

Finite Group Lattice Gauge Theories for Quantum Simulation

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INTRODUCTION

CHAPTER ONE

INTRODUCTION TO LATTICE GAUGE THEORIES

1.1 REVIEW OF YANG-MILLS THEORY

A Yang-Mills theory is a gauge field theory on Minkowski space $\mathbb{R}^{1,d}$, where the gauge group $U(1)$ or $SU(N)$, with matter fields, which are defined by a representation of the gauge group. For example, *Quantum Chromodynamics* (QCD) is an $SU(3)$ gauge theory with Dirac spinors in the fundamental representation. Keeping in mind the example of QCD, the Lagrangian of the theory is

$$\mathcal{L} = -\frac{1}{2g^2} \text{tr}(F_{\mu\nu}F^{\mu\nu}) + \bar{\psi}(i\gamma^\mu D_\mu - m)\psi, \quad (1.1)$$

where the fermions ψ are taken in the fundamental representation of $SU(N)$ and the covariant derivative is $D_\mu = \partial_\mu - iA_\mu$. We choose the convention where the Lie algebra generators T^a are Hermitian and $[T^a, T^b] = if^{abc}T^c$, with real structure constants f^{abc} . Furthermore, the generators are such that $\text{tr}(T^a T^b) = \frac{1}{2}\delta^{ab}$. The strength-field tensor $F_{\mu\nu}$ is given by

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu - i[A_\mu, A_\nu] \quad (1.2)$$

and transforms in the adjoint representation of $SU(N)$. Both the gauge field A_μ and the curvature tensor $F_{\mu\nu}$ live in the Lie algebra $\mathfrak{su}(N)$.

Under a gauge transformation given by a group-valued function $g(x) \in SU(N)$, such that $\psi \mapsto g(x)\psi(x)$, then

$$A_\mu(x) \mapsto g(x)A_\mu(x)g(x)^{-1} + ig(x)\partial_\mu g(x)^{-1}, \quad (1.3)$$

so that $D_\mu\psi(x) \mapsto g(x)D_\mu\psi(x)$, while

$$F_{\mu\nu} \mapsto g(x)F_{\mu\nu}g(x)^{-1}, \quad (1.4)$$

leaving the action invariant.

The action of the theory in $d + 1$ dimensions is the given by

$$S[A, \psi, \bar{\psi}] = \int d^{d+1} \mathcal{L} \quad (1.5)$$

and the path integral

$$Z = \int \mathcal{D}A \mathcal{D}\bar{\psi} \mathcal{D}\psi e^{iS[A, \psi, \bar{\psi}]}. \quad (1.6)$$

1.1.1 EUCLIDEAN FIELD THEORY

In order to work in a Euclidean space-time, we need first to perform a *Wick rotation*, where the time coordinate x_0 is mapped a forth space coordinate x_4 , through $x_0 = -ix_4$. This has the effect of changing the path-integral integrand from e^{iS} , which is oscillatory, to e^{-S} , which is positive and can be interpreted as a probability distribution of the configurations of the fields.

The Euclidean path integral is

$$Z_E = \int \mathcal{D}A \mathcal{D}\bar{\psi} \mathcal{D}\psi e^{-S}, \quad (1.7)$$

so that the Minkowski action and the Euclidean action satisfy $iS_M = -S_E$. The respective actions are given by

$$S_M = \int d^{d+1} x_M \mathcal{L}_M, \quad S_E = \int d^{d+1} x_E \mathcal{L}_E. \quad (1.8)$$

The rotation $x_0 = -ix_4$ leads to $\mathcal{L}_E = -\mathcal{L}_M$.

The Wick rotation does not change the form of the gauge kinetic term, i.e.,

$$-\frac{1}{2g^2} \text{tr}(F_{\mu\nu} F^{\mu\nu}). \quad (1.9)$$

The sum is a simple Euclidean sum, where there are no minus signs when raising or lowering indices and $\mu = 1, \dots, d + 1$.

Considering now the fermionic part of the Yang-Mills Lagrangian, we need to perform the Wick rotation on the Dirac operator. In Minkowski space-time

$$\bar{\psi}(i\gamma_M^\mu D_\mu - m)\psi = \bar{\psi}(i\gamma_M^\mu \partial_\mu + \gamma_M^\mu A_\mu - m)\psi, \quad (1.10)$$

where γ_M^μ are the gamma matrices of the Clifford algebra, and they satisfy $\{\gamma_M^\mu, \gamma_M^\nu\} = 2\eta^{\mu\nu}$. In the Euclidean Clifford algebra instead the gamma matrices γ_E^μ satisfy $\{\gamma_E^\mu, \gamma_E^\nu\} = 2\delta^{\mu\nu}$. Given the fact that we have $\partial_0 = i\partial_4$ and $A_0 = iA_4$, in order to obtain the correct form we have to put $\gamma_M^0 = \gamma_E^4$. This procedure yields

$$\bar{\psi}(i\gamma_M^\mu \partial_\mu + \gamma_M^\mu A_\mu - m)\psi = -\bar{\psi}(\gamma_E^\mu \partial_\mu + i\gamma_E^\mu A_\mu + m)\psi. \quad (1.11)$$

Since $\mathcal{L}_E = -\mathcal{L}_M$, we finally arrive at

$$\mathcal{L}_E = \frac{1}{2g^2} \text{tr}(F_{\mu\nu} F^{\mu\nu}) + \bar{\psi}(\gamma^\mu D_\mu)\psi, \quad (1.12)$$

where the indices are all Euclidean and $D_\mu = \partial_\mu + iA_\mu$.

1.1.2 HAMILTONIAN FORMULATION

The Hamiltonian formulation of a Yang-Mills theory can be tricky, especially the part about the gauge field. Usually, one has to proceed by computing the conjugate momenta and performing a Legendre transform. The main issue here is that the gauge field component A_0 , does not have a conjugate momentum:

$$\frac{\partial \mathcal{L}}{\partial \dot{A}_0} = 0. \quad (1.13)$$

Hence, the transformation is not invertible. The easiest way to remedy to the situation is to impose the gauge condition $A_0 = 0$, which is called *canonical gauge* or *temporal gauge*. With this condition, the kinetic term for the gauge fields can be written as

$$\mathcal{L} = -\frac{1}{2g^2} \text{tr}(F_{\mu\nu} F^{\mu\nu}) = \frac{1}{g^2} (\mathbf{E}^2 - \mathbf{B}^2) = \frac{1}{g^2} (E_i^a E_i^a - B_i^a B_i^a), \quad (1.14)$$

where \mathbf{E} and \mathbf{B} are, respectively, the “chromoelectric” and the “chromomagnetic” fields. In the temporal gauge $\mathbf{E} = \dot{\mathbf{A}}$, the time derivative of the spatial components \mathbf{A} of the gauge field A_μ , while \mathbf{B} corresponds to the spatial components of the strength-field tensor $F^{\mu\nu}$ and does not involve any time derivative. From the Legendre transformation of (1.14) we obtain the Hamiltonian density:

$$\mathcal{H} = \frac{1}{g^2} E_i^a \dot{A}_i^a - \frac{1}{2g^2} (E_i^a E_i^a - B_i^a B_i^a) = \frac{1}{2g^2} \text{tr}(\mathbf{E}^2 + \mathbf{B}^2), \quad (1.15)$$

hence the Hamiltonian in d spatial dimensions is

$$H = \int d^d x \frac{1}{2g^2} \text{tr}(\mathbf{E}^2 + \mathbf{B}^2). \quad (1.16)$$

In the Hamiltonian formulation, the fields \mathbf{A} and \mathbf{E} are now operators, satisfying the commutation relations

$$\begin{aligned} [A_i^a(x), E_j^b(y)] &= ig^2 \delta_{ij} \delta_{ab} \delta(x - y) \\ [E_i^a(x), E_j^b(y)] &= [A_i^a(x), A_j^b(y)] = 0. \end{aligned} \quad (1.17)$$

The \mathbf{B} operator is defined from \mathbf{A} .

In order to impose the canonical gauge $A_0 = 0$, the equation of motion for A_0 has to be satisfied. This leads [\[citation?\]](#) to the fact that one must have

$$D_i E_i = 0, \quad (1.18)$$

where D_i are the spatial components of the covariant derivative and E_i the spatial components of the chromoelectric field. Unfortunately, the equation above is inconsistent with the commutation relation (1.17) and so it cannot be implemented as an

operator equation. The easiest solution is to impose it on states that are considered *physical*:

$$D_i E_i |\psi_{\text{phys}}\rangle = 0, \quad (1.19)$$

for each component of $D_i E_i$. The constraints select a subspace of the overall Hilbert space, which will be labeled as the *physical Hilbert space*. The condition above for a $U(1)$ theory reduces to the well known *Gauss law* $\nabla \cdot \mathbf{E} = 0$.

1.1.3 THE SIGN PROBLEM

Even though interesting phases have been predicted for QCD in the μ - T plane^[citation?], such as quark-gluon plasma^[citation?] or superconductivity^[citation?], detailed quantitative analysis has been limited to the $\mu = 0$ region only. This is due mainly to the difficulty of studying QCD in the low energy regime, where the perturbative approach fails^[citation?]. Moreover, even *lattice gauge theories*, a non-perturbative approach to Yang-Mills theory (and gauge theories in general), is not applicable for $\mu \neq 0$ due to the *sign problem*.

In the Hamiltonian formulation, the chemical potential in the same manner as statistical mechanics. If \hat{H} is the Hamiltonian operator and \hat{N} a number operator, then one can simply replace \hat{H} with $\hat{H} - \mu \hat{N}$. In the case of a Yang-Mills theory, the number operator corresponds to fermion number, i.e., $\hat{N} = \psi^\dagger \psi$.

In the path integral formalism, fermions are Grassmann variables, which we can integrate over:

$$\int \mathcal{D}\psi \mathcal{D}\bar{\psi} \exp\left(-\int d^{d+1}\bar{\psi} K \psi\right) = \det K, \quad (1.20)$$

where K is the kinetic operator for the fermions. If the fermions are coupled to gauge field, as one would expect from a Yang-Mills theory, then K has some complicated dependence on the fields A_μ . If one includes the chemical potential term $\mu \bar{\psi}^\dagger \psi$ in the Lagrangian, then the fermion determinant $\det K$ turns out to be complex^[citation?], with a non-trivial phase factor.

As a result, the integrand of the path-integral is no longer positive, and it cannot be interpreted as a probability distribution. This is the infamous *sign problem* and poses severe limitation to Monte Carlo simulations in the finite μ region.

1.2 LATTICE FIELD THEORY

Starting from the path integral formulation, the first step in the formulation of a *lattice field theory* (LFT) is the discretization of space-time, where a discrete $d + 1$ -dimensional lattice substitutes the continuum space-time. The simplest choice in this regard is a hypercubic lattice with lattice spacing a , but in theory an LFT can be defined on any type of lattice. An immediate advantage of using a lattice instead of a continuum is the natural ultraviolet cutoff given by the inverse of the lattice spacing.

Formally a lattice Λ is defined as

$$\Gamma = \left\{ x \in \mathbb{R}^4 : x = \sum_{\mu=1}^{d+1} a n_{\mu} \hat{\mu} \quad n_{\mu} \in \mathbb{Z} \right\}, \quad (1.21)$$

where $\mu = 1, \dots, d+1$ and $\hat{\mu}$ is the unit vector in the direction μ . The edges will be labeled by a pair $(x, \hat{\mu})$, meaning that we are referring to the edge in the $\hat{\mu}$ direction from the vertex x . It is important to fix an orientation for each direction in the lattice. The most natural choice is to choose $+\hat{\mu}$ for each $\hat{\mu}$. So, even though $(x, \hat{\mu})$ and $(x + \hat{\mu}, -\hat{\mu})$ refers to the same link, the former is traversed in the positive direction while the latter in the negative direction.

In an LFT, both the vertices and edges (also called links) host degrees of freedom (d.o.f). In particular, the matter fields lives on the vertices while the gauge fields live on the links between vertices. However, the definition of these d.o.f. will need some care, because we have two main requirements, especially if we are interested in Yang-Mills theory:

- The lattice action should reduce to the continuum action in the continuum limit, i.e., $a \rightarrow 0$;
- The lattice action should respect the gauge symmetry.

Lorentz invariance is naturally broken on a lattice but we expect to recover it in the continuum limit.

1.2.1 GAUGE FIELDS ON A LATTICE

[spiegazione perché usiamo il gruppo di Lie e non l'algebra di Lie] Considering a general group G , we associate an element $U_{\mu}(x) \in G$ to each link (x, μ) . If one traverse the link in the opposite direction, one should obtain the inverse element U^{-1} . In the case of $SU(N)$, we take $U_{\mu}(x)$ to be the matrices in the fundamental representation. We can obtain a vector potential in the continuum limit by writing

$$U_{\mu}(x) = e^{iaA_{\mu}(x)}, \quad (1.22)$$

where a is the lattice spacing.

It is necessary to discuss about gauge invariance before moving to the dynamics of these gauge fields. A gauge transformation is described by a group-valued function $g(x)$ (in the appropriate representation), which acts on the vertices x . The variable $U_{\mu}(x)$ sits in the middle of the site x and $x + \hat{\mu}$, then it is reasonable to think its transformation is

$$U_{\mu}(x) \mapsto g(x)U_{\mu}(x)g(x + a\hat{\mu})^{\dagger}. \quad (1.23)$$

In order to introduce a dynamics for the gauge fields $U_{\mu}(x)$, we need to define their action which need two satisfy two requirements: it has to be gauge-invariant and reduce to the pure gauge Yang-Mills action in the continuum limit. From (1.23),

we can immediately deduce that taking the product of $U_\mu(x)$ along a closed curve will yield a gauge-invariant quantity. The simplest close curve we can consider is a *plaquette*, i.e., the smallest square face. Hence, on a plaquette \square we introduce W_\square to put in the action, defined as

$$W_\square = U_\mu(x)U_\nu(x + a\hat{\mu})U_\mu(x + a\hat{\nu})^\dagger U_\nu(x)^\dagger. \quad (1.24)$$

Notice that we do not have any sum in the indices μ and ν because they are not Lorentz indices. The quantity in (1.24) is called a single plaquette *Wilson loop*.

We can only have scalar quantities in the action, so we need to take the trace of W_\square . Then, our lattice action will be defined as the sum over the plaquettes of $\text{tr } W_\square$ (and its Hermitian conjugate):

$$\mathcal{S} = -\frac{1}{g^2} \sum_{\square} \left(\text{tr } W_\square + \text{tr } W_\square^\dagger \right). \quad (1.25)$$

This is known as the *Wilson action* [\[citation?\]](#). The quantity $\text{tr } W_\square$ behaves as expected in the continuum limit, where we have to work with the strength field $F^{\mu\nu}$:

$$\text{tr } W_\square \approx N - \frac{a^4}{2} \text{tr } F_{\mu\nu} F^{\mu\nu} + \mathcal{O}(a^6), \quad (1.26)$$

The lattice action is not unique. The Wilson action in (1.25) is the simplest choice that one can make that satisfy our requirement. Some other modification, for example, can include other types of closed loops and these modification can have their place. However, they will not be considered here.

Obviously, in a path-integral formulation of lattice gauge theories we need to define the path integral, which is immediate. The partition function is given by

$$Z = \int \prod_{(x,\hat{\mu})} dU_\mu(x) e^{-\mathcal{S}}, \quad (1.27)$$

where the integration measure $dU_\mu(x)$ is understood to be the Haar measure. In case of a compact group, like $SU(N)$, it is well defined and yields a finite value. Now that the path integral measure has been defined, the average of an observable \mathcal{O} can be computed as

$$\langle \mathcal{O} \rangle = \frac{1}{Z} \int \prod_{(x,\hat{\mu})} dU_\mu(x) \mathcal{O} e^{-\mathcal{S}} \quad (1.28)$$

1.2.2 ORDER PARAMETERS AND GAUGE INVARIANCE

The Wilson formulation of lattice gauge theories can resemble spin models studied in statistical mechanics. The link variables $U_\mu(x)$ can be thought as some sort of generalization of the spin degrees of freedom. They are distributed in a crystal-like structure and interact with their nearest neighbours, in this case through a four-body interaction (the plaquette action), instead of two-body interaction (like the

Ising model). If one wants to pursue this analogy, then it is reasonable to look at order parameters that behaves like the spontaneous magnetization, where a non-vanishing expectation value signals a phase transition. The analogue of such an order parameter in lattice gauge theory would be something like

$$\langle U_\mu(x) \rangle \neq 0, \quad (1.29)$$

but it has been shown^[citation?] [teorema di Elitzur] that this is impossible in Wilson theory.

In standard spin models, a non-zero magnetization represents a spontaneous breaking of the global symmetry of the system. Consider the simplest case of the classical Ising model, where the degrees of freedom are binary variables $\sigma = \pm 1$. Without an external field, the energy is given by the interaction of nearest neighbouring spins, i.e., $\sigma_i \sigma_j$. This system has an obvious global \mathbb{Z}_2 symmetry, that corresponds to the inversion $\sigma_i \mapsto -\sigma_i$ off all the spins. A ferromagnetic phase is, by definition, signaled by $\langle \sigma \rangle \neq 0$, which necessarily breaks the global \mathbb{Z}_2 symmetry of the model. Once a direction is selected by $\langle \sigma \rangle \neq 0$, it remains stable under thermal fluctuations because they cannot coherently shift of the magnetization of an infinite (or rather large) number of spins.

In a lattice gauge theory, an expectation value like (1.29) would *breaks* gauge invariance, which is a *local symmetry*, not a global one. As explained previously, gauge invariance means that the action is unchanged under local arbitrary “rotations” of the link variables $U_\mu(x)$, see (1.23). Hence, thermal fluctuations will induce such rotations and in the long run it will average on all the possible gauges. This leads to

$$\langle U_\mu \rangle = \int dU_\mu U_\mu = 0 \quad (1.30)$$

if U_μ contains only non-trivial irreducible representations of the group. This means that “magnetization” is always vanishing in a lattice gauge theory and gauge invariance cannot be spontaneously broken, which is the contents of the Elitzur theorem^[citation?].

The conclusion of this brief discussion may seem rather grim, as magnetization in spin models is the most convenient and used order parameter. But this does not mean that there are no other order parameters in a lattice gauge theory. We have just showed that the problem when considering something like $\langle U_\mu \rangle$ is gauge invariance. So, the most reasonable step forward is to consider *gauge-invariant quantities*. We have already seen that tracing over a product of U_μ variables along a closed curve is a gauge-invariant quantity, called *Wilson loop*.

In so far, we have considered only single plaquette loops but nothing restrains us from considering arbitrary large loops, indeed it serves as a *confinement test* for pure gauge theories. It has been shown^[citation?] that confinement is equivalent to the *area law* behaviour of Wilson loops, i.e.,

$$\langle W(C) \rangle \sim \exp(-\sigma A(C)), \quad (1.31)$$

where $A(\mathcal{C})$ is the minimal area inside the closed path \mathcal{C} and σ the *string tension* (the coefficient of the linear potential between two quarks). On the other hand, in the absence of confinement one finds instead the *perimeter law*

$$\langle W(\mathcal{C}) \rangle \sim \exp(-kP(\mathcal{C})), \quad (1.32)$$

where $P(\mathcal{C})$ is the perimeter of the curve \mathcal{C} and k just some constant.

The reason behind this behaviour can be seen with a simple qualitative picture. A closed timelike Wilson loop basically represents a process in which a quark-antiquark pair is produced, moved along the sides of the loop and annihilated. If we are in a confining phase we can then expect a linear potential between the quark and antiquark. We can imagine a flux tube *binding* the two charges, which swoops the whole inside the loop. Then, it is easy to image that the energy of this whole process will necessarily depend on the area of the loop. On the other hand, if we are in a deconfined phase then there is no potential binding the two quarks. In this case the energy of the whole process depends only on the self-energy of quarks, which move along the sides of the loop. Therefore, the leading energy contribution of this process depends on the perimeter, instead of the area. Obviously, this picture is no longer valid when dynamical matter is involved. In a confining phase, pair production is always preferred when separating two quarks at large distances. [[inserire immagine](#)]

From (1.31) and (1.32), we can deduce that the string tension σ can be used as an order parameter. It is non-zero for a confining phase, while it vanishes for a deconfined phase. But it is *non-local* in nature, as it involves the asymptotic behaviour of potential, and therefore of the correlation functions of the theory.

1.2.3 FERMIONS ON A LATTICE

Defining fermions is not an easy task due to the known *doubling problem*. In simple terms, when introducing fermions on a lattice, instead of a continuous space, it leads to a extra spurious fermions, which are just lattice artifacts.

In order to briefly see this, consider the correlation function for a single fermionic species. If K is the kinetic matrix for the fermions, then $G = K^{-1}$ gives their correlation matrix. One finds [\[citation?\]](#) that the correlation function between two sites x and y has the form

$$(G)_{x,y} = \frac{1}{a^d L^d} \sum_k \tilde{G}_k e^{2\pi i k \cdot (x-y)/L}, \quad (1.33)$$

where a is the lattice spacing, L^d the total volume and \tilde{G}_k the correlation function in momentum space:

$$\tilde{G}_k^{-1} = m + \frac{i}{a} \sum_{\mu} \gamma_{\mu} \sin(2\pi k_{\mu}/L). \quad (1.34)$$

It involves a trigonometric function because the derivative term involves nearest neighbouring sites. One can then take the model to a large lattice, which justifies in

substituting the discrete sums with integrals:

$$\frac{2\pi k_\mu}{La} \rightarrow q_\mu \quad \text{and} \quad \frac{1}{a^d L^d} \sum_k \rightarrow \int \frac{d^d q}{(2\pi)^d}, \quad (1.35)$$

where the q_μ 's are continuous momentum variables. This substitution maps (1.34) into

$$\tilde{G}_k^{-1} = m + \frac{i}{a} \sum_\mu \gamma_\mu \sin(aq_\mu). \quad (1.36)$$

One can naively think of taking the limit $a \rightarrow 0$ and expand $\sin(aq_\mu)$ around the zero and obtain something that look like the correct continuum limit:

$$\tilde{G}_\mu^{-1} = m + i\not{q} + \mathcal{O}(a^2). \quad (1.37)$$

But one should not be fooled by this sloppy procedure just because it appears to give the wanted result. Each component q_μ takes values in the region $[-\pi/a, +\pi/a]$, hence we have to integrate on the whole volume $[-\pi/a, +\pi/a]^d$. Looking at (1.36), it is clear that the major contributions to G in (1.33) comes from the zeros of \tilde{G}_k^{-1} . This, not only vanishes in the region $q_\mu \sim 0$ but also for large momentum $q_\mu \sim \pi/a$. The propagator has no suppression of momentum values near π/a . We can isolate the large momenta region by considering

$$\tilde{q}_\mu = q_\mu - \pi/a \quad (1.38)$$

for each direction in space. In this way, we de facto half the integration region,

$$\int_{-\pi/a}^{\pi/a} dq_\mu \rightarrow \int_{-\pi/2a}^{\pi/2a} (dq_\mu + d\tilde{q}_\mu) \quad (1.39)$$

and now the limit $a \rightarrow 0$ can be taken safely, but it comes with a price to pay. For each direction in space, we have two independent regions that gives a free fermion contribution to the propagator in the continuum limit. We have effectively *doubled* the number of fermions for each direction. In a d -dimensional lattice we end up with 2^d independent fermions, even though we initially started with just one. [qualcosa a che fare con la chiralità]

There are many solutions to this fermion doubling problem^[citation?], but we will focus only on one in this manuscript: *the staggered fermions*^[citation?]. We have seen that these fictitious fermions come from the large momenta regions, where $q_\mu \sim \pi/a$. Brutally cutting out this large momenta regions spoils the completeness of the Fourier transform, so it is not a solution, but a smarter solution can give out the same effect. The idea is to spread the fermionic degrees of freedom over multiple lattice sites, reducing effectively the momenta space. For example, in two dimensions it would correspond to placing the *particles* on *even* sites and the *antiparticles* on *odd* sites. A site is considered even or odd when $(-1)^x = (-1)^{x_1 + \dots + x_d} = +1$ or -1 .

To obtain a staggered fermion, we define a new fermionic species $\chi(x)$ such that

$$\psi(x) = \prod_{\mu} (\gamma^{\mu})^{n_{\mu}} \chi(x), \quad (1.40)$$

where $x_{\mu} = an_{\mu}$. Now, if we want to express the discretized covariant derivative, the term $\gamma^{\mu}\psi(x + an_{\mu})$ have two extra powers of γ^{μ} compared to $\bar{\psi}$. Since $(\gamma^{\mu})^2 = \pm 1$, we have therefore

$$\bar{\psi}(x)\gamma^{\mu}\psi(x) = (-1)^{\eta_{\mu}(x)}\chi(x)^{\dagger}\chi(x + an_{\mu}), \quad (1.41)$$

where $\eta_{\mu}(x)$ is some sign function depending on the site x . This function can be obtained from the commutation relations of the gamma matrices. In particular, in two dimensions we have

$$\eta_1(x) = 1 \quad \text{and} \quad \eta_2(x) = (-1)^{n_1}, \quad (1.42)$$

while in four (Euclidean) dimension we have instead [[Citare Tong, gauge theories](#)]

$$\eta_1(x) = 1, \quad \eta_2(x) = (-1)^{n_1}, \quad \eta_3(x) = (-1)^{n_1+n_2}, \quad \eta_4(x) = (-1)^{n_1+n_2+n_3}. \quad (1.43)$$

A similar reasoning can be applied to the mass term $m\bar{\psi}(x)\psi(x)$, where it becomes

$$m\bar{\psi}(x)\psi(x) = (-1)^{\eta(x)}\chi(x)^{\dagger}\chi(x), \quad (1.44)$$

for some sign function $\eta(x)$ that can be obtained from the commutation relations of the gamma functions.

CHAPTER TWO

LATTICE GAUGE THEORIES FOR QUANTUM SIMULATION

2.1 QUANTUM SIMULATION

2.2 FINITE-GROUP APPROACH

CHAPTER THREE

DUALITIES IN LATTICE GAUGE THEORIES

- 3.1 BOND-ALGEBRA APPROACH TO DUALITIES**
- 3.2 GAUGE-REDUCING DUALITIES**
- 3.3 DUALITIES OF TWO-DIMENSIONAL PURE LGTS**
- 3.4 DUALITIES OF LADDER LGTS**