrams contain five fermionic lines with two interaction, yielding -in principle- $5! = 120$ Wick contractions, each represented by a diagram. However, some of these diagrams will be identical to other ones.

From these cases, it is clear that some redundancies arise in the diagrammatic expansion of the Dyson series, since many contractions appear to correspond to the same graph.

8.3.2.1. Redundancies and Simplifications in the Diagrammatic Expansions. There are multiple causes for redundancies in the Feynmann expansion. For example, a third-order expansion of the Green function has, in principle, $7! = 5040$ possible Wick contractions. An identification of diagrammatic redundancies among these contractions is then of the utmost importance. Some ways in which to account for redundancies are treated as follows,

- **Permutation of interaction lines:** first, consider three interaction lines which occur at t_1, t_2, t_3 -times. It is obvious that there $3! = 6$ equivalent Wick contractions, equivalent in the sense that these can be obtained from permutations. Such contractions will yield the exact same product of Green functions with identical diagrams but with different labels. But, since integrations are carried out over all internal lines, albeit at different orders, their interchange is irrelevant. Therefore, in general, in an n -line diagram, a $\frac{1}{n!}$ -factor is associated with terms of order n , namely

$$S(\infty, -\infty) = \sum_{n=0}^{\infty} \frac{i^n}{n!} \mathcal{T} \int_{\mathbb{R}} dt_1 \cdot \int_{\mathbb{R}} dt_n \mathbf{H}_I(t_1) \cdots \mathbf{H}_I(t_n).$$

This $\frac{1}{n!}$ thus cancels itself out with the $n!$ symmetry factor for each diagram.

- **Reversal symmetry of an interaction line:** Each individual interaction line, in the momentum-time representation, has a direction, with two possible orientations. In the position-time representation, internal variables $\mathbf{x}_i, \mathbf{x}_j$, at the vertices, may be interchanged. In either case, such operation will yield diagrams having the same topology. Integrating over these internal variables will yield the same value for the contribution of each diagrams. This entails that, for each interaction line, there is a factor of $V(\mathbf{q})$, instead of $\frac{V(\mathbf{q})}{2}$ that appears in the Hamiltonian.

- Multiplicative factors: the fact that any contraction term of order n has n interaction lines, and $2n + 1$ Green functions, implies that the corresponding n th-order diagram must be associated with a factor i^n .
- Transposition of fermionic operators: Correct Green function operator ordering requires operator switching and gives rise to a factor of -1 . Furthermore, there is a $(-1)^L$ -factor associated with diagrams with L closed fermionic loops.

So far the denominator of [equation \(8.3.1.2\)](#) has been ignored. Upto first order in the interaction Hamiltonian it gives,

(8.3.2.7)

$$\begin{aligned} \langle \Psi_0 | S(\infty, -\infty) | \Psi_0 \rangle &\sim \langle \Psi_0 | \Psi_0 \rangle - \frac{i}{\hbar} \int_{\mathbb{R}} dt_1 \langle \Psi_0 | T \mathbf{H}_I(t_1) | \Psi_0 \rangle \\ &= 1 - \frac{i}{\hbar} \int_{\mathbb{R}} d^4x_1 d^4x_2 V_{\sigma_1\sigma_2\sigma'_1\sigma'_2}^{(x,x')} \langle \Psi_0 | \mathcal{T} \psi_{\sigma_1}^\dagger(x_1) \psi_{\sigma_2}^\dagger(x_2) \psi_{\sigma'_2}(x_2) \psi_{\sigma'_1}(x_1) | \Psi_0 \rangle = F_0 + F_1 \end{aligned}$$

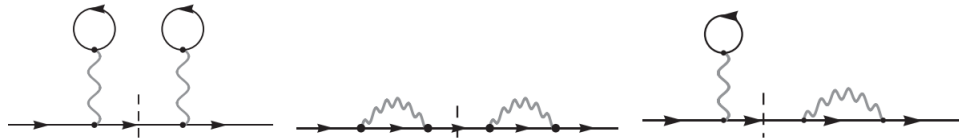
These diagrams do not have external fermionic lines, and consequently its diagrams are closed ones. Terms in the series $\langle \Psi_0 | S(\infty, -\infty) | \Psi_0 \rangle$ are called **vacuum polarization terms**, which have corresponding vacuum polarization diagrams of the closed type. Some contractions of the interaction Hamiltonian in the n th-order perturbative expansion do not link the external fermionic points $\mathbf{c}(t)$ and $\mathbf{c}^\dagger(t')$. Such contractions correspond to diagrams having disconnected closed components. It can be shown that there exists a factorization that separates the disconnected closed-type components from connected ones. This cancellation theorem demonstrates that, in such factorization, the closed diagrams in the denominator exactly cancel the disconnected counterpart diagrams of the perturbation expansion in the numerator. This means that the Green function is then given by the sum of all connected diagrams as

(8.3.2.8)

$$i\mathfrak{G}_{\sigma\sigma'}^{(x,x')} = \sum_{n=0}^{\infty} \frac{1}{n!} \left(\frac{-i}{\hbar} \right)^n \int_{\mathbb{R}} dt_1 \cdots \int_{\mathbb{R}} dt_n \langle \Phi_0 | \mathcal{T} [\mathbf{H}_I(t_1) \cdots \mathbf{H}_I(t_n) \psi_{H\sigma}(\mathbf{x}, t) \psi_{H\sigma'}^\dagger(\mathbf{x}', t_1)] | \Phi_0 \rangle_{\text{connected}}.$$

8.3.3. Infinite Sums: Dyson Series and Self-Energy. It turns out that some diagrams can be constructed by linking diagrams which appear in lower-order expansion with a fermionic line $\mathfrak{G}^{(0)}$. In other words, there are second- and higher-order diagrams which can be decomposed into first-order diagrams. This suggests that identifying building blocks, or irreducible or proper diagrams is useful, since all other diagrams can be generated from this one.

An irreducible diagram can be defined as one which cannot be disconnected into two parts by cutting a single internal electron line. If a diagram is not irreducible, it is said to be reducible. Some examples of second-order diagrams reducible diagrams are



which consist of linked first-order direct and exchange diagrams. All of these can be cut into two diagrams along the dashed line.

An irreducible self-energy diagram (ISED) is defined as the diagram obtained by removing external lines from an irreducible diagram. Equivalently, the ISED's mathematical expression is obtained by deleting two factors of $\mathfrak{G}^{(0)}$ due to external lines. The total ISED sum can be represented as follows

$$\Sigma^* = \text{self-energy diagrams} + \text{Hartree diagram} + \text{Fock diagram} + \dots$$

Then, the complete one-particle interacting Green function can be recovered by reconstructing the clipped diagrams as

$$\text{clipped } G = G^{(0)} + G^{(0)} \Sigma^* G$$

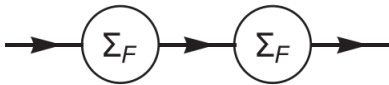
which is the Dyson equation. The appearance of the double line on the right-hand side indicates the self-consistent nature of this equation. The diagrammatic expression of the Dyson equations' main advantage is that it can be terminated at any order and still only include certain classes of diagrams. In many cases, the application of perturbation theory requires summation over whole classes of diagrams in order to obtain meaningful and reliable results, eg. divergences of perturbation terms in the interacting electron gas can be removed by performing an infinite sum of the terms arising from two particular diagrams, the Hartree or direct diagram and the exchange or Fock diagram.

In this case, the Hartree term in particular has $\mathbf{q} = 0$ due to momentum conservation and is cancelled by the background interaction. The Fock-term gives rise to a non-trivial contribution,

(8.3.3.1)

$$i\mathfrak{G}_{\text{Fock}}^{(1\text{-st order})} = i\mathfrak{G}^{(0)}(\mathbf{k}, \omega) \int \frac{d\varepsilon}{2\pi} \frac{1}{\Omega} \sum_{\mathbf{q}} V(\mathbf{q}) \mathfrak{G}^{(0)}(\mathbf{k} - \mathbf{q}, \omega - \varepsilon) \mathfrak{G}^{(0)}(\mathbf{k}, \omega) \equiv \Sigma_F^*(\mathbf{k}, \omega) = \text{Fock diagram}$$

$$i\mathfrak{G}_{\text{Fock}}^{(2\text{-nd order})} \approx \mathfrak{G}^{(0)}(\mathbf{k}, \omega) \Sigma_F^*(\mathbf{k}, \omega) \mathfrak{G}^{(0)}(\mathbf{k}, \omega) \Sigma_F^*(\mathbf{k}, \omega) \mathfrak{G}^{(0)}(\mathbf{k}, \omega) = \text{Fock diagram with two insertions}$$



Then, in order to calculate this exact contribution all higher-order terms must be included. These higher-order terms include repetitions of this precise sequence, generating an infinite series, the **Dyson series**, which involves terms as

$$\begin{aligned} \text{clipped } G(\mathbf{k}, \omega) &= G^{(0)}(\mathbf{k}, \omega) + G^{(0)} \Sigma_F G^{(0)} + G^{(0)} \Sigma_F G^{(0)} \Sigma_F G^{(0)} + \dots \\ &= \mathfrak{G}^{(0)}(\mathbf{k}, \omega) \sum_{n=0}^{\infty} \left(\Sigma_F^*(\mathbf{k}, \omega) \mathfrak{G}^{(0)}(\mathbf{k}, \omega) \right)^n \\ &= \frac{\mathfrak{G}^{(0)}(\mathbf{k}, \omega)}{1 - \Sigma_F^*(\mathbf{k}, \omega) \mathfrak{G}^{(0)}(\mathbf{k}, \omega)} = \frac{1}{[\mathfrak{G}^{(0)}(\mathbf{k}, \omega)]^{-1} - \Sigma_F^*(\mathbf{k}, \omega)}, \end{aligned}$$

which, by inserting the explicit form of the non-interactive Green function in [equation \(8.2.2.25\)](#)

$$\text{namely, } \mathfrak{G}^{(0)} = (\omega - \varepsilon(\mathbf{k}))^{-1}, \text{ yields } \mathfrak{G}(\mathbf{k}, \omega) = \frac{1}{\omega - \varepsilon(\mathbf{k}) - \Sigma_F^*(\mathbf{k}, \omega) \pm i\eta}.$$

$\Sigma_F^*(\mathbf{k}, \omega)$ is the self-energy of the interacting electron gas, obtained by carrying out the summation, a process known as renormalization. In general, as in [section 8.3.3](#), the self-energy is identified with the sum of all proper diagrams coming from all orders of perturbation, namely the ones which cannot be split into parts by breaking

$$\begin{aligned}
\langle \mathcal{T}[\Psi(\mathbf{x}, t) \Psi^\dagger(\mathbf{x}_1, t_1)] \rangle_{\Psi_0} &= \sum_{\mathbf{k} \mathbf{k}_1} \phi_{\mathbf{k}}(\mathbf{x}) \phi_{\mathbf{k}_1}^*(\mathbf{x}_1) \langle \mathcal{T}[\mathbf{c}_{\mathbf{k}}(t) \mathbf{c}_{\mathbf{k}_1}^\dagger(t_1)] \rangle_{\Phi_0} \\
(8.3.4.4) \quad &= \sum_{\mathbf{k} \mathbf{k}_1} \phi_{\mathbf{k}}(\mathbf{x}) \phi_{\mathbf{k}_1}^*(\mathbf{x}_1) \langle \mathcal{T}[\mathbf{c}_{\mathbf{k}} \mathbf{c}_{\mathbf{k}_1}^\dagger] \rangle_{\Phi_0} e^{-i(\varepsilon_{\mathbf{k}} t - \varepsilon_{\mathbf{k}_1} t_1)} \\
&= \sum_{\mathbf{k}} \phi_{\mathbf{k}}(\mathbf{x}) \phi_{\mathbf{k}}^*(\mathbf{x}_1) \delta_{\mathbf{k} \mathbf{k}_1} e^{-i\varepsilon_{\mathbf{k}}(t-t_1)} = \delta(\mathbf{x} - \mathbf{x}_1) \delta(t - t_1).
\end{aligned}$$

- With the equal space-time Green function being

$$(8.3.4.5) \quad i\bar{\mathfrak{G}}^{(0)}(x_2, x_2^+) = -\mathbf{n}^{(0)}(x_2),$$

ie. the density per spin of the non-interacting system.

Then,

$$\begin{aligned}
(8.3.4.6) \quad i\mathfrak{G}_{\sigma\sigma'}^{(1)} &= \frac{i}{\hbar} \sum_{\substack{\sigma_1 \sigma'_1 \\ \sigma_2 \sigma'_2}} V_{\sigma_1 \sigma'_1}^{\sigma_2 \sigma'_2}(x_1, x_2) i\mathfrak{G}_{\sigma\sigma_1}^{(x_1, x_2); (0)} \left[i\mathfrak{G}_{\sigma'_1 \sigma'_2}^{(x_1, x_2); (0)} \mathfrak{G}_{\sigma'_2 \sigma'}^{(x_2, x'); (0)} - i\mathfrak{G}_{\sigma'_1 \sigma'}^{(x_1, x'); (0)} i\mathfrak{G}_{\sigma'_2 \sigma_2}^{(x_2, x_2); (0)} \right] \\
&= \frac{i}{\hbar} \sum_{\substack{\sigma_1 \sigma'_1 \\ \sigma_2 \sigma'_2}} V_{\sigma_1 \sigma'_1}^{\sigma_2 \sigma'_2}(x_1, x_2) i\mathfrak{G}_{\sigma\sigma_1}^{(x_1, x_2); (0)} \delta_{\sigma\sigma_1} \left[i\mathfrak{G}_{\sigma'_1 \sigma'_2}^{(x_1, x_2); (0)} \delta_{\sigma'_1 \sigma'_2} \mathfrak{G}_{\sigma'_2 \sigma'}^{(x_2, x'); (0)} \delta_{\sigma'_2 \sigma'} - i\mathfrak{G}_{\sigma'_1 \sigma'}^{(x_1, x'); (0)} \delta_{\sigma'_1 \sigma'} i\mathfrak{G}_{\sigma'_2 \sigma_2}^{(x_2, x_2); (0)} \delta_{\sigma'_2 \sigma_2} \right] \\
&= \frac{i}{\hbar} \sum_{\substack{\sigma_1 \sigma'_1 \\ \sigma_2 \sigma'_2}} \left[\begin{aligned} &V(|\mathbf{x}_1 - \mathbf{x}_2|) \delta_{\sigma_1 \sigma_2} \delta_{\sigma'_1 \sigma'_2} i\delta(\mathbf{x}_1 - \mathbf{x}_2) \delta(t_1 - t_2) \delta_{\sigma\sigma_1} i\mathfrak{G}_{\sigma'_1 \sigma'_2}^{(x_1, x_2); (0)} \mathfrak{G}_{\sigma'_2 \sigma'}^{(x_2, x'); (0)} \delta_{\sigma'_2 \sigma'} \\ &- V(|\mathbf{x}_1 - \mathbf{x}_2|) \delta_{\sigma_1 \sigma_2} \delta_{\sigma'_1 \sigma'_2} i\delta(\mathbf{x}_1 - \mathbf{x}_2) \delta(t_1 - t_2) \delta_{\sigma\sigma_1} i\mathfrak{G}_{\sigma'_1 \sigma'}^{(x_1, x'); (0)} \delta_{\sigma'_1 \sigma'} i\mathfrak{G}_{\sigma'_2 \sigma_2}^{(x_2, x_2); (0)} \delta_{\sigma'_2 \sigma_2} \end{aligned} \right] \\
&= \frac{i}{\hbar} \left[\begin{aligned} &V(|\mathbf{x}_1 - \mathbf{x}_2|) \delta_{\sigma\sigma'} \delta(\mathbf{x}_1 - \mathbf{x}_2) \delta(t_1 - t_2) i\bar{\mathfrak{G}}^{(x_2, x_2^+); (0)} i\mathfrak{G}^{(x_2, x'); (0)} \\ &- V(|\mathbf{x}_1 - \mathbf{x}_2|) \delta_{\sigma\sigma'} i\delta(\mathbf{x}_1 - \mathbf{x}_2) \delta(t_1 - t_2) i\mathfrak{G}^{(x_1, x'); (0)} i\mathfrak{G}^{(x_2, x_2); (0)} \end{aligned} \right]
\end{aligned}$$

where $2S + 1$ accounts for the sum over spin. Consequently, the integral in the first self-energy term is just the Coulomb interaction of the fermionic test particle with all the particles of the system. The unphysical self-interaction is cancelled by its counterpart in the Fock term. The self-consistency built into the Hartree-Fock theory allows for the replacing of non-interacting Green functions $\mathfrak{G}^{(0)}$ with similar diagrams for the interacting Green function \mathfrak{G}

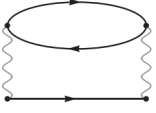
$$(8.3.4.7) \quad \Sigma_{\text{HF}}^*(x, x') = -\frac{1}{\hbar} \delta(t - t') \left[-\delta(\mathbf{x} - \mathbf{x}') (2S + 1) \int_{\mathbb{R}^3} d\mathbf{x}_2 \mathbf{n}(x_2) V(|\mathbf{x} - \mathbf{x}_2|) - V(|\mathbf{x} - \mathbf{x}'|) i\mathfrak{G}(\mathbf{x}, t; \mathbf{x}', t') \right]$$

$$\Sigma_{\text{HF}} = \text{diagram 1} + \text{diagram 2},$$

Diagram 1: A circle with a vertical wavy line extending downwards from its bottom. Diagram 2: A horizontal wavy line with a solid dot at each end.

Note that Σ_{HF}^* is frequency independent, since it depends on $\delta(t - t_1)$. It is purely real, which manifests the Hartree-Fock approximation as a \mathbf{k} -dependent shift in the non-interacting energy poles.

8.3.5. The Two-Particle Green Function and RPA. The Hartree-Fock approach, as a first-order approximation, does not give an adequate account of the density-density response function for the interacting electron gas. These shortfalls are the consequence of the long interaction range of the Coulomb potential, which gives rise to an infrared divergence because of the $\frac{1}{q^2}$ -behaviour. To make progress, higher-order diagrams and their approximations are to be considered. Moreover, screening effects that reduce the long-range nature of the Coulomb interaction arise.



In these higher-order diagrams, a new element appears: the fermionic bubble, also called a polarization loop.

8.3.5.1. A digression on Diagram Scaling. Consider the self-energy diagrams of n -th order in the Fourier (\mathbf{k}, ω) representation. They contain

- n interaction lines,
- $2n$ vertices
- and $2n - 1$ fermionic lines.

All of these quantities have distinct energy scalings, which can be written in terms of k_F as follows:

Scalings of different diagram elements		
Quantity	Per component	Total
Fermionic lines	k_F^{-2}	$(k_F^{-2})^{2n-1}$
Internal lines	k_F^5	$(k_F^5)^n$
Coulomb interaction	k_F^{-2}	$(k_F^{-2})^n$

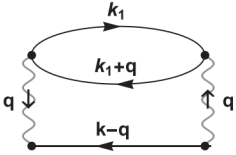
Integration is carried out over n internal momenta, each with three components, and n internal energies. However, given the conservation of momenta at each vertex, there are $2n$ constraints, reducing the number of independent variables to n .

Scaling up all energies and momenta in terms of k_F , for a diagram of order n yields,

$$(8.3.5.1) \quad (k_F^5)^n \times (k_F^{-2})^{2n-1} \times (k_F^{-2})^n = k_F^{-(n-2)} \propto r_s^{(n-2)}$$

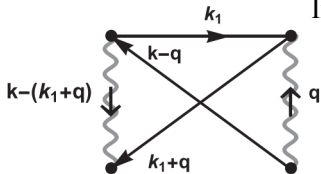
Thus, if $r_s < 1$, it follows that $\Sigma^{*(n+1)} \leq \Sigma^{*(n)}$.

As previously stated, quasiparticles can only be defined near the Fermi surface and that is precisely where Green functions assume large amplitudes, namely, for $k \simeq k_F$. Thus, configurations where all momenta appearing in the self-energy fermionic lines are close to k_F contribute significantly. Consider the following two diagrams.



First, consider this second-order self-energy diagram. For small-enough q , it turns out $V^2(q) \approx q^{-4}$, and that the summation over \mathbf{k}_1 , with $k_1 \simeq k_F$, is unconstrained over a shell of width $2q$. Such a second-order diagram diverges.

In contrast,



in this second-order diagram, for small-enough q and $|\mathbf{k} - \mathbf{q} - \mathbf{k}_1|$ in the interaction lines, the phase space for $k_1 \simeq k_F$ becomes severely restricted.

It can thus be inferred that, because of this relatively large difference of phase space volume over which $k_1 \simeq k_F$, the contribution of the first diagram dwarfs that of the second diagram, ie. that $\frac{\Sigma^*(a)}{\Sigma^*(b)} \gg 1$. This argument can be extended to higher-order diagrams and ascertain that the dominant diagram, in any order, will be the one with maximum phase space volume for all internal momenta to have $k_i \simeq k_F$, and that will be the case when the internal momenta are least constrained to be in the vicinity of the Fermi surface. These diagrams have the highest-degree of divergence, namely, that the highest power of $q^{-2} \rightarrow q^{-2n}$. These dominant diagrams are called **bubble** or **ring** diagrams.

Note that, although the contribution of each individual diagram diverges, their Dyson infinite sum in the high-density regime $r_s < 1$ gives a finite result. This inference describes the screening of the Coulomb potential by

free electrons, giving rise to a Yukawa potential, and is the basis for the random phase approximation.

8.3.5.2. Density-Density Correlation Function. Consider the density-density Green function, ie. the two-particle Green function, defined as

$$(8.3.5.2) \quad \Pi(\mathbf{q}, t) = i \left\langle \text{GS} \left| \mathcal{T} \left[\rho_\sigma(\mathbf{q}, t) \rho_\sigma(-\mathbf{q}, 0) \right] \right| \text{GS} \right\rangle, \text{ where } \begin{array}{l} |\text{GS}\rangle \text{ is the ground state} \\ \rho(\mathbf{q}, t) = \sum_{\mathbf{k}\sigma} \mathbf{c}_{\mathbf{k}\sigma}^\dagger(t) \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}(t) \end{array}$$

$$= i \sum_{\substack{\mathbf{k}\mathbf{k}' \\ \sigma\sigma'}} \langle \text{GS} | \mathcal{T} \left[\mathbf{c}_{\mathbf{k}\sigma}^\dagger(t) \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}(t) \mathbf{c}_{\mathbf{k}'\sigma'}^\dagger(0) \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}(0) \right] | \text{GS} \rangle.$$

This quantity describes the propagation of a particle-hole pair created at time 0:

- the particle is created in the state $|\mathbf{k}'\sigma'\rangle$,
- and the hole is created by annihilating a particle in the $|\mathbf{k}' - \mathbf{q}\sigma'\rangle$.
- The pair is ultimately annihilated in states $|\mathbf{k}\sigma\rangle$ and $|\mathbf{k} + \mathbf{q}\sigma\rangle$ at time t .

By analogy to the case of the single-particle Green function, this propagator will therefore have poles at energies corresponding to stable particle-hole excitations of the system. The summation over \mathbf{k}, \mathbf{k}' describes the propagation of a superposition of many particle-hole excitations. Such superpositions are, in fact, collective excitations. This is, in essence, a density fluctuation, so that [equation \(8.3.5.2\)](#) is also called the density-fluctuation propagator. For an interacting system, the perturbative expansion reads

$$(8.3.5.3) \quad \Pi(\mathbf{q}, t) = i \frac{\langle \Phi_0 | \mathcal{T} \left[\rho_\sigma(\mathbf{q}, t) \rho_\sigma(-\mathbf{q}, 0) \mathcal{S}(\infty, -\infty) \right] | \Phi_0 \rangle}{\langle \Phi_0 | \mathcal{S}(\infty, -\infty) | \Phi_0 \rangle}$$

with all quantities appearing in the interaction representation, and all expectation values being taken with respect to the non-interacting ground state, allowing for analog diagrammatic rules previously derived for any order in the perturbation expansion, including vacuum elimination.

8.3.5.3. Excitations spectrum of the Non-Interacting System. Consider the zero-th-order in the perturbative expansion given by [equation \(8.3.5.3\)](#), $\Pi_0(\mathbf{q}, t)$, which is just the two-particle correlation function of the non-interacting system. $\Pi_0(\mathbf{q}, t)$ provides information about the excitation spectrum of the Sommerfeld-Fermi gas, with a fixed number of particles N , which can be done using Wick's theorem. There are two possible contractions

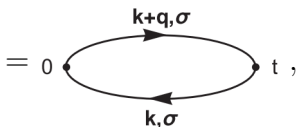
$$(8.3.5.4) \quad \langle \mathcal{T} \left[\overline{\mathbf{c}_{\mathbf{k}\sigma}^\dagger(t) \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}(t)} \overline{\mathbf{c}_{\mathbf{k}'\sigma'}^\dagger(0) \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}(0)} \right] \rangle, \quad \langle \mathcal{T} \left[\overline{\mathbf{c}_{\mathbf{k}\sigma}^\dagger(t) \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}(t) \mathbf{c}_{\mathbf{k}'\sigma'}^\dagger(0) \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}(0)} \right] \rangle$$

The first contraction requires $\mathbf{q} = 0$ for it to be non-zero, yielding

$$(8.3.5.5) \quad \left(\sum_{\mathbf{k}\sigma} \langle \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \rangle \right) \left(\sum_{\mathbf{k}'\sigma'} \langle \mathbf{c}_{\mathbf{k}'\sigma'}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'} \rangle \right) \propto N^2.$$

It describes the self-correlation of the uniform electron density and can be ignored. Thus, there is only one possible contraction. Using the definition $\mathfrak{G}^{(0)} = -i \langle \mathcal{T} c c^\dagger \rangle$ and performing three fermionic swaps, the zero-th-order density-density correlation function reads

$$(8.3.5.6) \quad \Pi_0(\mathbf{q}, t) = i \sum_{\mathbf{k}\sigma} \mathfrak{G}_\sigma^{(0)}(\mathbf{k} + \mathbf{q}, t) \mathfrak{G}_\sigma^{(0)}(\mathbf{k}, -t)$$



where the \mathbf{q} -momentum is an external one, whereas the \mathbf{k} -momentum is an internal momentum, which together with an internal spin are summed over.

Note that a hole can be thought of as an electron propagating backward in time, which is the reason such diagram is known as a particle-hole loop, or a polarization loop. Now, consider the Fourier transform with respect

to time in order to work out the momentum-energy representation of the density-density correlator, which will yield a convolution in energy

$$\begin{aligned}
 \Pi_0(\mathbf{q}, \omega) &= \int_{-\infty}^{\infty} dt e^{i\omega t} \Pi_0(\mathbf{q}, \omega) = i \sum_{\mathbf{k}, \sigma} \int \frac{d\varepsilon}{2\pi} \mathfrak{G}_{\sigma}^{(0)}(\mathbf{k}, \varepsilon) \mathfrak{G}_{\sigma}^{(0)}(\mathbf{k} + \mathbf{q}, \varepsilon + \omega) \\
 (8.3.5.7) \quad &= \sum_{\mathbf{k}, \sigma} \int \frac{d\varepsilon}{2\pi i} \frac{1}{\varepsilon - \varepsilon_{\mathbf{k}} + i\eta \text{sgn}(\varepsilon_{\mathbf{k}})} \frac{1}{\varepsilon + \omega - \varepsilon_{\mathbf{k}+\mathbf{q}} + i\eta \text{sgn}(\varepsilon_{\mathbf{k}+\mathbf{q}})} \\
 &= -2 \sum_{\mathbf{k}, \sigma} \left[\frac{\Theta(-\varepsilon_{\mathbf{k}+\mathbf{q}}) \Theta(\varepsilon_{\mathbf{k}})}{\omega - \varepsilon_{\mathbf{k}+\mathbf{q}} + \varepsilon_{\mathbf{k}} + i\eta} - \frac{\Theta(\varepsilon_{\mathbf{k}+\mathbf{q}}) \Theta(-\varepsilon_{\mathbf{k}})}{\omega - \varepsilon_{\mathbf{k}+\mathbf{q}} + \varepsilon_{\mathbf{k}} - i\eta} \right],
 \end{aligned}$$

where in the last step the integral was carried out in the complex plane. Note that both $\varepsilon_{\mathbf{k}+\mathbf{q}}$ and $\varepsilon_{\mathbf{k}}$ have the same sign, so that they are both positive/negative, the poles in the complex ε -plane are either in the lower or upper half-plane. Taking the contour in the pole-free half, and noting the asymptotic value of the integral is $\mathcal{O}(\varepsilon^{-2})$, yields zero. The integral is non-zero if $\varepsilon_{\mathbf{k}+\mathbf{q}}$ and $\varepsilon_{\mathbf{k}}$ have opposite signs, thus yielding the last line.

The physical interpretation of the Θ -functions is straightforward,

- For $\Theta(-\varepsilon_{\mathbf{k}})$, it follows that $n_p = \mathbf{n}_{\mathbf{k}}^0$ ie. a particle in the Fermi sea.
- In contrast, for $\Theta(\varepsilon_{\mathbf{k}})$, it follows that $n_p = 1 - \mathbf{n}_{\mathbf{k}}^0$ ie. a hole outside the Fermi sea.

Consequently, [equation \(8.3.5.7\)](#)'s first term is the retarded propagator whilst the second term is the advanced propagator. Furthermore the retarded propagator of a particle-hole pair, with $|\mathbf{k} + \mathbf{q}| > k_F$ and $|\mathbf{k}| < k_F$, correspond to a positive excitation energy, $\omega_{\mathbf{q}} = \varepsilon_{\mathbf{k}+\mathbf{q}} - \varepsilon_{\mathbf{k}} > 0$.

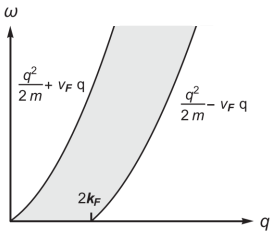
8.3.5.4. The Dynamical Form Factor. The imaginary part of Π_0 is the dynamical form factor,

$$\begin{aligned}
 \mathcal{S}_0(\mathbf{q}, \omega) &= \frac{1}{\pi} \Im \Pi_0(\mathbf{q}, \omega) \\
 (8.3.5.8) \quad &= 2 \int \frac{d\mathbf{k}}{(2\pi)^3} \mathbf{n}^0(\mathbf{k}) (1 - \mathbf{n}^0(\mathbf{k} + \mathbf{q})) \delta(\omega - \varepsilon_{\mathbf{k}+\mathbf{q}} + \varepsilon_{\mathbf{k}}).
 \end{aligned}$$

Hence, a non-zero \mathbf{n}^0 implies that the $|\mathbf{k} \dots\rangle$ -state is occupied and that an excitation of the particle in it will leave behind a hole, while a null $\mathbf{n}^0(\mathbf{k} + \mathbf{q})$ indicates an empty state that can be filled with a particle of $(\mathbf{k} + \mathbf{q})$ -momentum. Conservation of energy thus implies that the dynamical form factor provides information about real or absorptive processes. This is just a manifestation of all possible particle-hole pairs that carry \mathbf{q} -momentum and satisfy energy conservation, since it defines the real excitations of the Fermi gas.

In order to determine the spectrum, one must inspect the energy-conserving delta function

$$(8.3.5.9) \quad 0 \leq \omega = \frac{q^2}{2m} + \mathbf{v}_{\mathbf{k}} \cdot \mathbf{q} \Rightarrow \begin{cases} \omega_{\max} = \frac{q^2}{2m} + v_F q & q \leq 2k_F \\ \omega_{\min} = \frac{q^2}{2m} - v_F q & q \geq 2k_F \end{cases}.$$



The electron-hole excitation spectrum is a strip between two parabolas, ie. the gray area shown in this figure. Then, at a fixed \mathbf{q} , the dynamical form factor gives a wide range of allowed energies. It is then clear that these are incoherent excitations for particle-hole events, an excitation in the $|\mathbf{q}, \omega\rangle$ -state can immediately decay into excitations with energies lower than ω , but with same \mathbf{q} -momentum for the Sommerfeld gas.

This is in clear contrast with the coherent single-particle excitation events manifest in the corresponding spectral function,

$$(8.3.5.10) \quad A(\mathbf{k}, \omega) = \delta[\omega - \varepsilon(\mathbf{k})].$$

8.3.5.5. Virtual Excitations and Lindhard functions. If the energy conservation condition is not realized, the corresponding contribution is reactive, not dissipative, and it only appears in the real part of the density-density correlation function. The intermediate state is only virtual. In effect, consider the real part of the zero-th-order density-density correlator function given by **equation (8.3.5.7)**, which can be written as

$$(8.3.5.11) \quad \Re \Pi_0(\mathbf{q}, \omega) = 2 \int \frac{d\mathbf{k}}{(2\pi)^3} \frac{\mathbf{n}^{(0)}(\mathbf{k}) - \mathbf{n}^{(0)}(\mathbf{k} + \mathbf{q})}{\omega - \varepsilon_{\mathbf{k}+\mathbf{q}} + \varepsilon_{\mathbf{k}}},$$

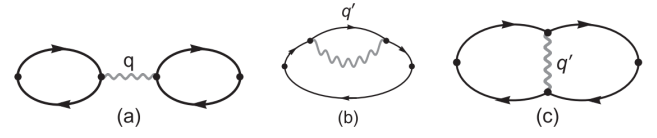
where $\mathcal{D}(\varepsilon_F)$ is the density of states at ε_F and where $F(\tilde{q}, \tilde{\omega})$ is the Lindhard function. The previous relationship carries information about the probability and phase of all possible virtual processes.

8.3.5.6. Interacting Systems and the Random Phase Approximation (RPA).

• Irreducible Polarization Bubble

By analogy to the treatment of the self-energy, an exact irreducible polarization bubble $\Pi(\mathbf{q}, \omega)$ can be defined as the sum over all diagrams that cannot be cut into two by splitting a single interaction line (an external line), ie.

The first-order diagrams for the density-density correlator are given by



(a) (b) (c)

$$(8.3.5.12) \quad \Rightarrow \quad \Pi(\mathbf{q}, \omega) = \text{bubble} = \text{loop} + \text{loop with wavy line} + \text{two loops} + \dots$$


Furthermore, the exact charge susceptibility, which is the response to external Coulomb perturbations, can be obtained from a Dyson series on the irreducible polarization as

$$(8.3.5.13) \quad \chi(\mathbf{q}, \omega) = \text{bubble} = \text{bubble} + \text{bubble} \text{---} \text{wavy line} \text{---} \text{bubble} + \dots$$

$$\chi(\mathbf{q}, \omega) = \frac{\Pi(\mathbf{q}, \omega)}{1 - V_c(\mathbf{q})\Pi(\mathbf{q}, \omega)}.$$

• Screened interaction

Moreover, an exact effective screened Coulomb interaction can be diagrammatically represented by a double wiggly line, and is given by a summation over concatenated irreducible polarization bubbles attached with internal line with several external interaction lines attached, namely



(8.3.5.14) or

$$V_{\text{eff}}^{\text{exact}}(\mathbf{q}, \omega) = \frac{V_c(\mathbf{q})}{1 - V_c(\mathbf{q}) \Pi(\mathbf{q}, \omega)}.$$

• Random Phase Approximation

Consider the first-order diagrams for the density-density correlator Π , given in the first line of **equation (8.3.5.12)**. A phase space analysis shows that the diagram (a) has the dominant contribution to the correlation function, where each of the two Π_0 bubbles has an unconstrained sum. By extension, it can be argued that the dominant contribution at any order consists of interacting Π_0 bubbles. A justifiable approximation can be made in terms of Dyson-like sums involving Π_0 bubbles, which is known as the **random phase approximation (RPA)**:

$$\begin{aligned}
\chi^{\text{RPA}}(\mathbf{q}, \omega) &= \text{Diagram 1} = \text{Diagram 2} + \text{Diagram 3} + \dots \\
(8.3.5.15) \quad &= \Pi_0(\mathbf{q}, \omega) + \Pi_0(\mathbf{q}, \omega)[V_c(\mathbf{q})\Pi_0(\mathbf{q}, \omega)] + \Pi_0(\mathbf{q}, \omega)[V_c(\mathbf{q})\Pi_0(\mathbf{q}, \omega)]^2 + \dots, \\
&= \frac{\Pi_0(\mathbf{q}, \omega)}{1 - V_c(\mathbf{q})\Pi_0(\mathbf{q}, \omega)}
\end{aligned}$$

Physically relevant information can be extracted from the imaginary part of χ^{RPA} , which is given by

$$(8.3.5.16) \quad \Im[\chi^{\text{RPA}}] = \frac{\Im\Pi_0}{(1 - V_c(\mathbf{q})\Re\Pi_0)^2 + V_c^2(\mathbf{q})(\Im\Pi_0)^2}$$

The numerator thus indicates that the electron-hole excitations spectrum carries over to the interacting case. Moreover, outside the electron-hole continuum ie. where $\Im[\Pi_0] = 0$, new excitations appear when $1 - V_c(\mathbf{q})\Re\Pi_0 = 0$. These excitations are plasmons, which manifest in Coulomb interacting systems. Said excitations are also manifest as zero-sound excitations in systems with short-range interactions.

The random phase approximation can be extended to the effective screened interaction, by including external fermionic lines, ie.

$$\begin{aligned}
&\text{Diagram 4} = \text{Diagram 5} + \text{Diagram 6} + \text{Diagram 7} + \dots \\
(8.3.5.17) \quad V^{\text{RPA}}(\mathbf{q}, \omega) &= V_c(\mathbf{q}) + V_c(\mathbf{q})[V_c(\mathbf{q})\Pi_0(\mathbf{q}, \omega)] + V_c(\mathbf{q})[V_c(\mathbf{q})\Pi_0(\mathbf{q}, \omega)]^2 \\
&= \frac{V_c(\mathbf{q})}{1 - V_c(\mathbf{q})\Pi_0(\mathbf{q}, \omega)} = \frac{V_c(\mathbf{q})}{\varepsilon(\mathbf{q}, \omega)}, \text{ where } \varepsilon^{-1}(\mathbf{q}, \omega) = 1 + V_c(\mathbf{q})\chi^{\text{RPA}},
\end{aligned}$$

where ε is the dielectric function. The effective interaction is ω -dependent, which means that it acquires a time-dependence, in contrast to the instantaneous character of the original non-relativistic Coulomb interaction. This is physically sound, since interactions via density fluctuations involve time delay.

In particular, consider static screening by examining the effective interaction in the ($\omega = 0$)-limit. Note that

$$(8.3.5.18) \quad \begin{aligned} \Im\Pi_0 &= 0 \\ \Re\Pi_0 &= \mathcal{D}(\varepsilon_F) = \frac{mk_F}{\pi^2\hbar^2} = \frac{3n}{2\varepsilon_F} \end{aligned} \Rightarrow 1 + V_c(\mathbf{q})\Pi_0 \underset{\omega, q \rightarrow 0}{=} 1 + \frac{6\pi ne^2}{q^2\varepsilon_F} \Rightarrow V^{\text{RPA}}(\mathbf{q}, \omega) \underset{\omega, q \rightarrow 0}{=} \frac{4\pi e^2}{q^2 + \kappa^2},$$

with $\kappa = \varepsilon_{\text{TF}}^{-1} = \sqrt{\frac{6\pi ne^2}{\varepsilon_F}}$ is the inverse Thomas-Fermi screening length. The effective interaction potential thus becomes a short-range Yukawa-type interaction, namely

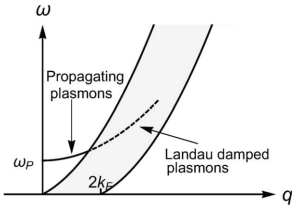
$$V_{\text{eff}}(q \rightarrow 0, 0) \approx \frac{e^{-\kappa r}}{r}.$$

Consequently, in the presence of the polarizable medium - ie. the electron sea - the long-range Coulomb interaction actually becomes short-range. To physically interpret this screening mechanism, consider the scenario where a negative point charge is introduced. Then, the electrons in its vicinity will be repelled from this region, creating a region around this point of lower-than-average density, namely a screening region of positive charge. At long distances, another electron sees not only the initial negative charge, but also the screening region around it, which *almost* cancel each other, hence making the Coulomb interaction effectively short-ranged.

8.3.5.7. Plasma Oscillations. Now, consider the ($\omega \neq 0$)-frequency. In particular, consider the domain of high-frequencies, long-wavelengths and low temperatures, and in particular the region outside the electron-hole continuum, wherein $\Im\Pi_0 = 0$. The dielectric function can be written down explicitly as

$$\begin{aligned}
 \epsilon^{\text{RPA}}(\mathbf{q}, \omega) &= 1 - \frac{4\pi e^2}{q^2} \Pi_0(\mathbf{q}, \omega) = 1 - \frac{4\pi e^2}{q^2} \int_{\mathbb{R}^3} \frac{2d\mathbf{k}}{(2\pi)^3} \frac{\mathbf{n}^0(\mathbf{k}) - \mathbf{n}^0(\mathbf{k} + \mathbf{q})}{\omega - \epsilon_{\mathbf{k}+\mathbf{q}} + \epsilon_{\mathbf{k}}} \\
 &\stackrel{\substack{\omega \gg \omega_F q \\ q \ll k_F \\ k_B T \ll k_F}}{=} 1 - \frac{4\pi e^2}{q^2} \int_{\mathbb{R}^3} \frac{2d\mathbf{k}}{(2\pi)^3} \frac{\mathbf{q} \cdot \nabla_{\mathbf{k}} \epsilon_{\mathbf{k}} \frac{dn(\epsilon)}{d\epsilon}}{\frac{\mathbf{q} \cdot \nabla_{\mathbf{k}} \epsilon_{\mathbf{k}}}{m}} \\
 (8.3.5.19) \quad &= 1 - \frac{4\pi e^2}{q^2} \int_{\mathbb{R}^3} \frac{2d\mathbf{k}}{(2\pi)^3} \frac{v_{\mathbf{k}} q \cos \theta \left. \frac{dn(\epsilon)}{d\epsilon} \right|_{\epsilon_F}}{\mathbf{q} \cdot \mathbf{v}_{\mathbf{k}}} \\
 &= 1 - \frac{4\pi e^2}{q^2} \int_{\mathbb{R}^3} \frac{2d\mathbf{k}}{(2\pi)^3} \frac{v_{\mathbf{k}} q \cos \theta \left. \frac{dn(\epsilon)}{d\epsilon} \right|_{\epsilon_F}}{v_{\mathbf{k}} q \cos \theta} \\
 &= 1 - \frac{4\pi e^2}{q^2} \frac{1}{2\pi^2} \int dk k^2 \delta(\epsilon_{\mathbf{k}} - \epsilon_F) \int_{-1}^{+1} dx \frac{v_{\mathbf{k}} q x}{\omega - v_{\mathbf{k}} q x} \\
 &= 1 - \frac{\omega_p^2}{\omega^2} \left[1 + \frac{3}{5} \left(\frac{q v_F}{\omega} \right)^2 \right].
 \end{aligned}$$

The plasmon dispersion is obtained from the zeros of the dielectric function, yielding



$$\epsilon^{\text{RPA}}(\mathbf{q}, \omega) = 0 \Rightarrow \omega(q) \simeq \omega_P + \frac{3}{10} \frac{v_F^2}{\omega_p} q^2.$$

Once the dispersion curve of plasmons enters the particle-hole continuum, wherein $\Im\Pi_0 \neq 0$, plasmons become unstable to decay into an electron-hole pair. Note that $\Im\chi(\mathbf{q}, \omega) \neq 0$ only in the shaded area.

8.3.6. *Finite Temperature Green Functions.* In a large number of application in many-body physics, the effects of finite temperature are to be included, wherein Green functions now read

$$(8.3.6.1) \quad \mathfrak{G}_{AB}(t) = \langle \mathbf{A}(t) \mathbf{B}(0) \rangle = -i \text{Tr} \left[\frac{1}{\mathcal{Z}} e^{-\beta \mathbf{H}} \mathbf{A}(t) \mathbf{B}(0) \right], \text{ where } \mathcal{Z} = \text{Tr} \left[e^{-\beta \mathbf{H}} \right], \quad \mathbf{H} \rightarrow \mathbf{H} - \mu \mathbf{N}$$

and where the $\mathbf{A}(t)$ - and $\mathbf{B}(t_1)$ - are in the Heisenberg picture.

The idea for finite-temperature many-body physics is to separate the Hamiltonian into a solvable part, \mathbf{H}_0 , and a perturbative quantity \mathbf{H}' , so that perturbation theory is applicable. Then, an expansion in powers of \mathbf{H}' can be done. Note that at zero temperature, the traces over ρ are replaced with ground-state averages, which allows to transform the problem to the interaction picture, thanks to the Gell-Mann-Low theorem and the scattering matrix, and to solve it perturbatively within the non-interaction framework.

8.3.6.1. Perturbative Expansion of ρ . The ρ -density operator satisfies the Bloch equation of motion

$$(8.3.6.2) \quad -\frac{\partial \rho}{\partial \beta} = \mathbf{H} \rho.$$

which is similar to the Schrödinger equation of motion, but with the time variable it replaced with β and with $\mathbf{H} \rightarrow \mathbf{H} - \mu \mathbf{N}$. Consequently, it is possible to perform a perturbative expansion in powers of \mathbf{H}' , with said expansion not being along the real-time axis but along the inverse temperature axis. This, a priori, entails a loss of the notion of time-ordered product of operators which evolve along two orthogonal axes.

8.3.6.2. Analytic Continuation. From a mathematical perspective, it would be desirable to include the thermal state $\rho = e^{-\beta\mathbf{H}}$ explicitly in the group of time-evolution operators. This requires the analytic continuation of the time-evolution group to complex times $t \rightarrow t + i\tau = z$ and to introduce new (imaginary)-time evolution operators $e^{iz\mathbf{H}}$. Consider the corresponding Green function

$$\begin{aligned} t > 0, \mathfrak{G}_{AB}(t) &= -\frac{i}{\mathcal{Z}} \text{Tr} \left[e^{-\beta\mathbf{H}} e^{i\mathbf{H}t} \mathbf{A} e^{-i\mathbf{H}t} \mathbf{B} \right] \\ &\rightarrow \mathfrak{G}_{AB}(z) = -\frac{i}{\mathcal{Z}} \text{Tr} \left[e^{-(\beta+\tau)\mathbf{H}} e^{i\mathbf{H}t} \mathbf{A} e^{-i\mathbf{H}t} \mathbf{B} \right] \end{aligned} \quad (8.3.6.3)$$

$$\begin{aligned} t < 0, \mathfrak{G}_{AB}(t) &= \pm \frac{i}{\mathcal{Z}} \text{Tr} \left[e^{-\beta\mathbf{H}} e^{-i\mathbf{H}t} \mathbf{A} e^{+i\mathbf{H}t} \mathbf{B} \right] \\ &\rightarrow \mathfrak{G}_{AB}(z) = \pm \frac{i}{\mathcal{Z}} \text{Tr} \left[e^{-(\beta-\tau)\mathbf{H}} e^{i\mathbf{H}t} \mathbf{A} e^{-i\mathbf{H}t} \mathbf{B} \right] \end{aligned}$$

where \pm corresponds to fermions and bosons, respectively. Furthermore, $e^{iz\mathbf{H}}$ is a well-defined operator on the strip given by

$$(8.3.6.4) \quad \begin{cases} -\beta \leq \tau \leq 0, t \in \mathbb{R} & t > 0 \\ 0 \leq \tau \leq \beta, t \in \mathbb{R} & t < 0 \end{cases},$$

thus defining the domains of complex time for which the preceding Green function is analytic. Note that this shows that the Green function of complex time is discontinuous on the whole real axis of time.

Physical Context

In an informal way, the $\pm\beta$ -boundaries along the imaginary time axis can be associated with decoherence induced by thermal fluctuations, namely the thermal fluctuations introduced by an uncertainty β^{-1} in energies, in which case $\tau_T = \beta^{-1}$ represents a characteristic time of a thermal fluctuation. Processes of duration longer than τ_T lose their phase coherence, so coherent quantum processes are limited within a world of finite temporal extent β .

8.3.6.3. Imaginary Time. The solution to this problem was first proposed by Takeo Matsubara and relies on taking advantage of the similarity between the operators $e^{-\beta\mathbf{H}}$ and $\mathbf{U}(t) = e^{-i\mathbf{H}t}$, but limiting the domain to the imaginary axis only, wherein

$$\rho = \frac{e^{-\beta\mathbf{H}}}{\text{Tr} e^{-\beta\mathbf{H}}} = \frac{\mathbf{U}(-i\beta)}{\text{Tr}\{\mathbf{U}\}(-i\beta)}.$$

The Matsubara approach is to treat time as an imaginary temperature. For time $t = -i\tau$, the Green functions are well-defined in the interval $-\beta \leq \tau \leq \beta$. Thus, the real-time evolution operators and elements can be replaced by its appropriate imaginary-time counterparts,

$$-i\tau, \tau \in \mathbb{R} : e^{-i\mathbf{H}t} \rightarrow e^{-i\mathbf{H}(-i\tau)} = e^{-\mathbf{H}\tau},$$

which evolves in imaginary- rather than real-time. This allows for the tractable expansion of the Green function in terms of these new imaginary-time operators. Note that these operators can then be analytically-continued to recover the real-time ones.

This approach renders finite-temperature many-body physics to be just as manageable as its zero-temperature counterpart.

8.3.6.4. Representations. To develop the imaginary-time formulation of the Heisenberg and interaction representations in terms of $t \rightarrow -i\tau$, consider the real-time Schrödinger equation,

$$(8.3.6.5) \quad \begin{aligned} i\partial_t |\Psi_S\rangle &= \mathbf{H} |\Psi_S\rangle \longrightarrow -\partial_\tau |\Psi_S\rangle = \mathbf{H} |\Psi_S\rangle \\ |\Psi_S(t)\rangle &= e^{i\mathbf{H}t} |\Psi_S(0)\rangle \longrightarrow |\Psi_S(\tau)\rangle = e^{-\mathbf{H}\tau} |\Psi_S(0)\rangle \\ \mathbf{O}_H(t) &= e^{i\mathbf{H}t} \mathbf{O}_s e^{-i\mathbf{H}t} \longrightarrow \mathbf{O}_H(\tau) = e^{\mathbf{H}\tau} \mathbf{O}_s e^{-\mathbf{H}\tau} \implies \partial_\tau \mathbf{O}_H = [\mathbf{H}, \mathbf{O}_H]. \end{aligned}$$

Example

Consider the free fermionic particle Hamiltonian

$$\mathbf{H} = \sum_{\mathbf{k}} \varepsilon_{\mathbf{k}} \mathbf{c}_{\mathbf{k}}^\dagger \mathbf{c}_{\mathbf{k}} \longrightarrow \begin{aligned} \partial_\tau \mathbf{c}_{\mathbf{k}} &= [\mathbf{H}, \mathbf{c}_{\mathbf{k}}] \rightarrow \mathbf{c}_{\mathbf{k}}(\tau) = e^{-\varepsilon_{\mathbf{k}}\tau} \mathbf{c}_{\mathbf{k}} \\ \partial_\tau \mathbf{c}_{\mathbf{k}}^\dagger &= [\mathbf{H}, \mathbf{c}_{\mathbf{k}}^\dagger] \rightarrow \mathbf{c}_{\mathbf{k}}^\dagger(\tau) = e^{\varepsilon_{\mathbf{k}}\tau} \mathbf{c}_{\mathbf{k}}^\dagger \end{aligned}$$

One should take note of an important difference between real-time and imaginary-time formalism, since in the latter the Heisenberg creation and annihilation operators are not hermitian conjugates, ie. $\mathbf{c}_{\mathbf{k}}^\tau = (\mathbf{c}_{\mathbf{k}}(-\tau))^\dagger \neq (\mathbf{c}_{\mathbf{k}}(\tau))^\dagger$.

8.3.6.5. Matsubara Functions. Green functions for operators with imaginary-time arguments are referred to as Matsubara functions, which are only used for thermal equilibrium calculations at finite temperature, and are defined as follows

$$(8.3.6.6) \quad \mathfrak{G}_{\mathbf{AB}}(\tau_1, \tau_2) = -\left\langle \mathcal{T}_\tau \mathbf{A}(\tau_1) \mathbf{B}(\tau_2) \right\rangle = -\frac{1}{\mathcal{Z}} \text{Tr} \left[e^{-\beta \mathbf{H}} \mathbf{A}(\tau_1) \mathbf{B}(\tau_2) \right],$$

where \mathcal{T}_τ is the chronological time ordering, whose action is defined as follows

$$\mathcal{T}_\tau \mathbf{A}(\tau_1) \mathbf{B}(\tau_2) = \Theta(\tau_1 - \tau_2) \mathbf{A}(\tau_1) \mathbf{B}(\tau_2) - \Theta(\tau_2 - \tau_1) \mathbf{B}(\tau_2) \mathbf{A}(\tau_1)$$

for fermions³².

An important property of the Matsubara functions is that it is periodic for bosons and antiperiodic for fermions over an interval β , ie.

$$(8.3.6.7) \quad \mathfrak{G}_{\mathbf{AB}}(\tau) = \begin{cases} -\mathfrak{G}_{\mathbf{AB}}(\tau \pm \beta), & \text{fermions} \\ \mathfrak{G}_{\mathbf{AB}}(\tau \pm \beta), & \text{bosons} \end{cases}, \text{ for } \tau > (<) 0$$

This antiperiodicity of the Matsubara-Green function for the fermionic case, can be used to extend the definition of the Matsubara function $\mathfrak{G}_{\mathbf{AB}}(\tau)$ to the entire imaginary time axis. An important type of Matsubara-Green function are the single-particle fermionic correlation functions \mathfrak{G} , defined as

³²Note that the Matsubara function is a function of the time difference alone, ie. $\mathfrak{G}_{\mathbf{AB}}(\tau_1, \tau_2) = \mathfrak{G}_{\mathbf{AB}}(\tau_1 - \tau_2, 0)$. In effect, note that for $\tau_1 > \tau_2$

$$\begin{aligned} \mathfrak{G}_{\mathbf{AB}}(\tau_1, \tau_2) &= -\frac{\Theta(\tau_1 - \tau_2)}{\mathcal{Z}} \text{Tr} \left[e^{-\beta \mathbf{H}} e^{\mathbf{H}\tau_1} \mathbf{A} e^{-\mathbf{H}\tau_1} e^{\mathbf{H}\tau_2} \mathbf{B} e^{-\mathbf{H}\tau_2} \right] + \frac{\Theta(\tau_2 - \tau_1)}{\mathcal{Z}} \text{Tr} \left[e^{-\beta \mathbf{H}} e^{\mathbf{H}\tau_2} \mathbf{A} e^{-\mathbf{H}\tau_2} e^{\mathbf{H}\tau_1} \mathbf{B} e^{-\mathbf{H}\tau_1} \right] \\ &= -\frac{\Theta(\tau_1 - \tau_2)}{\mathcal{Z}} \text{Tr} \left[e^{-\beta \mathbf{H}} e^{\mathbf{H}(\tau_1 - \tau_2)} \mathbf{A} e^{-\mathbf{H}(\tau_1 - \tau_2)} \mathbf{B} \right] + \frac{\Theta(\tau_2 - \tau_1)}{\mathcal{Z}} \text{Tr} \left[e^{-\beta \mathbf{H}} \mathbf{B} e^{\mathbf{H}(\tau_1 - \tau_2)} \mathbf{A} e^{-\mathbf{H}(\tau_1 - \tau_2)} \right] \\ &= \mathfrak{G}_{\mathbf{AB}}(\tau_1 - \tau_2, 0) = \mathfrak{G}_{\mathbf{AB}}(\tau), \end{aligned}$$

with a similar result holding for $\tau_2 > \tau_1$, with the condition $-\beta < \tau < \beta$ needed in order to assure the convergence.

$$(8.3.6.8) \quad \mathfrak{G}_{\sigma\sigma'}(\mathbf{x}, \tau; \mathbf{x}', \tau') = -\left\langle \mathcal{T}_\tau \left[\psi_\sigma(\mathbf{x}, \tau) \psi_{\sigma'}^\dagger(\mathbf{x}', \tau') \right] \right\rangle \text{ in real space.}$$

Consider for instance the Hamiltonian of a non-interacting fermionic system, which reads

$$(8.3.6.9) \quad \mathbf{H} = \sum_{\mu} \varepsilon_{\mu} \mathbf{c}_{\mu}^{\dagger} \mathbf{c}_{\mu}, \quad \text{where } \varepsilon_{\lambda} = E_{\lambda} - \mu \text{ is the one-particle energy shifted by the chemical potential.}$$

Its corresponding Matsubara function is

$$(8.3.6.10) \quad \begin{aligned} \mathfrak{G}_{\lambda\lambda'}(\tau - \tau') &= -\left\langle \mathcal{T}_\tau \mathbf{c}_{\lambda}(\tau) \mathbf{c}_{\lambda'}^{\dagger}(\tau') \right\rangle \\ &= -\frac{1}{\mathcal{Z}} \text{Tr} e^{-\beta \mathbf{H}} \left[\Theta(\tau - \tau') \mathbf{c}_{\lambda}(\tau) \mathbf{c}_{\lambda'}^{\dagger}(\tau') - \Theta(\tau' - \tau) \mathbf{c}_{\lambda'}^{\dagger}(\tau') \mathbf{c}_{\lambda}(\tau) \right], \quad \text{with } \begin{aligned} \mathbf{c}_{\lambda}(\tau) &= e^{-\varepsilon_{\lambda} \tau} \mathbf{c}_{\lambda}(0) \\ \mathbf{c}_{\lambda}^{\dagger}(\tau) &= e^{\varepsilon_{\lambda} \tau} \mathbf{c}_{\lambda}^{\dagger}(0) \end{aligned} \\ &= -\frac{1}{\mathcal{Z}} \text{Tr} \left[\Theta(\tau - \tau') e^{\varepsilon_{\lambda'} \tau' - \varepsilon_{\lambda} \tau} \mathbf{c}_{\lambda} \mathbf{c}_{\lambda'}^{\dagger} - \Theta(\tau' - \tau) e^{\varepsilon_{\lambda'} \tau' - \varepsilon_{\lambda} \tau} \mathbf{c}_{\lambda'}^{\dagger} \mathbf{c}_{\lambda} \right] \\ &= -\frac{1}{\mathcal{Z}} \text{Tr} e^{-\beta \mathbf{H}} \left[\Theta(\tau - \tau') e^{\varepsilon_{\lambda'} \tau' - \varepsilon_{\lambda} \tau} \mathbf{c}_{\lambda} \mathbf{c}_{\lambda'}^{\dagger} - \Theta(\tau' - \tau) e^{\varepsilon_{\lambda'} \tau' - \varepsilon_{\lambda} \tau} \mathbf{c}_{\lambda'}^{\dagger} \mathbf{c}_{\lambda} \right], \\ &\quad \langle \mathbf{c}_{\lambda'}^{\dagger} \mathbf{c}_{\lambda} \rangle = \delta_{\lambda\lambda'} \mathbf{n}_{\mathfrak{F}}(\varepsilon_{\lambda}) = \frac{1}{e^{\beta \varepsilon_{\lambda}} + 1} \\ &\quad \text{where} \\ &\quad \langle \mathbf{c}_{\lambda} \mathbf{c}_{\lambda'}^{\dagger} \rangle = \delta_{\lambda\lambda'} (1 - \mathbf{n}_{\mathfrak{F}}(\varepsilon_{\lambda})) \\ &= -\left[\Theta(\tau - \tau') e^{\varepsilon_{\lambda'} \tau' - \varepsilon_{\lambda} \tau} \langle \mathbf{c}_{\lambda} \mathbf{c}_{\lambda'}^{\dagger} \rangle - \Theta(\tau' - \tau) e^{\varepsilon_{\lambda'} \tau' - \varepsilon_{\lambda} \tau} \langle \mathbf{c}_{\lambda'}^{\dagger} \mathbf{c}_{\lambda} \rangle \right] \\ &= -\delta_{\lambda\lambda'} e^{\varepsilon_{\lambda}(\tau' - \tau)} \left[\Theta(\tau - \tau') \langle \mathbf{c}_{\lambda} \mathbf{c}_{\lambda}^{\dagger} \rangle - \Theta(\tau' - \tau) \langle \mathbf{c}_{\lambda}^{\dagger} \mathbf{c}_{\lambda} \rangle \right] \\ &= -\delta_{\lambda\lambda'} e^{\varepsilon_{\lambda}(\tau' - \tau)} \left[(1 - \mathbf{n}_{\mathfrak{F}}(\varepsilon_{\lambda})) \Theta(\tau - \tau') - \mathbf{n}_{\mathfrak{F}}(\varepsilon_{\lambda}) \Theta(\tau' - \tau) \right]. \end{aligned}$$

$$(8.3.6.11)$$

This result can be generalized to include the bosonic non-interacting case as well, yielding

$$\mathfrak{G}_{\lambda\lambda'}(\tau - \tau') = \delta_{\lambda\lambda'} \mathfrak{G}_{\lambda}(\tau - \tau'),$$

$$\mathfrak{G}_{\lambda}(\tau) = e^{\varepsilon_{\lambda} \tau} \left[(1 \pm \mathbf{n}_{\mathfrak{F}, \mathfrak{B}}(\varepsilon_{\lambda})) \Theta(\tau) \pm \mathbf{n}_{\mathfrak{F}, \mathfrak{B}}(\varepsilon_{\lambda}) \Theta(-\tau) \right]$$

For further references³³, check ????

8.3.6.6. Fourier Transform of the Matsubara Functions: Fourier series in Imaginary Time. The Fourier transform of the Matsubara-Green function $\mathfrak{G}_{\mathbf{AB}}(\tau)$ to the frequency domain. It can be proven that the Matsubara-Green function is periodic with period 2β on the interval $\tau \in \mathbb{R}_{[-\beta, \beta]}$. In effect, using the Matsubara-Green function's symmetry property twice yields

$$(8.3.6.13) \quad \mathfrak{G}_{\mathbf{AB}}(-\beta) = \mp \mathfrak{G}_{\mathbf{AB}}(0) = (\mp)^2 \mathfrak{G}_{\mathbf{AB}}(\beta) = \mathfrak{G}_{\mathbf{AB}}(\beta).$$

Hence, a discrete Fourier series can be written,

$$(8.3.6.12) \quad \begin{aligned} \langle \mathbf{c}_{\lambda'}^{\dagger} \mathbf{c}_{\lambda} \rangle &= \frac{\text{Tr} \exp(-\beta \mathbf{H}) \mathbf{c}_{\lambda'}^{\dagger} \mathbf{c}_{\lambda}}{\text{Tr} \exp(-\beta \mathbf{H})} = \frac{\text{Tr} \mathbf{c}_{\lambda} \exp(-\beta \mathbf{H}) \mathbf{c}_{\lambda'}^{\dagger}}{\text{Tr} \exp(-\beta \mathbf{H})} = \frac{\text{Tr} \mathbf{c}_{\lambda} \exp(-\beta \mathbf{H}) \mathbf{c}_{\lambda'}^{\dagger} \exp(-\beta \mathbf{H}) \exp(\beta \mathbf{H})}{\text{Tr} \exp(-\beta \mathbf{H})} \\ &= \frac{\text{Tr} \exp(-\beta \mathbf{H}) \mathbf{c}_{\lambda} \mathbf{c}_{\lambda'}^{\dagger}(\beta)}{\text{Tr} \exp(-\beta \mathbf{H})} = e^{-\beta \varepsilon_{\lambda}} \langle \mathbf{c}_{\lambda} \mathbf{c}_{\lambda'}^{\dagger} \rangle = e^{-\beta \varepsilon_{\lambda}} \left(\delta_{\lambda\lambda'} \mp \langle \mathbf{c}_{\lambda'}^{\dagger} \mathbf{c}_{\lambda} \rangle \right) \Rightarrow \langle \mathbf{c}_{\lambda'}^{\dagger} \mathbf{c}_{\lambda} \rangle = \frac{1}{e^{\beta \varepsilon_{\lambda}} \pm 1} = \mathbf{n}_{\mathfrak{F}, \mathfrak{B}}, \end{aligned}$$

ie. yielding the Fermi and Bose statistics, respectively.

³³In effect, these identities can be easily derived for both the fermionic and bosonic case as follows,

$$\mathfrak{G}_{\text{AB}}(\tau) = \frac{1}{\beta} \sum_{n \in \mathbb{Z}} e^{-\frac{in\pi\tau}{\beta}} \mathfrak{G}_{\text{AB}}(n) \iff \mathfrak{G}_{\text{AB}}(n) = \frac{1}{2} \int_{-\beta}^{\beta} d\tau e^{\frac{in\pi\tau}{\beta}} \mathfrak{G}_{\text{AB}}(\tau)$$

Employing the previous symmetry property yields a simplified transform,

$$\begin{aligned} \mathfrak{G}_{\text{AB}}(n) &= \frac{1}{2} \int_{-\beta}^{\beta} d\tau e^{\frac{in\pi\tau}{\beta}} \mathfrak{G}_{\text{AB}}(\tau) = \frac{1}{2} \int_{-\beta}^0 d\tau e^{\frac{in\pi\tau}{\beta}} \mathfrak{G}_{\text{AB}}(\tau) + \frac{1}{2} \int_0^{\beta} d\tau e^{\frac{in\pi\tau}{\beta}} \mathfrak{G}_{\text{AB}}(\tau) \\ (8.3.6.14) \quad &= e^{-in\pi} \frac{1}{2} \int_0^{\beta} d\tau e^{\frac{in\pi\tau}{\beta}} \mathfrak{G}_{\text{AB}}(\tau + \beta) + \frac{1}{2} \int_0^{\beta} d\tau e^{\frac{in\pi\tau}{\beta}} \mathfrak{G}_{\text{AB}}(\tau) \\ &= \frac{1}{2} (1 \mp e^{-in\pi}) \int_0^{\beta} d\tau e^{\frac{in\pi\tau}{\beta}} \mathfrak{G}_{\text{AB}}(\tau) \\ &\Rightarrow \mathfrak{G}_{\text{AB}}(n) = \int_0^{\beta} d\tau e^{\frac{in\pi\tau}{\beta}} \mathfrak{G}_{\text{AB}}(\tau), \begin{cases} n \in 2\mathbb{N} \text{ for fermions,} \\ n \in 2\mathbb{N} + 1 \text{ for bosons.} \end{cases} \end{aligned}$$

Furthermore, it is convenient to separate the Fourier transforms of the Green-Matsubara functions for bosons and fermions as follows,

$$\begin{aligned} \mathfrak{G}_{\text{AB}}(i\Omega_n) &= \int_0^{\beta} d\tau e^{i\Omega_n\tau} \mathfrak{G}_{\text{AB}}(\tau), \quad \Omega_n = \frac{(2n+1)\pi}{\beta} \text{ for fermions.} \\ (8.3.6.15) \quad \mathfrak{G}_{\text{AB}}(i\omega_n) &= \int_0^{\beta} d\tau e^{i\omega_n\tau} \mathfrak{G}_{\text{AB}}(\tau), \quad \omega_n = \frac{2n\pi}{\beta} \text{ for bosons.} \end{aligned}$$

It is custom to call both $i\Omega_n, i\omega_n$ as Matsubara frequencies. The temperature appears only in the Matsubara frequencies via β .

8.3.6.7. Analytic Continuation for the Matsubara and Green Functions. It turns out that the Matsubara and Green functions are manifestations of the same analytic function $\mathcal{F}_{\text{AB}}(z)$ in the complex-frequency plane: in the upper half-plane, it is equal to $\mathfrak{G}_{\text{AB}}(i\omega_n)$ on the imaginary axis and $\mathfrak{G}_{\text{AB}}^R(\omega)$ on the real axis, while in the lower complex-plane, the advanced Green function emerges on the real axis.

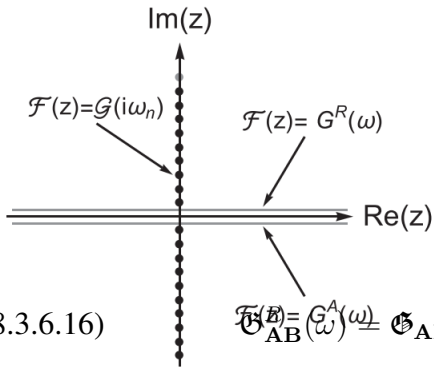
Hence, determination of one of the two guarantees that the other must follow as well by analytic continuation, which amounts to stating

$$(8.3.6.16) \quad \mathfrak{G}_{\text{AB}}^A(\omega) = \mathfrak{G}_{\text{AB}}(i\omega_n \rightarrow \omega + i\eta), \text{ where } \eta \text{ is a positive and infinitesimal constant.}$$

Some interesting properties of the Matsubara function are enumerated, as follows

- 1) $\mathfrak{G}_{\text{AB}}(\tau)$ is periodic along the imaginary-time axis,
- 2) $\mathfrak{G}_{\text{AB}}(\tau)$ does not have any statistical factor, eg. $n_{\mathfrak{B}}(\omega)$ nor $n_{\mathfrak{F}}(\omega)$, on the imaginary-frequency axis. The information on statistical probabilities implicitly enters through the so-called Matsubara frequency sums.
- 3) The Matsubara function is based on a thermal equilibrium theory. We note that, in some cases, the $(T \rightarrow 0)$ -limit of a thermal equilibrium can be different from results obtained from a $T = 0$ analysis. Note that the $(T \rightarrow 0)$ -formalism is designed to get the true ground state, including the prediction of equal thermal occupation for degenerate states, whereas a $T = 0$ analysis may identify only one of the available degenerate states, or because adiabatic turning on of the interaction may miss state crossings or a phase transition and land in an excited state.

The Matsubara approach provides an expedient recipe for obtaining its Green function counterparts. It is convenient to use the Lehmann representation to establish the relation between the two functions $\mathfrak{G}_{\text{AB}}(\tau)$ and $\mathfrak{G}_{\text{AB}}^R(\tau)$, whereby earlier results for the zero-temperature retarded single-particle Green function can be extended to finite temperatures as



$$\begin{aligned}
(8.3.6.17) \quad \mathfrak{G}_{\mathbf{AB}}(\omega) &= -\frac{i}{\mathcal{Z}} \int_{\mathbb{R}} dt e^{i\omega t} \mathfrak{G}_{\mathbf{AB}}(t) = -\frac{i}{\mathcal{Z}} \int_{\mathbb{R}} dt e^{i\omega t} \left\langle e^{-\beta \mathbf{H}} \mathcal{T}[\mathbf{A}(t) \mathbf{B}(0)] \right\rangle \\
&= -\frac{i}{\mathcal{Z}} \int_{\mathbb{R}} dt e^{i\omega t} \left\langle e^{-\beta \mathbf{H}} \left[\Theta(t) \mathbf{A}(t) \mathbf{B}(0) - \Theta(-t) \mathbf{B}(0) \mathbf{A}(t) \right] \right\rangle
\end{aligned}$$

$$\begin{aligned}
& -\frac{i}{\mathcal{Z}} \int_{\mathbb{R}} dt e^{i\omega t} \left\langle e^{-\beta \mathbf{H}} \Theta(t) \mathbf{A}(t) \mathbf{B}(0) \right\rangle & \frac{i}{\mathcal{Z}} \int_{\mathbb{R}} dt e^{i\omega t} \left\langle e^{-\beta \mathbf{H}} \Theta(-t) \mathbf{B}(0) \mathbf{A}(t) \right\rangle & = \\
& = -\frac{i}{\mathcal{Z}} \int_{\mathbb{R}} dt e^{i\omega t} \left\langle e^{-\beta \mathbf{H}} \Theta(t) e^{\frac{i}{\hbar} \mathbf{H} t} \mathbf{A}(0) e^{-\frac{i}{\hbar} \mathbf{H} t} \mathbf{B}(0) \right\rangle & = \frac{i}{\mathcal{Z}} \int_{\mathbb{R}} dt e^{i\omega t} \left\langle e^{-\beta \mathbf{H}} \Theta(-t) \mathbf{B}(0) e^{\frac{i}{\hbar} \mathbf{H} t} \mathbf{A}(0) e^{-\frac{i}{\hbar} \mathbf{H} t} \right\rangle \\
& = -\frac{i}{\mathcal{Z}} \int_{\mathbb{R}^+} dt e^{i\omega t} \text{Tr} \left[e^{-\beta \mathbf{H}} e^{\frac{i}{\hbar} \mathbf{H} t} \mathbf{A} e^{-\frac{i}{\hbar} \mathbf{H} t} \mathbf{B} \right] & = \frac{i}{\mathcal{Z}} \int_{\mathbb{R}^-} dt e^{i\omega t} \text{Tr} \left[e^{-\beta \mathbf{H}} \mathbf{B} e^{\frac{i}{\hbar} \mathbf{H} t} \mathbf{A} e^{-\frac{i}{\hbar} \mathbf{H} t} \right] \\
& = -\frac{i}{\mathcal{Z}} \int_{\mathbb{R}^+} dt e^{i\omega t} \sum_{\lambda \lambda'} \langle \lambda | e^{-\beta \mathbf{H}} e^{\frac{i}{\hbar} \mathbf{H} t} \mathbf{A} | \lambda' \rangle \langle \lambda' | e^{-\frac{i}{\hbar} \mathbf{H} t} \mathbf{B} | \lambda \rangle & = \frac{i}{\mathcal{Z}} \int_{\mathbb{R}^-} dt e^{i\omega t} \sum_{\lambda \lambda'} \langle \lambda | e^{-\beta \mathbf{H}} \mathbf{B} e^{\frac{i}{\hbar} \mathbf{H} t} | \lambda' \rangle \langle \lambda' | e^{-\frac{i}{\hbar} \mathbf{H} t} \mathbf{A} | \lambda \rangle \\
& = -\frac{i}{\mathcal{Z}} \int_{\mathbb{R}^+} dt e^{i\omega t} \sum_{\lambda \lambda'} e^{-\beta \varepsilon_{\lambda}} e^{\frac{i}{\hbar} \varepsilon_{\lambda} t} \mathbf{A}_{\lambda \lambda'} e^{-\frac{i}{\hbar} \varepsilon_{\lambda'} t} \mathbf{B}_{\lambda' \lambda} & = \frac{i}{\mathcal{Z}} \int_{\mathbb{R}^-} dt e^{i\omega t} \sum_{\lambda \lambda'} e^{-\beta \varepsilon_{\lambda'}} e^{\frac{i}{\hbar} \varepsilon_{\lambda'} t} \mathbf{B}_{\lambda \lambda'} e^{-\frac{i}{\hbar} \varepsilon_{\lambda} t} \mathbf{A}_{\lambda' \lambda} \\
& = \frac{1}{\mathcal{Z}} \sum_{\lambda \lambda'} \frac{\langle \lambda | \mathbf{A} | \lambda' \rangle \langle \lambda' | \mathbf{B} | \lambda \rangle}{\omega + \varepsilon_{\lambda} - \varepsilon_{\lambda'} + i\eta} e^{-\beta \varepsilon_{\lambda}} & = \frac{1}{\mathcal{Z}} \sum_{\lambda \lambda'} \frac{\langle \lambda | \mathbf{B} | \lambda' \rangle \langle \lambda' | \mathbf{A} | \lambda \rangle}{\omega + \varepsilon_{\lambda} - \varepsilon_{\lambda'} + i\eta} e^{-\beta \varepsilon_{\lambda'}}
\end{aligned}$$

$$(8.3.6.18) \quad \Rightarrow \mathfrak{G}_{\mathbf{AB}}(\omega) = \frac{1}{\mathcal{Z}} \sum_{\lambda \lambda'} \frac{\mathbf{A}_{\lambda \lambda'} \mathbf{B}_{\lambda' \lambda}}{\omega + \varepsilon_{\lambda} - \varepsilon_{\lambda'} + i\eta} (e^{-\beta \varepsilon_{\lambda}} + e^{-\beta \varepsilon_{\lambda'}})$$

As for the Matsubara function, for $\tau > 0$,

$$\begin{aligned}
(8.3.6.19) \quad \mathfrak{G}_{\mathbf{AB}}(\tau) &= -\frac{1}{\mathcal{Z}} \text{Tr} \left[e^{-\beta \mathbf{H}} e^{\mathbf{H} \tau} \mathbf{A} e^{-\mathbf{H} \tau} \mathbf{B} \right] = -\frac{1}{\mathcal{Z}} \sum_{\lambda \lambda'} \mathbf{A}_{\lambda \lambda'} \mathbf{B}_{\lambda' \lambda} e^{\tau(\varepsilon_{\lambda} - \varepsilon_{\lambda'})} \\
\mathfrak{G}_{\mathbf{AB}}(i\omega_n) &= \frac{1}{\mathcal{Z}} \sum_{\lambda \lambda'} \frac{\mathbf{A}_{\lambda \lambda'} \mathbf{B}_{\lambda' \lambda}}{i\omega_n + \varepsilon_{\lambda} - \varepsilon_{\lambda'} + i\eta} (e^{-\beta \varepsilon_{\lambda}} + e^{-\beta \varepsilon_{\lambda'}})
\end{aligned}$$

Hence, $\mathfrak{G}_{\mathbf{AB}}^R(\omega)$ have identical expressions $\mathfrak{G}_{\mathbf{AB}}(i\omega_n)$ and, moreover, both can be obtained from a common function

$$(8.3.6.20) \quad \mathfrak{F}_{\mathbf{AB}}(z) = \frac{1}{\mathcal{Z}} \sum_{\lambda \lambda'} \frac{\mathbf{A}_{\lambda \lambda'} \mathbf{B}_{\lambda' \lambda}}{z + \varepsilon_{\lambda} - \varepsilon_{\lambda'} + i\eta} (e^{-\beta \varepsilon_{\lambda}} + e^{-\beta \varepsilon_{\lambda'}}),$$

which is analytic in the upper and lower half-planes, but has a series of poles at $\varepsilon_{\lambda} - \varepsilon_{\lambda'}$, along the real axis.

Knowledge of $\mathfrak{G}_{\mathbf{AB}}(i\omega_n)$ implies knowledge of the $\mathfrak{F}_{\mathbf{AB}}$ -function of the complex z -variable on an infinite set of points on the imaginary axis $z_n = i\omega_n$, which is sufficient to determine the full $\mathfrak{F}_{\mathbf{AB}}$ -function in the complex plane, with the proviso that, since $\mathfrak{F}_{\mathbf{AB}}(z) = \mathfrak{G}_{\mathbf{AB}}(z)$, it has to be analytic in the upper-half-plane, leading to the prescription

$$\mathfrak{G}_{\mathbf{AB}}^R(\omega) = \mathfrak{G}_{\mathbf{AB}}(i\omega_n \rightarrow \omega + i\eta).$$

For example, consider again the fermionic Hamiltonian of a non-interacting theory given in [equation \(8.3.6.9\)](#). Its fermionic Matsubara function in the frequency domain can be obtained as the Fourier transform of [equation \(8.3.6.10\)](#), yielding

$$(8.3.6.21) \quad \begin{aligned} \mathfrak{G}_\lambda^{(0)}(i\Omega_n) &= \int_0^\beta d\tau e^{i\Omega_n \tau} \mathfrak{G}_\lambda^{(0)}(\tau), \quad \text{where } \Omega_n = \frac{(2n+1)\pi}{\beta}, \tau > 0, \\ &= \frac{1}{i\Omega_n - \varepsilon_\lambda}, \end{aligned}$$

since $e^{i\Omega_n \tau} = -1$ and since $(1 - \mathbf{n}(\varepsilon_\lambda)) = (1 + e^{-\beta\varepsilon_\lambda})^{-1}$. Hence, using the previously explained relationships between the Green and Matsubara functions, it follows that

$$(8.3.6.22) \quad \mathfrak{G}_\lambda^{(0)}(\omega) = \frac{1}{\omega - \varepsilon_\lambda + i\eta}.$$

8.3.6.8. Matsubara Approach to the Two-Particle Green Function. As an application of the Matsubara approach for finite temperature quantum field theory, one can calculate the two-particle correlation function χ at finite temperatures. Said correlation function reads

$$(8.3.6.23) \quad \begin{aligned} \chi(\mathbf{q}, \tau) &= - \sum_{\substack{\mathbf{k}, \mathbf{k}' \\ \sigma, \sigma'}} \left\langle \mathcal{T}_\tau [\mathbf{c}_{\mathbf{k}, \sigma}^\dagger(\tau + \varepsilon) \mathbf{c}_{\mathbf{k}+\mathbf{q}, \sigma}(\tau) \mathbf{c}_{\mathbf{k}', \sigma'}^\dagger(\varepsilon) \mathbf{c}_{\mathbf{k}'-\mathbf{q}, \sigma'}(0)] \right\rangle \\ &= \sum_{\substack{\mathbf{k}, \mathbf{k}' \\ \sigma, \sigma'}} \left\langle \mathcal{T}_\tau \left[\overbrace{\mathbf{c}_{\mathbf{k}, \sigma}^\dagger(\tau + \varepsilon) \mathbf{c}_{\mathbf{k}+\mathbf{q}, \sigma}(\tau) \mathbf{c}_{\mathbf{k}', \sigma'}^\dagger(\varepsilon) \mathbf{c}_{\mathbf{k}'-\mathbf{q}, \sigma'}(0)} - \overbrace{\mathbf{c}_{\mathbf{k}, \sigma}^\dagger(\tau + \varepsilon) \mathbf{c}_{\mathbf{k}+\mathbf{q}, \sigma}(\tau) \mathbf{c}_{\mathbf{k}', \sigma'}^\dagger(\varepsilon) \mathbf{c}_{\mathbf{k}'-\mathbf{q}, \sigma'}(0)} \right] \right\rangle \\ &= \sum_{\substack{\mathbf{k}, \mathbf{k}' \\ \sigma, \sigma'}} \left\langle \left[\begin{aligned} &\mathcal{T}_\tau \left(\mathbf{c}_{\mathbf{k}+\mathbf{q}, \sigma}(\tau) \mathbf{c}_{\mathbf{k}', \sigma'}^\dagger(0) \right) \mathcal{T}_\tau \left(\mathbf{c}_{\mathbf{k}'-\mathbf{q}, \sigma'}(0) \mathbf{c}_{\mathbf{k}, \sigma}^\dagger(\tau) \right) \\ &- \mathcal{T}_\tau \left(\mathbf{c}_{\mathbf{k}, \sigma}^\dagger(\tau + \varepsilon) \mathbf{c}_{\mathbf{k}+\mathbf{q}, \sigma}(\tau) \right) \mathcal{T}_\tau \left(\mathbf{c}_{\mathbf{k}', \sigma'}^\dagger(\varepsilon) \mathbf{c}_{\mathbf{k}'-\mathbf{q}, \sigma'}(0) \right) \end{aligned} \right] \right\rangle \\ &= \sum_{\substack{\mathbf{k}, \mathbf{k}' \\ \sigma, \sigma'}} \left[\mathfrak{G}_{\mathbf{k}+\mathbf{q}, \sigma}^{(0)}(\tau) \mathfrak{G}_{\mathbf{k}, \sigma}^{(0)}(-\tau) - \rho_{\mathbf{q}, \sigma}^{(0)} \rho_{-\mathbf{q}, \sigma}^{(0)} \right], \end{aligned}$$

where ζ is an infinitesimal positive constant that places the creation operators to the left of annihilation ones for equal time ordering. Note that the second term in the last line, does not depend on τ . Hence, its Fourier transform is done with respect to $\tau \sim \delta_{\omega_n, 0}$ and will thus disappear in the process of the analytical continuation $i\omega_n \rightarrow \omega + i\eta$.

Note that the density operators are physical observables, they are hermitian and commute, thus having bosonic character. This character should be reflected in the correlation function as well, i.e. it should be a function of a bosonic Matsubara frequency. Noting that the Fourier transform of a product in time-domain is a convolution in frequency-domain, the Fourier transform of -the first in- the correlation function is

$$(8.3.6.24) \quad \begin{aligned} \chi(\mathbf{q}, i\omega_n) &= \frac{1}{\beta} \sum_{i\Omega_n} \sum_{\mathbf{k}, \sigma} \mathfrak{G}_{\mathbf{k}+\mathbf{q}, \sigma}^{(0)}(i\Omega_m + i\omega_n) \mathfrak{G}_{\mathbf{k}, \sigma}^{(0)}(i\Omega_m) \\ &= \frac{2}{\beta} \sum_{\mathbf{k}, i\Omega_m} \left[\left(i(\Omega_m + \omega_n) - (\varepsilon_{\mathbf{k}+\mathbf{q}} - \mu) \right) \left(i\Omega_m - (\varepsilon_{\mathbf{k}} - \mu) \right) \right]^{-1}, \end{aligned}$$

wherein ω_n and Ω_m are bosonic and fermionic Matsubara frequencies, respectively. Performing the summation over Ω_m can be carried out by writing the sum as an integral over the lower-half-plane contour in the complex plane. The actual details of the calculations can be found in ?? p.509-511.

9. QUANTUM FIELD-THEORETIC APPROACH TO CONDENSED MATTER PHYSICS

9.1. Mean field theory. From "Many Body Quantum Theory in Condensed Matter Physics" by Henrik Bruus and Karsten Flensberg.

Physical Context

The physics of interacting particles is often very complicated because the motion of the individual particles depend on the position of all other, thus these variables become correlated. This is clearly the case for a system of charged particles interacting by Coulomb forces, such as eg. the electron gas. There, the probability to find two electrons in close proximity is expected to be small due to the strong repulsive interaction. Consequently, due to these correlation effects, there is a suppressed density in the neighbourhood of every electron, and there is a correlation hole. Nevertheless, in spite of these complications, there are a number of cases where a more crude treatment, not fully including the correlations, gives a good physical model. In these cases, it suffices to include correlations on the average. This means that the effect of the other particles is included as a mean density (or mean field), leaving an effective single particle problem, which is solvable. These mean fields are chosen so that they minimize the free energy, which in turn, ensures that the method is self-consistent. This method is called mean field theory, which allows to neglect detailed dynamics and time-dependent second quantization methods.

First, consider a system with two types of particles, described by operators $\mathbf{a}_\nu, \mathbf{b}_\mu$, respectively. In this model, the only interactions present are those between different kind of particles. The Hamiltonian then reads

$$(9.1.0.1) \quad \mathbf{H} = \mathbf{H}_0 + \mathbf{V}_{\text{int}}, \text{ with } \begin{aligned} \mathbf{H}_0 &= \sum_\nu \varepsilon_\nu^a \mathbf{a}_\nu^\dagger \mathbf{a}_\nu + \sum_\mu \varepsilon_\mu^b \mathbf{b}_\mu^\dagger \mathbf{b}_\mu, \\ \mathbf{V}_{\text{int}} &= \sum_{\nu\nu', \mu\mu'} \mathcal{V}_{\nu\mu, \nu'\mu'} \mathbf{a}_\nu^\dagger \mathbf{b}_\mu^\dagger \mathbf{b}_{\mu'} \mathbf{a}_{\nu'}. \end{aligned}$$

Now it may be expected that the density operator $\mathbf{a}_\nu^\dagger \mathbf{a}_{\nu'}$ and $\mathbf{b}_\mu^\dagger \mathbf{b}_{\mu'}$, deviate only little from their averages values, $\langle \mathbf{a}_\nu^\dagger \mathbf{a}_{\nu'} \rangle$ and $\langle \mathbf{b}_\mu^\dagger \mathbf{b}_{\mu'} \rangle$, respectively. It is natural then to use this deviation as a small parameter and perform an expansion. Let the deviation operators be

$$(9.1.0.2) \quad \mathbf{d}_{\nu\nu'} = \mathbf{a}_\nu^\dagger \mathbf{a}_{\nu'} - \langle \mathbf{a}_\nu^\dagger \mathbf{a}_{\nu'} \rangle, \quad \mathbf{e}_{\mu\mu'} = \mathbf{b}_\mu^\dagger \mathbf{b}_{\mu'} - \langle \mathbf{b}_\mu^\dagger \mathbf{b}_{\mu'} \rangle.$$

Then, in terms of these new operators, the Hamiltonian reads

$$\mathbf{H}_{\text{MFT}} = \mathbf{H}_0 + \mathbf{V}_{\text{MFT}} + \sum_{\nu\nu', \mu\mu'} \mathbf{d}_{\nu\nu'} \mathbf{e}_{\mu\mu'}, \text{ where } \mathbf{V}_{\text{MFT}} = \sum_{\nu\nu', \mu\mu'} \mathcal{V}_{\nu\mu, \nu'\mu'} \left[\begin{array}{l} \mathbf{a}_\nu^\dagger \mathbf{a}_{\nu'} \langle \mathbf{b}_\mu^\dagger \mathbf{b}_{\mu'} \rangle \\ + \mathbf{b}_\mu^\dagger \mathbf{b}_{\mu'} \langle \mathbf{a}_\nu^\dagger \mathbf{a}_{\nu'} \rangle \\ - \langle \mathbf{a}_\nu^\dagger \mathbf{a}_{\nu'} \rangle \langle \mathbf{b}_\mu^\dagger \mathbf{b}_{\mu'} \rangle \end{array} \right] + \mathcal{O}(de).$$

The mean field Hamiltonian \mathbf{H}_{MFT} thus contains only single-particle operators, and thus the original many-body problem has been reduced to a single-particle problem, which can always be solved³⁴

Mean field theory may be formulated in a different way: if an interaction term involving the product of two operators, \mathbf{A}, \mathbf{B} , ie. $\mathbf{H}_{AB} = \mathbf{A}\mathbf{B}$, then the mean field approximation is given by

³⁴In effect, any quadratic Hamiltonian may be diagonalized since it can be written as $\mathbf{H} = \sum_{ij} a_i^\dagger \mathcal{H}_{ij} a_j^\dagger$. Since this \mathcal{H} is hermitian, then there exists a transformation

$$\alpha_i \rightarrow \sum_j U_{ij} a_j, \quad U \in U(N),$$

such that \mathcal{H} is diagonal: $\sum_{kk'} U_{ik}^\dagger \mathcal{H}_{kk'} U_{k'j} = \delta_{ij} \varepsilon_i$. Then, in terms of this new basis, the Hamiltonian then becomes $\mathbf{H} = \sum_i \alpha_i^\dagger \varepsilon_i \alpha_i$, and $\{\varepsilon_i\}$ are thus \mathcal{H} 's eigenvalues.

$$\mathbf{H}_{AB}^{\text{MFT}} = \mathbf{A}\langle\mathbf{B}\rangle + \langle\mathbf{A}\rangle\mathbf{B} - \langle\mathbf{A}\rangle\langle\mathbf{B}\rangle.$$

The question is then, how to find the averages $\langle\mathbf{a}_\nu^\dagger\mathbf{a}_{\nu'}\rangle$ and $\langle\mathbf{b}_\mu^\dagger\mathbf{b}_{\mu'}\rangle$. In principle, there are two possible routes.

- Method 1: these averages are to be determined self-consistently, ie. when calculating the averages

$$(9.1.0.3) \quad \bar{\mathbf{n}}_{\nu\nu'}^a = \langle\mathbf{a}_\nu^\dagger\mathbf{a}_{\nu'}\rangle, \quad \bar{\mathbf{n}}_{\mu\mu'}^b = \langle\mathbf{b}_\mu^\dagger\mathbf{b}_{\mu'}\rangle,$$

using the new mean field Hamiltonian, the same answer should be re-obtained. This entails the following prescriptions for $\bar{\mathbf{n}}_{\nu\nu'}^a$ and $\bar{\mathbf{n}}_{\mu\mu'}^b$,

$$(9.1.0.4) \quad \begin{aligned} \bar{\mathbf{n}}_{\nu\nu'}^a &= \langle\mathbf{a}_\nu^\dagger\mathbf{a}_{\nu'}\rangle_{\text{MFT}} = \frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left(e^{-\beta\mathbf{H}_{\text{MFT}}} \mathbf{a}_\nu^\dagger\mathbf{a}_{\nu'} \right) \\ \bar{\mathbf{n}}_{\mu\mu'}^b &= \langle\mathbf{b}_\mu^\dagger\mathbf{b}_{\mu'}\rangle_{\text{MFT}} = \frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left(e^{-\beta\mathbf{H}_{\text{MFT}}} \mathbf{b}_\mu^\dagger\mathbf{b}_{\mu'} \right), \end{aligned} \quad \text{where } \mathcal{Z}_{\text{MFT}} = \text{Tr} \left(e^{-\beta\mathbf{H}_{\text{MFT}}} \right).$$

Said equations are the self-consistency equations, since the averages are given in terms of the mean field hamiltonian and the mean field partition function, which themselves depend on these averages.

- Method 2: the alternative route consists on using the $\mathbf{n}_{\nu\nu'}$ which minimizes the free energy F_{MFT} of the mean field Hamiltonian. Then,

$$(9.1.0.5) \quad \begin{aligned} 0 &= \frac{d}{d\bar{\mathbf{n}}_{\nu\nu'}^a} F_{\text{MFT}} = \frac{d}{d\bar{\mathbf{n}}_{\nu\nu'}^a} \left(-\frac{\log \mathcal{Z}_{\text{MFT}}}{\beta} \right) = \frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left(e^{-\beta\mathbf{H}_{\text{MFT}}} \frac{d}{d\bar{\mathbf{n}}_{\nu\nu'}^a} \mathbf{H}_{\text{MFT}} \right) \\ &= \frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left(e^{-\beta\mathbf{H}_{\text{MFT}}} \left(\sum_{\nu\nu',\mu\mu'} \mathcal{V}_{\nu\mu,\nu'\mu'} (\mathbf{b}_\mu^\dagger\mathbf{b}_\mu - \bar{\mathbf{n}}_{\mu\mu'}^b) \right) \right) \\ &= \frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left(e^{-\beta\mathbf{H}_{\text{MFT}}} \sum_{\nu\nu',\mu\mu'} \mathcal{V}_{\nu\mu,\nu'\mu'} \mathbf{b}_\mu^\dagger\mathbf{b}_\mu \right) - \frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left(e^{-\beta\mathbf{H}_{\text{MFT}}} \sum_{\nu\nu',\mu\mu'} \mathcal{V}_{\nu\mu,\nu'\mu'} \bar{\mathbf{n}}_{\mu\mu'}^b \right) \\ &= \sum_{\mu\mu'} \frac{1}{\mathcal{Z}_{\text{MFT}}} \mathcal{V}_{\nu\mu,\nu'\mu'} \text{Tr} \left(e^{-\beta\mathbf{H}_{\text{MFT}}} \mathbf{b}_\mu^\dagger\mathbf{b}_\mu \right) - \sum_{\mu\mu'} \mathcal{V}_{\nu\mu,\nu'\mu'} \bar{\mathbf{n}}_{\mu\mu'}^b \frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left(e^{-\beta\mathbf{H}_{\text{MFT}}} \right) \\ &= \sum_{\mu\mu'} \mathcal{V}_{\nu\mu,\nu'\mu'} \left(\langle\mathbf{b}_\mu^\dagger\mathbf{b}_{\mu'}\rangle_{\text{MFT}} - \bar{\mathbf{n}}_{\mu\mu'}^b \right), \end{aligned}$$

which should hold for any pair (ν, ν') and hence, the last parenthesis should vanish. Thus, the self-consistency equation is obtained as well for the second method, making both approaches equivalent.

Consider the average interaction energy, $\langle\mathbf{V}_{\text{int}}\rangle$, a natural approximation would be to evaluate the expectation value of the a and b operators separately, this is

$$\langle\mathbf{V}_{\text{int}}\rangle \approx \sum_{\nu\nu',\mu\mu'} \mathcal{V}_{\nu\mu,\nu'\mu'} \langle\mathbf{a}_\nu^\dagger\mathbf{a}_{\nu'}\rangle \langle\mathbf{b}_\mu^\dagger\mathbf{b}_{\mu'}\rangle,$$

which is equivalent to assuming that the a and b operators are uncorrelated. This is, in essence, the approximation done in the mean-field approach. In effect, the interaction term can be evaluated using the mean field Hamiltonian

$$\langle\mathbf{V}_{\text{int}}\rangle_{\text{MFT}} = \frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left((e^{-\beta\mathbf{H}_{\text{MFT}}} \mathbf{V}_{\text{int}}) \right).$$

Since, the mean field Hamiltonian can be separated into a part containing only a-operators and another term containing only b-operators, $\mathbf{H}_{\text{MFT}} = \mathbf{H}_{\text{MFT}}^a + \mathbf{H}_{\text{MFT}}^b$, the average factorization yields

$$(9.1.0.6) \quad \langle\mathbf{V}_{\text{int}}\rangle_{\text{MFT}} = \sum_{\mu\mu'} \mathcal{V}_{\nu\mu,\nu'\mu'} \langle\mathbf{a}_\nu^\dagger\mathbf{a}_{\nu'}\rangle_{\text{MFT}} \langle\mathbf{b}_\mu^\dagger\mathbf{b}_{\mu'}\rangle_{\text{MFT}}.$$

The mean field theory approach hence provides a consistent a physically sensible method to study interacting system where correlation are less important. Mean field theory is only valid provided d is actually small. This can be checked by calculating $\langle \mathbf{d} \rangle$, using the neglected term in the Hamiltonian (that concerning the cross term of $\mathbf{d}_{\nu\nu'} \mathbf{e}_{\mu\mu'}$) as a perturbation, and then compare the result to $\langle \mathbf{a}_{\nu}^{\dagger} \mathbf{a}_{\nu'} \rangle$. If the difference is not negligible, then either the wrong mean field parameter has been chosen or the method simply fails and other more adequate tools must be employed.

In practice, it is necessary to assume something about the averages $\langle \mathbf{a}_{\nu}^{\dagger} \mathbf{a}_{\nu'} \rangle$ and $\langle \mathbf{b}_{\mu}^{\dagger} \mathbf{b}_{\mu'} \rangle$. Even though, there is a prescription for which to find which averages are important, there are simply too many combinations. Suppose there are N different quantum numbers, then there are in principle N^2 combinations, which then yields N^2 coupled non-linear equations, which is only tractable for small systems. Often, symmetry arguments can help in reducing the number of parameters. Suppose for example that the system's Hamiltonian has translational symmetry, such that the momentum space is a natural choice. For a system of particles described by (c, c^{\dagger}) -operators, then

$$(9.1.0.7) \quad \langle \mathbf{c}_{\mathbf{k}}^{\dagger} \mathbf{c}_{\mathbf{k}'} \rangle = \frac{1}{\mathcal{V}} \int_{\mathbb{R}^6} d\mathbf{x} d\mathbf{x}' e^{-i\mathbf{k}' \cdot \mathbf{x}'} e^{i\mathbf{k} \cdot \mathbf{x}} \langle \Psi^{\dagger}(\mathbf{x}) \Psi(\mathbf{x}') \rangle.$$

It is natural to assume that the system is homogenous, which means that

$$(9.1.0.8) \quad \langle \Psi^{\dagger}(\mathbf{x}) \Psi(\mathbf{x}') \rangle = f(\mathbf{x} - \mathbf{x}') \Rightarrow \langle \mathbf{c}_{\mathbf{k}}^{\dagger} \mathbf{c}_{\mathbf{k}'} \rangle = \langle \mathbf{n}_{\mathbf{k}} \rangle \delta_{\mathbf{k}\mathbf{k}'}.$$

The homogeneity assumption is however not always true, since in some cases the system's symmetry is lower than that of the Hamiltonian. Eg. if the system spontaneously orders into a state with a spatial density variation, like a wave formation, then the average $\langle \Psi^{\dagger}(\mathbf{x}) \Psi(\mathbf{x}') \rangle$ is not a function of $\mathbf{x} - \mathbf{x}'$ only. Instead it has a lower and more restricted symmetry, namely that

$$\langle \Psi^{\dagger}(\mathbf{x}) \Psi(\mathbf{x}') \rangle = h(\mathbf{x}, \mathbf{x}') \text{ with } h(\mathbf{x}, \mathbf{x}') = h(\mathbf{x} + \mathbf{R}, \mathbf{x}' + \mathbf{R}), \mathbf{R} \text{ a lattice vector.}$$

Physical Context

This kind of crystal structure exists in nature and when it happens, the system undergoes a broken symmetry phenomena. It is important to realize that this solution cannot be found by taking the homogeneity assumption but instead by calculating $\langle \mathbf{c}_{\mathbf{k}}^{\dagger} \mathbf{c}_{\mathbf{k}+\mathbf{Q}} \rangle$, leading to the possibility of this quantity being finite, where $\mathbf{R} \cdot \mathbf{Q} = 2\pi$. Thus, the choice of the proper mean field parameters requires physical motivation about which possible states one expects.

9.1.0.1. Hartree-Fock theory. Mean field theory's methods, upto now, concern only interaction between different particles, wherein correlation terms are replaced by their average values plus small corrections. For interactions between identical particles this, however, does not exhaust the possibilities and only includes the so-called Hartree term, requiring a more general approximation scheme, given by the Hartree-Fock approximation. Consider a quantum system of interacting particles, whose Hamiltonian is given by

$$(9.1.0.9) \quad \mathbf{H} = \mathbf{H}_0 + \mathbf{V}_{\text{int}}, \text{ with } \mathbf{H}_0 = \sum_{\nu} \varepsilon_{\nu} \mathbf{c}_{\nu}^{\dagger} \mathbf{c}_{\nu} \\ \mathbf{V}_{\text{int}} = \frac{1}{2} \sum_{\nu\nu', \mu\mu'} \mathcal{V}_{\nu\mu, \nu'\mu'} \mathbf{c}_{\nu}^{\dagger} \mathbf{c}_{\mu}^{\dagger} \mathbf{c}_{\mu'} \mathbf{c}_{\nu'}.$$

The theory's main idea is that the operator $\rho_{\mu\mu'} = \mathbf{c}_{\mu}^{\dagger} \mathbf{c}_{\mu'}$ is large only when the average $\langle \rho_{\mu\mu'} \rangle$ is non-zero. For most $(\nu\nu')$ -combinations, the average $\langle \rho_{\nu\nu'} \rangle$ is zero. Therefore, the four-body interaction term may be rewritten in terms of a deviation from the average value

$$\mathbf{c}_{\nu}^{\dagger} \left(\mathbf{c}_{\mu}^{\dagger} \mathbf{c}_{\mu'} - \langle \mathbf{c}_{\mu}^{\dagger} \mathbf{c}_{\mu'} \rangle \right) \mathbf{c}_{\nu'} + \mathbf{c}_{\nu}^{\dagger} \mathbf{c}_{\nu'} \langle \mathbf{c}_{\mu}^{\dagger} \mathbf{c}_{\mu'} \rangle.$$

If $\nu' \neq \mu$, then c_ν operator commutes with the parenthesis, which is true upto zero-measure sets. With this assumption, then

(9.1.0.10)

$$\begin{aligned} c_\nu^\dagger \left(c_\mu^\dagger c_{\mu'} - \langle c_\mu^\dagger c_{\mu'} \rangle \right) c_{\nu'} + c_\nu^\dagger c_{\nu'} \langle c_\mu^\dagger c_{\mu'} \rangle &\stackrel{\mu \neq \nu'}{=} c_\nu^\dagger c_{\nu'} \left(c_\mu^\dagger c_{\mu'} - \langle c_\mu^\dagger c_{\mu'} \rangle \right) + c_\nu^\dagger c_{\nu'} \langle c_\mu^\dagger c_{\mu'} \rangle \\ &= \left(c_\nu^\dagger c_{\nu'} - \langle c_\nu^\dagger c_{\nu'} \rangle \right) \left(c_\mu^\dagger c_{\mu'} - \langle c_\mu^\dagger c_{\mu'} \rangle \right) + c_\nu^\dagger c_{\nu'} \langle c_\mu^\dagger c_{\mu'} \rangle - c_\mu^\dagger c_{\mu'} \langle c_\nu^\dagger c_{\nu'} \rangle - \langle c_\mu^\dagger c_{\mu'} \rangle \langle c_\nu^\dagger c_{\nu'} \rangle. \end{aligned}$$

The first term is proportional to the deviation squared and then may be neglected, yielding the Hartree approximation for interactions

(9.1.0.11)

$$\mathbf{V}_{\text{int}}^{\text{HF}} = \frac{1}{2} \sum_{\nu\nu', \mu\mu'} \mathcal{V}_{\nu\mu, \nu'\mu'} \left[\bar{n}_{\mu\mu'} c_\nu^\dagger c_{\nu'} + \bar{n}_{\nu\nu'} c_\mu^\dagger c_{\mu'} - \bar{n}_{\nu\nu'} \bar{n}_{\mu\mu'} \right],$$

which is the same result that would be obtained if the $(\nu\nu')$ and $(\mu\mu')$ -operators were considered to be from different kinds of particles. However, this is not yet the full result, since the μ and ν -labelled operators represent identical, indistinguishable particles and there is therefore one combination not accounted for by the Hartree approximation, the so-called Fock exchange term. This new term appears since the product of four one-body operators also gives a large contribution if $\langle c_\nu^\dagger c_{\nu'} \rangle$ finite. To derive the mean field contribution from this possibility the $\langle c_\nu^\dagger c_{\mu'} \rangle$ term must first be replaced by its average value and do the same with the combination of $\langle c_\mu^\dagger c_{\nu'} \rangle$, thus yielding

(9.1.0.12)

$$\mathbf{V}_{\text{int}}^{\text{Fock}} = -\frac{1}{2} \sum_{\nu\nu', \mu\mu'} \mathcal{V}_{\nu\mu, \nu'\mu'} \left[\bar{n}_{\nu\mu'} c_\mu^\dagger c_{\nu'} + \bar{n}_{\mu\nu'} c_\nu^\dagger c_{\mu'} - \bar{n}_{\mu\nu'} \bar{n}_{\nu\mu'} \right].$$

Therefore, the final mean field Hamiltonian within the Hartree-Fock approximation is

(9.1.0.13)

$$\mathbf{H}^{\text{HF}} = \mathbf{H}_0 + \mathbf{V}_{\text{int}}^{\text{Hartree}} + \mathbf{V}_{\text{int}}^{\text{Fock}}.$$

9.1.0.2. Homogeneous electron gas. Consider the example of a translationally invariant homogeneous electron gas, wherein the expectation value $\langle c_{\mathbf{k}}^\dagger c_{\mathbf{k}'} \rangle$ is diagonal. From the Hartree-Fock approximation, the Hamiltonian for this system reads

$$\begin{aligned} \mathbf{H}^{\text{HF}} &= \sum_{\mathbf{k}\sigma} \varepsilon_{\mathbf{k}}^{\text{HF}} c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}\sigma}, \quad \text{where} \\ \varepsilon_{\mathbf{k}}^{\text{HF}} &= \varepsilon_{\mathbf{k}} + \sum_{\mathbf{k}'\sigma'} [V(0) - \delta_{\sigma\sigma'} V(\mathbf{k} - \mathbf{k}')] n_{\mathbf{k}'\sigma'} \\ &= \varepsilon_{\mathbf{k}} + V(0)N - \sum_{\mathbf{k}'\sigma'} V(\mathbf{k} - \mathbf{k}') n_{\mathbf{k}'\sigma'}. \end{aligned}$$

The second term is the interaction with the average electron charge. In condensed matter systems, it is normally cancelled out by an equally large term due to the positively charged ionic background. The third term is precisely the exchange correlation term.

The Hartree-Fock approach depends crucially on what averages are assumed to be finite, and these assumptions must be based on physical knowledge or clever guess-work. For example, in the previous example, spin symmetry is assumed to be maintained, which then implies that $\langle c_{\mathbf{k}\downarrow} c_{\mathbf{k}\downarrow} \rangle = \langle c_{\mathbf{k}\uparrow} c_{\mathbf{k}\uparrow} \rangle$. Should this not be the case, a ferromagnetic solution may be obtained.

9.1.0.3. Broken Symmetry. Mean field theory is often used to study phase transitions and thus changes of symmetry. For a given Hamiltonian with some symmetry (eg. translational symmetry, rotational symmetry in real space or in spin space), there exists an operator which reflects this symmetry and therefore commutes with the Hamiltonian (eg. translation operator, rotation operator in either real or spin space). Since the operator and the Hamiltonian commute, a common set of eigenstates exists. Consider for example, the case of a liquid of particles,

wherein the Hamiltonian has translational symmetry. Thus, if $\mathcal{T}(\mathbf{R})$ is the translation operator which displaces all particles coordinates by an \mathbf{R} -vector, then $[\mathbf{H}, \mathcal{T}(\mathbf{R})] = 0$. Then, given that this translation operator is in reality a semigroup and given Stone's theorem, then $\mathcal{T} = e^{i\mathbf{R} \cdot \mathbf{P}}$, where \mathbf{P} is the total momentum operator. The total momentum operator is then a conserved quantity, given by

$$\mathbf{P} = \hbar \sum_{\mathbf{k}\sigma} \mathbf{k} c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}\sigma}.$$

An orthonormal basis of states $|\mathbf{P}\rangle$ with definite total momentum may be chosen. This fact can be used to "prove" the unphysical nature of a density wave. In effect, a \mathbf{Q} -density wave can be written in terms of its Fourier components as

$$\rho(\mathbf{Q}) = \sum_{\mathbf{k}\sigma} c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}+\mathbf{Q}\sigma}.$$

Said Fourier transform has a finite expectation value. However, it holds as well that

$$\langle c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}+\mathbf{Q}\sigma} \rangle = \frac{1}{\mathcal{Z}} \sum_{\mathbf{P}} e^{-\beta E_{\mathbf{P}}} \langle \mathbf{P} | c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}+\mathbf{Q}\sigma} | \mathbf{P} \rangle = 0,$$

since $c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}+\mathbf{Q}\sigma} |\mathbf{P}\rangle$ has momentum $\mathbf{P} - \mathbf{Q}$ and is thus orthogonal to \mathbf{P} . Then, according to this result, crystals do not exist. In the same way, magnetism, superconductivity and other physical phenomenon do not exist as well. What is wrong? The previous result breaks down if the sum of states in the thermodynamical average is restricted. Even though crystals with different spatial reference points (or ferromagnets with magnetizations in different directions) have formally the same energy, they are effectively decoupled due to the large energy barrier it takes to melt and then recrystallize into a new state with a shifted distance (or direction of magnetization). In those cases, where many states of the system are degenerate but separated by large energy barriers, it does not make sense to include them on equal footing in the statistical average, since they correspond to macroscopically totally different configurations. Therefore, the ergodicity postulate does not hold, giving rise to a phase space of physically separated sections, which is often illustrated by the double barrier model of phase transitions.

When at some critical temperature, the thermodynamical state of the system develops a non-zero expectation value of some macroscopic quantity, which has less symmetries than the original Hamiltonian, it is called spontaneous symmetry breaking. The quantity which signals that a phase transition has occurred is the order parameter. Some examples of order parameters are presented in [table 1](#).

Examples of Spontaneous Symmetry Breaking		
Phenomena	Order parameter physical	Order parameter mathematical
crystal	density wave	$\sum_{\mathbf{k}} \langle c_{\mathbf{k}}^\dagger c_{\mathbf{k}+\mathbf{Q}} \rangle$
ferromagnet	magnetization	$\sum_{\mathbf{k}} \langle c_{\mathbf{k}\uparrow}^\dagger c_{\mathbf{k}\uparrow} - c_{\mathbf{k}\downarrow}^\dagger c_{\mathbf{k}\downarrow} \rangle$
Bose-Einstein condensate	population of $\mathbf{k} = 0$ -state	$\langle a_{\mathbf{k}=0}^\dagger \rangle$
superconductor	pair condensate	$\langle c_{\mathbf{k}\uparrow}^\dagger c_{-\mathbf{k}\downarrow} \rangle$

TABLE 1. Typical examples of spontaneous symmetry breakings and their correspondent order parameters.

In order to avoid the paradox obtained before with the wave density, the possibility of phase transitions must be explicitly included into the theory. In the mean field approach, the trick is to include the order parameter in the choice of finite mean fields, and, of course, show that the resulting mean field Hamiltonian leads to a self-consistent finite result.

9.1.0.4. Ferromagnetism: The Heisenberg model. In ionic magnetic crystals, the interaction between the magnetic ions is due to the exchange interactions originating from the Coulomb interactions. The theory's effective Hamiltonian is given by

$$(9.1.0.15) \quad \mathbf{H} = -2 \sum_{i \sim j} J_{ij} \mathbf{S}_i \cdot \mathbf{S}_j, \text{ where } i \sim j \text{ if } \|i - j\|_1 = 1.$$

The interaction, which depends only on the inter-ion distance, is generally short ranged and assumed to be simply $J_{ij} = J_0 \delta_{\|i-j\|_1, 1}$, ie. is only non-zero for nearest-neighbours.

If $J < 0$, the spins tend to become anti-parallel whereas for $J > 0$, it is energetically favourable for the spins to be parallel, the first case corresponds to the antiferromagnetic case while the latter to the ferromagnetic case. For now, consider only the ferromagnetic case, $J > 0$.

As the model stands, it is inherently complicated and cannot be solved in general, the spins of the individual ions being strongly correlated. However, a mean field solution can be readily be found to give sensible physical results. The mean field decomposition reads

$$(9.1.0.16) \quad \mathbf{H} \approx \mathbf{H}_{\text{MF}} = -2 \sum_{i \sim j} J_{ij} \left(\langle \mathbf{S}_i \rangle \mathbf{S}_j + \mathbf{S}_i \langle \mathbf{S}_j \rangle - \langle \mathbf{S}_i \rangle \langle \mathbf{S}_j \rangle \right).$$

From symmetry arguments it is expected that the expectation value of $\langle \mathbf{S}_i \rangle$ is zero, since all directions are equivalent. But this is not the right answer, since a symmetry has been broken, and has to be presumed to be non-zero. Furthermore, since the system has translational symmetry, this average is expected to be position coordinate independent. This would not be the case for an antiferromagnetic solution, wherein position is important since the spins point on opposite directions on even and odd sites. In the ferromagnetic case, a finite but spatially independent average spin polarization is studied. In particular, if the z -axis is chosen to lie along the direction of the magnetization, then the mean field assumption is

$$\langle \mathbf{S}_i \rangle = \langle \mathbf{S}_z \rangle \mathbf{e}_z,$$

and each spin's magnetic moment \mathbf{m} , which by assumption is equal for all sites, reads

$$\mathbf{m} = 2 \sum_j J_{ij} \langle \mathbf{S}_j \rangle \mathbf{e}_z = 2nJ_0 \langle \mathbf{S}_z \rangle \mathbf{e}_z, \text{ where } n \text{ is the number of neighbours.}$$

For a square lattice, $\Lambda \subseteq \mathbb{Z}^d$ and $n = 2d$. Then mean field Hamiltonian is then

$$\mathbf{H}_{\text{MFT}} = -2 \sum_i \mathbf{m} \cdot \mathbf{S}_i + mN \langle \mathbf{S}_z \rangle,$$

which is now diagonal in the site index and hence easily solved. If the ions are spin- $\frac{1}{2}$ particles, with this simplification, the mean field partition function is

$$\mathcal{Z}_{\text{MFT}} = \left(e^{\beta m} + e^{-\beta m} \right)^N e^{\beta N m \langle \mathbf{S}_z \rangle} = \left[\left(e^{\beta m} + e^{-\beta m} \right) e^{\frac{\beta m^2}{2nJ_0}} \right]^N,$$

with one term for each possible spin projection.

The self-consistency equation is found by minimizing the free energy

$$\frac{\partial F_{\text{MFT}}}{\partial m} = -\frac{1}{\beta} \frac{\partial \log \mathcal{Z}_{\text{MFT}}}{\partial m} = -N \left(\frac{e^{\beta m} - e^{-\beta m}}{e^{\beta m} + e^{-\beta m}} \right) + N \frac{m}{nJ_0} = 0 \Rightarrow \alpha = \tanh(b\alpha), \text{ where } \alpha = \frac{m}{nJ_0},$$

$$b = nJ_0\beta.$$

For small α , the previous equation reads

$$\alpha \approx b\alpha - \frac{(b\alpha)^3}{3},$$

from which it is evident that there is no solution for $b < 1$. Thus, the critical temperature T_c can be found since at this temperature magnetism disappears. From the $(b_c = 1)$ -condition, it is found that $k_B T_c = n J_0$. Furthermore, for small α , the solution for the magnetization can be found

$$(9.1.0.17) \quad \alpha \approx \frac{1}{b} \sqrt{\frac{3(b-1)}{b}} \Rightarrow m \approx n J_0 \sqrt{3 \frac{T_c - T}{T_c}}, \text{ valid for } T \rightarrow T_c.$$

At $T = 0$, where $t \rightarrow \infty$, the solution is $\alpha = 1$ and hence $m = n J_0$. For the functional form of the magnetization in the entire range of temperature, the transcendental equation must be solved numerically.

9.1.0.5. The Stoner-Hubbard of metallic ferromagnets. In magnets where the electrons both generate magnetic moments and also form conduction bands, the Heisenberg model proves to be inadequate given that the spins are not localized. Metallic magnetism occurs eg. in transition metals, where the conduction bands are formed by the narrower d or f orbitals. The interaction between two particles in those orbitals is stronger than between electrons occupying the more spread out s or p orbitals and hence give a larger correlation between electrons. Typical metals where correlations between conduction band electrons are important for Fe and Ni.

Since the short-range of the interaction is important, the physical model -the so-called Hubbard model- must incorporate a simple physical fact. Namely, that the Coulomb interaction between electrons is taken to be point-like in real space, and hence constant in momentum space

$$(9.1.0.18) \quad \mathbf{H}_{\text{Hub}} = \sum_{\mathbf{k}\sigma} \varepsilon_{\mathbf{k}} \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} + \frac{U}{2\mathcal{V}} \sum_{\substack{\mathbf{k}, \mathbf{k}', \mathbf{q} \\ \sigma\sigma'}} \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}^\dagger \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'} \mathbf{c}_{\mathbf{k}\sigma}.$$

Then, the Hartree-Fock approximation scheme may be employed on this model, by searching for a ferromagnetic solution which allows expectation values to depend on the direction of the spin. The mean field parameters are

$$(9.1.0.19) \quad \langle \mathbf{c}_{\mathbf{k}\uparrow}^\dagger \mathbf{c}_{\mathbf{k}'\uparrow} \rangle = \delta_{\mathbf{k}\mathbf{k}'} \bar{n}_{\mathbf{k}\uparrow} \quad \langle \mathbf{c}_{\mathbf{k}\downarrow}^\dagger \mathbf{c}_{\mathbf{k}'\downarrow} \rangle = \delta_{\mathbf{k}\mathbf{k}'} \bar{n}_{\mathbf{k}\downarrow},$$

and the mean-field interaction Hamiltonian becomes

$$\begin{aligned} \mathbf{V}_{\text{int}}^{\text{MFT}} &= \frac{U}{\mathcal{V}} \sum_{\substack{\mathbf{k}, \mathbf{k}', \mathbf{q} \\ \sigma\sigma'}} \left[\mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}^\dagger \langle \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'} \rangle \mathbf{c}_{\mathbf{k}\sigma} - \langle \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'} \rangle \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \right. \\ &\quad \left. - \frac{1}{2} \left(\langle \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \rangle \langle \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'} \rangle - \langle \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'} \rangle \langle \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \rangle \right) \right] \\ &= \frac{U}{\mathcal{V}} \sum_{\substack{\mathbf{k}, \mathbf{k}', \mathbf{q} \\ \sigma\sigma'}} \left[\mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \delta_{\mathbf{k}'-\mathbf{q}\mathbf{k}} \bar{n}_{\mathbf{k}'\sigma'} - \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \delta_{\mathbf{k}+\mathbf{q}\mathbf{k}'} \delta_{\sigma\sigma'} \bar{n}_{\mathbf{k}'\sigma'} \right. \\ &\quad \left. - \frac{1}{2} \left(\delta_{\mathbf{k}+\mathbf{q}\mathbf{k}} \bar{n}_{\mathbf{k}\sigma} \delta_{\mathbf{k}'-\mathbf{q}\mathbf{k}'} \bar{n}_{\mathbf{k}'\sigma'} - \delta_{\mathbf{k}+\mathbf{q}\mathbf{k}'} \delta_{\sigma\sigma'} \bar{n}_{\mathbf{k}'\sigma'} \delta_{\mathbf{k}'-\mathbf{q}\mathbf{k}} \delta_{\sigma\sigma'} \bar{n}_{\mathbf{k}\sigma} \right) \right] \\ &= U \sum_{\substack{\mathbf{k}, \mathbf{k}' \\ \sigma\sigma'}} \left[\mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \frac{\bar{n}_{\mathbf{k}'\sigma'}}{\mathcal{V}} - \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \delta_{\sigma\sigma'} \frac{\bar{n}_{\mathbf{k}'\sigma'}}{\mathcal{V}} - \frac{1}{2} \left(\bar{n}_{\mathbf{k}\sigma} \frac{\bar{n}_{\mathbf{k}'\sigma'}}{\mathcal{V}} - \delta_{\sigma\sigma'} \bar{n}_{\mathbf{k}\sigma'} \frac{\bar{n}_{\mathbf{k}\sigma}}{\mathcal{V}} \right) \right] \\ &= U \sum_{\mathbf{k}\sigma\sigma'} \left[\mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} (\bar{n}_{\sigma'} - \delta_{\sigma\sigma'} \bar{n}_\sigma) - \mathcal{V} \bar{n}_\sigma \bar{n}_{\sigma'} + \mathcal{V} \bar{n}_\sigma^2 \right], \text{ where } \bar{n}_\sigma = \frac{1}{\mathcal{V}} \sum_{\mathbf{k}'} \langle \mathbf{c}_{\mathbf{k}'\sigma}^\dagger \mathbf{c}_{\mathbf{k}'\sigma} \rangle \text{ is the spin density.} \end{aligned}$$

The full Hamiltonian reads

$$(9.1.0.20) \quad \mathbf{H}_{\text{MFT}} = \sum_{\mathbf{k}\sigma} \varepsilon_{\mathbf{k}\sigma}^{\text{MFT}} \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} - \frac{U\mathcal{V}}{2} \sum_{\sigma\sigma'} \left(\bar{n}_\sigma \bar{n}_{\sigma'} - \bar{n}_\sigma^2 \delta_{\sigma\sigma'} \right), \text{ with } \varepsilon_{\mathbf{k}\sigma}^{\text{MFT}} = \varepsilon_{\mathbf{k}} + U(n_\uparrow + \bar{n}_\downarrow - \bar{n}_\sigma) = \varepsilon_{\mathbf{k}} + U \bar{n}_{\bar{\sigma}}.$$

The mean field solution is found via minimization,

$$\begin{aligned}
0 &= \frac{d}{d\bar{n}_\sigma} F_{\text{MFT}} = \frac{d}{d\bar{n}_\sigma} \left(-\frac{\log \mathcal{Z}_{\text{MFT}}}{\beta} \right) = \frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left(e^{-\beta \mathbf{H}_{\text{MFT}}} \frac{d}{d\bar{n}_\sigma} \mathbf{H}_{\text{MFT}} \right) \\
&= -\frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left[e^{-\beta \mathbf{H}_{\text{MFT}}} \left(\sum_{\mathbf{k}\sigma} \frac{d\varepsilon_{\mathbf{k}\sigma}^{\text{MFT}}}{d\bar{n}_\sigma} \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} - \frac{U\mathcal{V}}{2} \sum_{\sigma\sigma'} (\bar{n}_{\sigma'} + \bar{n}_\sigma \delta_{\sigma\sigma'} - 2\bar{n}_{\sigma'} \delta_{\sigma\sigma'}) \right) \right] \\
&= -\frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left[e^{\beta \mathbf{H}_{\text{MFT}}} \left(\sum_{\mathbf{k}\sigma} -U \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} - \frac{U\mathcal{V}}{2} \sum_{\sigma\sigma'} (\bar{n}_{\sigma'} - \bar{n}_{\sigma'} \delta_{\sigma\sigma'}) \right) \right] \\
&= U \sum_{\mathbf{k}\sigma} \langle \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \rangle_{\text{MFT}} - \frac{U\mathcal{V}}{2} \sum_{\sigma\sigma'} (\bar{n}_{\sigma'} - \bar{n}_{\sigma'} \delta_{\sigma\sigma'}) = U \langle \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \rangle_{\text{MFT}} + U\mathcal{V}\bar{n}_\sigma \\
&= \frac{U}{\mathcal{Z}_{\text{MFT}}} \sum_{\mathbf{k}\sigma} \text{Tr} \left[e^{-\beta \mathbf{H}_{\text{MFT}}} \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \right] + \frac{U\mathcal{V}}{2} \sum_{\sigma\sigma'} (\bar{n}_{\sigma'} - \bar{n}_{\sigma'} \delta_{\sigma\sigma'}) \frac{1}{\mathcal{Z}_{\text{MFT}}} \text{Tr} \left[e^{-\beta \mathbf{H}_{\text{MFT}}} \mathbf{1} \right] \\
&= \frac{U}{\mathcal{V}} \sum_{\mathbf{k}} \langle \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \rangle_{\text{MFT}} - \frac{U\mathcal{V}}{2} \left(2 \sum_{\sigma'} \bar{n}_{\sigma'} - \sum_{\sigma} \bar{n}_\sigma \right) = \sum_{\sigma} \left[\sum_{\mathbf{k}} \frac{U}{\mathcal{V}} \langle \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \rangle_{\text{MFT}} + \frac{U\mathcal{V}}{2} \bar{n}_\sigma \right] \\
&\Rightarrow \bar{n}_\sigma = \frac{1}{\mathcal{V}} \sum_{\mathbf{k}} \langle \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} \rangle_{\text{MFT}} = \frac{1}{\mathcal{V}} \sum_{\mathbf{k}} n_F(\varepsilon_{\mathbf{k}\sigma}^{\text{MFT}}).
\end{aligned}$$

At zero temperature,

$$\begin{aligned}
(9.1.0.21) \quad \bar{n}_\uparrow &= \int_{\mathbb{R}^3} \frac{d\mathbf{k}}{(2\pi)^3} \Theta \left(\mu - \frac{\hbar^2 k^2}{2m} - U\bar{n}_\downarrow \right) = \frac{k_{F\uparrow}^3}{6\pi^2} \\
\bar{n}_\downarrow &= \int_{\mathbb{R}^3} \frac{d\mathbf{k}}{(2\pi)^3} \Theta \left(\mu - \frac{\hbar^2 k^2}{2m} - U\bar{n}_\uparrow \right) = \frac{k_{F\downarrow}^3}{6\pi^2}, \quad \text{where } \frac{\hbar^2 k_{F\sigma}}{2m} + U\bar{n}_{(\sigma)^*} = \mu.
\end{aligned}$$

Thus, yielding two equations

$$(9.1.0.22) \quad \frac{\hbar^2}{2m} (6\pi)^{\frac{2}{3}} \bar{n}_\uparrow^{\frac{2}{3}} + U\bar{n}_\downarrow = \mu, \quad \frac{\hbar^2}{2m} (6\pi)^{\frac{2}{3}} \bar{n}_\downarrow^{\frac{2}{3}} + U\bar{n}_\uparrow = \mu.$$

Let,

$$\zeta = \frac{\bar{n}_\uparrow - \bar{n}_\downarrow}{\bar{n}}, \quad \gamma = \frac{2mU\bar{n}^{\frac{1}{3}}}{(3\pi^2)^{\frac{3}{2}}\hbar^2}, \quad \bar{n} = \bar{n}_\uparrow + \bar{n}_\downarrow,$$

then, by subtracting the two equations,

$$(9.1.0.23) \quad \bar{n}_\uparrow^{\frac{2}{3}} - \bar{n}_\downarrow^{\frac{2}{3}} = \frac{2mU}{\hbar^2} (6\pi)^{-\frac{2}{3}} (\bar{n}_\uparrow - \bar{n}_\downarrow) \Rightarrow \gamma\zeta = (1 + \zeta)^{\frac{2}{3}} - (1 - \zeta)^{\frac{2}{3}},$$

which has three types of solutions:

- $\gamma < \frac{4}{3}$: isotropic solution: $\zeta = 0$, which corresponds to a normal state.
- $\frac{4}{3} < \gamma < 2^{\frac{2}{3}}$: Partial polarization $0 < \zeta < 1$, which corresponds to a weak ferromagnet.
- $\gamma > 2^{\frac{2}{3}}$: Full polarization $\zeta = 1$, which corresponds to a strong ferromagnet.

The possibility for a magnetic solution can be traced back to the spin dependent eigenenergies, where it is clear that the mean field energy of a given spin direction depends on the occupation of the opposite spin direction, whereas the energy does not depend on the occupation of the same spin direction. This is a consequence of two things: the short range interaction and the exchange term. Both can be understood from the Pauli principle, which ensures that electrons with the same spin never occupy the same spatial orbitals and therefore, if the interaction is short-range, they cannot interact. This leaves interactions between opposite spins as the only possibility. Thus, the interaction energy is lowered by having a polarized ground state, which on the other hand for a fixed density

costs kinetic energy. The competition between the kinetic and potential energy gives rise to the phase transition.

The Stoner-Hubbard model gives a reasonable account for metallic magnets and is capable of qualitatively explaining the properties of excitations in the spin polarized states.

9.1.0.6. Superconductivity: BCS theory. One of the most striking examples of symmetry breaking is the superconducting phase transition. Below the critical temperature the metal which turns superconducting, has no resistance whatsoever, and it exhibits perfect diamagnetism, called Meisner effect, which means that magnetic fields are totally expelled from the interior of the material. These astonishing phenomena result from the superconducting state having a new form of symmetry breaking, namely loss of global gauge invariance. Both the new type of phase and the appearance of an excitation gap are explained by the Bardeen-Cooper-Schrieffer theory.

Breaking of global gauge symmetry and its consequences. What kind of broken symmetry may give rise to a superconducting state? the relevant symmetry is the global gauge symmetry, ie. all electrons may be added the same constant extra phase without altering the Hamiltonian. The ferromagnetic analog is that all spins can be rotated by some angle without changing the Hamiltonian, due to a global $SO(3)$ -symmetry. In that case, a broken $SO(3)$ -symmetry, entails that the expectation value $\langle \mathbf{S} \rangle$ is not invariant under rotation, since it will change the magnetization's direction. In the same venue, a phase rotation changes the superconducting order parameter, which is of the form $\langle c_\mu c_{\mu'} \rangle$ ie. the expectation value of two annihilation operators. Of course, the number of particles is conserved, but this is not a problem in our setting for we are only interested in superconductors connected to electron reservoirs.

While it is clear why the finite expectation value $\langle \mathbf{S} \rangle$ gives a magnetization in the case of a ferromagnet, it is not so clear why broken symmetry in the superconducting case leads to a system without resistance. Moreover, even though the superconducting phase is sensitive to a change in global phase, it is clear that a constant phase cannot have any measurable effect, since all expectation values are given in terms of the absolute square of the wave function. However, phase gradients can have an effect.

Consider a superconducting state, which is ascribed by a phase which depends on the position, $\phi(\mathbf{x})$, and varies extremely slowly, such that it takes a macroscopic distance to see any significant changes in $\phi(\mathbf{x})$. For any other non-superconducting system, it would not make sense to talk about a quantum mechanical phase difference over macroscopic distances, simply because quantum coherence is destroyed by all sorts of scattering events on rather short lengths ($\sim 10nm$ for metals). To argue that the superconducting state depends on phase differences over macroscopic distances is unusual and implies that superconductivity is a macroscopic quantum phenomenon.

Consider the following unitary transformation, which changes the phase

$$U = \exp \left(i \int_{\mathbb{R}^3} d\mathbf{x} \rho(\mathbf{x}) \phi(\mathbf{x}) \right), \text{ where } \rho(\mathbf{x}) = \Psi^\dagger(\mathbf{x}) \Psi(\mathbf{x}).$$

When applied to quantum field operators it has the following properties

$$(9.1.0.24) \quad \tilde{\Psi}(\mathbf{x}) = U \Psi(\mathbf{x}) U^{-1} = \exp \left(i \int_{\mathbb{R}^3} d\mathbf{x} \rho(\mathbf{x}) \phi(\mathbf{x}) \right) \Psi(\mathbf{x}) \exp \left(- i \int_{\mathbb{R}^3} d\mathbf{x} \rho(\mathbf{x}) \phi(\mathbf{x}) \right)$$

$$(9.1.0.25) \quad = \Psi(\mathbf{x}) e^{-i\phi(\mathbf{x})}$$

$$(9.1.0.26) \quad \tilde{\Psi}^\dagger(\mathbf{x}) = U \Psi^\dagger(\mathbf{x}) U^{-1} = \Psi^\dagger(\mathbf{x}) e^{i\phi(\mathbf{x})}.$$

These equations follow from the differential equation problem

$$\begin{cases} \frac{\delta \tilde{\Psi}(\mathbf{x})}{\delta \phi(\mathbf{x}')} = i U [\rho(\mathbf{x}'), \Psi(\mathbf{x})] U^{-1} = i \tilde{\Psi}(\mathbf{x}) \delta(\mathbf{x} - \mathbf{x}') \\ \tilde{\Psi}(\mathbf{x}) \Big|_{\phi=0} = \Psi(\mathbf{x}) \end{cases}$$

The transformed partition function is given by

$$\tilde{Z} = \text{Tr}'[U e^{-\beta \mathbf{H}} U^{-1}] = \text{Tr}'[e^{-\beta \tilde{\mathbf{H}}}],$$

note that, if the cyclic property of the trace had been used, U would have disappeared all-together. However, when dealing with systems with broken symmetry, the sum-over-states has to be restricted so that the cyclic property does not hold. In other words, the trace is restricted due to the spontaneous symmetry breaking. The transformation only changes the kinetic term of the Hamiltonian, since the Coulomb interaction term, the impurity scattering term and the phonon coupling term only involves the, phase-invariant, electron density operator $\rho(\mathbf{x})$. The kinetic term reads

$$\begin{aligned} \tilde{\mathbf{H}} &= \frac{1}{2m_e} \int_{\mathbb{R}^3} d\mathbf{x} \Psi^\dagger(\mathbf{x}) e^{i\phi(\mathbf{x})} \left(\frac{\hbar}{i} \nabla + e\mathbf{A} \right)^2 e^{-i\phi(\mathbf{x})} \Psi(\mathbf{x}) \\ (9.1.0.27) \quad &= \frac{1}{2m_e} \int_{\mathbb{R}^3} d\mathbf{x} \Psi^\dagger(\mathbf{x}) \left(\frac{\hbar}{i} \nabla + e\mathbf{A} - \hbar \nabla \phi(\mathbf{x}) \right)^2 \Psi(\mathbf{x}) \\ &= \mathbf{H} - \hbar \int_{\mathbb{R}^3} d\mathbf{x} \nabla \phi(\mathbf{x}) \cdot \mathbf{J}(\mathbf{x}) - \frac{\hbar}{2m_e} \rho(\mathbf{x}) (\nabla \phi(\mathbf{x}))^2. \end{aligned}$$

Therefore, the phase is a macroscopic quantity. The lowest free energy condition can be found by minimizing the free energy with respect to the phase. It is clear that the energy does not depend with ϕ itself, but on its gradient, hence

$$\begin{aligned} (9.1.0.28) \quad \frac{\delta F}{\delta \phi(\mathbf{x})} &= -\hbar \int_{\mathbb{R}^3} d\mathbf{x} \mathbf{J}(\mathbf{x}) - \frac{\hbar}{m_e} \rho(\mathbf{x}) \nabla \phi(\mathbf{x}) \\ &= -\langle \mathbf{J}(\mathbf{x}) \rangle + \frac{\hbar}{m_e} \langle \rho(\mathbf{x}) \rangle \nabla \phi(\mathbf{x}) = 0 \Rightarrow \langle \mathbf{J} \rangle = \frac{\hbar \rho_0}{m_e} \nabla \phi. \end{aligned}$$

Physical Context

The previous equation's physical meaning is that, by forcing a phase gradient onto the system, it minimizes its energy by carrying a current even in thermodynamical equilibrium, giving rise to a non-equilibrium current. In reality, the electron density should only include those electrons participating in the superconducting condensate, the so-called superfluid. In this derivation, it is assumed all electrons participate.

In the normal state of metals, a current is always associated with a non-equilibrium state, where energy is constantly dissipated from the driving source and absorbed in the conductor. Of course, there is an energy cost for the system to carry the current, but as long as this cost is smaller than the alternative which is to go out of the superconducting state, the current carrying state is chosen. The critical current is reached when the energies are equal, and then the superconductor goes into the normal state. Thus, when a phase gradient is imposed to a system, wherein the energy is assumed to depend on phase differences on a macroscopic scale, it unavoidably leads to the conclusion that the system will carry a dissipationless current in order to minimize the energy cost of the phase gradient. Finally, it should be noted that the appearance of the excitation gap is not the reason for the superconductivity itself. The superconductivity arises from the lack of gauge invariance, and in fact gapless superconductors do exist.

Microscopic theory:

An alternative picture of superconductivity can be given, the so-called Bardeen-Cooper-Schrieffer theory.

Historical Context:

The understanding that superconductivity was closely related to the electron-phonon coupling was clear from the 1950's onwards, where the isotope effect was discovered. Also the idea that superconductivity was somehow similar to the Bose-Einstein condensate, with the bosons being electron pairs, had been tried and in fact was the underlying idea of London's theory in 1935. However, only in 1956 Cooper showed that the Fermi surface of the normal metal state was unstable towards formation of bound pairs of electrons. Subsequently in 1957, when the superconducting state was derived using a variation wavefunction by BCS were the principles fully understood. It turns out that the phonon mediated electron-electron interaction, in fact has a range in frequency and momentum space where it is negative, ie. attractive. This happens for energy exchanges of order the Debye energy, ω_D , for metals is much smaller than the Fermi energy $\omega_D \ll E_F$. Furthermore, from the Cooper instability, it follows that the phonon-mediated interactions tends to pair electrons with opposite spin and momentum.

Let a Cooper pair operator be

$$(9.1.0.29) \quad \mathbf{b}_k = \mathbf{c}_{k\downarrow} \mathbf{c}_{-k\uparrow}.$$

Then, the previous two physical inputs, led BCS to suggest the following remarkably successful model Hamiltonian to explain the superconducting state. The BCS effective Hamiltonian model is

$$(9.1.0.30) \quad \mathbf{H}_{\text{BCS}} = \sum_{\mathbf{k}\sigma} \varepsilon_{\mathbf{k}} \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} + \sum_{\mathbf{k}\mathbf{k}'} \mathfrak{V}_{\mathbf{k}\mathbf{k}'} \mathbf{c}_{\mathbf{k}\uparrow}^\dagger \mathbf{c}_{-\mathbf{k}\downarrow}^\dagger \mathbf{c}_{-\mathbf{k}'\downarrow} \mathbf{c}_{\mathbf{k}'\uparrow},$$

where $\mathfrak{V}_{\mathbf{k}\mathbf{k}'} = -V_0 \left(\Theta(\varepsilon_{\mathbf{k}} + \varepsilon_F + \omega_D) - \Theta(\varepsilon_{\mathbf{k}} - \varepsilon_F - \omega_D) \right) \delta_{\mathbf{k}\mathbf{k}'} \neq 0$ if $\varepsilon_{\mathbf{k}} \in (\varepsilon_F - \omega_D, \varepsilon_F + \omega_D)$. The interaction includes only pair interactions and the remaining interaction is supposed to be included in $\varepsilon_{\mathbf{k}}$ via a Hartree-Fock term. The attractive interaction is due to the phenomenon of an electron propagating through a crystal, thereby attracting the positive ions and thus effectively creating a positive trace behind it, which is felt by the other electrons as an attractive interaction. It turns out that this effective interaction is most important for electrons occupying time reversed states, and in fact they can form a bounded state which is a Cooper pair. The Cooper pairs is thus a bounded state of an electron in state $\psi_{\mu}(\mathbf{x})$ and another electron in state $\psi_{\mu}^*(\mathbf{x})$, or in the homogeneous case, electrons in states \mathbf{k} and $-\mathbf{k}$.³⁵

The mean field assumption made by BCS is that the pair operator has a finite expectation value, which deviates only a little from its average values. Therefore, the BCS mean field Hamiltonian is derived in full analogy with Hartree-Fock's mean field theory, yielding

$$(9.1.0.31) \quad \mathbf{H}_{\text{BCS}}^{\text{MFT}} = \sum_{\mathbf{k}\sigma} \varepsilon_{\mathbf{k}} \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} - \sum_{\mathbf{k}} \Delta_{\mathbf{k}} \mathbf{c}_{\mathbf{k}\uparrow}^\dagger \mathbf{c}_{-\mathbf{k}\downarrow}^\dagger - \sum_{\mathbf{k}} \Delta_{\mathbf{k}}^* \mathbf{c}_{\mathbf{k}\downarrow} \mathbf{c}_{-\mathbf{k}\uparrow} + \sum_{\mathbf{k}\mathbf{k}'} \mathfrak{V}_{\mathbf{k}\mathbf{k}'} \langle \mathbf{c}_{\mathbf{k}\uparrow}^\dagger \mathbf{c}_{-\mathbf{k}\downarrow}^\dagger \rangle \langle \mathbf{c}_{\mathbf{k}\downarrow} \mathbf{c}_{-\mathbf{k}\uparrow} \rangle, \text{ where } \Delta_{\mathbf{k}} = - \sum_{\mathbf{k}'} \mathfrak{V}_{\mathbf{k}\mathbf{k}'} \langle \mathbf{c}_{-\mathbf{k}'\downarrow} \mathbf{c}_{\mathbf{k}'\uparrow} \rangle,$$

³⁵Note that, since the typical energy exchange due to the attractive interaction is the Debye energy ω_D one could naively expect that the energy scale for the superconducting transition temperature would be of $\mathcal{O}(\frac{\omega_D}{k})$ -order. This is however far from the truth, since $\frac{\omega_D}{k}$ is typically of the order of several hundred kelvin, while the critical temperatures found in conventional low superconductors are, in most cases, less than 10K. Thus, it seems that a new energy scale is needed.

which is quadratic in electron operators and should be readily solvable. It is however unusual in that terms $c^\dagger c^\dagger$ and cc both appear. Said Hamiltonian can then be solved via a Bogoliubov transformation. Note that $H_{\text{BCS}}^{\text{MFT}}$ may be rewritten in \mathbb{C}^2 -matrix form

$$(9.1.0.32) \quad \begin{aligned} H_{\text{BCS}}^{\text{MFT}} &= \sum_{\mathbf{k}} \begin{pmatrix} c_{\mathbf{k}\uparrow}^\dagger & c_{-\mathbf{k}\downarrow} \end{pmatrix} \begin{pmatrix} \varepsilon_{\mathbf{k}} & \Delta_{\mathbf{k}} \\ \Delta_{\mathbf{k}}^* & -\varepsilon_{\mathbf{k}} \end{pmatrix} \begin{pmatrix} c_{\mathbf{k}\downarrow}^\dagger \\ c_{-\mathbf{k}\uparrow} \end{pmatrix} + \sum_{\mathbf{k}} \varepsilon_{\mathbf{k}} + \sum_{\mathbf{k}\mathbf{k}'} \mathfrak{V}_{\mathbf{k}\mathbf{k}'} \langle c_{\mathbf{k}\uparrow}^\dagger c_{-\mathbf{k}\downarrow}^\dagger \rangle \langle c_{\mathbf{k}\downarrow} c_{-\mathbf{k}\uparrow} \rangle \\ &= \sum_{\mathbf{k}} \mathbf{A}_{\mathbf{k}}^\dagger \mathbf{H}_{\mathbf{k}} \mathbf{A}_{\mathbf{k}} + \text{tr}\{\mathbf{H}\}, \end{aligned}$$

where $(\mathbf{A}_{\mathbf{k}})^T = \begin{pmatrix} c_{\mathbf{k}\uparrow}^\dagger & c_{-\mathbf{k}\downarrow} \end{pmatrix}$. Then, the Hamiltonian can be rewritten in diagonal form, via a redefinition of operators, said redefinition being given by a unitary transformation,

$$(9.1.0.33) \quad \mathbf{B}_{\mathbf{k}} = \begin{pmatrix} \gamma_{\mathbf{k}\uparrow} \\ \gamma_{-\mathbf{k}\downarrow}^\dagger \end{pmatrix} = \mathbf{U}_{\mathbf{k}}^{-1} \mathbf{A}_{\mathbf{k}} \text{ where } \mathbf{U}_{\mathbf{k}} = \begin{pmatrix} u_{\mathbf{k}} & -v_{\mathbf{k}} \\ v_{\mathbf{k}}^* & u_{\mathbf{k}}^* \end{pmatrix},$$

which diagonalizes the Hamiltonian provided

$$(9.1.0.34) \quad \begin{aligned} \mathbf{U}_{\mathbf{k}}^\dagger \mathbf{H}_{\mathbf{k}} \mathbf{U}_{\mathbf{k}} &= \begin{pmatrix} u_{\mathbf{k}}^* & v_{\mathbf{k}} \\ -v_{\mathbf{k}}^* & u_{\mathbf{k}} \end{pmatrix} \\ &\Rightarrow \begin{pmatrix} u_{\mathbf{k}}^* & -v_{\mathbf{k}}^* \\ v_{\mathbf{k}} & u_{\mathbf{k}}^* \end{pmatrix} \begin{pmatrix} \varepsilon_{\mathbf{k}} & \Delta_{\mathbf{k}} \\ \Delta_{\mathbf{k}}^* & -\varepsilon_{\mathbf{k}} \end{pmatrix} \begin{pmatrix} u_{\mathbf{k}} & -v_{\mathbf{k}} \\ v_{\mathbf{k}}^* & u_{\mathbf{k}}^* \end{pmatrix} = \begin{pmatrix} u_{\mathbf{k}}^* & v_{\mathbf{k}} \\ -v_{\mathbf{k}}^* & u_{\mathbf{k}} \end{pmatrix} \begin{pmatrix} \varepsilon_{\mathbf{k}} u_{\mathbf{k}} + \Delta_{\mathbf{k}} v_{\mathbf{k}}^* & -\varepsilon_{\mathbf{k}} v_{\mathbf{k}} + \Delta_{\mathbf{k}} u_{\mathbf{k}}^* \\ \Delta_{\mathbf{k}}^* u_{\mathbf{k}} - \varepsilon_{\mathbf{k}} v_{\mathbf{k}}^* & -\Delta_{\mathbf{k}}^* v_{\mathbf{k}} - \varepsilon_{\mathbf{k}} u_{\mathbf{k}}^* \end{pmatrix} \\ &= \begin{pmatrix} \varepsilon_{\mathbf{k}} |u_{\mathbf{k}}|^2 + \Delta_{\mathbf{k}} u_{\mathbf{k}}^* v_{\mathbf{k}}^* + \Delta_{\mathbf{k}}^* u_{\mathbf{k}} v_{\mathbf{k}} - \varepsilon_{\mathbf{k}} |v_{\mathbf{k}}|^2 & -\varepsilon_{\mathbf{k}} u_{\mathbf{k}}^* v_{\mathbf{k}} + \Delta_{\mathbf{k}} |u_{\mathbf{k}}|^2 - \Delta_{\mathbf{k}}^* v_{\mathbf{k}}^2 - \varepsilon_{\mathbf{k}} u_{\mathbf{k}}^* v_{\mathbf{k}} \\ -\varepsilon_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}}^* - \Delta_{\mathbf{k}} |v_{\mathbf{k}}|^2 + \Delta_{\mathbf{k}}^* u_{\mathbf{k}}^2 - \varepsilon_{\mathbf{k}} u_{\mathbf{k}} v_{\mathbf{k}}^* & \varepsilon_{\mathbf{k}} |v_{\mathbf{k}}|^2 - \Delta_{\mathbf{k}} u_{\mathbf{k}}^* v_{\mathbf{k}}^* - \Delta_{\mathbf{k}}^* v_{\mathbf{k}} u_{\mathbf{k}} - \varepsilon_{\mathbf{k}} |u_{\mathbf{k}}|^2 \end{pmatrix} \\ &\Rightarrow \begin{aligned} |u_{\mathbf{k}}|^2 &= \frac{1}{2} \left(1 + \frac{\varepsilon_{\mathbf{k}}}{E_{\mathbf{k}}} \right) \\ |v_{\mathbf{k}}|^2 &= \frac{1}{2} \left(1 - \frac{\varepsilon_{\mathbf{k}}}{E_{\mathbf{k}}} \right) \end{aligned} \quad E_{\mathbf{k}} = \sqrt{\varepsilon_{\mathbf{k}}^2 + |\Delta_{\mathbf{k}}|^2} = -\tilde{E}_{\mathbf{k}}. \end{aligned}$$

The new fermionic operators, which diagonalize the Hamiltonian, are fermionic quasiparticles, frequently called bogoliubons and are the superposition of electrons and holes. This new particle arises due to a lack of particle conservation in the mean field Hamiltonian. Note that there are two different bogoliubons inherited from the two-fold spin degeneracy. The quasiparticles operators are given by

$$(9.1.0.35) \quad \begin{aligned} \begin{pmatrix} \gamma_{\mathbf{k}\uparrow} \\ \gamma_{-\mathbf{k}\downarrow}^\dagger \end{pmatrix} &= \begin{pmatrix} u_{\mathbf{k}}^* & v_{\mathbf{k}} \\ -v_{\mathbf{k}}^* & u_{\mathbf{k}} \end{pmatrix} \begin{pmatrix} c_{\mathbf{k}\uparrow} \\ c_{-\mathbf{k}\downarrow}^\dagger \end{pmatrix} \Leftrightarrow \begin{pmatrix} c_{\mathbf{k}\uparrow} \\ c_{-\mathbf{k}\downarrow}^\dagger \end{pmatrix} = \begin{pmatrix} u_{\mathbf{k}} & -v_{\mathbf{k}} \\ v_{\mathbf{k}}^* & u_{\mathbf{k}}^* \end{pmatrix} \begin{pmatrix} \gamma_{\mathbf{k}\uparrow} \\ \gamma_{-\mathbf{k}\downarrow}^\dagger \end{pmatrix} \\ \begin{pmatrix} \gamma_{\mathbf{k}\uparrow} \\ \gamma_{-\mathbf{k}\downarrow}^\dagger \\ \gamma_{\mathbf{k}\uparrow}^\dagger \\ \gamma_{-\mathbf{k}\downarrow} \end{pmatrix} &= \begin{pmatrix} u_{\mathbf{k}}^* c_{\mathbf{k}\uparrow} + v_{\mathbf{k}} c_{-\mathbf{k}\downarrow}^\dagger \\ -v_{\mathbf{k}}^* c_{\mathbf{k}\uparrow} + u_{\mathbf{k}} c_{-\mathbf{k}\downarrow}^\dagger \\ u_{\mathbf{k}} c_{\mathbf{k}\uparrow}^\dagger + v_{\mathbf{k}}^* c_{-\mathbf{k}\downarrow} \\ -v_{\mathbf{k}} c_{\mathbf{k}\uparrow}^\dagger + u_{\mathbf{k}} c_{-\mathbf{k}\downarrow} \end{pmatrix} \Leftrightarrow \begin{pmatrix} c_{\mathbf{k}\uparrow} \\ c_{-\mathbf{k}\downarrow}^\dagger \\ c_{\mathbf{k}\uparrow}^\dagger \\ c_{-\mathbf{k}\downarrow} \end{pmatrix} = \begin{pmatrix} u_{\mathbf{k}} \gamma_{\mathbf{k}\uparrow} - v_{\mathbf{k}} \gamma_{-\mathbf{k}\downarrow}^\dagger \\ v_{\mathbf{k}}^* \gamma_{\mathbf{k}\uparrow} + u_{\mathbf{k}}^* \gamma_{-\mathbf{k}\downarrow}^\dagger \\ u_{\mathbf{k}}^* \gamma_{\mathbf{k}\uparrow}^\dagger - v_{\mathbf{k}}^* \gamma_{-\mathbf{k}\downarrow}^\dagger \\ v_{\mathbf{k}} \gamma_{\mathbf{k}\uparrow}^\dagger + u_{\mathbf{k}} \gamma_{-\mathbf{k}\downarrow}^\dagger \end{pmatrix}. \end{aligned}$$

It is necessary to inquire in the bogoliubons' nature. Namely, it they obey fermionic anticommutation rules,

$$\begin{aligned}
\{\gamma_{\mathbf{k}\uparrow}^\dagger, \gamma_{\mathbf{k}'\uparrow}\} &= \{u_{\mathbf{k}}c_{\mathbf{k}\uparrow}^\dagger + v_{\mathbf{k}}^*c_{-\mathbf{k}\downarrow}, u_{\mathbf{k}'}^*c_{\mathbf{k}'\uparrow} + v_{\mathbf{k}'}c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&= u_{\mathbf{k}}u_{\mathbf{k}'}^*\{c_{\mathbf{k}\uparrow}^\dagger, c_{\mathbf{k}'\uparrow}\} + u_{\mathbf{k}}v_{\mathbf{k}'}^*\{c_{\mathbf{k}\uparrow}^\dagger, c_{-\mathbf{k}'\downarrow}\} \\
&\quad + u_{\mathbf{k}'}v_{\mathbf{k}}\{c_{-\mathbf{k}\downarrow}, c_{\mathbf{k}'\uparrow}^\dagger\} + v_{\mathbf{k}}v_{\mathbf{k}'}^*\{c_{-\mathbf{k}\downarrow}, c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&= u_{\mathbf{k}}u_{\mathbf{k}'}^*\delta_{\mathbf{k}\mathbf{k}'} + v_{\mathbf{k}}v_{\mathbf{k}'}^*\delta_{\mathbf{k}\mathbf{k}'} = \delta_{\mathbf{k}\mathbf{k}'} \\
\{\gamma_{-\mathbf{k}\downarrow}^\dagger, \gamma_{-\mathbf{k}'\downarrow}\} &= \{-v_{\mathbf{k}}^*c_{\mathbf{k}\uparrow} + u_{\mathbf{k}}c_{-\mathbf{k}\downarrow}^\dagger, -v_{\mathbf{k}'}^*c_{\mathbf{k}'\uparrow} + u_{\mathbf{k}'}c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&= v_{\mathbf{k}}v_{\mathbf{k}'}^*\{c_{\mathbf{k}\uparrow}, c_{\mathbf{k}'\uparrow}^\dagger\} - u_{\mathbf{k}}v_{\mathbf{k}'}^*\{c_{-\mathbf{k}\downarrow}^\dagger, c_{\mathbf{k}'\uparrow}^\dagger\} \\
&\quad - u_{\mathbf{k}'}v_{\mathbf{k}}^*\{c_{-\mathbf{k}\downarrow}^\dagger, c_{\mathbf{k}'\uparrow}^\dagger\} + v_{\mathbf{k}}v_{\mathbf{k}'}^*\{c_{-\mathbf{k}\downarrow}^\dagger, c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&= u_{\mathbf{k}}u_{\mathbf{k}'}^*\delta_{\mathbf{k}\mathbf{k}'} + v_{\mathbf{k}}v_{\mathbf{k}'}^*\delta_{\mathbf{k}\mathbf{k}'} = \delta_{\mathbf{k}\mathbf{k}'}
\end{aligned}$$

$$\begin{aligned}
\{\gamma_{\mathbf{k}\uparrow}^\dagger, \gamma_{-\mathbf{k}'\downarrow}^\dagger\} &= \{u_{\mathbf{k}}c_{\mathbf{k}\uparrow}^\dagger + v_{\mathbf{k}}^*c_{-\mathbf{k}\downarrow}, -v_{\mathbf{k}'}^*c_{\mathbf{k}'\uparrow} + u_{\mathbf{k}'}c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&= -u_{\mathbf{k}}v_{\mathbf{k}'}^*\{c_{\mathbf{k}\uparrow}^\dagger, c_{\mathbf{k}'\uparrow}\} + u_{\mathbf{k}}u_{\mathbf{k}'}\{c_{\mathbf{k}\uparrow}^\dagger, c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&\quad - v_{\mathbf{k}}v_{\mathbf{k}'}\{c_{-\mathbf{k}\downarrow}, c_{\mathbf{k}'\uparrow}^\dagger\} + v_{\mathbf{k}}^*u_{\mathbf{k}'}\{c_{-\mathbf{k}\downarrow}, c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&= -u_{\mathbf{k}}v_{\mathbf{k}'}^*\delta_{\mathbf{k}\mathbf{k}'} + u_{\mathbf{k}}v_{\mathbf{k}'}^*\delta_{\mathbf{k}\mathbf{k}'} = 0 \\
\{\gamma_{\mathbf{k}\uparrow}, \gamma_{-\mathbf{k}'\downarrow}\} &= \{u_{\mathbf{k}}^*c_{\mathbf{k}\uparrow} + v_{\mathbf{k}}c_{-\mathbf{k}\downarrow}^\dagger, -v_{\mathbf{k}'}^*c_{\mathbf{k}'\uparrow} + u_{\mathbf{k}'}c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&= -u_{\mathbf{k}}^*v_{\mathbf{k}'}\{c_{\mathbf{k}\uparrow}, c_{\mathbf{k}'\uparrow}^\dagger\} + u_{\mathbf{k}}^*u_{\mathbf{k}'}\{c_{\mathbf{k}\uparrow}, c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&\quad - v_{\mathbf{k}}^*v_{\mathbf{k}'}\{c_{-\mathbf{k}\downarrow}^\dagger, c_{\mathbf{k}'\uparrow}^\dagger\} + u_{\mathbf{k}'}v_{\mathbf{k}}\{c_{-\mathbf{k}\downarrow}^\dagger, c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&= -u_{\mathbf{k}}v_{\mathbf{k}'}^*\delta_{\mathbf{k}\mathbf{k}'} + u_{\mathbf{k}}v_{\mathbf{k}'}^*\delta_{\mathbf{k}\mathbf{k}'} = 0
\end{aligned}$$

$$\begin{aligned}
\{\gamma_{\mathbf{k}\uparrow}, \gamma_{\mathbf{k}'\uparrow}\} &= \{u_{\mathbf{k}}^*c_{\mathbf{k}\uparrow} + v_{\mathbf{k}}c_{-\mathbf{k}\downarrow}^\dagger, u_{\mathbf{k}'}^*c_{\mathbf{k}'\uparrow} + v_{\mathbf{k}'}c_{-\mathbf{k}'\downarrow}^\dagger\} \\
&= u_{\mathbf{k}}^*u_{\mathbf{k}'}\{c_{\mathbf{k}\uparrow}, c_{\mathbf{k}'\uparrow}\} + u_{\mathbf{k}}^*v_{\mathbf{k}'}\{c_{\mathbf{k}\uparrow}, c_{-\mathbf{k}'\downarrow}^\dagger\} + u_{\mathbf{k}'}^*v_{\mathbf{k}}\{c_{-\mathbf{k}\downarrow}^\dagger, c_{\mathbf{k}'\uparrow}\} + v_{\mathbf{k}}v_{\mathbf{k}'}\{c_{-\mathbf{k}\downarrow}^\dagger, c_{-\mathbf{k}'\downarrow}^\dagger\} = 0,
\end{aligned}$$

thus confirming the bogoliubon's fermionic character.

Then, the Hamiltonian, in terms of the bogoliubons, reads

$$\begin{aligned}
(9.1.0.36) \quad H_{\text{BCS}}^{\text{MFT}} &= \sum_{\mathbf{k}} \mathbf{A}_{\mathbf{k}}^\dagger \mathbf{H}_{\mathbf{k}} \mathbf{A}_{\mathbf{k}} + \text{tr}\{\mathbf{H}\} = \sum_{\mathbf{k}} \mathbf{B}_{\mathbf{k}}^\dagger \mathbf{U}_{\mathbf{k}}^\dagger \mathbf{H}_{\mathbf{k}} \mathbf{U}_{\mathbf{k}} \mathbf{B}_{\mathbf{k}} + \text{tr}\{\mathbf{H}\} \\
&= E_{\mathbf{k}} \left(\gamma_{\mathbf{k}\uparrow}^\dagger \gamma_{\mathbf{k}\uparrow} + \gamma_{\mathbf{k}\downarrow}^\dagger \gamma_{\mathbf{k}\downarrow} \right) + c
\end{aligned}$$

From the new Hamiltonian and its eigenvalues, it is clear that there are no fermion excitations possible with energy less than $|\Delta|$. The mean field parameter provides an energy gap, which is called the superconducting gap, which has a number of important consequences. The self-consistent solution is found from the gap equation, given by [equation \(9.1.0.31\)](#), by calculating $\Delta_{\mathbf{k}}$'s expectation value using the diagonalized Hamiltonian, which is equivalent to minimizing the free energy with respect to the mean field parameter $\langle \mathbf{b}_{\mathbf{k}} \rangle$; thus yielding

$$\begin{aligned}
(9.1.0.37) \quad \Delta_{\mathbf{k}} &= - \sum_{\mathbf{k}'} \mathfrak{V}_{\mathbf{k}\mathbf{k}'} \langle c_{-\mathbf{k}'\downarrow} c_{\mathbf{k}'\uparrow} \rangle = - \sum_{\mathbf{k}'} \mathfrak{V}_{\mathbf{k}\mathbf{k}'} \langle (v_{\mathbf{k}'} \gamma_{\mathbf{k}'\uparrow}^\dagger + u_{\mathbf{k}'} \gamma_{-\mathbf{k}'\downarrow}) (u_{\mathbf{k}'} \gamma_{\mathbf{k}'\uparrow} - v_{\mathbf{k}'} \gamma_{-\mathbf{k}'\downarrow}^\dagger) \rangle \\
&= - \sum_{\mathbf{k}'} \mathfrak{V}_{\mathbf{k}\mathbf{k}'} \left(v_{\mathbf{k}'} u_{\mathbf{k}'} \langle \gamma_{\mathbf{k}'\uparrow}^\dagger \gamma_{\mathbf{k}'\uparrow} \rangle - v_{\mathbf{k}'} u_{\mathbf{k}'} \langle \gamma_{-\mathbf{k}'\downarrow} \gamma_{-\mathbf{k}'\downarrow}^\dagger \rangle \right) \\
&= - \sum_{\mathbf{k}'} \mathfrak{V}_{\mathbf{k}\mathbf{k}'} v_{\mathbf{k}'} u_{\mathbf{k}'} \left(1 - n_{\mathfrak{F}\mathfrak{D}}(E'_{\mathbf{k}}) \right),
\end{aligned}$$

where, in the last step, the bogoliubons' fermionic nature was used. Since these are free fermionic particles, they obey Fermi-Dirac statistics. Now, given that $\mathfrak{V}_{\mathbf{k}\mathbf{k}'}$ is finitely-non zero only for $\varepsilon_{\mathbf{k}}, \varepsilon_{\mathbf{k}'} \in [-\omega_D, \omega_D]$, and given that $\omega_D \ll E_F$, such that the density of states is constant in the energy interval in question, thus yielding

$$(9.1.0.38) \quad |\Delta| = V_0 \frac{d(E_F)}{2} \int_{-\omega_D}^{\omega_D} d\varepsilon \frac{|\Delta|}{2E} \left(1 - n_{\mathfrak{F}\mathfrak{D}}(E'_{\mathbf{k}}) \right) \Rightarrow \frac{2}{V_0 d(E_F)} = \int_0^{\omega_D} d\varepsilon \frac{\tanh\left(\frac{\beta}{2} \sqrt{\varepsilon^2 + |\Delta|^2}\right)}{\sqrt{\varepsilon^2 + |\Delta|^2}},$$

which can be numerically solved. The critical temperature is found by setting $\Delta = 0$, yielding an approximate formula

$$kT_c = 1.13\omega_D e^{-\frac{2}{V_0}d(E_F)} \text{ and } \Delta_0 \approx 2\omega_D e^{-\frac{2}{V_0}d(E_F)}.$$

Note that for metals, $V_0 d(E_F)$ is a very small number. Combining both equations yield the BCS prediction

$$\frac{2\Delta_0}{kT_c} = 3.53,$$

which is very close to the experimental result. Note that both the gap and the critical temperature are reduced by an exponential factor, compared to the bare energy scale of the interaction ω_D . This strong renormalization cannot be obtained via perturbation theory. Moreover, no phase transition can be obtained from a perturbative approach alone, since eg. functions like $\exp(1/x)$ have no Taylor expansion at $x = 0$, ruling out any analytic connection between different phases.

9.2. Coherent States. In Quantum Field Theory, the basic ingredients are fields, operators $\phi(\mathbf{x})$ defined at each point in space. Thus, it is necessary to work in a basis that diagonalizes field operators, thus the coherent states arise. Moreover, coherent states comprise states with varying particle numbers, providing an extremely useful (non-orthonormal) basis for the Fock basis.

9.2.0.1. Bosonic coherent states. For a bosonic state with annihilation operator \mathbf{b} , the coherent states are defined as \mathbf{b} 's eigenstates, namely

$$\mathbf{b} |z\rangle = z |z\rangle, \text{ where } z = |z|e^{i\phi}.$$

Every state in the Hilbert space can thus be written as

$$|z\rangle = \mathcal{B}(\mathbf{b}^\dagger) |0\rangle, \quad \text{where the function } \mathcal{B}(\mathbf{b}^\dagger) \text{ has to satisfy } [\mathbf{b}, \mathcal{B}(\mathbf{b}^\dagger)] = z\mathcal{B}(\mathbf{b}^\dagger)$$

$$\text{Given that } [\mathbf{b}, \mathcal{B}(\mathbf{b}^\dagger)] = \frac{\partial \mathcal{B}(\mathbf{b}^\dagger)}{\partial \mathbf{b}^\dagger} \Rightarrow \frac{\partial \mathcal{B}(\mathbf{b}^\dagger)}{\partial \mathbf{b}^\dagger} = z\mathcal{B}(\mathbf{b}^\dagger) \Rightarrow \mathcal{B}(\mathbf{b}^\dagger) = Ce^{z\mathbf{b}^\dagger}.$$

Thus, every \mathbf{b} -eigenstate may be written as

$$\begin{aligned} |z\rangle &= Ce^{z\mathbf{b}^\dagger} |0\rangle, \text{ with } \langle z|z\rangle = |C|^2 \langle 0| e^{z^*\mathbf{b}} e^{z\mathbf{b}^\dagger} |0\rangle \Rightarrow |z\rangle = e^{-\frac{|z|^2}{2}} e^{z\mathbf{b}^\dagger} |0\rangle \\ &= |C|^2 e^{z^*\mathbf{b} + z\mathbf{b}^\dagger + \frac{1}{2}[z^*\mathbf{b}, z\mathbf{b}^\dagger]} \\ &= |C|^2 e^{z^*z} = 1 \Rightarrow |C|^2 = e^{-z^*z} = e^{-\frac{|z|^2}{2}} \sum_{n=0}^{\infty} \frac{z^n}{n!} |n\rangle. \end{aligned}$$

For an N -boson system, the most general coherent state may be written as

$$(9.2.0.1) \quad |\mathbf{z}\rangle = \prod_{i=1}^N e^{-\frac{|z_i|^2}{2}} e^{z_i \mathbf{b}_i^\dagger} |0\rangle = e^{-i \sum_{i=1}^N \frac{|z_i|^2}{2}} e^{\sum_{i=1}^N z_i \mathbf{b}_i^\dagger} |0\rangle.$$

The bosonic coherent states have some interesting consequences, namely

- The coherent states are not pair-wise orthogonal to each other:

$$(9.2.0.2) \quad \begin{aligned} \langle z_1 | z_2 \rangle &= e^{-\frac{|z_1|^2 + |z_2|^2}{2}} \langle 0 | e^{z_1^* \mathbf{b}} e^{z_2^* \mathbf{b}^\dagger} |0\rangle \\ &= e^{-\frac{|z_1|^2 + |z_2|^2}{2}} e^{z_1^* z_2} \end{aligned}$$

- The coherent states form an over complete basis

$$\text{id}_{\mathcal{F}} = \int_{\mathbb{C}} \frac{dz^* dz}{\pi} |z\rangle \langle z|.$$

The trace of an operator acting on the bosonic coherent states can be written as

$$\text{Tr}\{\mathbf{A}\} = \sum_n \langle n | \mathbf{A} | n \rangle = \int \frac{dz^* dz}{\pi} |z\rangle \langle z| e^{-|z|^2} \langle z | \mathbf{A} | z \rangle.$$

- The action of \mathbf{b}^\dagger on a coherent state $|z\rangle$ reads

$$(9.2.0.3) \quad \begin{aligned} \mathbf{b}^\dagger |z\rangle &= \mathbf{b}^\dagger \sum_{n=0}^{\infty} \frac{z^n}{n!} (\mathbf{b}^\dagger)^n |0\rangle = \sum_{n=0}^{\infty} \frac{z^n}{n!} (\mathbf{b}^\dagger)^{n+1} |0\rangle \\ &= \sum_{n=0}^{\infty} \frac{z^n}{(n+1)!} (\mathbf{b}^\dagger)^{(n+1)} |0\rangle = \sum_{n=1}^{\infty} \frac{z^{n-1}}{(n-1)!} (\mathbf{b}^\dagger)^{(n)} |0\rangle \\ &= \frac{\partial}{\partial z} |z\rangle. \end{aligned}$$

- Consider normal-ordered operator's matrix elements. Note that

$$\langle z_1 | \mathbf{b}^{\dagger n} \mathbf{b}^m | z_2 \rangle = z_1^{*n} z_2^m \langle z_1 | z_2 \rangle = z_1^{*n} z_2^m e^{z_1^* z_2},$$

which allows the matrix element of an operator to be written as

$$\mathbf{A} = \sum_{mn} a_{mn} \mathbf{b}^{\dagger n} \mathbf{b}^m \text{ such that } \langle z_1 | \sum_{mn} a_{mn} \mathbf{b}^{\dagger n} \mathbf{b}^m | z_2 \rangle = e^{z_1^* z_2} \sum_{mn} a_{mn} z_1^{*n} z_2^m.$$

- Any Fock state may be expanded in terms of coherent states, as

$$|\psi\rangle = \int \frac{dz^* dz}{\pi} |z\rangle \langle z | \psi \rangle = \int \frac{dz^* dz}{\pi} \phi(z^*) |z\rangle.$$

- Action of \mathbf{b}^\dagger and \mathbf{b} in coherent state representation reads

$$\begin{aligned} \langle z | \mathbf{b}^\dagger | z \rangle &= z^* \psi(z^*) \Rightarrow \mathbf{b}^\dagger \rightarrow z^* \\ \langle z | \mathbf{b} | z \rangle &= \frac{\partial}{\partial z^*} \psi(z^*) \Rightarrow \mathbf{b} \rightarrow \frac{\partial}{\partial z^*}. \end{aligned}$$

Furthermore, the coherent state variables satisfy the bosonic commutators

$$[z_i^*, z_j^*] = \left[\frac{\partial}{\partial z_i^*}, \frac{\partial}{\partial z_j^*} \right] = 0 \text{ and } \left[\frac{\partial}{\partial z_i^*}, z_j^* \right] = \delta_{ij}.$$

- Consider the coherent-state representation of the Schrödinger equation, which reads

$$\mathbf{H}(\mathbf{b}^\dagger, \mathbf{b}) |\psi\rangle = E |\psi\rangle.$$

Then, acting from the left by $\langle z |$ will yield

$$\mathbf{H}\left(z^*, \frac{\partial}{\partial z^*}\right) \phi(z^*) = E \phi(z^*),$$

which for a typical many-body Hamiltonian reads

$$\sum_{ij} \varepsilon_{ij} z_i^* \frac{\partial}{\partial z_j^*} + \frac{1}{2} \sum_{ijkl} \langle ij | V | kl \rangle z_i^* z_j^* \frac{\partial}{\partial z_i^*} \frac{\partial}{\partial z_j^*}.$$

9.2.0.2. Fermionic coherent states. Fermionic coherent states may be defined in a similar way to bosonic ones. However, when considering fermionic systems, the Pauli exclusion principle must be taken into account in this formulation. The eigenstates of a fermionic annihilation should satisfy the equation

$$\mathbf{f} |\zeta\rangle = \zeta |\zeta\rangle.$$

Since, for fermions there are only two possible state occupations, the most general state may be written as $|\zeta\rangle = \alpha |0\rangle + \beta |1\rangle$, which yields

$$\mathbf{f} |\zeta\rangle = \beta |1\rangle = \zeta (\alpha |0\rangle + \beta |1\rangle) \Rightarrow |\zeta\rangle = |0\rangle \text{ with eigenvalue } 0.$$

Therefore, to obtain nontrivial fermionic coherent states, the anticommutation of annihilation operators must be considered. For two operators $\mathbf{f}_i, \mathbf{f}_j$ with coherent eigenstates

$$\mathbf{f}_i |\zeta\rangle = \zeta_i |\zeta\rangle \text{ and } \mathbf{f}_j |\zeta\rangle = \zeta_j |\zeta\rangle \Rightarrow \text{if } \{\mathbf{f}_i, \mathbf{f}_j\} = 0 \Leftrightarrow \{\zeta_i, \zeta_j\} = 0.$$

This equation implies that $\zeta_i^2 = 0$ as well. This implies that the ζ_i -numbers are no ordinary numbers but are the elements of a Grassmann algebra, and such objects are identified as Grassmann numbers.

Grassmann Algebras

Formally, let \mathbf{V} be an n -dimensional complex vector space with basis $\{\theta_i\}_{i=1}^n$. The Grassmann algebra, whose Grassman variables are $\{\theta_i\}_{i=1}^n$, is defined as the exterior algebra of \mathbf{V} , namely

$$\begin{aligned}
\Lambda(\mathbf{V}) &= \mathbb{C} \oplus \mathbf{V} \oplus (\mathbf{V} \wedge \mathbf{V}) \oplus (\mathbf{V} \wedge \mathbf{V} \wedge \mathbf{V}) \oplus \cdots \oplus (\wedge^n \mathbf{V}) \\
(9.2.0.4) \quad &= \mathbb{C} \oplus \bigoplus_{k=1}^n \wedge^k \mathbf{V},
\end{aligned}$$

where \wedge is the exterior product. The individual elements of this algebra are the Grassmann numbers, which in their most general form may be written as

$$\begin{aligned}
(9.2.0.5) \quad z &= c_0 + \sum_{k=1}^n \sum_{i_1^\uparrow, \dots, i_k} c_{i_1, i_2, \dots, i_k} \theta_{i_1} \theta_{i_2} \cdots \theta_{i_k}, \\
&= c_0 + \sum_i c_i \theta_i + \sum_{i < j} c_i c_j \theta_i \theta_j + \sum_{i < j < k} c_i c_j c_k \theta_i \theta_j \theta_k + \cdots
\end{aligned}$$

where $i^\uparrow = (i_1, i_2, \dots, i_k)$ are strictly increasing k -tuples, with $1 \leq i \leq n$, $1 \leq j \leq k$, and where c_{i_1, i_2, \dots, i_k} are complex, completely antisymmetric tensors of rank k . Again, the θ_i , the $\theta_i \wedge \theta_j = \theta_i \theta_j$ (subject to $i < j$) and larger finite products, can be seen playing the role of a basis vector of Λ -subspaces. The Grassmann algebra generated by n linearly independent Grassmann variables has dimension 2^n , since the $(n+1)$ -fold product of variables must vanish. The dimension of $\wedge^k \mathbf{V}$ is given by the n choose k binomial coefficient. In the special case of a infinite dimensional \mathbf{V} -space, the previous series does not terminate, thus yielding a $\wedge^\infty \mathbf{V}$.

For Grassmann algebras with an even number of generators, $n = 2p$, a conjugation can be defined by selecting p generators θ_i and associating θ^* to each of them, such that

$$(\theta_i)^* = \theta_i^* \quad (\theta_i^*)^* = \theta_i.$$

In a fermionic field theory, the Grassmann algebra is defined by the set of its generating elements $\{1, \zeta_1, \zeta_1^*, \dots, \zeta_n, \zeta_n^*\}$, which satisfy

$$\begin{aligned}
(9.2.0.6) \quad \{\zeta_i, \zeta_j\} &= \{\zeta_j^*, \zeta_i^*\} = \{\zeta_i, \zeta_j^*\} = \zeta_i^2 = \zeta_i^{*2} = 0, \\
\{\zeta_i, \mathbf{f}_i\} &= \{\zeta_i^*, \mathbf{f}_i\} = \{\zeta_i, \mathbf{f}_i^\dagger\} = \{\zeta_i^*, \mathbf{f}_i^\dagger\} = 0.
\end{aligned}$$

Then, the basis set of the Grassmann algebra consists of all distinct products of its generators, namely $\{1, \zeta_1, \dots, \zeta_1 \zeta_2, \dots, \zeta_n\}$.

For a two-dimensional Grassmann algebra, with generators $\{\zeta, \zeta^*\}$, all analytic functions of a Grassmann variables reduce to a first degree polynomial $\phi \in \mathbb{C}[x]$, since $\zeta^2 = 0$,

$$\psi(\zeta) = \psi_0 + \psi_1 \zeta \Leftrightarrow \phi^*(\zeta) = \phi_0^* + \phi_1^* \zeta^*.$$

The most general bilinear function, acting on both ζ and ζ^* , is given by

$$F(\zeta, \zeta^*) = f_0 + f_1 \zeta + f_1^* \zeta^* + f_{12} \zeta^* \zeta.$$

Grassmann calculus

- Differentiation of Grassmann variables:

Since analytic functions of Grassmann variables have such a simple structure, differentiation is just as simple and is given by the usual complex derivative. In effect, let $\psi(\zeta) = \psi_0 + \psi_1 \zeta$. Then, the derivatives are defined to be

$$\partial_\zeta \psi(\zeta) = \psi_1, \quad \partial_{\zeta^*} \psi^*(\zeta) = \phi_1^*, \quad \partial_\zeta (\zeta^* \zeta) = -\bar{\zeta},$$

where in the last equality, $\partial_{\zeta^*} (\zeta^* \zeta) = -\partial_{\zeta^*} (\zeta \zeta^*) = -\zeta^*$. Now let $F(\zeta, \zeta^*) = f_0 + f_1 \zeta + f_1^* \zeta^* + f_{12} \zeta^* \zeta$, then

$$\partial_\zeta F(\zeta, \zeta^*) = f_1 - f_{12}\zeta^*, \quad \partial_{\zeta^*} F(\zeta, \zeta^*) = f_1 + f_{12}\zeta, \quad \partial_{\zeta^*} \partial_\zeta F(\zeta, \zeta^*) = -f_{12} = -\partial_\zeta \partial_{\zeta^*} F(\zeta, \zeta^*),$$

which in turn implies that $\{\partial_{\zeta^*}, \partial_\zeta\} = 0$. The chain rule for differentiation then reads

$$\partial_\zeta F(\psi) = \partial_\zeta \psi \partial_\psi F,$$

wherein, contrary to ordinary differentiation, the order of the terms matters.

- **Integration over Grassmann variables: the Berezin integral.**

Integration and differentiation are identical for Grassmann variables, wherein

$$\partial_{\zeta_i} F(\zeta) = \int d\zeta_i F(\zeta).$$

This definition thus ensures that two most fundamental of properties of ordinary integrals over functions vanishing at infinity are satisfied

– The integral of an exact differential form is zero:

$$\int d\zeta_i \partial_{\zeta_i} F(\zeta) = 0.$$

– The integral over ζ_i of $F(\zeta)$ does not depend on ζ_i , thus having a vanishing ∂_{ζ_i} -derivative,

$$\partial_{\zeta_i} \int d\zeta_i F(\zeta) = 0.$$

Both properties follow from the Grassmann derivative and its nilpotence. Another immediate consequence is that

$$\int d\zeta 1 = 0.$$

Thus, for a Grassmann algebra with a single generator,

$$(9.2.0.7) \quad F(\zeta) = c_0 + c_1\zeta \Rightarrow \int d\zeta F(\zeta) = c_1 = \partial_\zeta F(\zeta).$$

Properties of Fermionic coherent states

A fermionic coherent state is then defined as

$$(9.2.0.8) \quad |\zeta\rangle = \mathcal{N} e^{-\sum_\alpha \zeta_\alpha \mathbf{f}_\alpha^\dagger} |0\rangle = \mathcal{N} \prod_\alpha (1 - \zeta_\alpha \mathbf{f}_\alpha^\dagger) |0\rangle$$

$$(9.2.0.9) \quad \langle\zeta| = \mathcal{N} \langle 0| e^{-\sum_\alpha \zeta_\alpha^* \mathbf{f}_\alpha} = \mathcal{N} \prod_\alpha \langle 0| (1 - \zeta_\alpha^* \mathbf{f}_\alpha), \text{ since}$$

$$(9.2.0.10) \quad \begin{aligned} \mathbf{f}_\alpha |\zeta\rangle &= \mathcal{N} \mathbf{f}_\alpha \prod_\beta (1 - \zeta_\beta \mathbf{f}_\beta^\dagger) |0\rangle \\ &= \mathcal{N} \prod_{\beta \neq \alpha} (1 - \zeta_\beta \mathbf{f}_\beta^\dagger) \mathbf{f}_\alpha (1 - \zeta_\alpha \mathbf{f}_\alpha^\dagger) |0\rangle \\ &= \mathcal{N} \prod_{\beta \neq \alpha} (1 - \zeta_\beta \mathbf{f}_\beta^\dagger) \zeta_\alpha (1 - \zeta_\alpha \mathbf{f}_\alpha^\dagger) |0\rangle \\ &= \mathcal{N} \zeta_\alpha \prod_\beta (1 - \zeta_\beta \mathbf{f}_\beta^\dagger) |0\rangle \\ &= \zeta_\alpha |\zeta\rangle, \end{aligned} \quad \begin{aligned} \mathbf{f}_\alpha^\dagger |\zeta\rangle &= \mathcal{N} \mathbf{f}_\alpha^\dagger (1 - \zeta_\alpha \mathbf{f}_\alpha^\dagger) \prod_{\beta \neq \alpha} (1 - \zeta_\beta \mathbf{f}_\beta^\dagger) |0\rangle \\ &= \mathcal{N} \mathbf{f}_\alpha^\dagger \prod_{\beta \neq \alpha} (1 - \zeta_\beta \mathbf{f}_\beta^\dagger) |0\rangle \\ &= -\mathcal{N} \frac{\partial}{\partial \zeta_\alpha} (1 - \zeta_\alpha \mathbf{f}_\alpha^\dagger) \prod_{\beta \neq \alpha} (1 - \zeta_\beta \mathbf{f}_\beta^\dagger) |0\rangle \\ &= -\frac{\partial}{\partial \zeta_\alpha} |\zeta\rangle. \end{aligned}$$

$$(9.2.0.11) \quad \begin{aligned} \mathbf{f}_\alpha |\zeta\rangle &= \zeta_\alpha |\zeta\rangle, & \mathbf{f}_\alpha^\dagger |\zeta\rangle &= -\partial_{\zeta_\alpha} |\zeta\rangle, \\ \Rightarrow \quad \langle \zeta | \mathbf{f}_\alpha^\dagger &= \langle \zeta | \zeta_\alpha^*, & \langle \zeta | \mathbf{f}_\alpha &= +\partial_{\zeta_\alpha^*} \langle \zeta|. \end{aligned}$$

For a Grassmann-variable, η , the associated inner product is defined as

$$(9.2.0.12) \quad \langle \zeta | \eta \rangle = \langle 0 | e^{-\zeta^* \mathbf{f}} e^{\eta \mathbf{f}^\dagger} | 0 \rangle = 1 + \zeta^* \eta = e^{\zeta^* \eta}.$$

In particular, if $\eta = \zeta$, the normalization factor is $\mathcal{N} = e^{-\frac{\zeta^* \zeta}{2}}$. Then, generalizing to N fermions, the coherent states read

$$|\zeta\rangle = \prod_{i=1}^N e^{-\frac{\zeta_i^* \zeta_i}{2}} e^{-\zeta_i \mathbf{f}_i^\dagger} |0\rangle = e^{-\sum_i \frac{\zeta_i^* \zeta_i}{2}} e^{-\sum_i \zeta_i \mathbf{f}_i^\dagger} |0\rangle, \text{ since } [\zeta_i \mathbf{f}_i^\dagger, \zeta_j \mathbf{f}_j^\dagger] = 0.$$

Fermions obey a closure relation as well.

$$(9.2.0.13) \quad \int \prod_{i=1}^N d\zeta_i^* d\zeta_i |\zeta\rangle \langle \zeta| = \mathbb{1}_{\mathcal{F}}.$$

Fermionic closure relation

In effect, said relation holds if and only if

$$\langle \alpha_1, \dots, \alpha_n | \int \prod_{i=1}^N d\zeta_i^* d\zeta_i |\zeta\rangle \langle \zeta | \beta_1, \dots, \beta_m \rangle = \langle \alpha_1, \dots, \alpha_n | \beta_1, \dots, \beta_m \rangle.$$

Said overlap is

$$\langle \alpha_1, \dots, \alpha_n | \beta_1, \dots, \beta_m \rangle = \delta_{nm} (-1)^p \delta_{\alpha_{i_1} \beta_1} \dots \delta_{\alpha_{i_n} \beta_m}, \quad \text{where } p \text{ is the parity of the permutation } P \left(\begin{matrix} \alpha_1, \dots, \alpha_n \\ \alpha_{i_1}, \dots, \alpha_{i_n} \end{matrix} \right).$$

Note that $\langle \alpha_1, \dots, \alpha_n | \zeta \rangle = \langle 0 | \mathbf{f}_{\alpha_n} \dots \mathbf{f}_{\alpha_1} | \zeta \rangle = \zeta_{\alpha_n} \dots \zeta_{\alpha_1}$, thus leading to the following matrix element

$$(9.2.0.14) \quad \begin{aligned} \langle \alpha_1, \dots, \alpha_n | \int \prod_{i=1}^N d\zeta_i^* d\zeta_i |\zeta\rangle \langle \zeta | \beta_1, \dots, \beta_m \rangle &= \int \prod_{i=1}^N d\zeta_i^* d\zeta_i \langle \alpha_1, \dots, \alpha_n | \zeta \rangle \langle \zeta | \beta_1, \dots, \beta_m \rangle \\ &= \int \prod_{i=1}^N d\zeta_i^* d\zeta_i \prod_{j=1}^n \langle \alpha_j | \zeta \rangle \prod_{k=1}^m \langle \zeta | \beta_k \rangle \\ &= \int \prod_{i=1}^N d\zeta_i^* d\zeta_i \prod_{j=1}^n \left(\langle 0 | \mathbf{f}_{\alpha_j} | \zeta \rangle \right) \prod_{k=1}^m \left(\langle \zeta | \mathbf{f}_{\beta_k} | 0 \rangle \right) \\ &= \int \prod_{i=1}^N d\zeta_i^* d\zeta_i \prod_{j=1}^n e^{-\sum_\ell \frac{\zeta_\ell^* \zeta_\ell}{2}} \left(\zeta_{\alpha_j} \langle 0 | \zeta \rangle \right) \prod_{k=1}^m e^{-\sum_{\ell'} \frac{\zeta_{\ell'}^* \zeta_{\ell'}}{2}} \left(\zeta_{\beta_k}^* \langle \zeta | 0 \rangle \right) \\ &= \int \prod_{i=1}^N d\zeta_i^* d\zeta_i e^{-\sum_\ell \zeta_\ell^* \zeta_\ell} \prod_{j=1}^n \zeta_{\alpha_j} \prod_{k=1}^m \zeta_{\beta_k}^* \\ &= \int \prod_{i=1}^N d\zeta_i^* d\zeta_i \sum_\ell \left(1 - \zeta_\ell^* \zeta_\ell \right) \prod_{j=1}^n \zeta_{\alpha_j} \prod_{k=1}^m \zeta_{\beta_k}^* \end{aligned}$$

with the following integrals arising for a particular $|\gamma\rangle$ -state

$$(9.2.0.15) \quad \int d\zeta_\gamma^* d\zeta_\gamma (1 - \zeta_\ell^* \zeta_\ell) \begin{Bmatrix} \zeta_\gamma^* \zeta_\gamma \\ \zeta_\gamma^* \\ \zeta_\gamma \\ 1 \end{Bmatrix} = \begin{Bmatrix} 1 \\ 0 \\ 0 \\ 1 \end{Bmatrix}.$$

Therefore, this integral is non-vanishing only if each $|\gamma\rangle$ -state is occupied in both $\langle\alpha_1, \dots, \alpha_n|$ and $\langle\beta_1, \dots, \beta_m|$ or unoccupied in both states. In particular, if $m = n$, α is a permutation of β .

Overcompleteness allows to define Grassmann coherent state representation for Fock states. A Fock-state $|\psi\rangle$ can be expanded in terms of fermionic coherent states as

$$|\psi\rangle = \int \prod_{i=1}^N d\zeta_i^* d\zeta_i \psi(\zeta)^* |\zeta\rangle, \text{ where } \psi(\zeta)^* = e^{-\sum_\alpha \zeta_\alpha^* \zeta_\alpha} \langle\zeta|\psi\rangle$$

wherein, the following matrix elements readily are found to be

$$(9.2.0.16) \quad \begin{aligned} \langle\zeta|\mathbf{f}_i|\psi\rangle &= \int \prod_{i=1}^N d\zeta_i^* d\zeta_i \langle\zeta|\psi(\zeta)^* \mathbf{f}_i|\zeta\rangle & \langle\zeta|\mathbf{f}_i^\dagger|\psi\rangle &= \int \prod_{i=1}^N d\zeta_i^* d\zeta_i \langle\zeta|\psi(\zeta)^* \mathbf{f}_i^\dagger|\zeta\rangle \\ &= \int \prod_{i=1}^N d\zeta_i^* d\zeta_i \partial_{\zeta_i^*} \psi(\zeta)^* \langle\zeta|\zeta\rangle & &= \int \prod_{i=1}^N d\zeta_i^* d\zeta_i \zeta_i^* \psi(\zeta)^* \langle\zeta|\zeta\rangle \\ &= \partial_{\zeta_i^*} \psi(\zeta)^* & &= \zeta_i^* \psi(\zeta)^* \end{aligned}$$

The matrix elements of a normal-ordered operator $\mathbf{A}(\mathbf{f}_i^\dagger, \mathbf{f}_i)$ is

$$\langle\zeta|\mathbf{A}(\{\mathbf{f}_i^\dagger\}, \{\mathbf{f}_i\})|\zeta\rangle = e^{-\frac{|\zeta|^2 + |\zeta'|^2}{2}} e^{\sum_i \zeta_i^* \zeta'_i} \mathbf{A}(\{\zeta_i^*\}, \{\zeta_i\}).$$

In the case of the fermionic number operator, $\mathbf{N} = \sum_i \mathbf{f}_i^\dagger \mathbf{f}_i$ gives its expectation value, $\frac{\langle\zeta|\mathbf{N}|\zeta\rangle}{\langle\zeta|\zeta\rangle}$.

9.2.1. *Gaussian Integrals.* Quantum mechanical actions, commonly encountered in many-body path integral formalism, frequently involve exponential or complex or Grassmann variables. For quadratic-form actions, these integrals are simple generalizations of Gaussian integrals.

• **Real variables:**

If the integral can be written as

$$(9.2.1.1) \quad \mathcal{Z}(J) = \int_{\Omega \subseteq \mathbb{R}^{3N}} \prod_{i=1}^N dx_i e^{-\frac{1}{2} \mathbf{x}^T \mathbf{A} \mathbf{x} + \mathbf{J}^T \mathbf{x}}, \text{ with } \mathbf{A} = \mathbf{A}^T, \sigma(\mathbf{A}) > 0$$

Thus, the $3N$ -dimensional real symmetric and positive defined matrix \mathbf{A} can be diagonalized via an orthonormal transformation $\mathbf{A} = \mathbf{O}^T \mathbf{D} \mathbf{O}$, with $\mathbf{O} \in \mathbf{O}(3N)$ and $D_{ii} > 0$. Now, using these facts, the exponentiated factor can be rewritten as a bilinear form, without the linear term, by "completing the square". In effect, let $\mathbf{y} = \mathbf{O}(\mathbf{x} - \mathbf{A}^{-1} \mathbf{J})$, or equivalently $\mathbf{x} = \mathbf{O}^{-1} \mathbf{y} + \mathbf{A}^{-1} \mathbf{J}$

$$(9.2.1.2) \quad \begin{aligned} -\frac{1}{2} \mathbf{x}^T \mathbf{A} \mathbf{x} + \mathbf{J}^T \mathbf{x} &= -\frac{1}{2} (\mathbf{O}^{-1} \mathbf{y} + \mathbf{A}^{-1} \mathbf{J})^T \mathbf{A} (\mathbf{O}^{-1} \mathbf{y} + \mathbf{A}^{-1} \mathbf{J}) + \mathbf{J}^T (\mathbf{O}^{-1} \mathbf{y} + \mathbf{A}^{-1} \mathbf{J}) \\ &= -\frac{1}{2} (\mathbf{y}^T \mathbf{O} + \mathbf{J}^T (\mathbf{A}^{-1})^T) \mathbf{A} (\mathbf{O}^{-1} \mathbf{y} + \mathbf{A}^{-1} \mathbf{J}) + \mathbf{J}^T (\mathbf{O}^{-1} \mathbf{y} + \mathbf{A}^{-1} \mathbf{J}) \\ &= -\frac{1}{2} [\mathbf{y}^T \mathbf{O} \mathbf{A} \mathbf{O}^{-1} \mathbf{y} + \mathbf{y}^T \mathbf{O} \mathbf{A} \mathbf{A}^{-1} \mathbf{J} + \mathbf{J}^T (\mathbf{A}^{-1})^T \mathbf{A} \mathbf{O}^{-1} \mathbf{y} + \mathbf{J}^T (\mathbf{A}^{-1})^T \mathbf{A} \mathbf{A}^{-1} \mathbf{J}] \\ &\quad + \mathbf{J}^T \mathbf{O}^{-1} \mathbf{y} + \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J} \\ &= -\frac{1}{2} [\mathbf{y}^T \mathbf{D} \mathbf{y} + \mathbf{y}^T \mathbf{O} \mathbf{J} + \mathbf{J}^T \mathbf{A}^{-1} \mathbf{A} \mathbf{O}^{-1} \mathbf{y} + \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J}] + \mathbf{J}^T \mathbf{O}^{-1} \mathbf{y} + \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J} \\ &= -\frac{1}{2} [\mathbf{y}^T \mathbf{D} \mathbf{y} + (\mathbf{O} \mathbf{J})^T \mathbf{y} + \mathbf{J}^T \mathbf{O}^{-1} \mathbf{y} + \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J}] + \mathbf{J}^T \mathbf{O}^{-1} \mathbf{y} + \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J} \\ &= -\frac{1}{2} [\mathbf{y}^T \mathbf{D} \mathbf{y} + (\mathbf{J}^T \mathbf{O}^T \mathbf{y})^T + \mathbf{J}^T \mathbf{O}^{-1} \mathbf{y} + \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J}] + \mathbf{J}^T \mathbf{O}^{-1} \mathbf{y} + \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J} \\ &= -\frac{1}{2} \mathbf{y}^T \mathbf{D} \mathbf{y} - \frac{1}{2} (\mathbf{J}^T \mathbf{O}^T \mathbf{y})^T + \frac{1}{2} \mathbf{J}^T \mathbf{O}^{-1} \mathbf{y} + \frac{1}{2} \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J} \\ &= -\frac{1}{2} \mathbf{y}^T \mathbf{D} \mathbf{y} + \frac{1}{2} \mathbf{J}^T \mathbf{O}^{-1} \mathbf{y} + \frac{1}{2} \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J} \end{aligned}$$

$$(9.2.1.3) \quad \Rightarrow \mathcal{Z}(J) = e^{\frac{1}{2} \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J}} \prod_{i=1}^N \int_{\mathbb{R}} dy_i e^{-\frac{y_i^2 D_{ii}}{2}} = \frac{1}{2} \mathbf{J}^T \mathbf{A}^{-1} \mathbf{J} \sqrt{\frac{(2\pi)^N}{\det A}}$$

• **Complex variables:**

Consider now

$$(9.2.1.4) \quad \begin{aligned} \mathcal{Z}(\mathbf{J}^*, \mathbf{J}) &= \int_{\Omega \subseteq \mathbb{C}^{6N}} \prod_{i=1}^N \frac{dz_i^* dz_i}{2\pi i} e^{-\sum_{ij} z_i^* \mathbf{A}_{ij} z_j + \sum_i J_i^* z_i + \text{h.c.}}, \text{ with } \mathbf{A} = \mathbf{A}^\dagger, \sigma(\mathbf{A}) > 0 \\ &= \int_{\Omega \subseteq \mathbb{C}^{6N}} \prod_{i=1}^N \frac{dz_i^* dz_i}{2\pi i} e^{-\mathbf{z}^{\dagger} \mathbf{A} \mathbf{z} + (\mathbf{J}^\dagger \mathbf{z} + \text{h.c.})} \end{aligned}$$

Thus, \mathbf{A} may be diagonalized via a unitary transformation, $\mathbf{D} = \mathbf{U}^\dagger \mathbf{A} \mathbf{U}$, which induce the following change of variables

$$(9.2.1.5) \quad \begin{aligned} \mathbf{z}' &= \mathbf{U}(\mathbf{z} - \mathbf{A}^{-1} \mathbf{J}), \\ \mathbf{z}'^\dagger &= \mathbf{U}^\dagger(\mathbf{z}^\dagger - \mathbf{J}^\dagger \mathbf{A}^{-1}) \end{aligned} \Rightarrow \mathcal{Z}(\mathbf{J}^*, \mathbf{J}) = e^{\mathbf{J}^\dagger \mathbf{A}^{-1} \mathbf{J}} \int_{\mathbb{C}} \prod_{i=1}^N \frac{dz_i^* dz_i}{2\pi i} e^{-\mathbf{z}'^\dagger \mathbf{D} \mathbf{z}'} = e^{\mathbf{J}^\dagger \mathbf{A}^{-1} \mathbf{J}} (\det A)^{-1}$$

• **Grassmannian integrals:**

A Gaussian integral involving a conjugate pair of Grassmann variables can be written as

$$\int d\zeta^* d\zeta e^{-\zeta^* a \zeta} = \int d\zeta^* d\zeta (1 - \zeta^* a \zeta) = a.$$

By analogy, Grassmann gaussian integrals may be written as

$$(9.2.1.6) \quad \mathcal{Z}(\zeta^*, \zeta) = \int \prod_i d\zeta_i^* d\zeta_i e^{-\zeta^* \mathbf{A} \zeta + (\varepsilon^* \zeta + \text{h.c.})}.$$

To solve this Grassmann integral, the \mathbf{A} -matrix must be diagonalized by a suitable change of Grassmann-variables, which themselves have not a trivial transformation.

Transformation of Grassmann variables

Consider the integral^a

$$(9.2.1.7) \quad \begin{aligned} \mathcal{I} &= \int d\zeta F(\zeta), \text{ subject to a transformation } \zeta \rightarrow \alpha \zeta' + \beta, \alpha, \beta \in \mathbb{C}. \\ \Rightarrow \mathcal{I} &= \partial_\zeta F(\zeta) = \partial_\zeta(\zeta') \partial_{\zeta'} F(\zeta') = \frac{1}{\alpha} \partial_{\zeta'} F(\zeta') = \frac{1}{\alpha} \int d\zeta F(\alpha \zeta' + \beta). \end{aligned}$$

Therefore, the jacobian of the transformation is $\frac{1}{\alpha}$, instead of α for ordinary (bosonic) transformation. This result can then be generalized to a multidimensional integral

$$(9.2.1.8) \quad \begin{aligned} \mathcal{I} &= \int \prod_i d\zeta_i F(\zeta) = \prod_i \partial_{\zeta_i} F(\zeta) \\ &= \prod_i \left(\sum_j (\partial_{\zeta_i} \zeta'_j) \partial_{\zeta'_j} \right) F(\zeta) \\ &\stackrel{\zeta' = \mathbf{M} \zeta}{=} \prod_i \left(\sum_j \mathbf{M}_{ji} \partial_{\zeta'_j} \right) F(\zeta) \end{aligned}$$

Given that $\zeta_{i_1} \cdots \zeta_{i_n} = \varepsilon_{i_1 \cdots i_n} \zeta'_{i_1} \cdots \zeta'_{i_n}$, then

$$(9.2.1.9) \quad \begin{aligned} &= \prod_i \left(\varepsilon^{i_1 \cdots i_n} \mathbf{M}_{1i_1} \cdots \mathbf{M}_{ni_n} \partial_{\zeta'_{i_1}} \cdots \partial_{\zeta'_{i_n}} \right) F(\zeta) \\ &= \det(\partial_{\zeta_j} \zeta'_i) \prod_i \partial_{\zeta'_j} \zeta'_i F(\zeta) \\ &\Rightarrow \prod_i d\zeta_i = \det(\partial_{\zeta_j} \zeta'_i) \prod_i d\zeta'_i, \end{aligned}$$

the jacobian being the inverse of the determinant $\det(\partial_{\zeta_j} \zeta'_i)$.

^aNote that, if θ is a Grassmann variable, then $\theta = a\eta$ is another Grassmann variable

$$1 = \int d\theta \theta = \int d(a\eta) a\eta \Rightarrow d(a\eta) = \frac{1}{a} d\eta.$$

With this in mind, a wide class of Grassmannian integrals may be solved, for example,

$$\begin{aligned}
(9.2.1.10) \quad \mathcal{Z}(\zeta^*, \zeta) &= \int \prod_i d\zeta_i^* d\zeta_i e^{-\zeta^* \mathbf{A} \zeta} = \det(\mathbf{A}) \int \prod_i d\zeta_i^* d\zeta_i e^{\zeta^* \zeta} \\
&= \det(\mathbf{A}) \int \prod_i d\zeta_i^* d\zeta_i (1 - \zeta^* \zeta) = \det(\mathbf{A}),
\end{aligned}$$

or, for the more general integral given by [equation \(9.2.1.6\)](#)

$$(9.2.1.11) \quad \mathcal{Z}(\zeta^*, \zeta) = \det(\mathbf{A}) e^{-\zeta^* \mathbf{A}^{-1} \zeta}.$$

9.3. Photons: coupling to electrons. Phonons are quantized vibrations and play a fundamental role in sound physics, specific heat, elasticity and electrical resistivity of solids. More surprisingly, the electron-phonon coupling is the cause of conventional superconductivity. There are two distinct models for electron-phonon coupling, the jellium model, where the ions are represented by a smeared-out continuous positive background, and the lattice model where the ions oscillate around their equilibrium positions forming a crystal lattice.

Phonons are basically harmonic oscillators, they are bosons. Moreover, these naturally occur at finite temperature, so the thermal distribution function for bosons will be used.

9.3.0.1. The Jellium Model. Consider the Hamiltonian

(9.3.0.1)

$$\begin{aligned}
\mathbf{H} &= \mathbf{H}_+ + \mathbf{H}_{+e} + \mathbf{H}_e, \text{ where } \mathbf{H}_+ = \int_{\mathbb{R}^6} \frac{d\mathbf{x} d\mathbf{x}'}{2} \rho_+(\mathbf{x}) \rho_+(\mathbf{x}') \frac{e^{-\mu|\mathbf{x}-\mathbf{x}'|}}{|\mathbf{x}-\mathbf{x}'|} = \left(\frac{eN}{\sqrt{2}\Omega} \right)^2 \int dr d\theta \int_{\mathbb{R}^3} \frac{e^{-\mu r}}{r} = \frac{4\pi e^2 N^2}{2\mu^2 \Omega} \\
\mathbf{H}_{+e} &= -e \sum_i \int_{\mathbb{R}^3} d\mathbf{x} \rho_+(\mathbf{x}) \frac{e^{-\mu|\mathbf{x}-\mathbf{x}_i|}}{|\mathbf{x}-\mathbf{x}_i|} = -\frac{e^2 N}{\Omega} \sum_i \int_{\mathbb{R}^3} d\mathbf{x} \frac{e^{-\mu|\mathbf{x}-\mathbf{x}_i|}}{|\mathbf{x}-\mathbf{x}_i|} = -\frac{e^2 N^2}{\Omega} \frac{4\pi}{\mu^2} \\
\mathbf{H}_e &= \sum_i \frac{\mathbf{p}_i^2}{2m_e} + \frac{e^2}{2} \sum_{i \neq j} \frac{e^{-\mu|\mathbf{x}_i-\mathbf{x}_j|}}{|\mathbf{x}_i-\mathbf{x}_j|}, \text{ 2nd Quant. } \mathbf{T} = \sum_i \frac{\hbar^2 k_i^2}{2m_e} \mathbf{c}_i^\dagger \mathbf{c}_i,
\end{aligned}$$

where \mathbf{H}_+ is the positive background Hamiltonian, where \mathbf{H}_{+e} is the electron-positive background Hamiltonian and \mathbf{H}_e is the electron Hamiltonian. Consider in particular the electron Hamiltonian \mathbf{H}_e , whose second-quantization form is so since the single-electron wavefunctions are just plane waves, with the matrix element for the two-electron repulsion given by

(9.3.0.2)

$$\begin{aligned}
\left\langle \mathbf{k}_4 \sigma_i, \mathbf{k}_3 \sigma_j \left| \frac{e^{-\mu|\mathbf{x}_i-\mathbf{x}_j|}}{|\mathbf{x}_i-\mathbf{x}_j|} \right| \mathbf{k}_2 \sigma_j, \mathbf{k}_1 \sigma_i \right\rangle &= \frac{1}{\Omega^2} \int_{\mathbb{R}^6} d\mathbf{x}_i d\mathbf{x}_j e^{i(\mathbf{k}_1-\mathbf{k}_4) \cdot \mathbf{x}_i} e^{i(\mathbf{k}_2-\mathbf{k}_3) \cdot \mathbf{x}_j} \frac{e^{-\mu|\mathbf{x}_i-\mathbf{x}_j|}}{|\mathbf{x}_i-\mathbf{x}_j|} \\
&= \frac{1}{\Omega} \delta_{\mathbf{k}_1+\mathbf{k}_2-\mathbf{k}_3-\mathbf{k}_4} \int_{\mathbb{R}^3} d(\mathbf{x}_i - \mathbf{x}_j) e^{-i\mathbf{q} \cdot (\mathbf{x}_i - \mathbf{x}_j)} \frac{e^{-\mu r}}{r} \\
&= \frac{1}{\Omega} \delta_{\mathbf{k}_1+\mathbf{k}_2-\mathbf{k}_3-\mathbf{k}_4} \frac{4\pi}{\mu^2 + q^2} \\
\Rightarrow \mathbf{H}_e &= \sum_{\substack{\mathbf{k} \in \Lambda \\ \sigma \in \mathbb{C}^2}} \frac{\hbar^2 k^2}{2m_e} \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} + \sum_{\substack{\mathbf{k}, \mathbf{k}', \mathbf{q} \in \Lambda \\ \sigma, \sigma' \in \mathbb{C}^2}} \frac{2\pi e^2}{\Omega(\mu^2 + q^2)} \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}^\dagger \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'} \mathbf{c}_{\mathbf{k}\sigma}.
\end{aligned}$$

The ($\mathbf{q} = 0$)-term is special,

$$\begin{aligned}
 (9.3.0.3) \quad \sum_{\substack{\mathbf{k}, \mathbf{k}' \in \Lambda \\ \sigma, \sigma' \in \mathbb{C}^2}} \frac{2\pi e^2}{\Omega \mu^2} \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'} \mathbf{c}_{\mathbf{k}\sigma} &= \sum_{\substack{\mathbf{k}, \mathbf{k}' \in \Lambda \\ \sigma, \sigma' \in \mathbb{C}^2}} \frac{2\pi e^2}{\Omega \mu^2} \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} (\mathbf{c}_{\mathbf{k}'\sigma'}^\dagger \mathbf{c}_{\mathbf{k}\sigma} - \delta_{\mathbf{k}\mathbf{k}'} \delta_{\sigma\sigma'}) \\
 &= \frac{2\pi e^2}{\mu^2 \Omega} (N^2 - N).
 \end{aligned}$$

The $\frac{N^2 e^2}{\mu}$ -terms all add up to zero while the second term $-\frac{2\pi N e^2}{\mu^2 \Omega} \xrightarrow{\Omega \rightarrow \infty} 0$, thus reducing the system's Hamiltonian to

$$(9.3.0.4) \quad \mathbf{H} = \sum_{\substack{\mathbf{k} \in \Lambda \\ \mathbf{q} \neq 0 \\ \sigma \in \mathbb{C}^2}} \frac{\hbar^2 k^2}{2m_e} \mathbf{c}_{\mathbf{k}\sigma}^\dagger \mathbf{c}_{\mathbf{k}\sigma} + \sum_{\substack{\mathbf{k}, \mathbf{k}', \mathbf{q} \in \Lambda \\ \sigma, \sigma' \in \mathbb{C}^2}} \frac{2\pi e^2}{\Omega q^2} \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}^\dagger \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'} \mathbf{c}_{\mathbf{k}\sigma},$$

where $\mu = 0$ since the ($q = 0$)-singularity has been removed. The first term is the Hamiltonian for the Sommerfeld gas and where the second term is a perturbation. Notice that the ground state has been expressed as a product of single particle wavefunctions and not as an antisymmetrized linear combination of such terms, since the antisymmetrization has already been taken care of by the creation and annihilation operator commutators. Then, the jellium ground-state is given by

$$|\Omega\rangle = \prod_{\mathbf{k}=0}^{k=k_F} |\mathbf{k} \uparrow\rangle \langle \mathbf{k} \downarrow|.$$

9.3.0.2. The Wigner crystal. In the small-density limit, the Coulomb energy becomes dominant and therefore it is more appropriate to start from the ground state of the interaction term than from that of the kinetic energy. Wigner argued that the free electron-like Fermi sphere ground state could become unstable to an insulating lattice of localized electrons when the electron density becomes sufficiently small. This leads to a purely classical problem. Namely, to calculate the lowest-energy configuration of charged particles immersed into a homogeneous background of opposite charge.

Historical Context

This problem was first addressed by Wigner in 1934 as that of an inverted alkali metal, and he argued that at low enough densities the electrons would form a crystal. A consistent theory has of course to take into account the kinetic energy, leading to zero-point fluctuations of the electrons around their equilibrium positions. As the lattice constant decreases, these fluctuations become more and more important, until the Wigner crystal melts. Numerical simulations indicate this happens for $r_s \sim 100$. In fact, a two-dimensional Wigner crystal with a triangular structure has actually been observed for a very low-density electron system ($r_s \sim 10^4$) dispersed over the surface of liquid helium.

The first-order energy correction is given by

$$E^{(1)} = \langle \Omega | \mathbf{H}' | \Omega \rangle = \left\langle \Omega \left| \sum_{\substack{\mathbf{k}, \mathbf{k}', \mathbf{q} \in \Lambda \\ \sigma, \sigma' \in \mathbb{C}^2}} \frac{2\pi e^2}{\Omega q^2} \mathbf{c}_{\mathbf{k}+\mathbf{q}\sigma}^\dagger \mathbf{c}_{\mathbf{k}'-\mathbf{q}\sigma'}^\dagger \mathbf{c}_{\mathbf{k}'\sigma'} \mathbf{c}_{\mathbf{k}\sigma} \right| \Omega \right\rangle.$$

For this matrix-elements to be non-zero, the states $\mathbf{k} + \mathbf{q}$, $\mathbf{k}' - \mathbf{q}$, \mathbf{k}' and \mathbf{k} must be contained in the ground-state wavefunction product. This means that the only terms that do not vanish will be those that do not change $|\Omega\rangle$'s occupation numbers, therefore

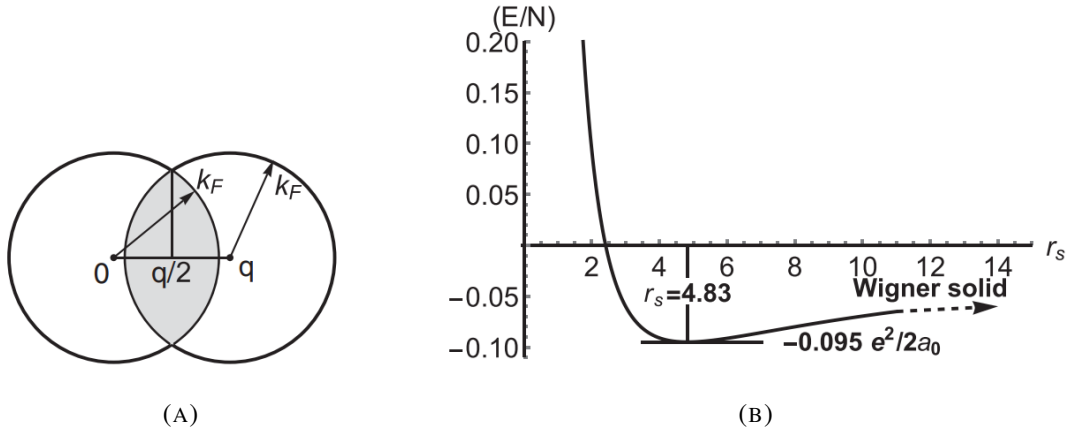


FIGURE 7. Fig. A: Geometry for the \mathbf{k} -integration is shown for arbitrary, but fixed, \mathbf{q} , which is just the intersection of two spheres.

Fig. B: Ground-state energy per electron, including first-order correction, for the jellium model.

$$(9.3.0.5) \quad \begin{cases} \mathbf{k} = \mathbf{k} + \mathbf{q}, & \mathbf{k}' = \mathbf{k}' - \mathbf{q} \Rightarrow \mathbf{q} = 0 \\ \mathbf{k}' = \mathbf{k} + \mathbf{q}, & \mathbf{k} = \mathbf{k}' - \mathbf{q} \end{cases}$$

This can be interpreted via [figure 8](#). The first configuration implies no scattering since $\mathbf{q} = 0$. The remaining term, the exchange scattering term, is represented by the second diagram in [figure 8](#) and depicts the case where the particle in $|\mathbf{k}\rangle$ -state is scattered into the $|\mathbf{k}'\rangle$ -state and vice-versa. The correction too E_0 is the exchange energy, given by

$$(9.3.0.6) \quad \begin{aligned} E^{(1)} &= \langle \Omega | \sum_{\substack{\mathbf{k}, \mathbf{k}', \mathbf{q} \in \Lambda \\ \sigma \in \mathbb{C}^2}} \frac{2\pi e^2}{\Omega q^2} \mathbf{c}_{\mathbf{k}+\mathbf{q} \sigma}^\dagger \mathbf{c}_{\mathbf{k}'-\mathbf{q} \sigma}^\dagger \mathbf{c}_{\mathbf{k}' \sigma} \mathbf{c}_{\mathbf{k} \sigma} | \Omega \rangle \\ &= \langle \Omega | \sum_{\substack{\mathbf{k}, \mathbf{q} \in \Lambda \\ \sigma \in \mathbb{C}^2}} \frac{2\pi e^2}{\Omega q^2} \mathbf{c}_{\mathbf{k}+\mathbf{q} \sigma}^\dagger \mathbf{c}_{\mathbf{k} \sigma}^\dagger \mathbf{c}_{\mathbf{k}+\mathbf{q} \sigma} \mathbf{c}_{\mathbf{k} \sigma} | \Omega \rangle \\ &= \langle \Omega | \sum_{\substack{\mathbf{k}, \mathbf{q} \in \Lambda \\ \sigma \in \mathbb{C}^2}} \frac{2\pi e^2}{\Omega q^2} (-N_{\mathbf{k}+\mathbf{q} \sigma} - N_{\mathbf{k} \sigma}) | \Omega \rangle \Rightarrow \varepsilon^{(1)}(\mathbf{k}) = -\frac{2\pi e^2}{\Omega} \sum_{\mathbf{q}} \frac{N_{\mathbf{k}+\mathbf{q}}}{q^2} \end{aligned}$$

(9.3.0.7)

$$\begin{aligned} \Rightarrow E^{(1)} &= -\frac{4\pi e^2}{\Omega} \frac{\Omega}{(2\pi)^6} \int_{\mathbb{R}^6} d\mathbf{k} d\mathbf{q} \frac{1}{q^2} \Theta(k_F - |\mathbf{k} + \mathbf{q}|) \Theta(k_F - k) \text{ with the integration region shown in [figure 7](#) (left)} \\ &= -\frac{4\pi e^2 \Omega}{(2\pi)^6} \int_{\mathbb{R}^3} \frac{d\mathbf{q}}{q^2} \int_{\mathbb{R}^3} d\mathbf{p} \Theta\left(k_F - \left|\mathbf{p} + \frac{1}{2}\mathbf{q}\right|\right) \Theta\left(k_F - \frac{1}{2}q\right) \\ &= -\frac{e^2}{2a_0} N \frac{3}{2\pi} \left(\frac{9\pi}{4}\right)^{\frac{1}{3}} \frac{1}{r_s}. \end{aligned}$$

This results, shown in [figure 7](#) (right), shows that the electron gas is stable when the repulsive Coulomb interaction is turned on. No external confinement potential is therefore needed to hold the electron gas in the ion jellium together. There exists an optimal density n^* , or interparticle distance r_s^* , which minimizes the energy and furthermore yields an $E^* < 0$. The negative exchange energy overcome the positive kinetic energy. This method of treating the electron gas is a special case of the Hartree-Fock approximation and is the simplest way to take into account interactions. The interesting point to note is that in this approximation the energy is reduced below that of the Sommerfeld gas, which is a manifestation of the exchange hole.

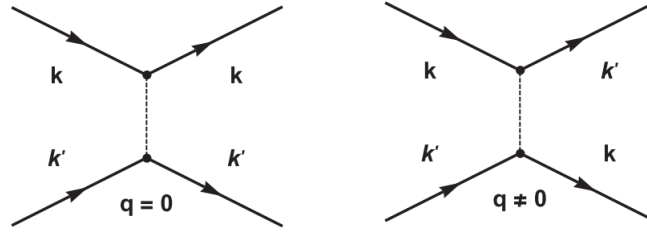


FIGURE 8. Diagrams for the jellium sea. The dashed line represents the interaction.

9.3.0.3. Hartree-Fock approximation as a Mean-Field Theory.

Consider the electronic Hamiltonian given by [equation \(9.3.0.2\)](#). The goal is to write an effective two-body Hamiltonian, which takes into account the average effects of the interactions. Therefore, the four-Fermi interaction is replaced with a sum of all possible two-body terms

$$(9.3.0.8) \quad \langle c_1^\dagger c_2^\dagger c_3 c_4 \rangle_{\text{MFT}} = -\langle c_1^\dagger c_3 \rangle c_2^\dagger c_4 - \langle c_2^\dagger c_4 \rangle c_1^\dagger c_3 + \langle c_1^\dagger c_4 \rangle c_2^\dagger c_3 + \langle c_2^\dagger c_3 \rangle c_1^\dagger c_4,$$

where $\langle c_{k\sigma}^\dagger c_{k'\sigma'} \rangle = n_{k\sigma} \delta_{kk'}^{\sigma\sigma'}$

The mean field terms can be understood as the average number of particles $n_{k\sigma}$ in the $|k\sigma\rangle$ -state, which will be weighted with the two-body interaction $V(\mathbf{q})$ to give the average interaction due to all other interaction terms. Then, the four-Fermi interaction in the mean-field approximation reads

$$(9.3.0.9) \quad \begin{aligned} \mathcal{V}_{\text{HF}} &= \frac{1}{2} \sum_{\substack{\mathbf{k}\mathbf{k}', \mathbf{q} \in \Lambda \\ \sigma\sigma' \in \mathbb{C}^2}} V(\mathbf{q}) \left[-\langle c_{k\sigma}^\dagger c_{k'\sigma'} \rangle c_{k'-\mathbf{q}\sigma'}^\dagger c_{k-\mathbf{q}\sigma} - \langle c_{k'-\mathbf{q}\sigma'}^\dagger c_{k-\mathbf{q}\sigma} \rangle c_{k\sigma}^\dagger c_{k'\sigma'} \right. \\ &\quad \left. + \langle c_{k\sigma}^\dagger c_{k-\mathbf{q}\sigma} \rangle c_{k'-\mathbf{q}\sigma'}^\dagger c_{k'\sigma'} + \langle c_{k'-\mathbf{q}\sigma'}^\dagger c_{k'\sigma'} \rangle c_{k\sigma}^\dagger c_{k-\mathbf{q}\sigma} \right] \\ &= - \sum_{\substack{\mathbf{k}\mathbf{q} \in \Lambda \\ \sigma \in \mathbb{C}^2}} V(\mathbf{q}) \langle c_{k\sigma}^\dagger c_{k\sigma} \rangle c_{k-\mathbf{q}\sigma}^\dagger c_{k-\mathbf{q}\sigma} + V(0) \sum_{\substack{\mathbf{k}\mathbf{k}' \in \Lambda \\ \sigma\sigma' \in \mathbb{C}^2}} \langle c_{k\sigma}^\dagger c_{k\sigma} \rangle c_{k'\sigma'}^\dagger c_{k'\sigma'} \\ &= \sum_{\substack{\mathbf{k} \in \Lambda \\ \sigma \in \mathbb{C}^2}} \left(- \sum_{\mathbf{q} \in \Lambda} n_{\mathbf{k}+\mathbf{q}\sigma} V(\mathbf{q}) + n V(0) \right) c_{k\sigma}^\dagger c_{k\sigma}, \\ \varepsilon_{\mathbf{k}}^{\text{HF}} &= \varepsilon_{\mathbf{k}} + \sum_{\substack{\mathbf{k}' \in \Lambda \\ \sigma' \in \mathbb{C}^2}} [V(0) - \delta_{\sigma\sigma'} V(\mathbf{k} - \mathbf{k}')] n_{\mathbf{k}'\sigma'} \\ \Rightarrow \mathbf{H}^{\text{HF}} &= \sum_{\substack{\mathbf{k} \in \Lambda \\ \sigma \in \mathbb{C}^2}} \varepsilon_{\mathbf{k}}^{\text{HF}} c_{k\sigma}^\dagger c_{k\sigma}, \text{ where } \varepsilon_{\mathbf{k}}^{\text{HF}} = \varepsilon_{\mathbf{k}} + V(0)N - \sum_{\substack{\mathbf{k}' \in \Lambda \\ \sigma' \in \mathbb{C}^2}} V(\mathbf{k} - \mathbf{k}') n_{\mathbf{k}'\sigma'}. \end{aligned}$$

Note that the Hartree or direct Coulomb term, which represents the average interaction energy of the electron $k\sigma$ with all the other electrons in the system, is a constant; when summed over all electron it cancels exactly with the constant arising from the sum of the self-energy of the positive background and the interaction energy of the electron gas with that of the background. The Fock, or exchange term, is a momentum-dependent shift.

Problems with the Hartree-Fock Theory

Although we argued that the Hartree-Fock approximation becomes a better approximation in the high-density limit for electron interacting via the Coulomb attraction, it never becomes exact. The Hartree-Fock energy correction is given by

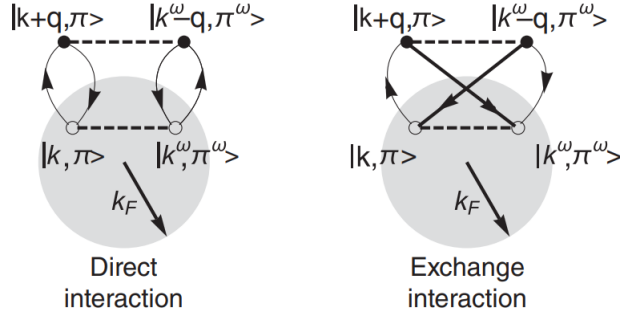


FIGURE 9. The only two possible processes in second-order perturbation theory for two states, $|\mathbf{k}_1, \sigma_1\rangle$ and $|\mathbf{k}_2, \sigma_2\rangle$, in the Fermi sea. The direct process gives a divergent contribution to $\frac{E}{N}$ while the exchange process gives a finite contribution.

$$\varepsilon^{(1)}(\mathbf{k}) = \frac{e^2 k_F}{4\pi^2} \left[\frac{k_F^2 - k^2}{k k_F} \log \left| \frac{k_F + k}{k_F - k} \right| + 2 \right]$$

The first term in parentheses has a logarithmic divergence in slope $k = k_F$. This means that while the energy shift might be small compared to the Fermi energy, the Fermi velocity $v_F = \nabla_{\mathbf{k}} \varepsilon(\mathbf{k})$, contains a non-finite term. This problem can be traced back to the long-range nature of the Coulomb force. Two electrons at large distances $\mathbf{x} - \mathbf{x}'$ do not feel the full $\frac{1}{|\mathbf{x} - \mathbf{x}'|}$ -interaction, but a screened version due to the presence of the intervening medium, namely the electron gas, rearranges itself to cancel out the long-range part of V .

Electron Interaction in Second-order Perturbation Theory

Trying to improve upon the first-order result by going to second-order perturbation theory yields a disastrous results for the matrix elements diverge. According to second-order perturbation theory,

$$(9.3.0.10) \quad \frac{E^{(2)}}{N} = \frac{1}{N} \sum_{\substack{|\psi\rangle \in \mathbb{H} \\ |\psi\rangle \neq |\Omega\rangle}} \frac{\langle \Omega | \mathbf{V}_c | \psi \rangle \langle \psi | \mathbf{V}_c | \Omega \rangle}{E^{(0)} - E_\psi}.$$

This, combined with the momentum-conserving Coulomb interaction, yields intermediate states where two particles are injected out of the Fermi sphere. From such an intermediate state, $|\Omega\rangle$, is restored by putting the excited electrons back into the holes they left behind. Therefore, there are only two types of possible processes: direct interaction and the exchange process, shown in [figure 9](#).

- 1) In the Direct interaction, it turns out that the direct interaction gives a divergent contribution $E_{\text{dir}}^{(2)}$ to $E^{(2)}$ due to the singular behaviour of the Coulomb interaction at small \mathbf{q} -momentum transfers. The $(|\psi\rangle = |\Omega\rangle)$ -constraint leads to the second p

$$|\psi\rangle = \Theta(|\mathbf{k}_1 + \mathbf{q}| - k_F) \Theta(|\mathbf{k}_2 - \mathbf{q}| - k_F) \Theta(k_F - |\mathbf{k}_2|) \Theta(k_F - |\mathbf{k}_1|) \mathbf{c}_{\mathbf{k}_1 + \mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}_2 - \mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}_2} \mathbf{c}_{\mathbf{k}_1} |\Omega\rangle.$$

To restore $|\Omega\rangle$, the same \mathbf{q} -momentum transfer must be involved in both $\langle \Omega | \mathbf{V}_c | - \rangle$ and $\langle \psi | \mathbf{V}_c | \Omega \rangle$, and writing $\mathbf{V}(\mathbf{q}) = \frac{4\pi e^2}{\mathbf{q}^2}$ yields

$$E_{\text{dir}}^{(2)} = \frac{1}{4\Omega^2} \sum_{\mathbf{q} \in \Lambda} \sum_{\substack{\mathbf{k}\mathbf{k}' \in \Lambda \\ \sigma\sigma' \in \mathbb{C}^2}} \frac{\mathbf{V}(\mathbf{q})^2}{E^{(0)} - E_{\nu}} \Theta(|\mathbf{k}_1 + \mathbf{q}| - k_F) \Theta(|\mathbf{k}_2 - \mathbf{q}| - k_F) \Theta(k_F - |\mathbf{k}_2|) \Theta(k_F - |\mathbf{k}_1|)$$

(9.3.0.11) Given that $\mathbf{V}(\mathbf{q})^2 \approx q^{-4} \Rightarrow E^{(0)} - E_{\psi} \stackrel{q \rightarrow 0}{\approx} \mathbf{k}_1^2 + \mathbf{k}_2^2 - (\mathbf{k}_1 + \mathbf{q})^2 - (\mathbf{k}_2 - \mathbf{q})^2 \approx q$

$$\sum_{\mathbf{q} \in \Lambda} \sum_{\substack{\mathbf{k}\mathbf{k}' \in \Lambda \\ \sigma\sigma' \in \mathbb{C}^2}} \cdots \Theta(|\mathbf{k}_2 - \mathbf{q}| - k_F) \Theta(k_F - |\mathbf{k}_2|) \Theta(k_F - |\mathbf{k}_1|) \stackrel{q \rightarrow 0}{\approx} q$$

$$\Rightarrow E_{\text{dir}}^{(2)} = \int_0 dq q^2 \frac{1}{q^4} \frac{1}{q} q q = \int_0 \frac{dq}{q} = \log q|_{0 \rightarrow \infty}.$$

- 2) The exchange process does not lead to a divergence since, in this case, the momentum transfer in the excitation part is \mathbf{q} , but in the relaxation part it is $\mathbf{k}_2 - \mathbf{k}_1 - \mathbf{q}$. Thus $\mathbf{V}^2(\mathbf{q})$ is replaced by $\mathbf{V}^2(\mathbf{q})\mathbf{V}(\mathbf{k}_2 - \mathbf{k}_1 - \mathbf{q}) \propto q^{-2}$ for $\mathbf{q} \rightarrow 0$, which is less singular than $\mathbf{V}^2(\mathbf{q}) \propto q^{-4}$.

Physically, the energy of an electron gas must be finite. The only hope for rescue lies in regularization of the divergent behaviour by taking higher-order perturbation terms into account. Consider now the ground-state energy of the jellium model of a three-dimensional electron gas. In the ($r_s < 1$)-regime, this energy can be expressed in terms of a power-series in r_s , as

$$E_0 = \frac{K}{r_s^2} [1 + br_s + cr_s^2 + \cdots],$$

where K, a, b, c are constants. This is not exactly right, since first-order perturbation theory gives a term of the (br_s)-form in the series. But, considering second-order perturbation theory, one finds a contribution that diverges like $\int_0 \frac{dq}{q}$, where q is the transfer-momentum in the Fourier transform \mathbf{V}_q of the Coulomb interaction, $\mathbf{V}_q \approx q^{-2}$. Thus, there is in fact a logarithmic divergence from the lower limit 0 of the momentum transfers. This divergence is associated with the long range of the Coulomb interaction. Furthermore, if one examines higher-order terms in the perturbation series, one finds that they diverge even more strongly.

On physical grounds, however, the energy of the interacting electron gas must be finite and well-defined, with no phase transitions occurring as the repulsive interactions are turned on. This failure of standard perturbation theory appears to signal that the energy does not have a standard power-series expansion in r_s . In 1957, Gell-Mann and Bruckner resolved this issue by developing many-body perturbation theory³⁶, in which the summation of the divergent terms in the series (an infinite number of terms) is performed before doing the momentum integrals, thus arriving to a finite and well-defined results. This is called the resummation of the perturbation series. They found that there was a term in the series for E_0 that is $\propto \log r_s$, and is non-analytic at $r_s = 0$. This holds so long the interaction doesn't cause any drastic changes to the system, such as a phase transition.

9.4. Random Phase Approximation. In the second quantization formalism, the fermionic field operators may be written in plane-wave basis as

$$\Psi(\mathbf{x}, t) = \frac{1}{\sqrt{\Omega}} \sum_{\mathbf{k} \in \Lambda} e^{i\mathbf{k} \cdot \mathbf{x}} \mathbf{c}_{\mathbf{k}}, \quad \Psi^\dagger(\mathbf{x}, t) = \frac{1}{\sqrt{\Omega}} \sum_{\mathbf{k} \in \Lambda} e^{-i\mathbf{k} \cdot \mathbf{x}} \mathbf{c}_{\mathbf{k}}^\dagger.$$

The particle density $\rho(\mathbf{x})$ at point \mathbf{x} is given by

$$\rho(\mathbf{x}) = \sum_{i \in \mathbb{N}} \delta(\mathbf{x} - \mathbf{x}_i) \Rightarrow \rho(\mathbf{x}) = \int_{\mathbb{R}^3} d\mathbf{x} \Psi^\dagger(\mathbf{x}) \left(\sum_{i \in \mathbb{N}} \delta(\mathbf{x} - \mathbf{x}_i) \right) \Psi(\mathbf{x}) = \sum_{i \in \mathbb{N}} \Psi^\dagger(\mathbf{x}) \Psi(\mathbf{x}).$$

³⁶Many-body perturbation expansions without diagrams. I. Normal states. Behnam Farid §

Its Fourier transform is given by $\rho_{\mathbf{q}} = \sum_{\mathbf{k} \in \Lambda} \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}}$. Its physical meaning is straightforward: it gives rise to particle-hole pairs charge fluctuations around the Fermi surface. Since $\rho^\dagger = \rho$, it follows that $\rho_{\mathbf{q}}^\dagger = \rho_{-\mathbf{q}}$.³⁷

Now, the following question naturally arises. Under what conditions an interacting system, with Hamiltonian

$$\mathbf{H} = \sum_{\mathbf{k} \in \Lambda} \varepsilon_{\mathbf{k}} \mathbf{c}_{\mathbf{k}}^\dagger \mathbf{c}_{\mathbf{k}} + \frac{1}{2} \sum_{\mathbf{q} \in \Lambda} \mathbf{V}_{\mathbf{q}} (\rho_{\mathbf{q}} \rho_{-\mathbf{q}} - N),$$

and ground-state $|\Omega\rangle$, $\mathbf{H}|\Omega\rangle = \varepsilon_0 |\Omega\rangle$, can support density excitations associated with $\rho_{\mathbf{q}}$? Namely, these excitations are

$$\mathbf{H} \rho_{\mathbf{q}}^\dagger |\Omega\rangle = (\varepsilon_0 + \hbar \omega_{\mathbf{q}}) \rho_{\mathbf{q}}^\dagger |\Omega\rangle$$

The previous equation may be rewritten as

$$\begin{aligned} [\mathbf{H}, \rho_{\mathbf{q}}^\dagger] &= \hbar \omega_{\mathbf{q}} \rho_{\mathbf{q}}^\dagger \Rightarrow \left[\sum_{\mathbf{k} \in \Lambda} \varepsilon_{\mathbf{k}} \mathbf{c}_{\mathbf{k}}^\dagger \mathbf{c}_{\mathbf{k}} + \frac{1}{2} \sum_{\mathbf{q} \in \Lambda} \mathbf{V}_{\mathbf{q}} (\rho_{\mathbf{q}} \rho_{-\mathbf{q}} - N), \rho_{\mathbf{q}}^\dagger \right] = \hbar \omega_{\mathbf{q}} \rho_{\mathbf{q}}^\dagger \\ &= \left[\sum_{\mathbf{k} \in \Lambda} \varepsilon_{\mathbf{k}} \mathbf{c}_{\mathbf{k}}^\dagger \mathbf{c}_{\mathbf{k}}, \rho_{\mathbf{q}}^\dagger \right] + \frac{1}{2} \left[\sum_{\mathbf{q} \in \Lambda} \mathbf{V}_{\mathbf{q}} (\rho_{\mathbf{q}} \rho_{-\mathbf{q}} - N), \rho_{\mathbf{q}}^\dagger \right] = \\ &= \sum_{\mathbf{k} \in \Lambda} \varepsilon_{\mathbf{k}} \left[\mathbf{c}_{\mathbf{k}}^\dagger \mathbf{c}_{\mathbf{k}}, \rho_{\mathbf{q}}^\dagger \right] + \frac{1}{2} \sum_{\mathbf{q} \in \Lambda} \mathbf{V}_{\mathbf{q}} \left[\cancel{\rho_{\mathbf{q}} \rho_{-\mathbf{q}}} , \rho_{\mathbf{q}}^\dagger \right] = \\ &= \sum_{\mathbf{k}, \mathbf{l} \in \Lambda} \varepsilon_{\mathbf{k}} \left[\mathbf{c}_{\mathbf{k}}^\dagger \mathbf{c}_{\mathbf{k}}, \mathbf{c}_{\mathbf{q}} \mathbf{c}_{\mathbf{l}+\mathbf{q}}^\dagger \right] = \\ &= \sum_{\mathbf{k}, \mathbf{l} \in \Lambda} \varepsilon_{\mathbf{k}} \left(\mathbf{c}_{\mathbf{k}}^\dagger [\mathbf{c}_{\mathbf{k}}, \mathbf{c}_{\mathbf{q}}] \mathbf{c}_{\mathbf{l}+\mathbf{q}}^\dagger + [\mathbf{c}_{\mathbf{k}}^\dagger, \mathbf{c}_{\mathbf{q}}] \mathbf{c}_{\mathbf{k}} \mathbf{c}_{\mathbf{l}+\mathbf{q}}^\dagger + \mathbf{c}_{\mathbf{q}} \mathbf{c}_{\mathbf{k}}^\dagger [\mathbf{c}_{\mathbf{k}}, \mathbf{c}_{\mathbf{l}+\mathbf{q}}^\dagger] + \mathbf{c}_{\mathbf{q}} [\mathbf{c}_{\mathbf{k}}^\dagger, \mathbf{c}_{\mathbf{l}+\mathbf{q}}^\dagger] \mathbf{c}_{\mathbf{k}} \right) = \\ &= \sum_{\mathbf{k} \in \Lambda} (\varepsilon_{\mathbf{k}+\mathbf{q}} - \varepsilon_{\mathbf{k}}) \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}} = \end{aligned} \tag{9.4.0.1}$$

If $\rho_{\mathbf{q}}^\dagger$ actually creates an excitation of energy $\hbar \omega$, then

³⁷For example, the following four-body term may be recast as

$$\sum_{\mathbf{k}, \mathbf{k}' \in \Lambda} \mathbf{c}_{\mathbf{k}-\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}'+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}'} \mathbf{c}_{\mathbf{k}} = - \sum_{\mathbf{k}' \in \Lambda} \mathbf{n}_{\mathbf{k}'} + \rho_{\mathbf{q}} \rho_{-\mathbf{q}}.$$

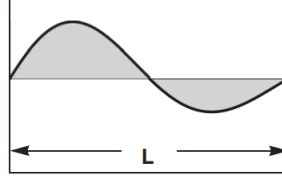


FIGURE 10. Consider a length- L box. Then, ρ_q measures the quantity approximately equal to the difference in the particle number in the two halves of the container.

(9.4.0.2)

$$\begin{aligned}
 [\mathbf{H}, [\mathbf{H}, \rho_q^\dagger]] &= (\hbar\omega)^2 \rho_q^\dagger \\
 \Rightarrow (\hbar\omega)^2 \rho_q^\dagger &= \sum_{\mathbf{k} \in \Lambda} (\varepsilon_{\mathbf{k}+\mathbf{q}} - \varepsilon_{\mathbf{k}}) [\mathbf{H}, \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}}] = \sum_{\mathbf{k} \in \Lambda} (\varepsilon_{\mathbf{k}+\mathbf{q}} - \varepsilon_{\mathbf{k}})^2 \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}} \\
 &\quad + \sum_{\mathbf{k}, \mathbf{q}' \in \Lambda} (\varepsilon_{\mathbf{k}+\mathbf{q}} - \varepsilon_{\mathbf{k}}) \frac{V_{\mathbf{q}'}}{2} \left[(\mathbf{c}_{\mathbf{k}+\mathbf{q}-\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}} - \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}+\mathbf{q}'}) \rho_{\mathbf{q}'}^\dagger + \rho_{\mathbf{q}'} (\mathbf{c}_{\mathbf{k}+\mathbf{q}+\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}} - \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}-\mathbf{q}'}) \right] \\
 &= \sum_{\mathbf{k} \in \Lambda} \left[\frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \right]^2 \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}} + \sum_{\mathbf{k}, \mathbf{q}' \in \Lambda} \frac{V_{\mathbf{q}'}}{2} \left[\frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) (\mathbf{c}_{\mathbf{k}+\mathbf{q}-\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}} - \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}+\mathbf{q}'}) \right. \\
 &\quad \left. - \frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) (\mathbf{c}_{\mathbf{k}+\mathbf{q}+\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}} - \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}-\mathbf{q}'}) \right] \\
 &= \sum_{\mathbf{k} \in \Lambda} \left[\frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \right]^2 \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}} + \sum_{\mathbf{q}' \in \Lambda} \frac{V_{\mathbf{q}'}}{2} \left[\sum_{\mathbf{k} \in \Lambda} \frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) (\mathbf{c}_{\mathbf{k}+\mathbf{q}-\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}} - \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}+\mathbf{q}'}) \right. \\
 &\quad \left. - \sum_{\mathbf{k} \in \Lambda} \frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) (\mathbf{c}_{\mathbf{k}+\mathbf{q}+\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}} - \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}-\mathbf{q}'}) \right] \\
 &= \sum_{\mathbf{k} \in \Lambda} \left[\frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \right]^2 \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}} \\
 &\quad + \sum_{\mathbf{q}' \in \Lambda} \frac{V_{\mathbf{q}'}}{2} \left[\sum_{\mathbf{k} \in \Lambda} \frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \mathbf{c}_{\mathbf{k}+\mathbf{q}-\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}} - \frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}+\mathbf{q}'} \right. \\
 &\quad \left. - \sum_{\mathbf{k} \in \Lambda} \frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \mathbf{c}_{\mathbf{k}+\mathbf{q}+\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}} - \frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}-\mathbf{q}'} \right] \\
 &= \sum_{\mathbf{k} \in \Lambda} \left[\frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \right]^2 \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}} \quad \text{let } \mathbf{k}' = \mathbf{k} + \mathbf{q}', \mathbf{k}'' = \mathbf{k} - \mathbf{q}'
 \end{aligned}$$

$$\begin{aligned}
 (9.4.0.3) \quad &+ \sum_{\mathbf{q}' \in \Lambda} \frac{V_{\mathbf{q}'}}{2} \left[\sum_{\mathbf{k} \in \Lambda} \frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \mathbf{c}_{\mathbf{k}+\mathbf{q}-\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}} - \sum_{\mathbf{k}' \in \Lambda} \frac{\hbar^2}{2m} (2(\mathbf{k}' - \mathbf{q}') \cdot \mathbf{q} + \mathbf{q}^2) \mathbf{c}_{\mathbf{k}'+\mathbf{q}-\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}'} \right. \\
 &\quad \left. - \sum_{\mathbf{k} \in \Lambda} \frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \mathbf{c}_{\mathbf{k}+\mathbf{q}+\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}} - \sum_{\mathbf{k}'' \in \Lambda} \frac{\hbar^2}{2m} (2(\mathbf{k}'' + \mathbf{q}') \cdot \mathbf{q} + \mathbf{q}^2) \mathbf{c}_{\mathbf{k}''+\mathbf{q}+\mathbf{q}'}^\dagger \mathbf{c}_{\mathbf{k}''} \right]
 \end{aligned}$$

$$(9.4.0.4) \quad (\hbar\omega)^2 \rho_q^\dagger = \sum_{\mathbf{k} \in \Lambda} \left[\frac{\hbar^2}{2m} (2\mathbf{k} \cdot \mathbf{q} + \mathbf{q}^2) \right]^2 \mathbf{c}_{\mathbf{k}+\mathbf{q}}^\dagger \mathbf{c}_{\mathbf{k}} + \sum_{\mathbf{q}' \in \Lambda} \frac{V_{\mathbf{q}'}}{2} \frac{\hbar^2 \mathbf{q} \cdot \mathbf{q}'}{m} \left(\rho_{\mathbf{q}'-\mathbf{q}} \rho_{\mathbf{q}'}^\dagger - \rho_{-\mathbf{q}'}^\dagger \rho_{-\mathbf{q}-\mathbf{q}'} \right).$$

Now, in order to study [equation \(9.4.0.4\)](#), note that ρ_0 is actually much larger than all the other components. It is just the average particle density in the system. Consider a box of length- L , shown in [figure 10](#), containing a large number of electrons of ρ_0 average density, then ρ_0 's first Fourier component is

$$\int_{\Omega \subset \mathbb{R}^3} d\mathbf{x} \rho(\mathbf{x}) e^{\frac{2i\pi x}{L}},$$

will be approximately equal to the difference between the particle number in the left and right hand sides of the box, the gray area. The number difference will then be small, compared to the total number and so ρ_0 will prevail in the summation over \mathbf{q}' . Since the $(\mathbf{q}' = 0)$ -term is omitted, ρ_0 will only appear when $\mathbf{q}' = \pm \mathbf{q}$. The neglect of all other terms, when $\mathbf{q}' \neq \pm \mathbf{q}$ is known as the **random phase approximation**³⁸.

Plasma oscillations

Consider the case where \mathbf{q} is small enough that, for the moment, the first term in [equation \(9.4.0.4\)](#) may be neglected and the second term approximated via RPA. In that case,

$$(\hbar\omega)^2 \rho_{\mathbf{q}}^\dagger = \frac{4\pi e^2}{q^2} \frac{\hbar^2 q^2}{m} (\rho_0 \rho_{\mathbf{q}}^\dagger + \rho_{\mathbf{q}}^\dagger \rho_0) \rightarrow \omega^2 = \frac{4\pi e^2 \rho_0}{m},$$

which gives just the classical plasma frequency ω_P . The contribution of the first term may be approximated by setting $k \sim k_F$ and $\langle 2\mathbf{k}_F \cdot \mathbf{q} \rangle \gg q^2$. Averaging over all angles yields

$$(\mathbf{k}_F \cdot \mathbf{q})^2 \simeq \frac{1}{4\pi} 4\pi^2 k_F^2 q^2 \int_{-1}^1 dx x^2 = 2\pi k_F^2 q^2,$$

which yield the plasmon dispersion relation

$$\omega^2(\mathbf{q}) = \omega_P^2 + \frac{2\pi \hbar^2 k_F^2}{3m^2} q^2 \simeq \omega_P^2 + \frac{\pi v_F^2}{3\omega_P} q^2.$$

Thus, the relevant long-wavelength excitations are collective motions of the electron gas. Bohm and Pines suggested that electrons act as interacting particles through a bare Coulomb potential $\frac{4\pi e^2}{q^2}$ only for length scales less than an effective screening length, which may be taken to be the Thomas-Fermi distance, $\sim q_{TF}^{-1}$. Below q_{TF} , the interaction contribute only to the plasma oscillations.

9.4.0.1. .

10. PATH INTEGRAL FORMALISM AND ITS APPLICATIONS IN CONDENSED MATTER THEORY

Path Integrals in CMT.tex Path Integrals in CMT.tex

10.1. Path Integrals.

10.1.0.1. **Quantum Propagators and Path Integrals.** The dynamical information in quantum mechanics is contained in the matrix elements of the time-evolution operator $\mathcal{U}(t_f; t_i)$, which for a time-independent Hamiltonian reads

$$(10.1.0.1) \quad \mathcal{U}(t_f; t_i) = e^{-\frac{i\mathbf{H}(t_i - t_f)}{\hbar}}.$$

The probability amplitude for a particle to start at \mathbf{x}_i -position and t_i -time and end at \mathbf{x}_f -position and t_f -time is written as

³⁸An alternative interpretation for the RPA is as follows. Consider the first-quantized version of the density operator $\rho_{\mathbf{q} \in \Lambda} = \frac{1}{\Omega} \sum_{n \in \mathbb{N}} e^{i\mathbf{q} \cdot \mathbf{x}_n}$, where \mathbf{x}_n is the n -th's particles position. Then, $\rho_{\mathbf{q} \pm \mathbf{q}'}$ may be written in first quantized form as

$$\rho_{\mathbf{q} \pm \mathbf{q}'} \frac{1}{\Omega} \sum_{n \in \mathbb{N}} e^{i(\mathbf{q} \pm \mathbf{q}') \cdot \mathbf{x}_n}.$$

Then, if the electrons are statistically randomly distributed in the system, the preceding sum only gives a finite result if $\mathbf{q} = \mp \mathbf{q}'$, since the sum involves random phases that add destructively. In RPA, [equation \(9.4.0.4\)](#) is approximated by keeping only terms with $\mathbf{q} = \mp \mathbf{q}'$.

$$(10.1.0.2) \quad \begin{aligned} \mathcal{U}(x_f, t_f; x_i, t_i) &= \langle x_f | \mathcal{U}(t_f; t_i) | x_i \rangle \\ &= \int_{x_i}^{x_f} \mathcal{D}[x] \exp \left[\frac{i}{\hbar} \int_{t_i}^{t_f} dt \mathcal{L}[\cdot, x] \right] = \int_{x_i}^{x_f} \mathcal{D}[x] \exp \left[\frac{i}{\hbar} \mathcal{S}[x(t)], \right] \end{aligned}$$

where \mathcal{L} is the Lagrangian and \mathcal{S} is the classical action associated with the $x(t)$ -trajectory and where the integration is to be understood as over all paths with starting point $x(t_i) = x_i$ and end-point $x(t_f) = x_f$. Hence, the path integral representation thus provides an elegant interpretation of the quantum-mechanical transition-amplitude, since the particle takes all possible trajectories with each $x(t)$ -trajectory contributing $\exp \left[-\frac{i}{\hbar} \mathcal{S}[x(t)] \right]$ to the amplitude. Moreover, the path integral is an exact representation of the evolution operator $\mathcal{U}(t) = \exp \left[\frac{\hbar}{i} \mathbf{H} t \right]$ without explicitly incorporating any quantum-mechanical operators.

10.1.0.2. Classical Limit and the Stationary Phase Approximation. It can be seen from [equation \(10.1.0.2\)](#) that all paths have the same absolute magnitude but with different phase: the phase of each $x(t)$ -path is its classical action $\mathcal{S}[x(t)]$ divided by its microscopic action \hbar . Hence, each path taken in isolation is equally important, i.e. the classical path is no more important than any arbitrarily complicated path.

However, given that the path integral is a coherent summation of paths, different paths interfere with one another. The important trajectories, hence, contribute significantly to the path integral must be those for which the action varies very slowly when the $x(t)$ -path is slightly deformed, which means that they leave the \mathcal{S} -action stationary, i.e.

$$(10.1.0.3) \quad \frac{\delta \mathcal{S}[x]}{\delta x} = 0.$$

By contrast, the contributions which rapidly oscillating exponential factor add up to zero. Thus, in the classical limit, only one trajectory $x_c(t)$ contributes, that which is the solution to the classical equations of motions.

In the classical limit, $\hbar \rightarrow 0$. In particular, consider a deviation from the classical path $x(t) = x_{cl}(t) + \eta(t)$, which allows for a second order expansion in this *small* deviation as follows

$$(10.1.0.4) \quad \begin{aligned} \mathcal{U}(x_f, x_i; t_f - t_i) &\simeq e^{\frac{i}{\hbar} \mathcal{S}_{cl}[x_{cl}]} \int_{\eta(t_i)=0}^{\eta(t_f)=0} \mathcal{D}[\eta] \exp \left[\frac{i}{2\hbar} \int_{t_i}^{t_f} dt \frac{\delta^{(2)} \mathcal{S}[x]}{\delta x(t) \delta x(t')} \eta(t) \eta(t') \right] \\ &= e^{\frac{i}{\hbar} \mathcal{S}_{cl}[x_{cl}]} \det \left[\frac{1}{2\pi\hbar} \frac{\delta^{(2)} \mathcal{S}[x]}{\delta x(t) \delta x(t')} \right]^{-\frac{1}{2}} \end{aligned}$$

The quantum correction to the classical action thus measures the strength of the quantum fluctuations. In general, the action for the quantum η -fluctuations is not the same as the action for the classical trajectory.

10.1.1. Path integrals in Statistical Mechanics. Path integrals offer an elegant and powerful technique for calculating observables, in particular correlation functions, for statistical mechanics systems, since -in this formalism- the partition function is a path integral. Hence, the finite-temperature quantum-mechanical properties of the system can be accessible from the path integral.

A statistical ensemble in equilibrium at finite β^{-1} -temperature is described by a partition function

$$(10.1.1.1) \quad \mathcal{Z} = \text{Tr} e^{-\beta \mathbf{H}}.$$

The partition function can be written in the eigenstate basis

$$(10.1.1.2) \quad \mathcal{Z} = \int dx \langle x | e^{-\beta \mathbf{H}} | x \rangle = \int dx \int dx_1 \langle x | x_1 \rangle \langle x_1 | e^{-\beta \mathbf{H}} | x \rangle,$$

where $\langle x_1 | e^{-\beta \mathbf{H}} | x \rangle$ can now be viewed as analytic continuation of the propagator [equation \(10.1.0.1\)](#) to imaginary time $\tau = it$, where the propagation is considered to start at $\tau = 0$ and ending at $\tau = \beta \hbar$, i.e. a Wick rotation which maps Minkowski \mathbb{M}^4 to a euclidean \mathbb{R}^4 . Performing these steps yields the following representation for the propagator

$$(10.1.1.3) \quad \langle x_f | e^{-\beta \mathbf{H}} | x_i \rangle = \mathcal{N} \int \mathcal{D}[x(\tau)] e^{-\frac{\mathcal{S}_E[x]}{\hbar}},$$

$$\mathcal{S}_E[x] = \int_0^{\hbar\beta} d\tau \mathbf{H}(x), \quad \text{is the Euclidean action.}$$

Note that, in the previous relations, the Hamiltonian has replaced the Lagrangian in the action, since $t = -i\tau$. The positive form of the Euclidean actions eliminates any convergence problems in the integral.

Given that the partition function is, by definition, a trace over states, boundary conditions must be imposed s.t. that the initial and final states are the same state $x(0) = x(\hbar\beta)$. Then, a sum is taken across all such states. In other words, the imaginary-time periodic boundary condition reads: $x(\tau) = x(\tau + \hbar\beta)$. Taking into account the previous considerations, the partition function now reads,

$$(10.1.1.4) \quad \mathcal{Z} = \mathcal{N} \int dx \int_{x(0)=x}^{x(\hbar\beta)=x} \mathcal{D}[x(\tau)] e^{-\frac{\mathcal{S}_E[x]}{\hbar}}.$$

Real-time dynamics and quantum-statistical properties can be obtained by the transformation $t \rightarrow -it$, a Wick rotation of $\frac{\pi}{2}$ in the complex-time plane.

10.1.1.1. Correlation Functions. Usually, one is interested in measuring thermal-averages of (products of) local observables. Analytically, this reads as

$$(10.1.1.5) \quad \langle \mathbf{O}_1 \cdots \mathbf{O}_n \rangle = \frac{1}{\mathcal{Z}} \text{Tr} \left(\mathbf{O}_1 \cdots \mathbf{O}_n e^{-\beta \mathbf{H}} \right).$$

Such averages can be readily obtained in an alternative way, by defining a modified partition function

$$(10.1.1.6) \quad \mathcal{Z}(h_1, \dots, h_n) = \text{Tr} e^{-\beta \mathbf{H} + \mathbf{H}_{\text{ext}}}, \quad \mathbf{H}_{\text{ext}} = \sum_i h_i \mathbf{O}_i,$$

wherein \mathbf{H}_{ext} represents the Hamiltonian of external sources h_i which couple to the relevant operators \mathbf{O}_i . This allows for the simple representation of an n -body correlator as follows

$$(10.1.1.7) \quad \langle \mathbf{O}_1 \cdots \mathbf{O}_n \rangle = \frac{1}{\mathcal{Z}[\{h_i = 0\}]} \left. \frac{\partial^n \mathcal{Z}[h_1, \dots, h_n]}{\partial h_1 \cdots \partial h_n} \right|_{\{h_i = 0\}}.$$

Now, the question arises as to the possibility of extending this thermal averaging process to correlation functions of the form,

$$(10.1.1.8) \quad \langle \mathbf{O}_1(\tau_1) \cdots \mathbf{O}_n(\tau_n) \rangle = \frac{\int_{x(0)=x(\beta)} \mathcal{D}[x(\tau)] \mathbf{O}_1(\tau_1) \cdots \mathbf{O}_n(\tau_n) e^{-\frac{\mathcal{S}_E}{\hbar}}}{\int_{x(0)=x(\beta)} \mathcal{D}[x(\tau)] e^{-\frac{\mathcal{S}_E}{\hbar}}}, \quad \text{where } \mathbf{O}_i(\tau_i) = \mathbf{O}_i(x(\tau_i)).$$

Remark. To interpret such a correlation function in imaginary time, one may start by considering the action of operators on a given state function $|\phi(t)\rangle$ at different times. Operator \mathbf{O}_1 acts at t_1 s.t.

$$(10.1.1.9)$$

Operator \mathbf{O}_1 acts at t_1 s.t. $\mathbf{O}_1 |\phi(t_1)\rangle = \mathbf{O}_1 \mathcal{U}(t_1, 0) |\phi(0)\rangle$

Operator \mathbf{O}_2 acts at $t_2 > t_1$ s.t. $\mathbf{O}_2 |\phi(t_2)\rangle = \mathbf{O}_2 \mathcal{U}(t_2, t_1) \mathbf{O}_1 \mathcal{U}(t_1, 0) |\phi(0)\rangle$

\Rightarrow The projection of this states onto the unperturbed state reads $\langle \phi(0) | \mathbf{O}_2 \mathcal{U}(t_2, t_1) \mathbf{O}_1 \mathcal{U}(t_1, 0) | \phi(0) \rangle = \langle \psi(0) | \mathbf{O}_{2H}(t_2) \mathbf{O}_{1H}(t_1) | \psi(0) \rangle$

Its analytic continuation to imaginary time reads

$$\frac{\text{Tr } e^{-\beta \mathbf{H}} \mathcal{T}_\tau [\mathbf{O}_{2H}(t_2) \mathbf{O}_{1H}(t_1)]}{\text{Tr } e^{-\beta \mathbf{H}}}.$$

where a time-ordering operator has been introduced since, now, $\mathbf{O}_H^\dagger(\tau) \neq [\mathbf{O}_H(\emptyset)]^\dagger$. It follows then, by using a complete set of basis functions, that

(10.1.1.10)

$$\begin{aligned} \text{Tr } e^{-\beta \mathbf{H}} \mathcal{T}_\tau [\mathbf{O}_{2H}(t_2) \mathbf{O}_{1H}(t_1)] &= \int_{\mathbb{R}^9} dx dx_1 dx_2 \langle x | e^{-(\beta-\tau_2)\mathbf{H}} | x_2 \rangle \langle x_2 | \mathbf{O}_2 e^{-(\tau_2-\tau_1)\mathbf{H}} | x_1 \rangle \langle x_1 | \mathbf{O}_1 e^{-\tau_1\mathbf{H}} | x \rangle \\ &= \int_{\mathbb{R}^9} dx dx_1 dx_2 \mathbf{O}_2(x_2) \mathbf{O}_1(x_1) \langle x | e^{-(\beta-\tau_2)\mathbf{H}} | x_2 \rangle \langle x_2 | e^{-(\tau_2-\tau_1)\mathbf{H}} | x_1 \rangle \langle x_1 | e^{-\tau_1\mathbf{H}} | x \rangle \\ &= \int_{\mathbb{R}^9} dx dx_1 dx_2 \mathbf{O}_2(x_2) \mathbf{O}_1(x_1) \left[\int_{x(\tau_2)=x_2}^{x(\beta)=x} \mathcal{D}[x(\tau)] \int_{x(\tau_1)=x_1}^{x(\tau_2)=x_2} \mathcal{D}[x'(\tau)] \int_{x(0)=x}^{x(\tau_1)=x_1} \mathcal{D}[x''(\tau)] \right. \\ &\quad \times e^{-\int_{\tau_2}^{\beta} d\tau \mathcal{L}_E(x(\tau))} e^{-\int_{\tau_1}^{\tau_2} d\tau \mathcal{L}_E(x(\tau))} e^{-\int_0^{\tau_1} d\tau \mathcal{L}_E(x(\tau))} \left. \right] \\ &= \int_{x(0)=x(\beta)} \mathcal{D}[x(\tau)] \mathbf{O}_2(x_2) \mathbf{O}_1(x_1) e^{-\int_0^{\beta} d\tau \mathcal{L}_E(x(\tau))}. \end{aligned}$$

Hence, the time-ordered correlation functions naturally manifest as moments of the path integral. Moreover, all operators in the path-integral disappear and are replaced by classical c-numbers. Furthermore, since all paths must satisfy the condition $x(0) = x(\beta)$, the correlation functions can be expanded in terms of a Matsubara Fourier series. The real-time correlation function is obtained then through analytic continuation of Matsubara sums. Thus, path integrals present very convenient tools to obtain time-ordered correlation functions.

10.1.1.2. Correlation Function in Terms of State Kets. Frequently, the transition amplitudes between two state kets $|\psi_i\rangle$ and $|\psi_f\rangle$ is desired, which can be written as

$$\begin{aligned} \langle \psi_f | \psi_i \rangle &= \int dx_f dx_i \langle \psi_f | x_f, t_f \rangle \langle x_f, t_f | x_i, t_i \rangle \\ (10.1.1.11) \quad &= N \int dx_f dx_i \psi_f^*(x_f, t_f) \psi_i(x_i, t_i) \int \mathcal{D}[x] e^{\frac{i}{\hbar} S[x]}. \end{aligned}$$

This definition can be extended to time-ordered correlation functions by writing

$$\begin{aligned} \langle \overline{\psi_f} | \psi_i \rangle &= \langle \psi_f | \mathcal{T}[\mathbf{O}(\tau_1) \cdots \mathbf{O}(\tau_n)] | \psi_i \rangle \\ (10.1.1.12) \quad &= N \int dx_f dx_i \psi_f^*(x_f, t_f) \psi_i(x_i, t_i) \int \mathcal{D}[x] \mathcal{T}[\mathbf{O}(\tau_1) \cdots \mathbf{O}(\tau_n)] e^{\frac{i}{\hbar} S[x]}. \end{aligned}$$

These correlation functions can be written by defining a generalized action, in the form

$$(10.1.1.13) \quad S[x, \mathbf{h}] = S[x] + \int_{t_i}^{t_f} dt \mathcal{O}[x(t)] \cdot \mathbf{h}, \quad \mathcal{O}[x(t)] = [\mathbf{O}_1(x(\tau_1)) \cdots \mathbf{O}_1(x(\tau_n))], \\ \mathbf{h}^T = [h_1 \cdots h_n].$$

From which, the correlation functions are given by

(10.1.1.14)

$$\langle \psi_f | \mathcal{T}[\mathbf{O}(\tau_1) \cdots \mathbf{O}(\tau_n)] | \psi_i \rangle = N(-i\hbar)^n \int dx_f dx_i \psi_f^*(x_f, t_f) \psi_i(x_i, t_i) \int \mathcal{D}[x] \frac{\delta^{(n)} S[x, \mathbf{h}]}{\delta h_1(t_1) \cdots \delta h_n(t_n)} \Big|_{\mathbf{h}=0} e^{\frac{i}{\hbar} S[x]}.$$

10.1.2. *Functional Integrals in Many-Body Theory.* Upto now, the path integral has been formulated in terms of the single-particle eigenstates, e.g. position and momentum eigenkets. However, in many body systems, the functional integrals will be written in terms, not of position/momentum eigenstates, but in terms of coherent states, which themselves are parameterized in terms of complex scalar variables (or Grassmann variables in the case of fermions). A functional integral approach will, then, be written by replacing operators with their corresponding scalar variables.

10.1.2.1. **Coherent State Functional Integral.** The grand-canonical partition function can be written in terms of the many-body coherent states $|\psi\rangle$ as follows,

(10.1.2.1)

$$\mathcal{Z} = \sum_n \int d(\bar{\psi}, \psi) e^{-\sum_\alpha \bar{\psi}_\alpha \bar{\psi}_\alpha} \langle n|\psi\rangle \langle\psi| e^{-\beta(\mathbf{H}-\mu\mathbf{N})} |n\rangle, \quad d(\bar{\psi}, \psi) = \begin{cases} \prod_\alpha \frac{d\bar{\psi}_\alpha d\psi_\alpha}{2\pi i}, & \text{bosons} \\ \prod_\alpha d\bar{\psi}_\alpha d\psi_\alpha, & \text{fermions.} \end{cases}$$

where ϕ_a is a c-number for bosons and a Grassmann variable for fermions, and where $d(\bar{\psi}, \psi)$ is the integration measure. In order to eliminate the sum over n , the matrix elements must be recast into a resolution of the identity on the $|n\rangle$ -eigenstates. This means that the $\langle n|\bar{\psi}\rangle$ must be commuted with $\langle\psi| e^{-\beta(\mathbf{H}-\mu\mathbf{N})} |n\rangle$. In the case of bosons, the commutation holds but in the case of fermions, the Grassmann variables anticommute, leading to

$$\langle n|\psi\rangle \langle\psi|n\rangle = \langle -\psi|n\rangle \langle n|\psi\rangle,$$

where $\langle -\psi| = \langle\emptyset| \exp[-\sum_\alpha \bar{\psi}_\alpha \mathbf{c}_\alpha^\dagger]$. Given that both the Hamiltonian and number operators contain even numbers of creation and annihilation operators, and thus are not susceptible to sign changes, it follows that

(10.1.2.2)

$$\mathcal{Z} = \int d(\bar{\psi}, \psi) e^{-\sum_\alpha \bar{\psi}_\alpha \psi_\alpha} \langle \zeta \psi | e^{-\beta(\mathbf{H}-\mu\mathbf{N})} | \psi \rangle,$$

with $\zeta = \pm 1$ for bosons and fermions, respectively. The interval β can be divided into N segments, with $\delta\beta = \beta/N$ s.t.

$$\langle \zeta \psi | e^{-\beta(\mathbf{H}-\mu\mathbf{N})} | \psi \rangle = \langle \zeta \psi | \bigwedge_{j=1}^N e^{-\delta\beta(\mathbf{H}-\mu\mathbf{N})} | \psi \rangle.$$

Then, it follows that

$$(10.1.2.3) \quad \mathcal{Z} = \int_{\substack{\bar{\psi}_N = \zeta \bar{\psi}_0 \\ \psi_N = \zeta \psi_0}} \mathcal{D}(\bar{\psi}, \psi) e^{-S[\bar{\psi}, \psi]}, \quad S[\bar{\psi}, \psi] = \int_0^\beta d\tau \left[\bar{\psi} \partial_\tau \psi + \mathbf{H}(\bar{\psi}, \psi) - \mu \mathbf{N}(\bar{\psi}, \psi) \right].$$

10.1.3. *Hubbard-Stratonovich Transformation.* Biquadratic, i.e. four-variable terms, in the Lagrangian cannot be readily integrated in the path integral. These terms can be decoupled via the Hubbard-Stratonovich transformation. The Hubbard-Stratonovich transformation (HST) is a general method for replacing a four-body interaction with two-body operators with external bosonic exchange fields, and is a generalization of the Gaussian Integral

(10.1.3.1)

$$e^{\mathbf{O}^2} = \int_{\mathbb{R}} dx \exp \left[-\pi x^2 - 2\sqrt{\pi} x \mathbf{O} \right],$$

where \mathbf{O} is an operator. Hence, if a biquadratic term appears in the partition function [equation \(10.1.2.3\)](#) - e.g. a term similar to the left-hand side of the previous equation- representing a two-body interaction, it can be replaced by the right-hand side, a Gaussian integral which only contains two-body operators. However, these single particles now interact with an external bosonic field that must be integrated.

In order to apply the HST to the case of a bilinear interaction of the form,

$$\mathbf{H}_I(\tau_i) = V_{12} \mathbf{O}_1(\tau_i) \mathbf{O}_2(\tau_i) = V_{12} \left[\left(\frac{\mathbf{O}_1 + \mathbf{O}_2}{2} \right)^2 - \left(\frac{\mathbf{O}_1 - \mathbf{O}_2}{2} \right)^2 \right],$$

valid for commuting \mathbf{O}_1 and \mathbf{O}_2 . Then, it follows that

$$\begin{aligned}
 (10.1.3.2) \quad e^{\mathbf{H}_I(\tau_i)} &= \int \mathcal{D}[\phi] \exp \left[-\pi \int_0^\beta d\tau \phi(\tau) V_{12}^{-1} \phi^*(\tau) \right] \\
 &\quad \times \mathcal{T}_\tau \exp \left[- \int_0^\beta d\tau \sqrt{\pi} \left\{ \mathbf{O}_1(\tau) \phi^*(\tau) - \mathbf{O}_2(\tau) \phi(\tau) \right\} \right],
 \end{aligned}$$

where $\phi(\tau)$ is a complex field. In most cases, $\mathbf{O}_2 = \mathbf{O}_1^\dagger$.

Remark.

Choosing the appropriated HST is always an educated guess, which may have profound physical ramifications. As a matter of fact, one can consider pairings of \mathbf{a} -operators in three distinct and physically meaningful ways, which are discussed as follows

- For Hartree or direct Coulomb interactions, the interaction Hamiltonian reads

$$(10.1.3.3) \quad \mathbf{H}_I = \frac{1}{2} \sum_{\substack{\mathbf{k}, \mathbf{k}', \mathbf{q} \\ \sigma, \sigma' \in \mathbb{C}^2}} \frac{4\pi e^2}{q^2} \mathbf{c}_{\mathbf{k}+\mathbf{q}, \sigma}^\dagger \mathbf{c}_{\mathbf{k}'-\mathbf{q}, \sigma'}^\dagger \mathbf{c}_{\mathbf{k}' \sigma'} \mathbf{c}_{\mathbf{k} \sigma} = \frac{1}{2} \sum_{\substack{\mathbf{q} \\ \sigma, \sigma' \in \mathbb{C}^2}} \frac{4\pi e^2}{q^2} \rho_\sigma(\mathbf{q}) \rho_{\sigma'}(-\mathbf{q}) + c.$$

By identifying $\mathbf{O}_1 = e \sum_{\sigma \in \mathbb{C}^2} \rho_\sigma(\mathbf{q})$ and $\mathbf{O}_2 = e \sum_{\sigma \in \mathbb{C}^2} \rho_\sigma(-\mathbf{q})$, the HST yields

$$(10.1.3.4) \quad -\frac{1}{8\pi} \sum_{\mathbf{q}} q^2 \phi_{-\mathbf{q}} \phi_{\mathbf{q}} + i \sum_{\mathbf{q}} \left[e \rho(\mathbf{q}) \phi_{-\mathbf{q}} - e \rho(-\mathbf{q}) \phi_{\mathbf{q}} \right]$$

The inverse Fourier transform of the first term³⁹ yields $\frac{1}{8\pi} (\nabla \phi(\mathbf{x}))^2 = E^2/8\pi$, which is the bosonic electric field. Furthermore, for a Coulomb-electron gas, the action, given in [equation \(10.1.2.3\)](#), can be written as

$$(10.1.3.5) \quad S = \int d\mathbf{x} d\tau \left[\bar{\psi}(\mathbf{x}) \left(\partial_\tau - \frac{p^2}{2m} + e\phi(\mathbf{x}) - \mu \right) \psi(\mathbf{x}) - \frac{1}{2} \nabla^2 \phi^2 \right].$$

- Consider, for instance, the contact interaction model, whose interaction Hamiltonian reads

$$(10.1.3.6) \quad g \sum_{s, s'} \Psi_s^\dagger(x) \Psi_{s'}^\dagger(x) \Psi_{s'}(x) \Psi_s(x) = -\mathbf{s}(x) \cdot \mathbf{s}(x),$$

where $\mathbf{s}(x) = \Psi_s^\dagger(x) \sigma_{ss'} \Psi_{s'}(x)$. By introducing an exchange field \mathbf{m} , the HST yields

$$(10.1.3.7) \quad \int \mathcal{D}[\mathbf{m}] \exp \left[- \int_0^\beta d\tau d\mathbf{x} (m^2 - 2\mathbf{m} \cdot \mathbf{s}) \right] = N \exp \left[\int_0^\beta d\tau d\mathbf{x} S^2(\mathbf{x}) \right]$$

with $S = \int_0^\beta d\mathbf{x} d\tau \left[\sum_{s \in \mathbb{C}^2} \bar{\psi}_s (\partial_\tau - \mathbf{H}_0 - \mu) \psi_s - 2g\mathbf{m} \cdot \mathbf{s} + gm^2 \right],$

wherein the \mathbf{m} -field is allowed to fluctuate and the integration is carried out over all possible paths of this field.

- Consider the pairing interaction Hamiltonian associated with superconductivity BCS theory,

$$(10.1.3.8) \quad \mathbf{H}_I = -g \sum_{\mathbf{k}, \mathbf{k}'} \mathbf{c}_{\mathbf{k}\uparrow}^\dagger \mathbf{c}_{-\mathbf{k}\downarrow}^\dagger \mathbf{c}_{-\mathbf{k}'\downarrow} \mathbf{c}_{\mathbf{k}'\uparrow} = -g \left(\sum_{\mathbf{k}} \mathbf{c}_{\mathbf{k}\uparrow}^\dagger \mathbf{c}_{-\mathbf{k}\downarrow}^\dagger \right) \left(\sum_{\mathbf{k}} \mathbf{c}_{-\mathbf{k}\downarrow} \mathbf{c}_{\mathbf{k}\uparrow} \right).$$

Applying the HST with a complex pairing field Δ yields

³⁹In effect, consider the definition of the inverse Fourier transform

$$\begin{aligned} -\nabla^2 \phi(\mathbf{x}) &= \rho(\mathbf{x}), \quad \text{let } \phi(\mathbf{x}) = \int d\mathbf{k} \hat{\phi}(\mathbf{k}) e^{i\mathbf{k} \cdot \mathbf{x}} \Rightarrow \nabla^2 \phi(\mathbf{x}) = \nabla^2 \int d\mathbf{k} \hat{\phi}(\mathbf{k}) e^{i\mathbf{k} \cdot \mathbf{x}} \\ &= \int d\mathbf{k} \hat{\phi}(\mathbf{k}) \nabla^2 e^{i\mathbf{k} \cdot \mathbf{x}} \\ &= \int d\mathbf{k} [-k^2 \hat{\phi}(\mathbf{k})] e^{i\mathbf{k} \cdot \mathbf{x}}. \end{aligned}$$

Then, it follows that $\mathcal{F}[\nabla^2 \phi(\mathbf{x})] = -k^2 \hat{\phi}(\mathbf{k})$.

$$(10.1.3.9) \quad \frac{\bar{\Delta}\Delta}{g} + \bar{\Delta} \left(\sum_{\mathbf{k}} \mathbf{c}_{\mathbf{k}\uparrow}^\dagger \mathbf{c}_{-\mathbf{k}\downarrow}^\dagger \right) + \Delta \left(\sum_{\mathbf{k}} \mathbf{c}_{-\mathbf{k}'\downarrow} \mathbf{c}_{\mathbf{k}'\uparrow} \right).$$

10.1.4. *Mean Field Theory and Saddle-Point Approximation.* So far, the treatment of four-body interactions and its variants is exact, no approximations being made anywhere along the way. The HST gives the exact expression

(10.1.4.1)

$$\begin{aligned} \mathcal{Z} &= \int \mathcal{D}[\phi^*, \phi] \mathcal{D}[\bar{\psi}, \psi] \exp \left[-\pi \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) \right] \exp \left[-\int_0^\beta d\tau (\bar{\psi} \partial_\tau \psi + \mathbf{H}_{\text{eff}}) \right] \\ &= \int \mathcal{D}[\phi^*, \phi] \mathcal{D}[\bar{\psi}, \psi] e^{-S[\phi^*, \phi, \bar{\psi}, \psi]}, \\ &\equiv \int \mathcal{D}[\phi^*, \phi] e^{-S_{\text{eff}}[\phi^*, \phi]}, \end{aligned}$$

$$\text{where } e^{-S_{\text{eff}}[\phi^*, \phi]} \equiv \int \mathcal{D}[\bar{\psi}, \psi] e^{-S[\phi^*, \phi, \bar{\psi}, \psi]}$$

$$e^{-S[\phi^*, \phi, \bar{\psi}, \psi]} = \exp \left[-\pi \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) \right] \exp \left[-\int_0^\beta d\tau (\bar{\psi} \partial_\tau \psi + \mathbf{H}_{\text{eff}}) \right]$$

$$\text{where } S_{\text{eff}} \text{ is an effective action given by } S_{\text{eff}}[\phi^*, \phi] = -\pi \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) - \int_0^\beta d\tau \left\{ \bar{\psi} (\partial_\tau - \mu) \psi + \mathbf{H}_{\text{eff}} \right\},$$

$$\text{and } \mathbf{H}_{\text{eff}} = \bar{\psi} \mathbf{H}_0 \psi + \sum_i (\bar{\mathbf{O}}_i \phi_i + \mathbf{O}_i \phi_i^*),$$

where the \mathbf{O}_i operators are bilinear in the $(\bar{\psi}, \psi)$ fields. In principle, one can carry out the functional integral in the previous equation, but the problem lies in that handling the functional integral over $\phi(x, t)$ is not easy nor known. Note that the four-body interactions become encoded in the dynamics of the bosonic ϕ -field, and hence the various approximations to treat it. The simplest way to treat it is to subject the action to a stationary phase treatment, which amounts to perform a saddle-point approximation for the functional integral over ϕ , by extremizing the action, as follows

$$(10.1.4.2) \quad \frac{\delta S_{\text{eff}}[\phi^*, \phi]}{\delta \phi_i^*} = \frac{\delta \log \mathcal{Z}}{\delta \phi_i^*} = \frac{1}{\mathcal{Z}} \int \mathcal{D}[\bar{\psi}, \psi] (V_i^{-1} \phi_i + \mathbf{O}_i) e^{-S[\phi^*, \phi, \bar{\psi}, \psi]} = 0.$$

This amounts to a classical approximation. It is clear that such a procedure will lead to a fixed configuration of the ϕ -field, namely

$$\phi_i^0 = -V_i \langle \mathbf{O}_i \rangle,$$

whereby the stationary auxiliary field is the average of the four-body interaction over the associated operator. This is just the mean-field approximation, where for a example a quartic term $\mathbf{c}_i^\dagger \mathbf{c}_j \mathbf{c}_k^\dagger \mathbf{c}_\ell$ becomes $\langle \mathbf{c}_i^\dagger \mathbf{c}_j \rangle \mathbf{c}_k^\dagger \mathbf{c}_\ell + \mathbf{c}_i^\dagger \mathbf{c}_j \langle \mathbf{c}_k^\dagger \mathbf{c}_\ell \rangle$.

Remark.

Moreover, it is worthwhile to remark, that both methods are equivalent, with the HST being systematic, as it might be easier to figure out how to go beyond the mean field as well contributions and combinations from different channels. In the case of superconductivity, $\mathbf{c}_i^\dagger \mathbf{c}_j \mathbf{c}_k^\dagger \mathbf{c}_\ell \mapsto \langle \mathbf{c}_i^\dagger \mathbf{c}_j \rangle \mathbf{c}_k^\dagger \mathbf{c}_\ell + \mathbf{c}_i^\dagger \mathbf{c}_j \langle \mathbf{c}_k^\dagger \mathbf{c}_\ell \rangle$. However, the HST are arbitrary, allowing for any arbitrary combination of any number of them, and the different emerging mean-field theories give different results.

Remark.

Spontaneous symmetry breaking occurs when the order parameter ϕ_i^0 or $\langle \mathbf{O}_i \rangle$ assumes a non-zero value. In both methods, order parameters can be then solved self-consistently to yield the decoupled Hamiltonian.

In the example of the Coulombic electron gas, this yields

$$(10.1.4.3) \quad \frac{\delta S_{\text{eff}}[\phi^*, \phi]}{\delta \phi_i^*} = \frac{1}{\mathcal{Z}} \int \mathcal{D}[\bar{\psi}, \psi] \left(e\rho(\mathbf{x}) + \nabla^2 \phi(\mathbf{x}) \right) e^{-S} = 0,$$

and the saddle-point solution satisfies the Gauss' law

$$(10.1.4.4) \quad -\nabla^2 \phi = \nabla \cdot \mathbf{E} = e\rho(\mathbf{x}).$$

Spatial and temporal homogeneity translates to a uniform potential ϕ_0 , and charge neutrality sets it to zero.

10.1.4.1. Effective Action and Fluctuations. One can, then, calculate systematic corrections to the lowest-order stationary-phase approximation ϕ^0 . If the integration on the $(\bar{\psi}, \psi)$ is carried, the properties of functional fluctuations on the bosonic field $\delta\phi = \phi - \phi^0$ can be studied. From [equation \(10.1.2.3\)](#) it follows that,

$$\mathcal{Z} = \int \mathcal{D}(\delta\phi^*, \delta\phi) e^{-S_{\text{eff}}[\delta\phi^*, \delta\phi]}, \text{ where } e^{-S_{\text{eff}}[\phi^*, \phi]} \equiv \int \mathcal{D}[\bar{\psi}, \psi] e^{-S[\phi^*, \phi, \bar{\psi}, \psi]}$$

$$\text{and where } S_{\text{eff}}[\phi^*, \phi] = S_{\text{int}}[\phi^*, \phi] + S^{(1)}[\phi^*, \phi]$$

$$= - \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) - \int_0^\beta d\tau \left\{ \bar{\psi}(\partial_\tau - \mu) \psi + \mathbf{H}^{\text{eff}} \right\},$$

The second term in the last line of the previous equation is not known. However, the "classical" partition function associated with it, is. In effect,

$$(10.1.4.5) \quad \begin{aligned} \mathcal{Z}^{(1)} &= \int \mathcal{D}[\bar{\psi}, \psi] e^{-S^{(1)}[\phi^*, \phi]} \\ &= \int \mathcal{D}[\bar{\psi}, \psi] \exp \left[- \int_0^\beta d\tau \left\{ \bar{\psi}(\partial_\tau - \mu) \psi + \mathbf{H}^{\text{eff}} \right\} \right] = \det^\zeta \left[\partial_\tau - \mu + \mathbf{H}^{\text{eff}} \right] \\ &\Rightarrow S^{(1)}[\bar{\psi}, \psi] = \log \mathcal{Z}^{(1)} = \log \det^\zeta \left[\partial_\tau - \mu + \mathbf{H}^{\text{eff}} \right] = \text{Tr} \log \left[\partial_\tau - \mu + \mathbf{H}^{\text{eff}} \right] \end{aligned}$$

Using this result in the first equation of this section, yields an exact and tractable expression for the effective action, i.e.

(10.1.4.6)

$$\begin{aligned} S_{\text{eff}}[\phi^*, \phi] &= S_{\text{int}}[\phi^*, \phi] + S^{(1)}[\phi^*, \phi] \\ &= - \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) - \int_0^\beta d\tau \left\{ \bar{\psi}(\partial_\tau - \mu) \psi + \mathbf{H}^{\text{eff}} \right\}, \\ &= - \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) - \zeta \text{Tr} \log \left[\partial_\tau - \mu + \mathbf{H}^{\text{eff}} \right] \\ &= - \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) - \zeta \text{Tr} \log \left[\partial_\tau + \mathbf{H}_0 - \mu + \sum_i (\bar{\mathbf{O}}_i \phi_i + \mathbf{O}_i \phi_i^*) \right] \end{aligned}$$

But notice that $\partial_\tau + \mathbf{H}_0 - \mu = (\mathfrak{G}^{(0)})^{-1}$, the non-interacting Green function s.t.

$$\begin{aligned} \exp[\text{Tr} \log -(\mathfrak{G}^{(0)})^{-1}] &= \exp[-\text{Tr} \log -(\mathfrak{G}^{(0)})] = \exp[\log \det(\mathfrak{G}^{(0)})] = \mathcal{Z}_0, \text{ then} \\ &= - \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) - \zeta \text{Tr} \log \left[-(\mathfrak{G}^{(0)})^{-1} + \sum_i (\bar{\mathbf{O}}_i \phi_i + \mathbf{O}_i \phi_i^*) \right] \\ &= - \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) - \zeta \text{Tr} \log \left\{ -(\mathfrak{G}^{(0)})^{-1} \times [1 - \mathfrak{G}^{(0)} \sum_i (\bar{\mathbf{O}}_i \phi_i + \mathbf{O}_i \phi_i^*)] \right\} \\ &= - \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) - \zeta \text{Tr} \log \left[-(\mathfrak{G}^{(0)})^{-1} \right] - \zeta \text{Tr} \log \left\{ [1 - \mathfrak{G}^{(0)} \sum_i (\bar{\mathbf{O}}_i \phi_i + \mathbf{O}_i \phi_i^*)] \right\} \end{aligned}$$

$$\begin{aligned}
&= - \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) - \zeta \operatorname{Tr} \log \left[- (\mathfrak{G}^{(0)})^{-1} \right] - \zeta \sum_{n=1}^{\infty} \frac{1}{n!} \operatorname{Tr} \left[\mathfrak{G}^{(0)} \sum_i (\bar{\mathbf{O}}_i \phi_i + \mathbf{O}_i \phi_i^*) \right]^n \\
&\Rightarrow \mathcal{Z} = \mathcal{Z}_0 \int \mathcal{D}(\delta\phi^*, \delta\phi) \exp \left[- \int_0^\beta d\tau \phi(\tau) V^{-1} \phi^*(\tau) - \zeta \sum_{n=1}^{\infty} \frac{1}{n!} \operatorname{Tr} \left[\mathfrak{G}^{(0)} \sum_i (\bar{\mathbf{O}}_i \phi_i + \mathbf{O}_i \phi_i^*) \right]^n \right]
\end{aligned}$$

where \mathcal{Z}_0 is a normalization factor stemming from the non-interacting theory.

Example

For example, consider the Coulomb gas, which is mediated by a scalar photon field ϕ , which is real and periodic $\phi(\tau + \beta) = \phi(\tau)$. A plane-wave basis expansion and Matsubara-Fourier series expansion in the fermionic and bosonic Matsubara frequencies Ω_n, ω_m , respectively, yields

$$\begin{aligned}
(10.1.4.7) \quad S_{\text{eff}}[\phi] &= \frac{1}{2} \sum_{\mathbf{q}} \left(\frac{q^2}{2\pi} - e^2 \Pi_0(\mathbf{q}) \right) |\phi(\mathbf{q})|^2 + \mathcal{O}(e^4) \\
&= \frac{1}{2} \sum_{\mathbf{q}} \frac{q^2}{2\pi} \left(1 - \frac{4\pi e^2}{q^2} \Pi_0(\mathbf{q}) \right) |\phi(\mathbf{q})|^2 = \sum_q \varepsilon(q) \frac{q^2}{4\pi},
\end{aligned}$$

which describes the screening of Coulomb field fluctuations by the electron gas.

10.2. . ref.bib