

\mathbb{Z}_2 gauge field coupled to a Fermion system

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In the whole article we use σ and s as shorthands of σ^z and s^z , unless confusion may be caused by the notation.

1 The model Hamiltonian

The model investigated in this article is shown in [1]. The model Hamiltonian is

$$H = H_Z - J \underbrace{\sum_{\langle \mathbf{i}, \mathbf{j} \rangle} \sigma_{\mathbf{i}\mathbf{j}} s_{\mathbf{i}} s_{\mathbf{j}}}_{H_{\text{Ising}}} + \sum_i h_i^x s_i^x - t \sum_{\langle \mathbf{i}, \mathbf{j} \rangle} \sigma_{\mathbf{i}\mathbf{j}} c_{\mathbf{i}}^\dagger c_{\mathbf{j}}, \quad (1)$$

where

$$H_Z = -g \sum_{\mathbf{l} \in \square_{\mathbf{i}^*}} \sigma_{\mathbf{i}\mathbf{j}}^z - h \sum_{\langle \mathbf{i}, \mathbf{j} \rangle} \sigma_{\mathbf{i}\mathbf{j}}^x, \quad (2)$$

where \mathbf{i}^* is a site in the dual lattice (i.e. a site placed at the center of a plaquette), $\square_{\mathbf{i}^*}$ is the plaquette whose center is \mathbf{i}^* , and \mathbf{l} denotes a certain bond of a plaquette. Such kind of model Hamiltonian usually emerges from orthogonal metals [1, 2].

In this section, we introduce every part of (1).

1.1 \mathbb{Z}_2 gauge field and its dual theory

1.1.1 The plaquette term

The plaquette term in the Hamiltonian of the \mathbb{Z}_2 gauge field is be a function of

$$F_{\mathbf{i}^*} = \prod_{\mathbf{l} \in \square_{\mathbf{i}^*}} \sigma_{\mathbf{l}}, \quad (3)$$

which is invariant under a \mathbb{Z}_2 gauge transformation

$$Q_{\mathbf{i}} = \prod_{\mathbf{l} \in +_{\mathbf{i}}} \sigma_{\mathbf{l}}^x. \quad (4)$$

A convenient convention is to let a plaquette share the same index with the site in its left bottom corner. An example of a \mathbb{Z}_2 gauge field configuration can be found in Figure 1 on page 1.

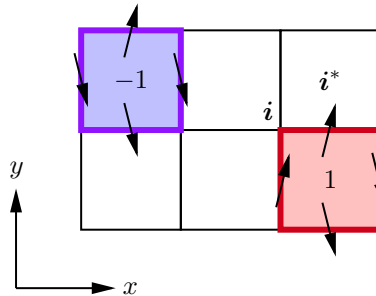


Figure 1: A \mathbb{Z}_2 gauge field configuration. The blue plaquette's $F_{\mathbf{i}^*}$ is $(-1) \times 1 \times (-1) \times (-1) = -1$, while the red plaquette's $F_{\mathbf{i}^*}$ is $1 \times 1 \times (-1) \times (-1) = 1$. Note that we assign the same index to a plaquette's center (labeled as \mathbf{i}^* at the right top of the diagram) and the plaquette's left bottom site \mathbf{i} . The terms “up” and “down” are defined in the Cartesian coordinates given in the diagram.

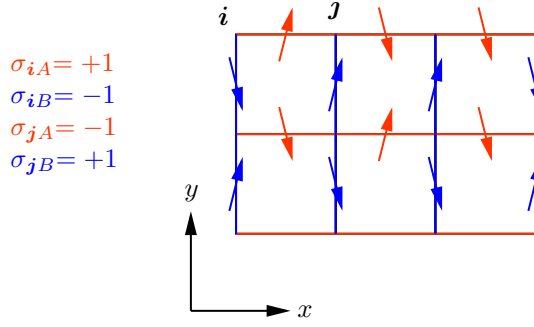


Figure 2: Dividing a gauge field configuration into two sublattices

1.1.2 The transverse field

We name the \mathbb{Z}_2 gauge theory with plaquette terms only as the **Wegner model**, i.e.

$$H_W = -g \sum_i F_{i^*}. \quad (5)$$

There is no quantum fluctuation in H_W . The existence of interaction between the \mathbb{Z}_2 field and the Ising field and the fermions introduces effective interaction channels between \mathbb{Z}_2 excitations, but all effective interaction between \mathbb{Z}_2 excitations are in terms of σ_{ij}^z , which commutes with H_Z and therefore do not bring in quantum fluctuation.

An idiomatic way to add quantum fluctuation is to add a transverse field. In this project we consider a transverse field Hamiltonian in the form of

$$H_h = -h \sum_{\langle i,j \rangle} \sigma_{ij}^x, \quad (6)$$

where the parameter h measures the quantum fluctuation. (6) obviously commutes with Q_i for every i , so it can be a term in a \mathbb{Z}_2 gauge invariant Hamiltonian. In the language of string-net condensation, H_h is a string tension term.

1.1.3 Equivalence between a \mathbb{Z}_2 gauge theory and a bundle of Ising chains

With the gauge choice

$$\sigma_{i, i+\hat{x}} = 1, \quad (7)$$

we have

$$F_{i^*} = \sigma_{i, i+\hat{y}} \sigma_{i+\hat{x}, i+\hat{x}+\hat{y}}, \quad (8)$$

which, by renaming $\sigma_{i, i+\hat{y}}$ into S_i , reads

$$F_{i^*} = S_i S_{i+\hat{x}}. \quad (9)$$

The hopping constant σ_{ij} , respectively, is

$$\sigma_{ij} = \begin{cases} 1, & j = i + \hat{x}, \\ S_i, & j = i + \hat{y}, \end{cases}, \quad \sigma_{ji} = \sigma_{ij}. \quad (10)$$

This, actually, means that the \mathbb{Z}_2 gauge field may also be transformed into a dual transverse field Ising field. For example consider Wegner model. With the definition of S_i s it is rephrased into

$$H_W = -g \sum_i S_i S_{i+\hat{x}}, \quad (11)$$

so the model is actually a bundle of 1D Ising spin chain.



Figure 3: Converting a \mathbb{Z}_2 gauge field theory into a bundle of Ising chains. There is no interaction between parallel Ising chains.

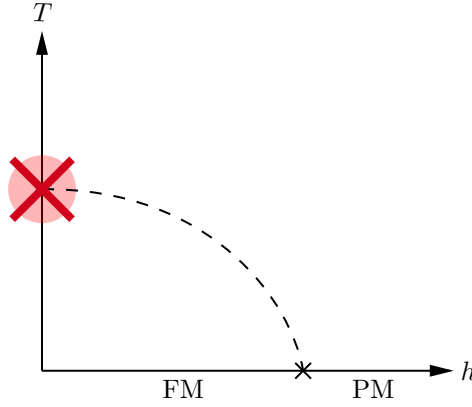


Figure 4: No deconfined phase when $T > 0$: if there is a deconfined phase when $T > 0$, there must be a thermal phase transition at the $h = 0$ line, which is forbidden by the fact that a classical 1D Ising chain does not show thermal phase transitions.

1.1.4 Absence of a finite temperature deconfined phase

When $T = 0$ and $h \ll g$, the \mathbb{Z}_2 gauge field exhibits a *deconfined* phase, i.e. there are freely moving \mathbb{Z}_2 excitations such as \mathbb{Z}_2 fluxes.

When $h = 0$, the absence of quantum fluctuation here means the model is indeed totally classical, so there is no thermal phase transition, because a 1D Ising chain does not show thermal phase transition. On the other hand, we know when $T = 0$ the model exhibits deconfined \mathbb{Z}_2 charges, and when a transverse field is introduced a zero temperature quantum phase transition can be observed from the deconfined phase into the confined phase as the transverse field grows. If for $T > 0$ there exists a deconfined phase, a simple analysis of the phase diagram tells us there must be a thermal phase transition when there is no transverse field. The fact that Wegner model is dual to a bundle of 1D Ising chain and therefore lacks thermal phase transition implies that no deconfined phase exists with finite non zero temperature.

Here the terminology may cause confusion: we do not actually know if the "finite temperature deconfined" phase shares behavior with the standard zero temperature deconfined phase in Wegner model, and the "confined" phase in the high temperature condition is not caused by a strong transverse field which coerces all σ degrees of freedom into \rightarrow or \leftarrow , but rather, by thermal fluctuation that erases all information of \mathbb{Z}_2 excitations. With the criteria of Wilson loops we know the "deconfined" phase obeys the perimeter law, agreeing with the ordinary deconfined phase in Wegner model, but whether the Wilson loop operator is a good detector in non-local models is still under question, because in local models, existence of long range correlation indicates exotic phenomena, while in non-local models long range correlation may be a trivial consequence of the non-locality. All these questions remain open, and in this project we simply use the terms "confined" and "deconfined" as a shorthand of the two phases in a \mathbb{Z}_2 gauge field theory with a dual theory of Dyson-Ising spin chains, without inquiring into their similarity and discrepancy with local \mathbb{Z}_2 theories such as Wegner model.

1.1.5 (Trivial) deconfined phase in models with long-range interaction

If we are to extend the deconfined phase to $T > 0$, at least some stronger correlation must be introduced. Without introducing quantum fluctuation, a reasonable proposal may be

$$H_Z = - \sum_i \sum_{r=1}^{\infty} J(r) \prod_{a=0}^{r-1} F_{i^*+a\hat{x}},$$

or in terms of S_i s,

$$H_Z = - \sum_{r=1}^{\infty} J(r) \sum_i S_i S_{i+r\hat{x}},$$

which is again a bundle of spin chains without quantum fluctuation but this time with long range interaction. A famous example is the **Dyson-Ising chain**, which is defined as

$$H_Z = -g \sum_i S_i S_{i+\hat{x}} - J_r \sum_i \sum_r \frac{S_i S_{i+r\hat{x}}}{r^\omega}.$$

For 1D Ising chain with long-range interaction, there exists a thermal phase transition with zero transverse field, and indeed we get a deconfined phase with finite temperature, and as the temperature goes up the deconfined phase switches to the confined phase.

However, the so-called deconfined phase of a long-range interacting model is not of particular interest, because if you put manually something long-range into a model, *of course* it exhibits some long-range behaviors, for example Wilson loops obeying the perimeter law. That is why we call such a “deconfined phase” a trivial one. Despite its triviality, “deconfined” phases in long-range models give us a hint that strong interaction between \mathbb{Z}_2 fluxes is important for a thermal deconfined phase. A natural question to ask is, if we introduce more things into the model to induce effective interaction channels between \mathbb{Z}_2 fluxes, what happens when $T > 0$?

1.2 Orthogonal metals, emergent fermions and an Ising field

1.2.1 Fermion fractionalization in orthogonal metals and the effective model

Orthogonal metal is a type of fractionalized electron systems where an electron is split into another fermion and an Ising spin, i.e.

$$c_{i\alpha}^\dagger = f_{i\alpha}^\dagger \sigma_i^z, \quad (12)$$

where f operators and σ^z operators commute and $f_{i\alpha}^\dagger$ fermions can move around freely. Since the Hilbert space spanned by the f fermions and Ising spins are larger than the original electronic system’s Hilbert space, certain constraints must be imposed, resulting in an emergent gauge field, gluing up f and σ fields.

That justifies the way we construct (1). Suppose we have a tight-binding electron models with certain interaction channels that are strong enough to create a non-Fermi liquid phase, the Hamiltonian of which is

$$H = -t \sum_{\langle i,j \rangle, \alpha} c_{i\alpha}^\dagger c_{j\alpha} + \sum_{i,j} V_{ij} n_i n_j. \quad (13)$$

Substituting (12) into (13) and doing necessary mean field approximations, we find that the Hamiltonian about f fermions in the model is also a tight-binding Hamiltonian with its hopping constants being the same as σ_{ij} , which couples the fermions with the \mathbb{Z}_2 gauge field, endowing the fermions \mathbb{Z}_2 charges.

The idea of orthogonal metals gives us an approach to find a local \mathbb{Z}_2 gauge model that shows deconfined phase at finite temperature. Our logic is the inverse of the derivation of orthogonal metals: if an orthogonal metal with \mathbb{Z}_2 gauge structure does exist, then it can be described by (1). That implies the existence of a model Hamiltonian in the form of (1) with a deconfined phase. Our goal is, therefore, to check under what condition (1) *can never* be in a deconfined phase.

1.2.2 Confined phase in orthogonal metals

When the \mathbb{Z}_2 charges get trapped into a confined phase, it can be expected that the fermions and the Ising spins are also confined. Note that what the composite particles generated by the gauge field as a glue in the confined phase are is not quite clear. One possibility is that the fermions are glued together, forming something like Cooper pairs, where the fermion excitations are now gapped. Nonetheless, note that orthogonal metals are generated in a strongly correlated *electron* systems, so it is highly likely that *one fermion and one Ising spin* are glued together, restoring the electrons. In this case, the confined phase is just an ordinary metal, with gapless fermions (electrons). At high temperature \mathbb{Z}_2 excitations are confined (in the sense defined in subsection 1.1.4), so (1) is just an ordinary metal.

2 Path integral formulation

We are going to study (1) numerically. In the partition function, the (non-normalized) weight of a configuration is $\langle n | e^{-\beta H} | n \rangle$, given that $\{|n\rangle\}$ is a basis. We use a discrete path integral method to evaluate these weights.

We introduce a Trotter decomposition with imaginary time step $\Delta\tau = \beta/m$, choose σ_{ij}, s_i as labels for the \mathbb{Z}_2 field and the Ising field, respectively, and integrate out the fermion hopping term. Now a configuration of the systems is a sequence of length m , and at each imaginary time point τ there is a \mathbb{Z}_2 field $\sigma_{ij}(\tau)$ and an Ising field $s_{ij}(\tau)$. The weight of a configuration is

$$\begin{aligned} W(\sigma, s) &= \text{tr}_{\text{fermion}} \prod_{\tau=1}^{m\Delta\tau} \langle (\sigma, s)(\tau + \Delta\tau) | e^{-\Delta\tau H_Z} e^{-\Delta\tau H_{\text{Ising}}} | (\sigma, s)(\tau) \rangle e^{\Delta\tau t \sum_{\langle i,j \rangle} \sigma_{ij}(\tau) c_i^\dagger c_j} \\ &= \prod_{\tau=1}^{m\Delta\tau} \langle (\sigma, s)(\tau + \Delta\tau) | e^{-\Delta\tau H_Z} e^{-\Delta\tau H_{\text{Ising}}} | (\sigma, s)(\tau) \rangle \det \left(1 + \prod_{\tau=1}^{m\Delta\tau} e^{\Delta\tau t \sigma(\tau)} \right) \\ &= \prod_{\tau=1}^{m\Delta\tau} \langle \sigma(\tau + \Delta\tau) | e^{-\Delta\tau H_Z} | \sigma(\tau) \rangle \langle s(\tau + \Delta\tau) | e^{-\Delta\tau H_{\text{Ising}}} | s(\tau) \rangle \det \left(1 + \prod_{\tau=1}^{m\Delta\tau} e^{\Delta\tau t \sigma(\tau)} \right). \end{aligned}$$

where $\sigma = \{\sigma_{ij}(\tau)\}$ and $s = \{s_i(\tau)\}$, and in the last line values of $\sigma(\tau)$ replaces the σ^z operators in H_{Ising} .

2.1 Quantum fluctuation introduced by H_h

With the presence of H_h , the H_Z factor of each imaginary time step reads

$$\begin{aligned} &\langle \sigma(\tau + \Delta\tau) | e^{-\Delta\tau H_{Z0}} e^{-\Delta\tau H_h} | \sigma(\tau) \rangle \\ &= e^{-\Delta\tau H_{Z0} | \sigma(\tau)} \sum_{\sigma^x} e^{-\Delta\tau H_h | \sigma^x} \langle \sigma^z(\tau + \Delta\tau) | \sigma^x \rangle \langle \sigma^x | \sigma^z(\tau) \rangle \\ &= e^{-\Delta\tau H_{Z0} | \sigma(\tau)} \sum_{\sigma^x} e^{h\Delta\tau \sum_{\langle i,j \rangle} \sigma_{ij}^x} \langle \sigma^z(\tau + \Delta\tau) | \sigma^x \rangle \langle \sigma^x | \sigma^z(\tau) \rangle \\ &= e^{-\Delta\tau H_{Z0} | \sigma(\tau)} \prod_{\text{bond } l} \sum_{\sigma_l^x = \pm 1} e^{h\Delta\tau \sigma_l^x} \langle \sigma_l^z(\tau + \Delta\tau) | \sigma_l^x \rangle \langle \sigma_l^x | \sigma_l^z(\tau) \rangle \\ &= e^{-\Delta\tau H_{Z0} | \sigma(\tau)} \prod_{\text{bond } l} \sum_{\sigma_l^x = \pm 1} e^{h\Delta\tau \sigma_l^x} \langle \sigma_l^z(\tau + \Delta\tau) | \sigma_l^x \rangle \langle \sigma_l^x | \sigma_l^z(\tau) \rangle. \end{aligned}$$

By the formula

$$\langle \sigma_l^x | \sigma_l^z(\tau) \rangle = \frac{1}{\sqrt{2}} e^{i\pi \frac{1-\sigma_l^x}{2} \frac{1-\sigma_l^z(\tau)}{2}},$$

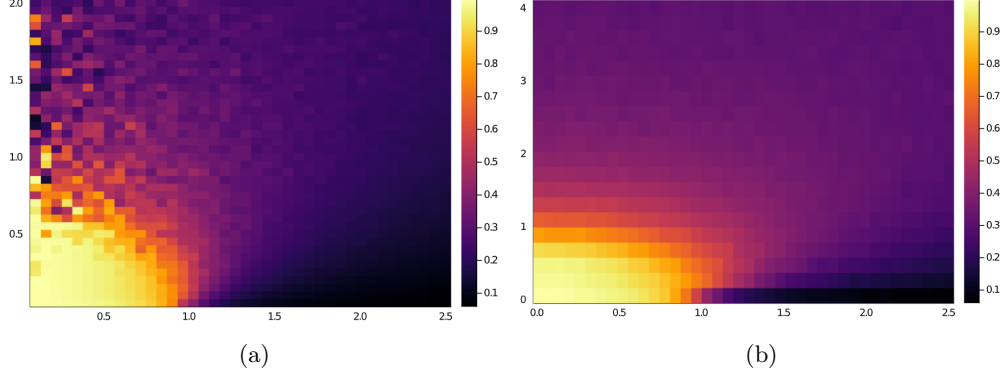


Figure 5: Phase diagrams of 1D transverse field Ising chain obtained with different updating algorithms. The x coordinate is h and the y coordinate is T . (a) Metropolis algorithm (b) Wolff cluster algorithm.

we have

$$\begin{aligned}
& \langle \sigma(\tau + \Delta\tau) | e^{-\Delta\tau H_{Z0}} e^{-\Delta\tau H_h} | \sigma(\tau) \rangle \\
&= e^{-\Delta\tau H_{Z0} | \sigma(\tau)} \prod_{\text{bond } l} \sum_{\sigma_l^x = \pm 1} e^{h\Delta\tau \sigma_l^x} \frac{1}{2} e^{i\pi \frac{1-\sigma_l^x}{2} \frac{1-\sigma_l^z(\tau)}{2}} e^{i\pi \frac{1-\sigma_l^x}{2} \frac{1-\sigma_l^z(\tau+\Delta\tau)}{2}} \\
&= \frac{1}{2^{2N}} e^{-\Delta\tau H_{Z0} | \sigma(\tau)} \prod_{\text{bond } l} \left(e^{h\Delta\tau} + e^{-h\Delta\tau} e^{i\pi \left(\frac{1-\sigma_l^z(\tau)}{2} + \frac{1-\sigma_l^z(\tau+\Delta\tau)}{2} \right)} \right) \\
&= \frac{1}{2^{2N}} e^{-\Delta\tau H_{Z0} | \sigma(\tau)} \prod_{\text{bond } l} (e^{h\Delta\tau} + e^{-h\Delta\tau} \sigma_l^z(\tau) \sigma_l^z(\tau + \Delta\tau)) \\
&= \frac{1}{2^{2N}} e^{-\Delta\tau H_{Z0} | \sigma(\tau)} \prod_{\text{bond } l} e^{J_\tau \sigma_l^z(\tau) \sigma_l^z(\tau + \Delta\tau)}.
\end{aligned}$$

where

$$\tanh J_\tau = e^{-2h\Delta\tau}. \quad (14)$$

The last few steps all use the fact that $\sigma_l^z = \pm 1$. With the gauge $\sigma_{i,i+\hat{x}} = 1$, the weight

3 Monte Carlo simulation

3.1 Examine the phase diagram of the pure \mathbb{Z}_2 gauge theory

Since J_x and J_y differ a lot, Metropolis algorithm is incapable for the simulation of the anisotropic Ising model. Cluster update methods - in this project Wolff cluster updating [3] - must be used. Figure 5 shows a comparison between Metropolis algorithm and Wolff algorithm, where Metropolis algorithm cannot update the system sufficiently when $h = 0$, since in that case the 2D classical Ising model corresponding to the 1D transverse field Ising chain degenerates into a classical 1D Ising chain due to the vanishing quantum fluctuation, so the coupling strength in the temporal direction approaches to infinite.

3.2 Get Fermions involved

References

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