NOTES FROM NSF-MSGI INTERNSHIP

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1. Introduction

We will explore the connection between statistical mechanics models and quantum Hamiltonians, primarily via examples.

More stuff here later. Basic objects/definitions. Mass gap/correlation length?

2. Ising Models

The Ising models are a family of statistical mechanics models with nearest-neighbor interaction. Given any lattice $\mathcal{L} \subset (a_i\mathbb{Z})^d$ with N total points for some scaling factors $a_i \in \mathbb{R}_{>0}$, define a *spin* at each site of the lattice via a variable $s \in \Omega_0 = \{\pm 1\}$. A *configuration* is a collection $\{s\} = \{s_\ell\}_{\ell \in \mathcal{L}} \in \Omega_0^N$. For convenience, we will usually work with cubic lattices of the form $\mathcal{L} = (a_i\{-n, -n+1, \ldots, n-1, n\})^d$ or $\mathcal{L}(= a_i\{1 \ldots, n\})^d$ and impose periodic boundary conditions. The *action* associated to this model is the function

$$\mathcal{S}(\{s\}, \tau, h) = -\beta_\tau(\tau) \sum_{i \sim j} s_i s_j - \beta(\tau) h \sum_i s_i,$$

where τ is the lattice spacing in the "time" direction, h is the strength of an external magnetic field, and the first sum is over nearest neighbor sites of the lattice. The coefficients β_τ and β may depend on the time direction lattice spacing and are otherwise chosen to reflect some physical situation (e.g. they may be taken to be proportional to the "inverse temperature" $\frac{1}{kT}$ where k is Boltzmann's constant).

We will be interested in the *partition function* Z defined as

$$Z(\tau, h) = \sum_{\{s\} \in \Omega_0^N} e^{-\mathcal{S}(\{s\}, \tau, h)}.$$

It turns out that we can use statistical mechanics to show that many other interesting physical quantities can be derived from the partition function [4, Ch. 3]. The *pressure in* $\mathcal{L}$ of the model is defined to be

$$\psi_\mathcal{L}(\tau, h) := \frac{1}{N} \log Z(\tau, h).$$

The *magnetization density in* $\mathcal{L}$ is by definition

$$m_\mathcal{L} := \frac{1}{N} \sum_{\ell \in \mathcal{L}} s_\ell.$$

The expected magnetization density, $\langle m_\mathcal{L} \rangle$, is related to the pressure and partition function:

$$\langle m_\mathcal{L} \rangle(\tau, h) = \frac{1}{\beta} \frac{\partial \psi_\mathcal{L}}{\partial h}(\tau, h) = \frac{\partial}{\partial h}\left(\frac{1}{\beta N} \log Z(\tau, h)\right).$$

2.1. Infinite Lattice Ising Models. The Ising model can be generalized to an *infinite volume model* by letting the lattice $\mathcal{L}$ grow to $\mathbb{Z}^d$ appropriately. We'll briefly sketch the ideas here - a detailed exposition can be found in [4] in chapters 3 and 6. In general, we say a sequence of lattices $\{\mathcal{L}_n\}$ converges to $\mathbb{Z}^d$, denoted $\mathcal{L}_n \Uparrow \mathbb{Z}^d$, if

(1) $\mathcal{L}_n$ is increasing, i.e. $\mathcal{L}_n \subset \mathcal{L}_{n+1}$ for all n,
(2) $\bigcup_n \mathcal{L}_n = \mathbb{Z}^d$,
(3) $\lim\limits_{n\to\infty} \frac{|\partial \mathcal{L}_n|}{|\mathcal{L}_n|} = 0$, where $\partial \mathcal{L}_n$ is the boundary of the lattice, i.e.

$$\partial \mathcal{L}_n := \{\ell \in \mathcal{L}_n \mid \ell \sim m \text{ for some } m \in \mathbb{Z}^d \setminus \mathcal{L}_n\}.$$

It turns out that there is a well-defined limit

$$\psi(\tau, h) := \lim_{\mathcal{L} \Uparrow \mathbb{Z}^d} \psi_{\mathcal{L}}(\tau, h).$$

The resulting function ψ is convex as a function from $\mathbb{R} \times \mathbb{R}_{>0}$ and even as a function of h. Because ψ is a convex function of h, the *average magnetization density* given by

$$m(\tau, h) = \lim_{\mathcal{L} \Uparrow \mathbb{Z}^d} \langle m_{\mathcal{L}} \rangle(\tau, h),$$

exists for all but a countable set of $h \in \mathbb{R}$. The points at which the average magnetization density fails to exist (always because the left and right derivatives do not equal each other) are called *first-order phase transitions.*

2.2. 1+0 Ising. In the 1+0 Ising model, we take a 1-dimensional statistical mechanics system and relate it to a 0-dimensional (point) quantum Hamiltonian. See [3, 10]. The definitions above simplify to the following. The lattice is now a set of N points $x_0, \ldots, x_{N-1}$ on a circle so that $s_0 = s_N$. The action is given by

$$\mathcal{S}(\{s\}, \tau, h) = -\beta_\tau(\tau) \sum_{i=0}^{N-1} s_i s_{i+1} - \beta(\tau)h \sum_{i=0}^{N-1} s_i.$$

It will be helpful to rewrite this expression to be symmetric and expressed in terms of sums and differences of the s_i. Up to a constant factor of $-\beta_\tau N$, we find

$$\mathcal{S}(\{s\}, \tau) = \frac{1}{2}\beta_\tau(\tau) \sum_{i=0}^{N-1} (s_i - s_{i+1})^2 - \frac{1}{2}\beta(\tau)h \sum_{i=0}^{N-1} (s_i + s_{i+1}).$$

Note that we needed periodicity in order to rewrite the second term in this way. Further, we now have $\mathcal{S} = 0$ when all of the spins are identical and $h = 0$.

Now the partition function becomes

$$Z = \sum_{\{s\} \in \Omega_0^N} e^{-\frac{1}{2}\beta_\tau(\tau) \sum_{i=0}^{N-1}(s_i - s_{i+1})^2 + \frac{1}{2}\beta(\tau)h \sum_{i=0}^{N-1}(s_i + s_{i+1})}$$

$$= \sum_{s_0 \in \{\pm 1\}} \cdots \sum_{s_{N-1} \in \{\pm 1\}} e^{-\frac{1}{2}\beta_\tau(s_0 - s_1)^2} e^{\frac{1}{2}\beta h(s_0 + s_1)} \cdot \ldots \cdot e^{-\frac{1}{2}\beta_\tau(s_{N-1} - s_0)^2} e^{\frac{1}{2}\beta h(s_{N-1} + s_0)}.$$

However, each factor in the product depends only on what the values of $s_i - s_{i+1}$ and $s_i + s_{i+1}$ are. We can construct the *transfer matrix*, T, to record this data, indexed by the possible spins at each site. So we have

$$T_{-1,-1} = e^{-\beta h}, \quad T_{-1,1} = T_{1,-1} = e^{-2\beta_\tau}, \quad \text{and } T_{1,1} = e^{\beta h}.$$

Thus

$$T = \begin{bmatrix} e^{\beta h} & e^{-2\beta_\tau} \\ e^{-2\beta_\tau} & e^{-\beta h} \end{bmatrix}.$$

Now we can write $Z = \operatorname{Tr} T^N$. From this result, we can derive the $N \to \infty$ and $\tau \to 0$ limits of the quantities of interest above, in particular the average magnetization. Since T is diagonalizable, we can write $T = SDS^{-1}$ and so

$$\operatorname{Tr} T^N = \operatorname{Tr}(SD^N S^{-1}) = \operatorname{Tr}(D^N) = \lambda_0^N + \lambda_1^N,$$

where λ_0, λ_1 are the eigenvalues of T.

Via a bit of algebra, we see that

$$\lambda_0 = \cosh(\beta h) + \sqrt{e^{-4\beta_\tau} + \sinh^2(\beta h)}$$

$$\lambda_1 = \cosh(\beta h) - \sqrt{e^{-4\beta_\tau} + \sinh^2(\beta h)}$$

Later, we will examine the Hamiltonian when $\beta_\tau = -\frac{1}{2}\log\tau$ and $\beta = \tau$. In this case,

$$\lambda_0 = \cosh(\tau h) + \sqrt{\tau^2 + \sinh^2(\tau h)}$$

$$\lambda_1 = \cosh(\tau h) - \sqrt{\tau^2 + \sinh^2(\tau h)}$$

Now the average magnetization is

$$m(\tau, h) = \lim_{N \to \infty} \frac{\partial}{\partial h}\left(\frac{1}{\beta N}\log(\lambda_0^N + \lambda_1^N)\right)$$

$$= \frac{\partial}{\partial h}\lim_{N \to \infty}\left(\frac{1}{\beta N}\left(\log(\lambda_0^N) + \log\left(1 + \left(\frac{\lambda_1}{\lambda_0}\right)^N\right)\right)\right)$$

$$= \frac{1}{\beta}\frac{\partial}{\partial h}\log(\lambda_0),$$

since $|\lambda_1/\lambda_0| < 1$. In our example case,

$$m(\tau, h) = \frac{1/\tau}{\cosh(\tau h) + \sqrt{\tau^2 + \sinh^2(\tau h)}}\left(\tau\sinh(\tau h) + \frac{\tau\sinh(\tau h)\cosh(\tau h)}{\sqrt{\tau^2 + \sinh^2(\tau h)}}\right)$$

$$= \frac{\sinh(\tau h)}{\sqrt{\tau^2 + \sinh(\tau h)}}.$$

We can now compute the limit of the expected average magnetization as we take the lattice distance τ to 0.

$$\lim_{\tau \to 0} m(\tau, h) = \lim_{\tau \to 0}\frac{\sinh(\tau h)}{\sqrt{\tau^2 + \sinh(\tau h)}} = \frac{h}{\sqrt{1 + h^2}}.$$

From these computations we can see that $m(\tau, h)$ is a smooth function of h and so the 1+0 Ising model does not have a (first-order) phase transition.

2.2.1. *Deriving the quantum Hamiltonian:* Our goal is to choose functions $\beta_\tau(\tau)$ and $\beta(\tau)$ so that the transfer matrix operator T has the form $T = e^{-\tau H} \approx I - \tau H$ for some quantum Hamiltonian H (independent of τ) acting on a 2-dimensional vector space. This is the τ-*continuum Hamiltonian* for the model. Here the idea is that statistical mechanics properties of the lattice action system (e.g. magnetization per site, average magnetization, two-point correlations, correlation length) will map to properties of the operator H related to its eigenvalues and eigenvectors.

For example, we could choose $\beta_\tau(\tau) = -\frac{1}{2}\log\tau$ and $\beta(\tau) = \tau$. Then we have

$$T = \begin{bmatrix} e^{\tau h} & \tau \\ \tau & e^{-\tau h} \end{bmatrix} \approx I_2 - \tau\begin{bmatrix} -h & -1 \\ -1 & h \end{bmatrix}.$$

From this it follows that

$$H = \begin{bmatrix} -h & -1 \\ -1 & h \end{bmatrix} = -\sigma_1 - h\sigma_3,$$

where σ_1, σ_2, and σ_3 are the Pauli matrices.

More generally, we can ask what constraints β and β_τ must satisfy to be able to do this. Analyzing each of the four matrix entries shows that for small τ we need

$$e^{\beta h} \approx 1 - \tau H_{0,0}, \quad e^{-2\beta_\tau} \approx -\tau H_{1,0} = -\tau H_{0,1}, \quad \text{and } e^{-\beta h} \approx 1 - \tau H_{1,1}.$$

Since $H_{1,0}$ and $H_{0,1}$ are constant with respect to τ we need $-2\beta_\tau \approx \log(\lambda\tau)$ for small τ and some $\lambda \in \mathbb{R}_{>0}$. Similarly, if β is small when τ is small we have $e^{\beta h} \approx 1 - \beta h \approx 1 - \tau H_{0,0}$ and $e^{-\beta h} \approx 1 + \beta h \approx 1 - \tau H_{1,1}$. So to first order we have $\beta \approx \tau$ and we see $H_{0,0} = -H_{1,1} = -\mu h$ for some $\mu \in \mathbb{R} \setminus \{0\}$.

Putting all of this together, we see that the general τ-continuum Hamiltonian for this model will have the form

$$H = \begin{bmatrix} -\mu h & -\lambda \\ -\lambda & \mu h \end{bmatrix} = -\lambda\sigma_1 - \mu h\sigma_3.$$

Now we could also attempt to find a second order approximation of T so that $T \approx I - \tau H + \frac{1}{2}\tau^2 H^2$ for an operator H not depending on τ. However, following the approach above will show us that there is no easy solution in this case. To get a higher order approximation (or the actual matrix logarithm solving $T = e^{-\tau H}$ for H), we may need to allow H to depend on τ.

In the second order case, we need

$$e^{\beta h} \approx 1 - \tau H_{0,0} + \frac{\tau^2}{2}(H_{0,0}^2 + H_{0,1}^2),$$

$$e^{-2\beta_\tau} \approx -\tau H_{0,1} + \frac{\tau^2}{2}(H_{0,0}H_{0,1} + H_{0,1}H_{1,1}),$$

$$e^{-\beta h} \approx 1 - \tau H_{1,1} + \frac{\tau^2}{2}(H_{0,1}^2 + H_{0,0}^2)$$

for small τ, where we have used the fact that H should be real and symmetric and so $H_{0,1} = H_{1,0}$.

From the first and third equations, we see that $H_{0,0}$ and $H_{1,1}$ must depend on h. But the second equation has no dependence on h, so the dependence of $H_{0,1}$ on h must be inversely related to the $H_{0,0}$ and $H_{1,1}$ dependence. On the other hand adding the first and third equations yields

$$\cosh(\beta h) \approx 1 - \frac{\tau}{2}(H_{0,0} + H_{1,1}) + \frac{\tau^2}{4}(H_{0,0}^2 + 2H_{0,1}^2 + H_{1,1}^2),$$

which naturally suggests taking $\beta = \mu\tau$, $H_{0,0} = -H_{1,1} = \mu h$, and $H_{0,1} = H1,0 = 0$. But this is a contradiction.

Let's return to the running example where our Hamiltonian is now $H = -\sigma_1 - h\sigma_3$. The eigenvalues of H are $\lambda_0 = -\sqrt{1 + h^2}$ and $\lambda_1 = \sqrt{1 + h^2}$. The corresponding normalized eigenvectors are

$$v_0 = \frac{1}{\sqrt{2(1 + h^2 + h\sqrt{1 + h^2})}}\begin{pmatrix} h + \sqrt{1 + h^2} \\ 1 \end{pmatrix} \text{ and } v_1 = \frac{1}{\sqrt{2(1 + h^2 - h\sqrt{1 + h^2})}}\begin{pmatrix} h - \sqrt{1 + h^2} \\ 1 \end{pmatrix}.$$

Now the ground state is $|gs\rangle = v_0$ and if we compute $\langle gs|\sigma_3|gs\rangle$, we get

$$\langle gs|\sigma_3|gs\rangle\rangle = \frac{1}{2(1 + h^2 + h\sqrt{1 + h^2})}\begin{pmatrix} h + \sqrt{1 + h^2} & 1 \end{pmatrix}\begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}\begin{pmatrix} h + \sqrt{1 + h^2} \\ 1 \end{pmatrix} = \frac{h}{\sqrt{1 + h^2}}.$$

This demonstrates the correspondence between properties of the original statistical mechanics model and observables derived from operators applied to the ground state of the Hamiltonian.

2.3. 1+1 Ising Model. We may analyze the $1+1$ Ising model similarly to the $1+0$ model. We are going to relate a statistical mechanics system on a 2-dimensional lattice (with one temporal and one spatial dimension) to a quantum mechanical Hamiltonian on a 1-dimensional system of interacting spins. See [3, 10].

As in the 1+0 Ising model, we will work with a lattice with periodic boundary conditions. Suppose $\mathcal{L}$ is a $N_x \times N_\tau$ lattice of points in $\mathbb{Z}^2$ with $\vec{\tau}$ a unit vector in the time direction and $\vec{x}$ a unit vector in the spatial direction. Then the action for the model is

$$\mathcal{S} = -\sum_\ell \beta_\tau s_\ell s_{\ell+\vec{\tau}} + \beta s_\ell s_{\ell+\vec{x}},$$

where the sum is over all lattice points $\ell \in \mathcal{L}$.

We want to write $\mathcal{S} = \sum_{j=1}^{N_\tau} L(j, j+1)$ for some $L(j, j+1)$ that describes the interaction between the spatial rows j and $j+1$. To do this, we first rewrite

$$\mathcal{S} = \sum_\ell \frac{\beta_\tau}{2}(s_\ell - s_{\ell+\vec{\tau}})^2 - \frac{1}{2}\beta(s_\ell s_{\ell+\vec{x}} + s_{\ell+\vec{\tau}} s_{\ell+\vec{\tau}+\vec{x}}),$$

using the periodic boundary conditions. Note that the new $\mathcal{S}$ differs from the previous one by a normalization constant so that the first term is 0 when all of the spins are aligned, rather than being $-N\beta_\tau$. We can now define

$$L(j, j+1) := \sum_{\ell=1}^{N_x} \frac{\beta_\tau}{2}(s_\ell - \tilde{s}_\ell)^2 - \frac{\beta}{2}(s_\ell s_{\ell+1} + \tilde{s}_\ell \tilde{s}_{\ell+1}),$$

where the sum is over the N_x indices in the spatial row, $\{s\}$ is the configuration of row j and $\{\tilde{s}\}$ is the configuration of row $j+1$.

Then the partition function is

$$Z = \sum_{\{s\}} e^{-L(1,2)} e^{-L(2,3)} \dots e^{-L(N_\tau, 1)}.$$

As in the previous section, we express Z as the trace of the N_τ^{th} power of a transfer matrix $\hat{T}$ describing the transition between rows. Since there are 2^{N_x} configurations for each row, $\hat{T}$ will be a $2^{N_x} \times 2^{N_x}$ matrix. The elements of $\hat{T}$ can be organized by the number of spin flips between configurations, since these determine the value of the first term of $L(j, j+1)$. We want to find β, β_τ so that for τ near 0, $\hat{T} \approx 1 - \tau \hat{H}$.

$$\hat{T}|_{0\,\text{flips}} = e^{\beta \sum_{\ell=1}^{N_x} s_\ell s_{\ell+1}},$$

$$\hat{T}|_{1\,\text{flip}} = e^{-2\beta_\tau} e^{\frac{\beta}{2} \sum_{\ell=1}^{N_x}(s_\ell s_{\ell+1} + \tilde{s}_\ell \tilde{s}_{\ell+1})},$$

$$\vdots \qquad\qquad \vdots$$

$$\hat{T}|_{k\,\text{flips}} = e^{-2k\beta_\tau} e^{\frac{\beta}{2} \sum_{\ell=1}^{N_x}(s_\ell s_{\ell+1} + \tilde{s}_\ell \tilde{s}_{\ell+1})}$$

The nicest solution to this is to choose $\beta = \lambda\tau$ and $\tau = e^{-2\beta_\tau}$ for some $\lambda \in \mathbb{R}_{>0}$. Then to first order approximation in τ,

$$\hat{T}|_{0\,\text{flips}} \approx 1 - \mu\lambda\tau,$$

$$\hat{T}|_{1\,\text{flip}} \approx \tau(1 + \kappa\lambda\tau) \approx \tau,$$

$$\vdots \qquad\qquad \vdots$$

$$\hat{T}|_{k\,\text{flips}} \approx \tau^k(1 + \kappa\lambda\tau) \approx 0,$$

where $\mu = \sum_{\ell=1}^{N_x} s_\ell s_{\ell+1}$ and $\kappa = \frac{1}{2}\sum_{\ell=1}^{N_x}(s_\ell s_{\ell+1} + \tilde{s}_\ell \tilde{s}_{\ell+1})$ are independent of τ.

We can now express $\hat{T}$ as $I - \tau\hat{H}$ where the Hamiltonian $\hat{H}$ is given in terms of the Pauli operators at each site ℓ, $\hat{\sigma}_1(\ell)$ and $\hat{\sigma}_3(\ell)$:

$$\hat{H} = -\lambda \sum_{\ell=1}^{N_x} \hat{\sigma}_3(\ell)\hat{\sigma}_3(\ell+1) - \sum_{\ell=1}^{N_x} \hat{\sigma}_1(\ell). \tag{1}$$

2.3.1. *Phase transition of 1+1 Ising model.* There are a variety of ways to see that the 1+1 Ising model has a phase transition. (e.g. duality)

2.3.2. *Jordan-Wigner transform and solution in terms of Fermionic operators.* A Hamiltonian of the form of Eq. (1) can be rewritten in terms of Fermion operators $\{a_j, a_j^\dagger\}_{j=1}^n$ that satisfy the *canonical commutation relations* (CCRs)

$$\{a_j, a_k^\dagger\} = \delta_{k,j}I; \qquad \{a_j, a_k\} = 0,$$

where $\{A, B\} = AB + BA$ is the anticommutator of two operators [10, 12, 14].

We can summarize this transformation as follows:

(1) Use duality to swap the roles of $\hat{\sigma}_1$ and $\hat{\sigma}_3$ in Eq. (1).
(2) Use raising and lowering operators to rewrite H in terms of fermion operators. (Jordan-Wigner transform)
(3) Convert the resulting operators and quadratic Hamiltonian to momentum space.
(4) Diagonalize the momentum Fermionic Hamiltonian.
(5) Determine the eigenvalues of the resulting Hamiltonian.

3. O(N) Model

These models are extensively treated in [3, 7, 10]. As in the Ising model, we have a lattice $\mathcal{L} \subset \mathbb{Z}^d$ of M points, usually cubic. The spins at each site i of the lattice are now $\vec{x}_i \in \Omega_0 = S^{N-1} = \{\vec{x} \in \mathbb{R}^N \mid \|\vec{x}\| = 1\}$. The action is

$$\mathcal{S} = \sum_{i,j} -J_{i,j}\vec{x}_i \cdot \vec{x}_j,$$

where the sum is over nearest-neighbor pairs in $\mathcal{L}$. Frequently we will take the interaction constants $J_{i,j}$ to depend only on which coordinate of the lattice the points differ in. As in the earlier examples, it may be helpful to renormalize the action so that when all the $\vec{x}_i$ are equal we get $\mathcal{S} = 0$. This results in a renormalized action

$$\mathcal{S} = \sum_{i,j} \frac{1}{2}J_{i,j}(\vec{x}_i - \vec{x}_j)^2.$$

As before, we are interested in thinking of our d-dimensional lattice as having 1 time dimension and $d-1$ spatial dimensions and finding a τ-continuum quantum mechanical Hamiltonian H that corresponds to the action as the lattice spacing τ goes to 0. Since the spin configuration space is now continuous, the partition function will now be an integral rather than a summation:

$$Z = \int_{(S^N)^M} d\vec{x}_1 \ldots dx\vec{x}_M e^{-\beta \sum_{i,j} \frac{1}{2}J_{i,j}(\vec{x}_i - \vec{x}_j)^2}.$$

3.1. O(2) model on a 1-dimensional lattice. First we consider the $O(2)$ model on a 1-dimensional lattice of N points with periodic boundary conditions and a magnetic field of strength

h. Since our configuration space is S^1, we can parameterize the spins $\vec{x_i} = (\cos(\theta_i), \sin(\theta_i))$ for $\theta \in [0, 2\pi)$. Under this parameterization, the action becomes

$$\mathcal{S} = -\beta_\tau \sum_{j=1}^N \cos(\theta_j)\cos(\theta_{j+1}) + \sin(\theta_j)\sin(\theta_{j+1}) - \beta h \sum_{j=1}^N \cos(\theta_j)$$

$$= -\beta_\tau \sum_{j=1}^N \cos(\theta_j - \theta_{j+1}) - \frac{\beta h}{2} \sum_{j=1}^N (\cos(\theta_j) + \cos(\theta_{j+1})).$$

The corresponding partition function is then

$$Z = \int_{[0,2\pi]^N} \left(\prod_{j=1}^N d\theta_j \right) e^{\beta_\tau \sum_{j=1}^N \cos(\theta_j - \theta_{j+1})} e^{\frac{\beta h}{2} \sum_{j=1}^N \cos(\theta_j) + \cos(\theta_{j+1})}.$$

We now wish find an analogue of the transfer matrix from the analysis of the Ising model. To do this, we need the concept of integral operators. Given any L^2 function $f(x, y)$ on a domain $E \times E$, we can define the operator $L_f : L^2(E) \to L^2(E)$ by

$$(L_f g)(x) = \int_E f(x, y) g(y) \, dy$$

for any $g \in L^2(E)$.

Define $f(\theta, \phi) : [0, 2\pi] \times [0, 2\pi]$ by

$$f(\theta, \phi) = e^{\beta_\tau \cos(\theta - \phi)} e^{\frac{\beta h}{2}(\cos(\theta) + \cos(\phi))}.$$

Then the integral operator $\hat{T} = L_f$ plays the same role as the transfer matrix above. Since f is in $L^2([0, 2\pi]^2)$, $\hat{T}$ is trace-class and we have

$$Z = \int_{[0,2\pi]^N} \left(\prod_{j=1}^N d\theta_j \right) f(\theta_1, \theta_2) \dots f(\theta_N, \theta_1) = \operatorname{Tr} \hat{T}^N.$$

Again, our new goal is to write $\hat{T} = I - \tau \hat{H}$ when τ is small for some Hamiltonian on the one site Hilbert space $L^2([0, 2\pi])$. Recall that the set $B = \{\psi_m := \frac{1}{\sqrt{2\pi}} e^{im\theta} \mid m \in \mathbb{Z}\}$ forms an orthonormal Hilbert basis for $L^2([0, 2\pi])$. We will approximate the action of $\hat{T}$ on elements ψ_m of this basis in order to find our approximate Hamiltonian $\hat{H}$.

$$(\hat{T}\psi_m)(\theta) = \frac{1}{\sqrt{2\pi}} \int_0^{2\pi} e^{\beta_\tau \cos(\theta - \phi)} e^{\frac{\beta h}{2}(\cos(\theta) + \cos(\phi))} e^{im\phi} \, d\phi$$

$$= \frac{1}{\sqrt{2\pi}} e^{\frac{1}{2}\beta h \cos(\theta)} \int_0^{2\pi} e^{\beta_\tau \cos(\theta - \phi)} e^{\frac{1}{2}\beta h \cos(\phi)} e^{im\phi} \, d\phi. \tag{2}$$

We can use a Fourier transform to rewrite the $\cos(\theta - \phi)$ portion of the exponential in Eq. (2).

$$e^{-\beta_\tau + \beta_\tau \cos(\theta - \phi)} = \sum_{\ell \in \mathbb{Z}} e^{i\ell(\theta - \phi)} I_\ell(\beta_\tau), \tag{3}$$

where $I_\ell(\beta)$ is the Bessel function of imaginary argument. In particular, we will let $\beta_\tau = \tau^{-1}$ so that β_τ is large when τ goes to 0. Then we can approximate $I_\ell(\beta_\tau)$ by the Gaussian $e^{-\ell^2/2\beta_\tau}$ and

Eq. (3) becomes

$$e^{\beta_\tau \cos(\theta-\phi)} \approx e^{\beta_\tau} \sum_{\ell\in\mathbb{Z}} e^{i\ell(\theta-\phi)} e^{-\ell^2/2\beta_\tau} \tag{4}$$

$$= e^{1/\tau} \sum_{\ell\in\mathbb{Z}} e^{i\ell(\theta-\phi)} e^{-\tau\ell^2/2}. \tag{5}$$

After substituting Eq. (5) into Eq. (2), we can use absolute convergence of the sum to rewrite the expression further.

$$(\hat{T}\psi_m)(\theta) \approx \frac{1}{\sqrt{2\pi}} e^{\frac{1}{2}\beta h \cos(\theta)} \int_0^{2\pi} e^{1/\tau} \left(\sum_{\ell\in\mathbb{Z}} e^{i\ell(\theta-\phi)} e^{-\tau\ell^2/2}\right) e^{\frac{1}{2}\beta h \cos(\phi)} e^{im\phi}\, d\phi$$

$$= \frac{1}{\sqrt{2\pi}} e^{\frac{1}{2}\beta h \cos(\theta)+1/\tau} \sum_{\ell\in\mathbb{Z}} e^{i\ell\theta-\tau\ell^2/2} \left(\int_0^{2\pi} e^{\frac{1}{2}\beta h \cos(\phi)} e^{i(m-\ell)\phi}\, d\phi\right). \tag{6}$$

At this point, we want $\beta\beta_\tau$ to remain finite as τ goes to 0, so $\beta = \lambda\tau$ for some constant λ. Thus $\frac{1}{2}\beta h \cos(\phi)$ goes to 0 as τ goes to 0 and we can expand the corresponding exponential in Eq. (6). We then use orthogonality to evaluate the resulting integral.

$$(\hat{T}\psi_m)(\theta) \approx \frac{1}{\sqrt{2\pi}} e^{\frac{\lambda\tau h}{2} \cos(\theta)+1/\tau} \sum_{\ell\in\mathbb{Z}} e^{i\ell\theta-\tau\ell^2/2} \left(\int_0^{2\pi} (1 + \left(\frac{\lambda\tau h}{2} \cos(\phi)\right) e^{i(m-\ell)\phi}\right) d\phi,$$

$$= \frac{1}{\sqrt{2\pi}} e^{\frac{\lambda\tau h}{2} \cos(\theta)+1/\tau} \sum_{\ell\in\mathbb{Z}} e^{i\ell\theta-\tau\ell^2/2} \left(2\pi\delta_{\ell,m} + \frac{\lambda\tau h}{2} \int_0^{2\pi} \left(\frac{e^{i\theta}+e^{-i\theta}}{2}\right) e^{i(m-\ell)\phi}\, d\phi\right),$$

$$= \frac{1}{\sqrt{2\pi}} e^{\frac{\lambda\tau h}{2} \cos(\theta)+1/\tau} \sum_{\ell\in\mathbb{Z}} 2\pi e^{i\ell\theta-\tau\ell^2/2} \left(\delta_{\ell,m} + \frac{\lambda\tau h}{4}(\delta_{m+1,\ell} + \delta_{m-1,\ell})\right),$$

$$= \sqrt{2\pi} e^{\frac{\lambda\tau h}{2} \cos(\theta)+1/\tau} (e^{im\theta-\tau m^2/2} + \frac{\lambda\tau h}{4}(e^{i(m+1)\theta-\tau(m+1)^2/2} + e^{i(m-1)\theta-\tau(m-1)^2/2})).$$

At this point, we need to rescale the problem to drop the factor of $e^{\frac{1}{\tau}}$, as this term goes to infinity as τ goes to 0. After we do this, we will write the action of $\hat{T}$ on ψ_m in terms of the operators J_z and $J_\pm$ defined on our basis and extended linearly by

$$J_z\psi_m = m \cdot \psi_m, \quad J_\pm\psi_m = \psi_{m\pm1}.$$

To do this we use the fact that τ is small to expand the remaining exponentials that do not depend on θ and drop all terms involving a power of τ greater than one.

$$(\hat{T}\psi_m)(\theta) \approx \sqrt{2\pi} \left(1 + \frac{1}{2}\lambda\tau h \cos(\theta)\right) \left(\left(1 - \frac{1}{2}\tau m^2\right) e^{im\theta} + \frac{\lambda\tau h}{4}\left(e^{i(m+1)\theta} + e^{i(m-1)\theta}\right)\right),$$

$$\approx \sqrt{2\pi} \left(\left(1 - \frac{1}{2}\tau m^2 + \frac{1}{2}\lambda\tau h \cos(\theta)\right) e^{im\theta} + \frac{\lambda\tau h}{4}\left(e^{i(m+1)\theta} + e^{i(m-1)\theta}\right)\right),$$

$$= \sqrt{2\pi} \left(\left(1 - \frac{1}{2}\tau m^2\right) e^{im\theta} + \frac{\lambda\tau h}{2}\left(e^{i(m+1)\theta} + e^{i(m-1)\theta}\right)\right),$$

$$= \left(2\pi\left(I - \tau\left(\frac{J_z^2}{2} - \frac{\lambda h}{2}(J_+ + J_-)\right)\right)\psi_m\right)(\theta). \tag{7}$$

From Eq. (7) we see that the desired Hamiltonian is

$$\hat{H} = 2\pi\left(\frac{1}{2}J_z^2 - \frac{\lambda h}{2}(J_+ + J_-)\right). \tag{8}$$

3.2. O(2) model on a 2-dimensional lattice. We can perform a similar analysis on a 2-dimensional $M \times N$ lattice where there are M spatial rows and N temporal columns. For simplicity, we will drop the magnetic field term and consider only the interactions between neighboring spins. We will fix all interaction constants in the temporal direction (with unit vector $\vec{\tau}$ and temporal lattice spacing τ) to be some $\beta_\tau(\tau)$. The interaction constants in the spatial direction (with unit vector $\vec{x}$) will be $\beta(\tau)$. For brevity we will omit the explicit dependence on τ.

The action then becomes

$$S = -\beta_\tau \sum_\ell \cos(\theta_\ell - \theta_{\ell+\vec{\tau}}) - \beta \sum_\ell \cos(\theta_\ell - \theta_{\ell+\vec{x}}).$$

It will be helpful later to adjust the action by a constant factor of $MN\beta_\tau - \frac{MN}{2}\log(2\pi\beta_\tau)$. We also write θ_j^k for the angle in row k and column j of the lattice and rewrite the second term to be symmetric between adjacent rows. The resulting action is

$$S = \sum_{k=1}^{M} L(\vec{\theta}^k, \vec{\theta}^{k+1}),$$

where

$$L(\vec{\theta}, \vec{\phi}) = -\frac{N}{2}\log(2\pi\beta_\tau) - \beta_\tau \sum_{j=1}^{N}(-1 + \cos(\theta_j - \phi_j)) - \frac{\beta}{2}\sum_{j=1}^{N}(\cos(\theta_j - \theta_{j+1}) + \cos(\phi_j - \phi_{j+1}))$$

Then

$$Z = \int d\vec{\theta}^1 \dots d\vec{\theta}^M e^{-L(\vec{\theta}^1, \vec{\theta}^2)} \dots e^{-L(\vec{\theta}^N, \vec{\theta}^1)} = \mathrm{Tr}\left(\hat{T}^M\right),$$

where $\hat{T}$ is the integral operator on the Hilbert space $L^2([0, 2\pi]^N) \cong L^2([0, 2\pi])^{\otimes N}$ defined by

$$(\hat{T}f)(\vec{\theta}) = \int_{[0, 2\pi]^N} e^{-L(\vec{\theta}, \vec{\phi})} f(\vec{\phi}) \, d\vec{\phi}.$$

As in the previous section, we will analyze the action of $\hat{T}$ on elements of the orthornomal Hilbert basis

$$B = \left\{ \psi_{\vec{m}}(\vec{\theta}) := \frac{1}{(2\pi)^{N/2}} \prod_{j=1}^{N} e^{im_j\theta_j} \mid \vec{m} \in \mathbb{Z}^N \right\}$$

of $L^2([0, 2\pi]^N)$. This will go very similarly and we will need the following approximations for small τ and $m \in \mathbb{Z}$. We will also take $\beta_\tau = 1/\tau$ and $\beta = \lambda\tau$ for some constant λ.

$$\int_0^{2\pi} e^{\frac{1}{\tau}(-1 + \cos(\theta - \phi))} e^{im\phi} \, d\phi \approx \frac{\sqrt{\tau}}{\sqrt{2\pi}} \int_0^{2\pi} \sum_{\ell \in \mathbb{Z}} e^{i\ell(\theta - \phi) - \ell^2\tau/2 + im\phi} \, d\phi = \sqrt{2\pi\tau} e^{-m^2\tau/2} e^{im\theta}. \qquad (9)$$

$$e^{\frac{1}{2}\lambda\tau \sum_{j=1}^{N} \cos(\phi_j - \phi_{j+1})} \approx 1 + \frac{\lambda\tau}{4} \sum_{j=1}^{N} (e^{i\phi_j} e^{-i\phi_{j+1}} + e^{-i\phi_j} e^{i\phi_{j+1}}) \qquad (10)$$

We will write our Hamiltonian in terms of the operators $J_z(j)$, $J_+(j)$, and $J_-(j)$ defined on the basis $\{\psi_{\vec{m}}\}$ by

$$(J_z(j)\psi_{\vec{m}})(\vec{\theta}) := m_j\psi_{\vec{m}},$$
$$(J_\pm(j)\psi_{\vec{m}})(\vec{\theta}) := e^{\pm i\theta_j}\psi_{\vec{m}}.$$

Now

$$(\hat{T}\psi_{\vec{m}})(\vec{\theta}) = \frac{1}{(2\pi)^{N/2}} \int_{[0,2\pi]^N} e^{-L(\vec{\theta},\vec{\phi})} e^{\sum_{j=1}^{N} im_j\phi_j} \, d\vec{\phi}.$$

$$= F(\vec{\theta}) \int e^{\frac{1}{\tau}\sum_{j=1}^{N}(-1+\cos(\theta_j-\phi_j)+im_j\phi_j)+\frac{\lambda\tau}{2}\sum_{j=1}^{N}\cos(\phi_j-\phi_{j+1})} \, d\vec{\phi},$$

where $F(\vec{\theta}) = \left(\frac{1}{\tau}\right)^{N/2} e^{\frac{\lambda\tau}{2}\sum_{j=1}^{N}\cos(\theta_j-\theta_{j+1})}$. We can now use Eq. (9) and Eq. (10). Whenever we take an expansion in powers of τ we will neglect any terms with order at least 2.

$$(\hat{T}\psi_{\vec{m}})(\vec{\theta}) \approx F(\vec{\theta}) \int e^{\frac{1}{\tau}\sum_{j=1}^{N}(-1+\cos(\theta_j-\phi_j)+im_j\phi_j)} \left(1 + \frac{\lambda\tau}{4}\sum_{j=1}^{N}(e^{i\phi_j}e^{-i\phi_{j+1}} + e^{-i\phi_j}e^{i\phi_{j+1}})\right) d\vec{\phi},$$

$$\approx F(\vec{\theta}) \left(\prod_{j=1}^{N}\sqrt{2\pi\tau}e^{-m_j^2\tau/2}e^{im_j\theta_j} + \sum_{\substack{1\leq k,\ell\leq N \\ k\neq\ell}}\prod_{j=1}^{N}\sqrt{2\pi\tau}e^{i(m_j+\delta_{k,j}-\delta_{\ell,j})\theta_j}e^{-\tau(m_j+\delta_{k,j}-\delta_{\ell,j})^2/2}\right),$$

$$\approx F(\vec{\theta}) \left(\prod_{j=1}^{N}\sqrt{2\pi\tau}(1-m_j^2\tau/2)e^{im_j\theta_j} + \sum_{k\neq\ell}\prod_{j=1}^{N}\sqrt{2\pi\tau}(1-\tau(m_j+\delta_{k,j}-\delta_{\ell,j})^2/2)e^{i(m_j+\delta_{k,j}-\delta_{\ell,j})\theta_j}\right),$$

$$\approx F(\vec{\theta})(2\pi\sqrt{\tau})^N \left(\left(I - \frac{\tau}{2}\sum_{j=1}^{N}J_z^2(j) + \frac{\lambda\tau}{4}\sum_{j=1}^{N}(J_+(j)J_-(j+1) + J_-(j)J_+(j+1))\right)\right)\psi_{\vec{m}}(\vec{\theta}),$$

$$\approx (2\pi)^N \left(I - \frac{\tau}{2}\sum_{j=1}^{N}J_z^2(j) - \lambda(J_+(j)J_-(j+1) + J_-(j)J_+(j+1))\right)\psi_{\vec{m}}(\vec{\theta}).$$

It now follows that up to a multiplicative constant the desired Hamiltonian is

$$\hat{H} := \frac{1}{2}\sum_{j=1}^{N}J_z^2(j) - \lambda(J_+(j)J_-(j+1) + J_-(j)J_+(j+1)).$$

4. (Mean) Spherical Model

Detailed analysis of the spherical model in various dimensions can be found in [8, 15]. We will sketch the approach of [8] to determining a quantum Hamiltonian for the spherical model with some minor corrections to notation and using a slightly different line of reasoning. Again we take a lattice $\mathcal{L}$ of N sites, usually cubic. Now our spins at each site are $x_i \in \Omega_0 := \mathbb{R}$ subject to periodic boundary conditions and the *mean spherical constraint* $\sum_i x_i^2 = N$. (This constraint is the origin of the name: in the mean, the spins have norm 1 i.e. lie on a unit circle). We can work with a *grand partition function* given by

$$Q_N(S) = \int_{\mathbb{R}^N} dx_1 \ldots d_{x_N} e^{-S\sum_i x_i^2 + \beta J \sum_{i,j} D_{i,j}x_ix_j + \beta h \sum_i x_i},$$

where β is the inverse temperature, J is the exchange energy between nearest neighbor spins, $D_{i,j}$ are relative strength of coupling constants, and h is the strength of an external magnetic field, and S implements the mean spherical constraint.

This grand partition function is related to the usual partition function and action by

$$Z = \frac{1}{2\pi i}\int_{\alpha_0-i\infty}^{\alpha_0+i\infty} e^{NS}Q_N(S)\,dS,$$

where α_0 is chosen so that any singularities of the integrand occur to the left of the line $s = \alpha_0$. The mean spherical constraint is equivalent to requiring

$$N = -\frac{\partial}{\partial S} \log Q_N(S).$$

Let n be the number of sites in the temporal direction and $\widetilde{N} = N/n$ be the number of sites in any hyperplane slice of $\mathcal{L}$ with constant temporal coordinate. Now we can take the hyperplane slices of the lattice across one dimension (the temporal dimension) and consider the integral operator $\hat{T}$ defined by

$$\hat{T}f(\vec{x}) = \int_{\mathbb{R}^{\widetilde{N}}} L(\vec{x}, \vec{y})f(\vec{y})\, d\vec{y},$$

where $\vec{x}, \vec{y} \in \mathbb{R}^{\widetilde{N}}$ and $L(\vec{x}, \vec{y})$ describes the interactions between hyperplanes. We define

$$L(\vec{x}, \vec{y}) = e^{-\frac{1}{2}S(\vec{x}^2 + \vec{y}^2) + \frac{1}{4}K(\vec{x}^T M \vec{x} + \vec{y}^T M \vec{y}) + K_\tau \vec{x}\vec{y} + \frac{1}{2}K'h \sum_{i=1}^{\widetilde{N}}(x_i + y_i)}, \tag{11}$$

where M is a symmetric matrix describing the spin-spin interactions and K, K_τ, K' are anisotropic coupling constants that may depend on the lattice spacing τ. This lets us write $Q_N(S) = \operatorname{Tr} T^n$. As in our earlier examples, we now want to extract a quantum Hamiltonian from this formulation by making appropriate choices of couplings and letting the temporal lattice spacing τ tend to 0.

We will need to make use of the operator identity

$$(e^{-t\hbar\Delta/(2m)}\psi)(\vec{x}_0) = \left(\frac{m}{2\pi t\hbar}\right)^{s/2} \int_{\mathbb{R}^s} \exp\left\{-\frac{m}{2t\hbar}|\vec{x}_1 - \vec{x}_0|^2\right\}\psi(\vec{x}_1)d\vec{x}_1 \tag{12}$$

which can be found, for example, in [5, Section 20.3].

Let $\varphi, \psi \in L^2(\mathbb{R}^{\widetilde{N}})$ and define

$$\widetilde{\varphi}(\vec{y}) = \varphi(\vec{y})e^{-\frac{1}{2}(S(\tau) - K_\tau(\tau))\vec{y}^2 + \frac{1}{4}K(\tau)\vec{y}^T M \vec{y} + \frac{1}{2}K'(\tau)h \sum_i y_i} \tag{13}$$

$$\widetilde{\psi}(\vec{x}) = \psi(\vec{x})e^{-\frac{1}{2}(S(\tau) - K_\tau(\tau))\vec{x}^2 + \frac{1}{4}K(\tau)\vec{x}^T M \vec{x} + \frac{1}{2}K'(\tau)h \sum_i x_i}. \tag{14}$$

We can then write

$$\langle \varphi | \hat{T}_\tau | \psi \rangle = \int d^{\widetilde{N}}\vec{x}\, d^{\widetilde{N}}\vec{y}\, \widetilde{\varphi}(\vec{y})e^{-\frac{1}{2}K_\tau(\tau)(\vec{x} - \vec{y})^2}\widetilde{\psi}(\vec{x}). \tag{15}$$

Applying Eq. (12) to the $\vec{y}$ integral we have

$$\langle \varphi | \hat{T}_\tau | \psi \rangle = \left(\frac{\pi}{2K_\tau}\right)^{\widetilde{N}/2} \int d^{\widetilde{N}}\vec{x}\, \widetilde{\varphi}(\vec{x})e^{\frac{1}{4K_\tau(\tau)}\Delta}(\vec{x})\widetilde{\psi}(\vec{x}) \tag{16}$$

$$= \left(\frac{\pi}{2K_\tau}\right)^{\widetilde{N}/2} \int d^{\widetilde{N}}\vec{x}\, \widetilde{\varphi}(\vec{x})(\mathbb{1} + \frac{1}{4K_\tau(\tau)}\Delta + \dots)\widetilde{\psi}(\vec{x}) \tag{17}$$

But this implies that $\hat{T}_\tau$ is actually multiplication by

$$\hat{T}_\tau(\vec{x}) = \left(\frac{\pi}{2K_\tau}\right)^{\widetilde{N}/2} e^{-\frac{1}{2}(S - K_\tau)\vec{x}^2 + \frac{1}{4}K\vec{x}^T M \vec{x} + \frac{1}{2}K'h \sum_i x_i} e^{\frac{1}{4K_\tau}\Delta}(\vec{x})e^{-\frac{1}{2}(S - K_\tau)\vec{x}^2 + \frac{1}{4}K\vec{x}^T M \vec{x} + \frac{1}{2}K'h \sum_i x_i}. \tag{18}$$

If we expand the exponentials, writing out explicitly the terms that are the product of only the degree 1 and degree 0 terms, this becomes

$$\hat{T}_\tau(\vec{x}) = \left(\frac{\pi}{2K_\tau}\right)^{\widetilde{N}/2} \left(1 - (S(\tau) - K_\tau(\tau))\vec{x}^2 + \frac{1}{2}K(\tau)\vec{x}^T M \vec{x} + \frac{1}{2}K'(\tau)h \sum_i x_i + \frac{1}{4K_\tau(\tau)}\Delta + \dots\right). \tag{19}$$

Each of the constants $K, K\tau$, and K' depend implicitly on β. Separating this out, we can now make a choice of the explicit dependence on τ so that the terms above are the only terms of the

transfer operator involving terms with a power of τ at most 1. In particular, for some constant ω, set

$$S(\tau) - 2K_\tau(\tau) = \frac{\omega^2}{8}K_\tau(\tau), \quad K(\tau) = \beta\tau, \quad K_\tau(\tau) = \frac{\beta}{4\tau}, \quad K'(\tau) = \beta\tau. \tag{20}$$

Then in the limit $\tau \to 0$ we have

$$\hat{T}_\tau(\vec{x}) = (2\pi\tau\beta^{-1})^{\tilde{N}/2}\left(1 - \frac{\omega^2\tau}{2\beta}\vec{x}^2 + \frac{1}{2}\tau\beta\vec{x}^T M\vec{x} + \frac{1}{2}\beta\tau h\sum_i x_i + \frac{\tau}{\beta}\Delta + O(\tau^2)\right). \tag{21}$$

After pulling out constants and approximating a Hamiltonian with $\hat{T} = \mathbb{1} - \tau\hat{H}$, we see

$$\hat{H} = -\Delta + \frac{1}{2}(\omega^2\vec{x}^2 - \beta^2\vec{x}^T M\vec{x} - \beta^2 h\sum_i x_i). \tag{22}$$

5. Lattice Gauge Theories

A *gauge theory* is a physical model which is left invariant under the action of some set of *gauge transformations* at points of the model. For example, the electric field is the gradient of electric potential and so is left invariant under translation of the potential by any fixed constant. Mathematically, a gauge is a choice of a section of some bundle and a gauge transformation is a map between two such sections. In some cases, it can be useful to fix a particular gauge condition to simplify the description of a problem.

5.1. Ising Lattice Gauge Theory.
Outline/add details to some material from [10], maybe discuss [11] plaquette solution technique?

The traditional Ising model possesses a global symmetry - if every spin in the lattice is flipped at once, then the value of the action S is preserved. On the other hand, if we flip only the spin associated to a single site in the lattice, the terms involving that spin all flip and so the value of S may not be preserved. However, we can define a very similar model that is invariant under local transformations as well as global ones.

Let $\mathcal{L}$ be a d-dimensional cubic lattice and label the links of the lattice by an adjacent site ℓ and a unit vector in the lattice μ. On each link we place an spin $\sigma_3(\ell, \mu) \in \{\pm 1\}$ and define a *local gauge transformation* $G(\ell)$ to be the operation of flipping the $2d$ spins connected to the site ℓ. Any action which is invariant under such transformations must be constructed from products of σ_3 on closed paths in the lattice so that any gauge transformation changes an even number of signs in each term of the action. For example, the most basic such action is

$$S = -J\sum_{\ell,\mu,\nu}\sigma_3(\ell, \mu)\sigma_3(\ell + \mu, \nu)\sigma_3(\ell + \mu + \nu, -\mu)\sigma_3(\ell + \nu, -\nu),$$

where ℓ ranges over all sites, and μ, ν over all unordered pairs of distinct unit vectors in the lattice. These fundamental squares are called *plaquettes*. Sometimes we may omit the indexing for brevity and write

$$S = -J\sum_\square \sigma_3\sigma_3\sigma_3\sigma_3$$

for the sum of the product of spins over all plaquettes in the lattice. Two different configurations s and $\tilde{s}$ of spins are called *gauge equivalent* if there is a sequence $G(\ell_1), G(\ell_2), \ldots$ of gauge transformations so that

$$\tilde{s} = s \cdot G(\ell_1)G(\ell_2)\ldots.$$

We will allow infinite products of gauge transformations when we work with an infinite lattice.

We may 'spend' the gauge symmetry to impose some condition on the configurations we work with and sum only over configurations satisfying this condition. As long as we are only interested

in observables which are gauge-invariant, this will not affect the resulting value of the observable, since we can write

$$\langle \hat{f} \rangle = \frac{1}{Z} \sum_{\{\sigma\}} \hat{f}(\{\sigma\}) e^{-\beta S(\{\sigma\})}$$

$$= \frac{\sum_{\{\sigma\}} \hat{f}(\{\sigma\}) e^{-\beta S(\{\sigma\})}}{\sum_{\{\sigma\}} e^{-\beta S(\{\sigma\})}}$$

$$= \frac{\sum_C \left(\sum_{\{\sigma\} \in C} \hat{f}(\{\sigma\}) e^{-\beta S(\{\sigma\})} \right)}{\sum_C \left(\sum_{\{\sigma\} \in C} e^{-\beta S(\{\sigma\})} \right)}$$

$$= \frac{\sum_C \left(\hat{f}(\{\sigma_C\}) e^{-\beta S(\{\sigma_C\})} \sum_{\{\sigma\} \in C} 1 \right)}{\sum_C \left(e^{-\beta S(\{\sigma_C\})} \sum_{\{\sigma\} \in C} 1 \right)}$$

$$= \frac{\sum_C \hat{f}(\{\sigma_C\}) e^{-\beta S(\{\sigma_C\})}}{\sum_C e^{-\beta S(\{\sigma_C\})}}.$$

Here C denotes the classes of equivalent configurations subject to the desired constraints. Every such class must have the same cardinality, justifying the cancellation in the final line.

If we are working with a lattice which is infinite or has free boundary conditions in the temporal direction, we can apply gauge transformations to pick a representative of any configuration class that satisfies the *temporal gauge*

$$\sigma_3(\ell, \tau) = 1$$

for every link in the temporal direction. Note that this does not yield a unique representative, since any product of gauge transformations that includes all sites with the same spatial coordinates (and varying temporal coordinate) will preserve the temporal gauge on that temporal hyperplane. In the temporal gauge, we can rewrite the Ising action as

$$S = -J \sum_{\square_t} \sigma_3(\ell, \mu) \sigma_3(\ell + \tau, \mu) - J \sum_{\square_x} \sigma_3 \sigma_3 \sigma_3 \sigma_3,$$

splitting the sum over the plaquettes involving temporal links and those involving only spatial links.

For example, on a two-dimensional $N_x \times N_t$ lattice with free boundary conditions and unit lattice vectors (μ, τ), imposing the temporal gauge yields the Ising gauge action

$$S = -J \sum_{\ell} \sigma_3(\ell, \mu) \sigma_3(\ell + \tau, \mu).$$

This is precisely the action of N_x non-interacting copies of a 1-dimensional standard Ising model with free boundary conditions and zero external magnetic field. The analysis of this model is very similar to that in Section 2.2 where the model was assumed to have periodic boundary conditions.

Careful: Under periodic boundary conditions in the temporal direction, gauge transformations can only increase the number of temporal links with $\sigma_3(\ell, \tau) = 1$ when applied to sites with $\sigma_3(\ell, \tau) = \sigma_3(\ell, -\tau) = -1$. Thus the temporal gauge can be imposed only on the configuration classes consisting of configurations with an *even* number of negative spins on temporal links.

5.2. Lattice $\mathbb{Z}_n$-QED Model in 1-d. Outline results from [2, 13, 16].

Fix E,V,U convention to be self-consistent/match numerical results in later section of [2]

After more thought, I don't think can have both consistency and agreement with results of [2].

The standard Lagrangian for QED is

$$\mathcal{L} = \psi^\dagger \gamma^0 [\gamma^\mu (i\partial_\mu + gA_\mu) - m]\psi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu}. \tag{23}$$

This can be discretized to get a Hamiltonian on a 1-dimensional spatial lattice [2, 13, 16].

$$H = -\frac{1}{2a} \sum_\ell (\psi_\ell^\dagger U_{\ell,\ell+1} \psi_{\ell+1} + H.c.) + m \sum_\ell (-1)^\ell \psi_\ell^\dagger \psi_\ell + \frac{g^2 a}{2} \sum_\ell E_{\ell,\ell+1}^2, \tag{24}$$

where ℓ labels the sites of a 1-dimensional lattice with N sites and lattice spacing a and H.c. denotes the Hermitian conjugate of the previous terms. We follow the references in using a staggered mass $(-1)^\ell m$ to avoid the fermion-doubling problem. Fermionic matter is represented by the one component spinor creation/annihilation operators $\psi_\ell^\dagger$ and ψ_ℓ defined on each site. Gauge fields are defined on the links $(\ell, \ell+1)$ of the lattice via electric potential variables $E_{\ell,\ell+1}$ and vector potential variables $A_{\ell,\ell+1}$. These give rise to comparators $U_{\ell,\ell+1}(\eta) = e^{-i\eta A_{\ell,\ell+1}}$ and $V_{\ell,\ell+1}(\xi) = e^{i\xi E_{\ell,\ell+1}}$ which commute at different links, are parameterized by $\eta, \xi \in \mathbb{R}$, and satisfy the canonical commutation relations

$$V_{\ell,\ell+1}(\xi) U_{\ell,\ell+1}(\eta) = e^{i\eta\xi} U_{\ell,\ell+1}(\eta) V_{\ell,\ell+1}(\xi). \tag{25}$$

This relation is the exponentiated version of the canonical commutation relation for the $E_{\ell,\ell+1}$ and $A_{\ell,\ell+1}$ variables $[E_{\ell,\ell+1}, A_{\ell,\ell+1}] = i$.

For simulation, we need to have finite quantities everywhere. For this reason, we develop an approximation of the $U(1)$ gauge group by $\mathbb{Z}_n$ via operators that preserve a discrete version of the CCR of Eq. (25). We will take a finite-dimensional representation of the two parameter projective unitary Weyl group $\{e^{i(\xi E_{\ell,\ell+1} - \eta A_{\ell,\ell+1})}\}_{\xi,\eta \in \mathbb{R}}$ with the particular choices of parameters $(\xi, \eta) = (0, \sqrt{2\pi/n}), (\sqrt{2\pi/n}, 0)$. This leads to operators

$$U_{\ell,\ell+1} = e^{-i\sqrt{\frac{2\pi}{n}} A_{\ell,\ell+1}} \qquad V_{\ell,\ell+1} = e^{i\sqrt{\frac{2\pi}{n}} E_{\ell,\ell+1}} \tag{26}$$

satisfying commutation relations for discrete powers

$$U_{\ell,\ell+1}^m V_{\ell,\ell+1}^k = e^{i\frac{2\pi}{n} km} V_{\ell,\ell+1}^k U_{\ell,\ell+1}^m \tag{27}$$

with $k, m \in \mathbb{Z}_n$.

We can implement the operators U and V by placing an n-dimensional Hilbert space $\mathcal{H} \cong \mathbb{C}^n$ on each link with orthonormal basis $\{|v_k\rangle\}_{0 \le k \le n-1}$. Then V is the diagonal operator given by

$$V |v_k\rangle = e^{-2\pi ik/n} |v_k\rangle. \tag{28}$$

We define U to be the canonical conjugate operator to V, which acts cyclically on the $|v_k\rangle$ basis:

$$U |v_k\rangle = |v_{k+1 \pmod n}\rangle. \tag{29}$$

We can check easily that this choice of U and V do satisfy Eq. (27). Now the definition of V together with Eq. (26) implies that E must also be diagonal in the $|v_k\rangle$ basis. Further since the eigenvalues of an exponentiated operator are the exponentials of the original eigenvalues, we have

$$e^{-2\pi ik/n} = e^{i\sqrt{2\pi/n}e_k}, \tag{30}$$

where $E |v_k\rangle = e_k |v_k\rangle$ for $k = 0, \ldots n-1$. Now Eq. (30) implies that

$$e_k = -k' \sqrt{2\pi/n} \tag{31}$$

for some $k' \equiv k \mod n$.

In [2], the authors claim that

$$e_k = \sqrt{\frac{2\pi}{n}} \left(k - \frac{n-1}{2} + \phi \right)$$

for some $\phi \in \mathbb{R}$. However, this leads to a contradiction unless $n = 2$.

For physical reasons, we need the states of our model to satisfy Gauss's law at each site, which here becomes

$$G_\ell := \psi_\ell^\dagger \psi_\ell + \frac{1}{2}((-1)^\ell - 1) - (E_{\ell,\ell+1} - E_{\ell-1,\ell}) \approx 0. \tag{32}$$

Now this gives us a condition on eigenvalues for physical states. For sites that are occupied by a fermion, $\psi_\ell^\dagger \psi_\ell$ has eigenvalue 1. Let $n_\ell = 1$ if the site ℓ is occupied and $n_\ell = 0$ if it is empty. Then after a rescaling of the energy to absorb the factor of $\sqrt{2\pi/n}$ into the coupling constants of the Hamiltonian, Eq. (32) yields the condition

$$n_\ell + \frac{1}{2}((-1)^\ell - 1) - (e_R - e_L) = 0, \tag{33}$$

where e_R and e_L are the values of E on the links $(\ell, \ell+1)$ and $(\ell-1, \ell)$ respectively. Restricting our attention to eigenstates that our physical states are constructed from so that e_R and e_L have the form k_r and k_L for some integers $k_r, k_L \in \mathbb{Z}_n$, we have four cases to consider.

If ℓ is an occupied even site, then

$$1 + k_L \equiv k_R \pmod{n}.$$

If ℓ is an empty even site, then

$$k_L \equiv k_R \pmod{n}.$$

If ℓ is an occupied odd site, then

$$k_L \equiv k_R \pmod{n}.$$

If ℓ is an empty odd site, then

$$k_L - 1 \equiv k_R \pmod{n}.$$

5.3. Lattice $\mathbb{Z}_n$-QED Model in 2-d. Generalize results from [2, 13] using general description from [16].

Think I have a correct interpretation of Hamiltonian/Gauss Law given in [16] using constructions from [2, 13]. Don't know about reliability given consistency problems of [2].

5.4. Lattice G-QCD Model in 1-d. Long term/ambitious: set up lattice Hamiltonian for QCD with SU(2) or SU(3) gauge group and find a sequence of finite groups G to approximate by.

Don't understand description in [9, 16] of the formulation of lattice Hamiltonian for SU(2) gauge group. In particular, not clear how upper indices of L, R are related to generators of SU(2) or what the defn of L,R ops in [16] is.

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