

1 Integrability from extra dimensions

sec:integrability

Throughout the thesis, we discuss correspondences between a certain class of supersymmetric gauge theories and integrable lattice models, which are schematically written as

$$\mathcal{I}_{\mathsf{T}_{4d}[\mathsf{L}_{2d}]} = Z_{\mathsf{L}_{2d}[\mathsf{T}_{4d}]}, \quad (1.1)$$

eq:correspondence

where the left-hand side is the supersymmetric index of a four-dimensional gauge theory T_{4d} and the right-hand side denotes the statistical partition function of the corresponding integrable lattice model L_{2d} specified by the four-dimensional theory T_{4d} . As it turns out, the correspondence emerges from topological quantum field theory (TQFT) with extra dimensions [1]. In this section we give a step-by-step explanation of the above correspondence from an elementary level. Besides, to explain these, we will start by clarifying the precise meaning of the terms TQFT, lattice model, and integrability.

1.1 Preliminaries

sec:preliminaries

In section 1.2, we will introduce lattice model as a discrete version of quantum field theory (QFT). To do so, let us start from the general argument of QFT. We first briefly review one-dimensional QFT, which is nothing but quantum mechanics, and then extend the discussion to general quantum field theory in $(d + 1)$ -dimensions. Though we do not have a complete definition of general QFT, a special class of QFT which is called *topological* QFT is axiomatized by Atiyah [2].

1.1.1 What is quantum field theory?

sec:what_is_qft

Let us begin with one-dimensional QFT, as known as quantum mechanics (QM). Suppose we have a one-dimensional manifold,¹ which may be thought of as “time” $M^1 = \text{interval or } S^1 \text{ (or } \mathbb{R})$, and data to specify a QM:

- \mathcal{H} : a vector space (Hilbert space/state space),
- \mathcal{O} : a set of self-adjoint operators (observables, acting on \mathcal{H}),
- H : an operator called *Hamiltonian*,

where $H \in \mathcal{O}$, and the state space \mathcal{H} is a finite or an infinite dimensional \mathbb{C} -vector space. We sometimes need to allow a set of operators \mathcal{O} include not self-adjoint operators, but for simplicity we do not care about that at this moment. In undergraduate, we have learned that one should consider the Schrödinger equation and find a good basis in \mathcal{H} which diagonalizes the Hamiltonian.

¹We assume all the manifold in this thesis is smooth and oriented, unless otherwise stated.

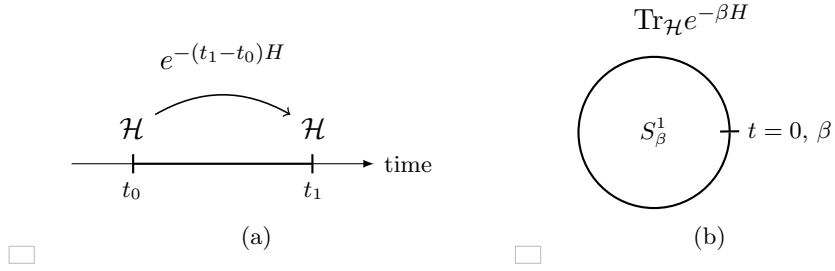


Figure 1: Two most crucial properties of quantum mechanics.

fig:qm_two_properties

As is well known, the procedure of solving the Schrödinger equation, at least formally, leads to an expression of time evolution of states by a linear map² e^{-tH} in quantum mechanics, and the partition function in many-body statistical mechanics. In turn, we have two most crucial properties of QM. Given an interval $M^1 = [t_0, t_1]$, we have state vectors at the endpoints of the interval and a time evolution between the states:

$$\begin{array}{|} \hline \\ \hline \end{array}_{t_0}^{t_1} \rightsquigarrow e^{-(t_1-t_0)H} : \mathcal{H} \longrightarrow \mathcal{H} , \quad (1.2)$$

eq:qm_linear_map

and given a circle $M^1 = S^1_\beta$, we get a number called *partition function*:

$$\begin{array}{c} \circ \\ S^1_\beta \end{array} \rightsquigarrow \text{Tr}_{\mathcal{H}} e^{-\beta H} . \quad (1.3)$$

eq:qm_number

In addition, from physical facts it is quite reasonable to assume that the time evolution and the partition function are compatible with cutting and gluing intervals and S^1 s. This simply means that a time evolution from time t_0 to t_1 followed by another time evolution from t_1 to t_2 is equal to a single time evolution from t_0 to t_2 , etc. Altogether, the properties of quantum mechanics, or one-dimensional QFT, are rephrased by the language of given one-dimensional manifold M^1 (see figure 1):

- Given an interval, QM produces state vectors at the endpoints and a linear map between them.
- Given an S^1 , QM produces a number.
- These two properties are compatible with cutting and gluing intervals and S^1 s.

Thus, in the abstract we conclude that quantum mechanics is a gadget satisfying the above properties for each given one-dimensional manifold M^1 .

From these observations of quantum mechanics, we wish to extend the discussion to general $(d+1)$ -dimensional QFT. The starting point of defining a $(d+1)$ -dimensional QFT is the choice of a $(d+1)$ -dimensional manifold M^{d+1} , which has d spatial directions and one

²We do not get into the argument of Euclideanization.

“time” direction. For most QFTs the manifold M^{d+1} is viewed as a Riemannian manifold with a smooth metric on it. We will be always considering a positive definite Riemannian metric, QFT on which is usually referred to as an Euclidean QFT, and hence precisely there is no notion of “time” in such a theory, at least globally. The manifold M^{d+1} may or may not have boundaries. In case it does have boundaries some additional information is needed at the boundaries to define the QFT. In addition to a Riemannian metric, depending on the situation one wants to consider, one often needs some more structures on the manifold M^{d+1} , e.g. conformal structure, spin structure, etc.

To obtain the data to specify a QFT, let us extend the observations seen in QM. Suppose we have a $(d+1)$ -dimensional manifold M^{d+1} , then we wish to “define” a QFT by a gadget Z , which should produce a vector when M^{d+1} has a boundary

$$Z(M^{d+1}) \in \mathcal{H}_{\text{bdy}}, \quad (1.4)$$

and should produce a number when M^{d+1} has no boundary

$$Z(M^{d+1}) \in \mathbb{C}. \quad (1.5)$$

For the case of M^{d+1} with a boundary, a vector space associated on the boundary is called the space of states or just Hilbert space in the physics literature. Meanwhile, if M^{d+1} does not have boundary, the number $Z(M^{d+1})$ is called the partition function. When the boundary of M^{d+1} consists of several disconnected components, namely the boundary is given by a disjoint union of simply connected d -dimensional closed manifolds $\{N_i^d\}$: $\partial M^{d+1} = \sqcup_i N_i^d$, Z should define a linear map among the vector spaces defined on the boundaries. In particular, in the case $M^{d+1} = N^d \times I$, where I is an interval of length T , M^{d+1} has two boundaries N^d and Z gives rise to a linear map $Z(N^d \times I) =: U(T)$,

$$U(T) : \mathcal{H}_N \longrightarrow \mathcal{H}_N. \quad (1.6)$$

From physical facts, Z also should be compatible with cutting and gluing of $(d+1)$ -dimensional manifolds, and thus we learn that the linear map given above satisfies $U(T_1)U(T_2) = U(T_1 + T_2)$. This really defines an operator H as the generator of U ,

$$U(T) = \exp(-TH). \quad (1.7)$$

H is called the Hamiltonian of the system, acting on the space of states \mathcal{H}_N . If one considers a manifold with Lorentzian metric, $U(T)$ is represented as

$$U(T) = \exp(-iTH), \quad (1.8)$$

and then the interval I is regarded as “physical time.”

1.1.2 Construction of Z

So far we have discussed only in an abstract way what quantum field theory is, or in other words what the gadget Z should satisfy. Here let us see a couple of examples how we specify Z to define a QFT. Broadly speaking, there are three kinds of constructions of Z . They are not totally independent and have many aspects of applicabilities. To the end of this subsection, M denotes a $(d + 1)$ -dimensional manifold with or without boundary, and N denotes a d -dimensional closed manifold without boundary.

Construct from an axiom

The first construction is in a sense the simplest one; we write down appropriate properties that a QFT must satisfy, and axiomatize them. Construction of Z is to give a mathematical formulation which satisfies the axioms as the data to specify the theory. This is actually the only way that one can define a QFT by mathematically rigorous procedures. Normally, such a theory is first studied by physicists as an ideal or a toy model from physical motivations, and then refined as rigorous mathematics by mathematicians.³ There are several kinds of such theories. We now introduce typical three examples.

The first one is *free theories* in any dimensions, which are toy models of field theories and play important roles as a probe of more complicated QFTs. In any spacetime dimensions, for the Riemannian manifolds with or without spin structure, we can rigorously define the free field theories. One of them is the free scalar field theory. Besides a $(d + 1)$ -dimensional closed manifold M , pick a G -bundle $P \rightarrow M$ with a connection A and a G -vector space V . Then we construct the associated vector bundle $P \times_G V$, whose covariant derivative is denoted by D_A . We now have the Laplacian Δ_A given by the covariant derivative D_A , and then the free scalar field theory is defined by the partition function

$$Z_{\text{scalar}}(M; A) := 1/\det \Delta_A. \quad (1.9)$$

eq:Zscalar

If the determinant of the Laplacian involves a divergence, it must be properly regularized.

Another free theory is the free fermion theory. In this case M also needs a spin (or spin^c) structure and the spin representations S, S' . Then we construct the Dirac operator \mathcal{D}_A on the spin bundles, and the partition function of the free fermion theory is given by

$$Z_{\text{fermion}}(M; A) := \det \mathcal{D}_A, \quad (1.10)$$

again the determinant is taken appropriately.⁴

Next, in two dimensions, we have another axiomatic quantum field theory, that is *conformal field theory* (CFT) [4]. To discuss two-dimensional CFT, we need a conformal (complex)

³The fact that a QFT can be treated in a mathematically rigorous way implies that the theory may have an enormous amount of symmetry, and the difficulties of the QFT are completely controlled by them. Even better, these theories can often be exactly solved in an appropriate sense.

⁴Fermion theories may have anomalies. If it is the case, the axiom needs to be somehow modified. In particular the partition function for a closed manifold is given by the eta invariant [3].

structure on M , i.e. M becomes a Riemann surface. CFT in two dimensions was formulated by Belavin, Polyakov, and Zamolodchikov [5] as a model of physical systems at critical points. They established the renowned Virasoro algebra as an infinite dimensional symmetry of the system and fully investigated the minimal model. The holomorphic part of the Virasoro algebra is captured by vertex operator algebras, and since then there are many mathematically rigorous discussions on them. Their kinematic behaviors on Riemann surfaces are governed by the conformal blocks, and its geometric meaning has been studied in [6]. In turn, the study of irrational CFTs [7, 8] has led to the AGT correspondence [9], which has brought to us large amount of developments in both physics and mathematics. We will give a general statement of AGT correspondence and apply it to our work in section ??.

The final example is *topological quantum field theories* (TQFTs). TQFT is one of the main focuses in this thesis. These theories originate from Witten’s proposals of topological field theories [10–12]. Inspired by Witten’s works, Atiyah and Segal axiomatized the topological QFT,⁵ and happily this also has brought to us a numerous amount of applications both to physics and to pure mathematics. For instance, two-dimensional TQFTs are known to be functorially equivalent to the category of commutative Frobenius algebras [13]. The three-dimensional Chern-Simons theory for a compact Lie group G is rigorously constructed by Turaev-Viro and Reshetikhin-Turaev using quantum groups [14, 15], and Kohno [16] using the monodromy representation of Knizhnik-Zamolodchikov equation [17]. The three-dimensional Chern-Simons theory has many applications to knot or link invariants and three-manifold invariants. Furthermore, both 2d and 3d TQFTs have applications to the mirror symmetry, or topological string theory (see e.g. [18]), those have led to fruitful interactions between physics and mathematics. We will take more time for TQFT to introduce the Atiyah’s axioms in some detail in section 1.1.3.

Path integrate a Boltzmann weight

Before going to Atiyah’s TQFT axioms, let us see two more constructions of Z . These are no longer mathematically rigorous, but rather familiar constructions for physicists. The second construction is to define a partition function by *path integral*. The general prescription is given as follows: one first introduces an action functional over the classical field configurations, which deduces the classical equation of motion by the variational principle. Then one exponentiates the action and integrate it over the space of fields.

The partition function for the free fields introduced above can also be defined by the path integral expression. For example, for the free complex scalar field theory consider a section ϕ of the vector bundle $P \times_G V$ over M , where V is a representation of a compact Lie group G . Then define an action functional

⁵Precisely speaking, what Segal axiomatized is the definition of conformal field theories [4]. Segal’s definition of CFT is functorial and quite similar to Atiyah’s TQFT. For the geometric aspects of CFT, see e.g. [6].

$$S : \Gamma(P \times_G V) \longrightarrow \mathbb{R}, \quad (1.11)$$

such that

$$S(\phi) = \int_M \frac{1}{2} D_A \phi \wedge * D_A \phi, \quad (1.12)$$

where $*$ is the Hodge star on M . D_A is again the covariant derivative given by the connection A . Using this action functional, physicists “define” its partition function by

$$Z_{\text{scalar}}(M; A) := \int_{\Gamma(P \times_G V)} \mathcal{D}\phi e^{-S(\phi)}. \quad (1.13)$$

The integrand e^{-S} of path integral is generically called the *Boltzmann weight*. For the free field theories, the path integral can be expressed as an infinite product of the Gaussian integral. Introducing an appropriate regularization, one can compute the exact partition function and it leads to the same result as in (1.9).

As another example, let us consider pure Yang-Mills theory. We introduce the kinetic term of the connection A , and define the action functional

$$S(A) = \int_M \frac{1}{4g^2} \text{Tr} F_A \wedge * F_A, \quad (1.14)$$

where F_A is the curvature of the G -connection A , and Tr is a Killing form of the Lie algebra \mathfrak{g} of G . Define the partition function of Yang-Mills theory by

$$Z_{\text{YM}}(M) = \int_{\mathcal{A}_M/\mathcal{G}} \mathcal{D}A e^{-S(A)}, \quad (1.15)$$

where the integral is taken over the space of connections on M modulo gauge transformations. Unlike the free scalar theory, making this path integral mathematically precise is extremely difficult. Although it is still ill-defined, physicists have been working on this expression in understanding of many properties of quantum gauge theories. Experimentally, we discretize the manifold M to a $(d+1)$ -dimensional lattice and then the path integral reduces to a well-defined integral. By putting it on a supercomputer, we can reproduce many experimental results to reasonable accuracy, at least numerically.

Deduce from string/M-theory

The final example of how to construct Z is to use string or M-theory. This is also our main tool to construct a QFT to realize the correspondence (1.1). Nonetheless, string or M-theory is less rigorous than path-integral expression of the construction of Z , and thus it is more hopeless to give a mathematically precise meaning to the construction. We here would like to show just two examples which “define” a class of quantum field theories through string/M-theory.

The first example is the so-called six-dimensional $\mathcal{N} = (2, 0)$ superconformal field theories (SCFTs) [19–21]. We start from a ten-dimensional string theory called the type IIB string

theory, which roughly assigns the partition function $Z_{\text{IIB}}(\tilde{M})$ to a ten-dimensional manifold \tilde{M} . Pick a finite subgroup Γ_G of type $G = A_n, D_n$ or $E_{6,7,8}$. We define a six-dimensional QFT Q_G by its partition function for a six-dimensional manifold M ,

$$Z_{Q_G}(M) = Z_{\text{IIB}}(M \times \mathbb{C}^2 / \Gamma_G). \quad (1.16)$$

They are examples of six-dimensional $\mathcal{N} = (2, 0)$ superconformal QFTs. These theories are known not to have a Lagrangian description, and hence their partition function cannot be obtained from the path integral formulation. They have another description via M-theory, as the low-energy dynamics of M5-branes. The construction from M5-branes leads to the AGT correspondence and the notion of class- \mathcal{S} theories, which we will explain in later sections.

The second example is the AdS/CFT correspondence [22–24]. This claims, originally, the equivalence between the type IIB string theory in ten-dimensional asymptotically AdS spacetime and four-dimensional $\mathcal{N} = 4$ super Yang-Mills theory living on the boundary of the AdS. The correspondence is schematically written as

$$Z_{\text{IIB}}(M \times Y) = Z[Y](\partial M), \quad (1.17)$$

where M denotes the five-dimensional AdS space and Y is a five-dimensional Sasaki-Einstein manifold. The right-hand side of the above equation is specified by Y , which is often called “internal space.” When $Y = S^5$ one finds the gauge theory side is give by $\mathcal{N} = 4$ super Yang-Mills theory: $Z[S^5] = Z_{\text{SYM}}$. The AdS/CFT correspondence is an extremely important subject in its own right. In this thesis, however, we do not get into the detail. For those who have interest in this direction, see e.g. [25].

1.1.3 Atiyah’s topological QFT

As already mentioned some times, in this thesis TQFT in extra dimensions plays a central role to realize the correspondence between supersymmetric gauge theories and integrable lattice models. So now let us pause here and introduce the Atiyah’s axioms of TQFT [2]. We first list the axioms, and right then give their physical meanings. The reader will notice that most of these axioms are physically quite natural and actually mathematical rephrasing of the properties that Z should satisfy, which we have already seen.

Axiom (Atiyah’s $(d + 1)$ -dimensional TQFT)

$(d + 1)$ -dimensional TQFT is defined by Z consisting of the following two data and five assignments:

- For each oriented d -dimensional closed manifold N , Z assigns a finite dimensional \mathbb{C} -vector space \mathcal{H}_N :

$$Z(N) = \mathcal{H}_N. \quad (1.18)$$

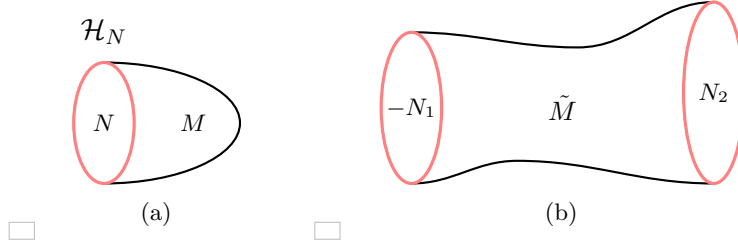


Figure 2: $(d+1)$ -dimensional manifolds M and \tilde{M} . They have a single boundary and two disconnected boundaries. For each case, Z gives a vector and a linear map, respectively.

fig:mfd_w_b
dys

This corresponds to (a half of) canonical quantization, or geometric quantization, known in the physics literature. Why we say it is “a half of” will be explained soon. In physics, especially for a field theory which has a Lagrangian description, we consider the phase space of fields, take a constant-time surface, and then perform canonical quantization by imposing canonical commutation relations on fields and their conjugate momenta. The constant-time surface is a codimension-1 hypersurface in $(d+1)$ -dimensional manifold M , which in this case is nothing but the d -dimensional closed manifold N . Thus, the manifold N is viewed as a collection of d spatial directions. One can think of the assigned vector space \mathcal{H}_N as “the space of functionals on the classical fields on N ,” usually called the space of states or Hilbert space of the system. The only major difference is that for topological theory the Hilbert space is of finite dimension.⁶

- For each oriented $(d+1)$ -dimensional manifold M with a boundary $\partial M = N$, Z assigns a vector

$$Z(M) = Z \left(\begin{array}{c} \text{red oval } N \\ \text{black oval } M \end{array} \right) \in \mathcal{H}_N. \quad (1.19)$$

This expresses the path integral quantization on M with boundary. Recall that if one has a spacetime manifold with boundary, one needs to impose a boundary condition on classical fields on the boundary, and it leads to the state vector associated to the boundary condition by path integral expression; let φ be a fixed field configuration on N ,

$$Z(M; \varphi) = \int_{X|_N = \varphi} \mathcal{D}X e^{-S(X)} \in \mathcal{H}_N. \quad (1.20)$$

In other words, the quantum state at a generic time t is given by the path integral of the Boltzmann weight over all the classical fields with time $< t$.

These two data are subjected to the following five assignments:

1. (involutory) Let $-N$ denote a manifold N with the opposite orientation, then

$$\mathcal{H}_{-N} = \mathcal{H}_N^*, \quad (1.21)$$

⁶One can remove the finiteness condition from the axiom. In TQFT we can in fact define a non-degenerate bilinear form from the other axioms, and from that as a corollary we can deduce the space of states is finite-dimensional.

where \mathcal{H}_N^* is the dual to \mathcal{H}_N .

2. (multiplicative) If N is a disjoint union of two d -dimensional closed manifolds N_1 and N_2 , then the vector space associated on it is factorized to a tensor product:

$$\mathcal{H}_{N_1 \sqcup N_2} = \mathcal{H}_{N_1} \otimes \mathcal{H}_{N_2}. \quad (1.22)$$

This is really natural in the physics point of view. This condition means that if we have two physical systems defined on spatially disjoint union $N_1 \sqcup N_2$, then the space of states is represented as the tensor product of each state space. Also, from the conditions so far one finds that if a $(d+1)$ -dimensional manifold \tilde{M} has two boundaries such that

$$\partial \tilde{M} = -N_1 \sqcup N_2, \quad (1.23)$$

then Z defines a linear map

$$Z(\tilde{M}) = Z \left(\begin{array}{c} \text{---} N_1 \text{---} \tilde{M} \text{---} N_2 \text{---} \end{array} \right) \in \text{Hom}_{\mathbb{C}}(\mathcal{H}_{N_1}, \mathcal{H}_{N_2}), \quad (1.24)$$

namely, for a cobordism between N_1 and N_2 , we have a linear map between \mathcal{H}_{N_1} and \mathcal{H}_{N_2} . This defines a time evolution, or transition amplitude, between the states in \mathcal{H}_{N_1} and \mathcal{H}_{N_2} . When the field theory has a Hamiltonian, $Z(\tilde{M})$ may correspond to the time evolution operator

$$Z(\tilde{M}) := e^{-tH} : \mathcal{H}_{N_1} \longrightarrow \mathcal{H}_{N_2}. \quad (1.25)$$

Since canonical quantization is a procedure to make an assignment of Hilbert spaces and the time evolution among them as we saw in QM, this association of linear maps gives the other half of canonical quantization. One can also explicitly express $Z(\tilde{M})$ in the path integral expression by the Feynman kernel (also as known as propagator or Green's function). Suppose we have an initial state $\Psi_0 \in \mathcal{H}_{N_1}$, then the state $\Psi_t \in \mathcal{H}_{N_2}$ at time t is expressed by

$$\begin{aligned} \Psi_t(\varphi_t) &= (e^{-tH} \Psi_0)(\varphi_t) \\ &= \int K(\varphi_t, \varphi_0) \Psi_0(\varphi_0) \mathcal{D}\varphi_0, \end{aligned} \quad (1.26)$$

where

$$K(\varphi_t, \varphi_0) = \int_{X|_{N_1}=\varphi_0}^{X|_{N_2}=\varphi_t} \mathcal{D}X e^{-S(X)}. \quad (1.27)$$

3. For two cobordisms

$$\partial M_1 = -N_1 \sqcup N_2, \quad \partial M_2 = -N_2 \sqcup N_3, \quad (1.28)$$

it follows that

$$Z(M_1 \cup_{N_2} M_2) = Z(M_2)Z(M_1), \quad (1.29)$$

where the cobordisms M_1 and M_2 are glued along N_2 by a certain diffeomorphism on N_2 . This asserts that the linear maps are transitive when we compose cobordisms. This is nothing but the physical requirement that the time evolution is compatible with the cutting and gluing the manifolds.

4. Given $N = \emptyset$ as an empty d -dimensional manifold, then the associated vector space is one-dimensional:

$$Z(\emptyset) = \mathbb{C}. \quad (1.30)$$

This is a non-triviality condition. From this, for each $(d+1)$ -dimensional manifold M without boundary, $\partial M = \emptyset$, Z assigns a number:

$$Z(M) \in \mathbb{C}. \quad (1.31)$$

This number associated to a closed $(d+1)$ -dimensional manifold is known as partition function. Not only is it an important quantity physically, but it also has mathematical applications, such as giving a topological invariant of the closed manifold M .

5. Let I be an interval, then for each cylinder $M = N \times I$, the linear map on \mathcal{H}_N is trivial:

$$Z(N \times I) = \text{id}_{\mathcal{H}_N}. \quad (1.32)$$

Since in a topological theory manifolds diffeomorphic to each other should be thought of as identical, each cobordism class defines a linear map. This means that the addition of a cylinder is a trivial operation; the associated linear map is the identity. This is another major difference from ordinary QFTs. For instance, for the three-dimensional Chern-Simons theory one can explicitly see its Hamiltonian is identically zero and hence the time evolution of Chern-Simons theory is trivial.

This is all the assignments for Atiyah's TQFT. Although we gave the path integral expressions at the explanation of the physical meanings for some conditions, the axioms themselves are really the basis for the rigorous mathematical definition of Z . One can actually take an equivalent definition of TQFT in slightly different axioms. For more details, see the original paper by Atiyah [2].

At first sight, this might look a little bit complicated and too abstract. This definition, however, is really natural and works quite well in the sense that it gives a kind of “homomorphism” between geometry and algebra:

$$Z : \text{‘geometry’} \longrightarrow \text{‘algebra’}.$$

In fact, all in all the axioms of Atiyah's TQFT define a functor from the category of $(d+1)$ -dimensional cobordisms to the category of finite-dimensional \mathbb{C} -vector spaces:⁷

$$Z : \text{Bord}_{d+1} \longrightarrow \text{Vect}_{\mathbb{C}}. \quad (1.33)$$

⁷What is more, the category of cobordisms and the category of vector spaces are endowed with the product structures, that is, taking disjoint union and taking tensor product, which are really symmetric operations. As such, TQFTs are keeping such structures and then called symmetric monoidal functor.

TQFT luckily has a mathematical precise definition. For a general QFT, i.e. not TQFT, in many cases there is no precise definition as we saw in the examples of the construction of Z . However, a QFT is generically expected to be characterized by some functor from the geometric structure of spacetime manifolds to the algebraic description of physical states and observables.

1.2 Lattice model as discrete QFT

Let us move on to the discussion on lattice model. First of all, let us begin with the question of “what is lattice model?” Probably there is no unique definition of lattice model, though. We here would like to characterize lattice model by a discrete version of QFT. What we have learned in the previous subsection is that a $(d+1)$ -dimensional QFT may be defined by an appropriate functor Z_Q^{d+1} , which produces a number for each $(d+1)$ -dimensional closed manifold M :

$$Z_Q^{d+1}(M) \in \mathbb{C}, \quad (1.34)$$

which is called the partition function of the model, and Z_Q^{d+1} satisfies additional some reasonable conditions for each QFT of interest.

We introduce lattice model in a same manner. For a $(d+1)$ -dimensional closed lattice L , a $(d+1)$ -dimensional lattice model is defined by Z_L^{d+1} , which produces a number

$$Z_L^{d+1}(L) \in \mathbb{C}, \quad (1.35)$$

which is again called the partition function, and Z_L^{d+1} satisfies additional some reasonable conditions. The typical example of lattice model in particle physics is lattice gauge theory as we mentioned before.⁸ For such theories, it is very clear that the model is defined by a discrete version of QFTs; one compactifies the spacetime manifold \mathbb{R}^4 to four-torus T^4 , and then discretize the theory and put it on a lattice on the torus.

In either case, therefore, “define a model” means to specify Z . In what follows, we will consider only $d = 0$ or 1 case of lattice model, and we refer each case as 1d or 2d lattice model. To prepare for the discussion of integrability, we take the prominent example of statistical lattice model known as the Ising model. We first briefly review the generality of the Ising model and compare with field theory, and then introduce transfer matrix which is essentially a time evolution operator in a discrete quantum system.

1.2.1 Prominent example – the Ising model

The Ising model is a very good introductory example to integrable lattice model. The 1d and 2d Ising model is known to be exactly solvable, whose meaning we will clarify in a moment,

⁸For sure, for any field theories when one computes a physical quantity on a computer, one needs to discretize the theory and put it on a lattice. In this sense, numerical analysis of field theory is always regarded as a lattice model.

and it is said to be integrable. To see the generality of the Ising model, we first define 2d lattice and spin configurations on it.

Define a 2d periodic lattice on two-torus T^2 by

$$\begin{aligned} L &:= \mathbb{Z}_m \times \mathbb{Z}_n \\ &= \{1, \dots, m\} \times \{1, \dots, n\}. \end{aligned} \quad (1.36)$$

A spin configuration on the lattice L is given by a map

$$\mathbf{s} : L \longrightarrow \{+1, -1\}, \quad (1.37)$$

where usually $+1$ is called “up spin” and -1 “down spin.” The map defines a configuration of up and down spins at each site of the lattice L . So it can also be thought of as an assignment of spins on all the sites of the lattice L . We often denote the spin at each site by its image $s_I := \mathbf{s}(I)$, $I \in L$. Let $S(L)$ be the set of all the spin configurations. Since spin configurations are defined on the $m \times n$ periodic lattice, the number of elements in $S(L)$ is 2^{mn} . In other words, the number of all the allowed configurations of spins on the lattice L is 2^{mn} .

For each spin configuration $\mathbf{s} \in S(L)$, define the energy functional of the Ising model by

$$\begin{aligned} E_{\text{Ising}}(\mathbf{s}) &= -J \sum_{\langle I, I' \rangle} s_I s_{I'} \\ &:= -J \left(\sum_{i,j=1}^{m,n} s_{ij} s_{i,j+1} + \sum_{i,j} s_{ij} s_{i+1,j} \right), \end{aligned} \quad (1.38)$$

where $s_{ij} = s_{i+m,j} = s_{i,j+n}$, and $J \in \mathbb{R}_{>0}$ is a constant parameter. This is one of the simplest spin systems, in which only the nearest-neighbor spins have simple interactions. The partition function of the Ising model is defined through this energy functional,

$$Z(L, E_{\text{Ising}}; \beta) := \sum_{\mathbf{s} \in S(L)} e^{-\beta E_{\text{Ising}}(\mathbf{s})}, \quad (1.39)$$

where $\beta \in \mathbb{R}_{>0}$ is the inverse temperature. The summand is generically called the Boltzmann weight as well as in field theory. In this expression, one recognizes that the right-hand side is a discrete version of path integral expression in field theory. The energy functional is corresponding to the action functional of a field theory, β^{-1} is the Planck’s constant, and the sum over all the spin configurations is the path integral over all the field configurations. Given a periodic lattice L , the partition function returns a number, which is a discrete version of the partition function of QFT which returns a number as well if a closed manifold M is given. For later use, define another quantity called the *free energy*,

$$f(\beta) := -\frac{1}{\beta} \frac{1}{mn} \log Z(\beta), \quad (1.40)$$

where $Z(\beta)$ is the partition function defined above.

According to the general story of statistical mechanics, the probability that a configuration \mathbf{s} with energy $E(\mathbf{s})$ is realized is given by the canonical ensemble⁹

$$p(\mathbf{s}) = \frac{1}{Z} e^{-\beta E(\mathbf{s})}, \quad (1.41)$$

where Z is the partition function of a statistical system. In a spin system, a physical observable is in general given by a functional on the space of spin configurations:

$$\mathcal{O} : S(L) \longrightarrow \mathbb{R}, \quad (1.42)$$

and its expectation value is given by the canonical ensemble as

$$\langle \mathcal{O} \rangle_p := \sum_{\mathbf{s} \in S(L)} \mathcal{O}(\mathbf{s}) p(\mathbf{s}) = \frac{1}{Z} \sum_{\mathbf{s} \in S(L)} \mathcal{O}(\mathbf{s}) e^{-\beta E(\mathbf{s})}. \quad (1.43)$$

In particular, an energy functional is an example of physical observable,

$$\frac{1}{mn} \langle E \rangle_p = \frac{1}{mn} \frac{1}{Z} \sum_{\mathbf{s} \in S(L)} E(\mathbf{s}) e^{-\beta E(\mathbf{s})}, \quad (1.44)$$

which is obtained from the derivative of the free energy,

$$\frac{1}{mn} \langle E \rangle_p = \frac{\partial}{\partial \beta} (\beta f(\beta)). \quad (1.45)$$

Generally speaking, the free energy provides us all the information of the system. For example, the state most likely to happen is governed by critical points of the free energy, and we can in principle compute physical quantities such as expectation value of energy, fluctuation, specific heat, and so on. To make a lattice model as a physically meaningful system, however, one needs to take the thermodynamic limit; $m, n \rightarrow \infty$. Therefore, in this sense, integrability of lattice model may be characterized by the calculability of an exact free energy at the thermodynamic limit. This can be rephrased as well by the calculability of the partition function. Based on these argument, we shall see the integrability of the Ising model, only for 1d case for simplicity.

1.2.2 Transfer matrix and integrability

It is well known that 1d and 2d Ising model is exactly solvable in the sense that one can exactly compute its free energy in the thermodynamic limit. Both 1d and 2d cases are solved by so-called the method of transfer matrix. To see this, for simplicity in this section we consider 1d Ising model.

The setup is almost the same as in the 2d case. A 1d lattice on torus and a spin configuration is defined by

$$L_{1d} = \mathbb{Z}_n = \{1, \dots, n\}, \quad (1.46)$$

⁹Only here we argue for a general energy functional, but the reader may assume the Ising model.

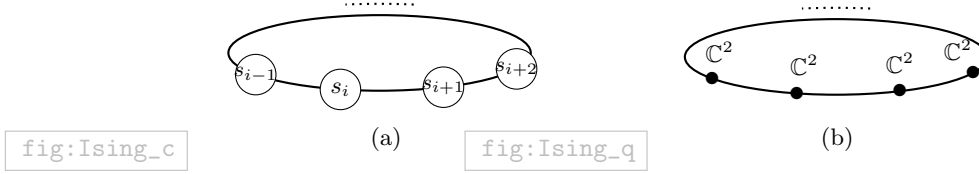


Figure 3: (a) Classical Ising spin chain. Spin variables s_i take values in $\{\pm 1\}$. They are just c numbers. (b) Quantum spins are vectors in \mathbb{C}^2 associated at each site. The discrete time evolution of quantum spins is given by the transfer matrix.

$$\mathbf{s} : L_{1d} \longrightarrow \{\pm 1\}. \quad (1.47)$$

Define the energy functional of 1d Ising model by

$$E_{1d \text{ Ising}}(\mathbf{s}) = -J \sum_{i=1}^n s_i s_{i+1}, \quad s_{n+1} = s_1. \quad (1.48)$$

Then the partition function of 1d Ising model is given as

$$\begin{aligned} Z(L_{1d}, E_{1d}; \beta) &:= \sum_{\mathbf{s}} e^{-\beta E_{1d}(\mathbf{s})} \\ &= \sum_{s_1, \dots, s_n = \pm 1} e^{K s_1 s_n} \dots e^{K s_3 s_2} e^{K s_2 s_1}, \end{aligned} \quad (1.49)$$

where $K := \beta J$.

Now introduce the *transfer matrix*,

$$T := \left(e^{K s s'} \right)_{s, s' = \pm 1} = \begin{pmatrix} e^K & e^{-K} \\ e^{-K} & e^K \end{pmatrix}, \quad (1.50)$$

where the indices of row and column are labeled by $+1$ and -1 , respectively. Using this matrix T , the partition function is rewritten as

$$Z(\beta) = \sum_{s_1, \dots, s_n = \pm 1} T_{s_1 s_n} \dots T_{s_3 s_2} T_{s_2 s_1}. \quad (1.51)$$

By definition of matrix multiplication, the partition function is eventually given by a trace of a 2×2 matrix:

$$\begin{aligned} Z(\beta) &= \sum_{s_1 = \pm 1} (T^n)_{s_1 s_1} \\ &= \text{Tr}_{\mathbb{C}^2} (T^n). \end{aligned} \quad (1.52)$$

eq:Ztransfer

This expression tells us that we have the quantum Hilbert space \mathbb{C}^2 at the boundary of each 1d segment (see figure 3), and the transfer matrix T sends a state to the adjacent site, which is the discrete “time evolution” of this system; $\log T \propto \text{Hamiltonian}$. This corresponds to the analogue of Hamiltonian/operator formalism in field theory. In operator formalism, spins are replaced by the Pauli matrices, and original classical up spin and down spin are replaced

by the eigenvalues and eigenvectors of the Pauli matrix σ^z . Further, the non-commutativity is now manifest in a way that the spins and their time evolution is given by matrices, this is the consequence of quantization.

For the expression (1.52) of the partition function, let us consider the thermodynamic limit $n \rightarrow \infty$. The eigenvalues of T are easily obtained and let them be $\lambda_0 > \lambda_1$, then we have the free energy

$$\begin{aligned} -\beta f(\beta) &= \lim_{n \rightarrow \infty} \frac{1}{n} \log Z(\beta) \\ &= \lim_{n \rightarrow \infty} \frac{1}{n} \log \lambda_0^n \left(1 + \left(\frac{\lambda_1}{\lambda_0} \right)^n \right) \\ &= \log \lambda_0. \end{aligned} \tag{1.53}$$

We conclude that the free energy in the thermodynamic limit is given by the largest eigenvalue of the transfer matrix. Finally, integrability of lattice model is rephrased as

One can exactly find the eigenvalues of transfer matrix.

\Rightarrow Can exactly compute the free energy.

\Rightarrow The system is exactly solvable.

The transfer matrices of 1d and 2d Ising model are diagonalizable and one can exactly find their eigenvectors and eigenvalues using, for example, algebraic Bethe ansatz. In this sense, 1d and 2d Ising model are said to be exactly solvable, or integrable. Then, a natural question arises; when can we diagonalize the transfer matrix of a lattice model? This actually leads to the most fundamental answer to the question of “what is integrable model.” A canonical answer is the *Yang-Baxter equation*. When the Boltzmann weight with spectral parameters, which take values in a Riemann surface, satisfies the Yang-Baxter equation, the lattice model is integrable. We are going to give an explanation of this definition of integrability in the following subsection in some detail, with its origin from TQFT with extra dimensions.

1.3 Integrability from TQFT in extra dimensions

We have seen that a QFT is ideally given by some functor Z , and lattice model is also given by some Z as a discrete version of QFT. Now, we would like to relate these two. The key concept is that “line operators in two-dimensional TQFT form a lattice model.” In particular, if the two-dimensional TQFT has extra dimensions, the corresponding lattice model naturally turns out to be integrable.

The purpose of this section is to provide a general prescription to relate four-dimensional supersymmetric gauge theories and integrable lattice models. We first give a step-by-step explanation from two-dimensional TQFT to integrable lattice model, and then put the argument into the setup of brane tilings. Branes in string theory are powerful enough to yield the systematic method to relate a class of supersymmetric gauge theories with integrable

sec:integrability_from_TQFT

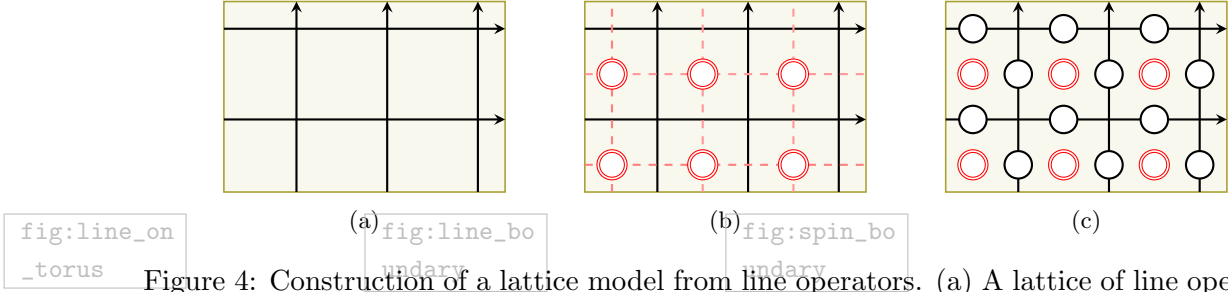


Figure 4: Construction of a lattice model from line operators. (a) A lattice of line operators on a torus. (b) The torus with holes obtained by gluing the pieces. (c) The corresponding spin model.

lattice models, as we will see. To begin with, we present a general prescription of brane constructions at a formal level. Concrete setups and more details will be given in later sections. The discussion in this subsection is mainly based on [1].

1.3.1 Lattice models from TQFTs with line operators

Suppose we have a two-dimensional TQFT T on a two-torus T^2 equipped with line operators $\mathcal{L}_i, i = 1, \dots, l$. Wrap them around one-cycles C_i on the torus in such a way that they form an $m \times n$ lattice; see figure 4(a). We wish to consider the correlation function of this configuration of line operators. To compute it, our strategy is to break up the torus into pieces and perform the path integral piecewise. Then the original correlation function is reconstructed by combining these pieces.¹⁰

Let us consider the following piece of surface which contains an intersection of two line operators \mathcal{L}_i and \mathcal{L}_j :

$$\begin{array}{c}
 a \quad d \\
 \swarrow \quad \searrow \\
 i \quad \uparrow \quad j \\
 \downarrow \quad \swarrow \\
 b \quad c
 \end{array}
 \quad . \quad (1.54)$$

If one thinks of the above picture as the world-sheet of two scattering open strings, one may regard it as follows: The endpoints of the open strings sweep out the double-lined arcs, on which D-branes are there and the boundary conditions (Chan-Paton factors) are specified by labels a, b, c, d . The line operators are the world-lines of particles associated with the open strings.

The path integral for the piece (1.54) produces a linear map (recall the property of quantum mechanics (1.2) and an assignment of TQFT), usually called S-matrix of the particle

¹⁰The original Atiyah's TQFT is referred to as “closed TQFT” nowadays. In the following, we will compute the correlation function by embedding the closed TQFT with line operators into an open/closed TQFT with line operators, or so-called defect TQFT; see e.g. [26–28].

scattering,

$$\begin{array}{c} a \\ \uparrow \\ i \text{---} \text{---} d \\ \downarrow \\ b \text{---} \text{---} c \\ j \end{array} \rightsquigarrow R_{ij} \begin{pmatrix} a & d \\ b & c \end{pmatrix} : V_{ab,i} \otimes V_{bc,j} \longrightarrow V_{ad,j} \otimes V_{dc,i}, \quad (1.55)$$

where $V_{ab,i}$ is the space of states of the open string propagating along \mathcal{L}_i with the boundary conditions on the left and the right ends specified by a, b , respectively. We call this map the R-matrix (or R-operator) associated with this decorated surface. To reconstruct the original configuration of line operators, we glue these pieces together. Gluing them amounts to the composition of R-matrices. For example, gluing two pieces horizontally gives

$$\begin{array}{c} a \quad d \quad e \\ \uparrow \quad \uparrow \\ i \text{---} \text{---} \text{---} \\ \downarrow \quad \downarrow \\ b \quad c \quad f \\ j \quad k \end{array} \rightsquigarrow R_{ik} \begin{pmatrix} d & e \\ c & f \end{pmatrix} \circ_{V_{dc,i}} R_{ij} \begin{pmatrix} a & d \\ b & c \end{pmatrix}. \quad (1.56)$$

The original configuration is thus obtained by gluing all the pieces. However, it still has holes assigned boundary conditions as in figure 4(b), which we must fill by summing over the boundary conditions. To do this, let us recall that the path integral on a finite-length cylinder with boundary condition a imposed on one end defines a closed string state $|a\rangle$ as known as a boundary state. Similarly, the path integral on a disk with no insertion of operators defines a state $|1\rangle$ on the boundary. Assume that we have chosen the set of boundary conditions B to be sufficiently large so that the state $|1\rangle$ can be written as a superposition of boundary states:

$$|1\rangle = \sum_{a \in B} c_a |a\rangle. \quad (1.57)$$

Then, the sum over the boundary conditions gives the state $|1\rangle$ on the boundary of each hole, which is replaced with a disk:

$$\sum_{a \in B} c_a \left(a \text{---} \text{---} \right) = \sum_{a \in B} c_a \left(|a\rangle \text{---} \text{---} \right) = |1\rangle \text{---} \text{---} = \text{---} \text{---} . \quad (1.58)$$

Thus the holes are filled and the original configuration on the torus is reconstructed.

Now that we have understood how to reconstruct the correlation function of line operators in the field theoretic point of view, let us reinterpret this procedure as an operation in statistical lattice model. Rephrasing the procedures so far, it is clear that the above configuration defines a partition function of a statistical spin model. The model has spins located at two kinds of sites, \bigcirc and \odot . They correspond to the open string states and the boundary states; see figure 4(c). A spin at \bigcirc takes values in the chosen basis for the relevant open string states $V_{ab,i}$, while that at \odot is valued in B . The Boltzmann weights for each spin configuration are given by the matrix elements of R-matrix and the coefficients of boundary states c_a . Thus, we conclude that the correlation function of line operators $\{\mathcal{L}_i\}$

on the torus coincides with the partition function of a statistical spin model defined on the periodic lattice formed by the line operators:

$$\left\langle \prod_{i=1}^l \mathcal{L}_i[C_i] \right\rangle_{\mathbb{T}, T^2} = Z_{\mathbb{L}(\mathbb{T})}(\{\mathcal{L}_i[C_i]\}). \quad (1.59)$$

$\mathbb{L}(\mathbb{T})$ denotes the lattice model arising from the TQFT \mathbb{T} with the line operators and $\{\mathcal{L}_i[C_i]\}$ is the lattice formed by line operators \mathcal{L}_i wrapping around C_i . In other words, the partition function of lattice model (1.35) is in this case defined by the lattice $\{\mathcal{L}_i[C_i]\}$, and specified by the left-hand side. By this construction, we can view typical examples of lattice models.

IRF model

If $\dim V_{ab,i} = 1$ for any a, b, i , in this case we can ignore the spins at \bigcirc . We just sum over the boundary conditions and such a spin system $\mathbb{L}(\mathbb{T})$ is called *interaction-round-a-face model*, or *IRF model* for short. The spins are placed on the faces of the lattice, and interaction takes place among four spins surrounding a vertex. The 2d Ising model actually can be formulated as an IRF model by taking vertex-face transformation.

Vertex model

If B consists of a single boundary condition, say a , we can simply write the open string state space by $V_i := V_{aa,i}$.¹¹ The R-matrix is represented only by a crossing of two lines:

$$R_{ij} = i \begin{array}{c} \uparrow \\ \hline \rightarrow \\ \hline \downarrow \\ j \end{array} , \quad (1.60)$$

$$R_{ij} : V_i \otimes V_j \longrightarrow V_j \otimes V_i . \quad (1.61)$$

The space of states V_i can also be thought of as the Hilbert space of a particle propagating on the line operator \mathcal{L}_i . In this case we can ignore the spins at \bigcirc since there is no summation over boundary conditions. This means that the lattice model $\mathbb{L}(\mathbb{T})$ is called a *vertex model*: spins are living on the edges of the lattice and interact with each other at the vertices.

We can always recast our lattice model into a vertex model by setting $V_i := \bigoplus_{a,b \in B} V_{ab,i}$, at least formally, and declaring that all newly introduced R-matrix elements, which correspond to scattering processes with inconsistent Chan-Paton factors, vanish. We can also absorb the coefficients c_a into the R-matrix elements by appropriate rescaling. In what follows we implicitly perform this reformulation, and restrict ourselves to vertex model.

A remarkable aspect of this construction of lattice models is that it allows us to understand integrability from a higher-dimensional point of view. This is crucial observation by Costello [29]. In our lattice model, consider a horizontal line operator \mathcal{L}_i intersecting

¹¹More precisely, for such a case there is actually no boundary condition at all.

vertical line operators \mathcal{L}_j , $j = 1, \dots, n$. Concatenating the R-matrices in this row, we get the row-to-row transfer matrix:

$$T_i = \begin{array}{c} \begin{array}{c} \uparrow \\ \text{---} \end{array} \begin{array}{c} \uparrow \\ \text{---} \end{array} \dots \begin{array}{c} \uparrow \\ \text{---} \end{array} \\ \begin{array}{c} \text{---} \\ \downarrow \end{array} \begin{array}{c} \text{---} \\ \downarrow \end{array} \dots \begin{array}{c} \text{---} \\ \downarrow \end{array} \end{array} = \text{Tr}_{V_i} (R_{i,n} \circ_{V_i} \dots \circ_{V_i} R_{i,1}). \quad (1.62)$$

The hooks on the horizontal line are to remind us that the periodic boundary condition is imposed. This gives an endomorphism of $\bigotimes_{j=1}^n V_j$ which maps a state just below \mathcal{L}_j to another state above it, namely the transfer matrix defines the discrete time evolution of this spin system and $\mathcal{H} := \bigotimes_{j=1}^n V_j$ is the total quantum Hilbert space. In terms of transfer matrices, the partition function is written as a trace,

$$Z_{\text{L}(\mathbb{T})}(\{\mathcal{L}_i[C_i]\}) = \text{Tr}_{\mathcal{H}} (T_{n+m} \dots T_{n+1}). \quad (1.63)$$

An important point is that the underlying field theory is a TQFT, and hence a state evolves trivially on a cylinder unless it hits something, line operators in the present case.

Now suppose that each line operator depends on a continuous parameter which is an element of some smooth manifold S . This parameter is called *spectral parameter* of the lattice model. Let u_i be the spectral parameter of \mathcal{L}_i , then the R-matrix and the transfer matrix are rewritten by

$$R_{ij} \longrightarrow R_{ij}(u_i, u_j), \quad (1.64)$$

$$T_i \longrightarrow T_i(u_i; u_1, \dots, u_n). \quad (1.65)$$

To avoid clutter, we fix u_1, \dots, u_n and suppress them below. A vertex model is said to be integrable if the transfer matrices $T_i(u_i)$ are meromorphic functions of u_i and commute with each other:

$$\begin{array}{c} \begin{array}{c} \uparrow \\ \text{---} \end{array} \begin{array}{c} \uparrow \\ \text{---} \end{array} \dots \begin{array}{c} \uparrow \\ \text{---} \end{array} \\ \begin{array}{c} \text{---} \\ \downarrow \end{array} \begin{array}{c} \text{---} \\ \downarrow \end{array} \dots \begin{array}{c} \text{---} \\ \downarrow \end{array} \end{array} = \begin{array}{c} \begin{array}{c} \uparrow \\ \text{---} \end{array} \begin{array}{c} \uparrow \\ \text{---} \end{array} \dots \begin{array}{c} \uparrow \\ \text{---} \end{array} \\ \begin{array}{c} \text{---} \\ \downarrow \end{array} \begin{array}{c} \text{---} \\ \downarrow \end{array} \dots \begin{array}{c} \text{---} \\ \downarrow \end{array} \end{array} \iff [T_i(u_i), T_j(u_j)] = 0, \quad u_i \neq u_j. \quad (1.66)$$

eq:commuting
gT

When the transfer matrices with different spectral parameters commute, we can find a series of mutually commuting operators on the total Hilbert space $\bigotimes_{j=1}^n V_j$ by Laurent expansion of the transfer matrix. In particular, they commute with the transfer matrix itself, which means they produce an infinite number of conserved quantities. In fact, the commutativity of transfer matrix implies that one can find the exact eigenvalues and eigenvectors of the transfer matrix, which was the underlying reason of integrability of the Ising model.

The situation considered here, namely the addition of spectral parameters and commuting transfer matrix, is naturally realized if the TQFT has “extra dimensions.” In this scenario, we really start with a higher-dimensional theory $\tilde{\mathbb{T}}$ formulated on the product space $S \times T^2$, which is topological on the torus T^2 but not on S . We wrap line operators \mathcal{L}_i around $\{u_i\} \times C_i$, where u_i are points on S . If one can see only the torus T^2 and is unaware of the extra dimensions S , the theory seems to be the previous two-dimensional TQFT $\mathbb{T} \cong \tilde{\mathbb{T}}[S]$ that has parameters taking values in S . One finds that line operators $\mathcal{L}_i(u_i)$ wrapping

around C_i carry continuous parameters u_i in the seemingly two-dimensional theory, and the correlation function for this configuration of line operators is given by the partition function of a lattice model $L(\tilde{T}[S])$ defined on the lattice $\{\mathcal{L}_i(u_i)[C_i]\}$.

For a generic choice of $\{u_i\}$, the transfer matrices of the lattice model $L(\tilde{T}[S])$ commute since the two horizontal line operators such as equation (1.66) can move freely and interchange their positions owing to the topological nature along T^2 ; no phase transition occurs when they pass each other as they do not meet in the full spacetime $S \times T^2$. Thus, integrability follows from the existence of extra dimensions, whose coordinates provide continuous spectral parameters.

In fact, we can deduce integrability from another point of view. By the same logic, we have the unitarity relation

$$\begin{array}{c} i \\ \diagdown \quad \diagup \\ j \end{array} = \begin{array}{c} i \\ \longrightarrow \\ j \end{array} \quad (1.67)$$

$$\Longleftrightarrow R_{ji}(u_j, u_i) R_{ij}(u_i, u_j) = \text{id}_{V_i \otimes V_j}, \quad (1.68)$$

and the Yang-Baxter equation

$$\begin{array}{c} i \\ \diagdown \quad \diagup \\ j \end{array} \begin{array}{c} \longrightarrow \\ \longrightarrow \\ k \end{array} = \begin{array}{c} i \\ \diagdown \quad \diagup \\ j \end{array} \begin{array}{c} \longrightarrow \\ \longrightarrow \\ k \end{array} \quad (1.69)$$

$$\Longleftrightarrow R_{ij}(u_i, u_j) R_{ik}(u_i, u_k) R_{jk}(u_j, u_k) = R_{jk}(u_j, u_k) R_{ik}(u_i, u_k) R_{ij}(u_i, u_j), \quad (1.70)$$

where the R-matrix R_{ij} acts on $V_i \otimes V_j$ as an intertwiner and trivial on V_k , etc. From these two relations we can reproduce the commutativity of the transfer matrix. We should emphasize that the Yang-Baxter equation is a local condition. When the Boltzmann weight locally satisfies the Yang-Baxter equation, it extends to the commutativity of the transfer matrix and hence to the integrability of the model. In this sense, the Yang-Baxter equation is the fundamental condition of integrability of a lattice model.

Before getting into the discussion of brane construction, we would like to generalize the above arguments to further higher-dimensional situations. First of all, replace the two-torus T^2 with a general two-dimensional surface Σ , along which line operators are wrapped. We now have $S \times \Sigma$, and similarly to the above a lattice model is defined on Σ by line operators with spectral parameters. Let us consider the case that we really have more extra dimensions and a higher-dimensional theory is formulated on $S \times M \times \Sigma$, where M is some smooth manifold.¹² In such a case, the line operators we had may descend from extended operators of dimension greater than one. Let T again be the new higher-dimensional theory, and suppose that it is topological on Σ and has extended operators \mathcal{E}_i whose codimension is greater than $\dim S$. Place \mathcal{E}_i on submanifolds of the form $\{u_i\} \times N_i \times C_i$. Since $T[S \times M]$

¹²Note that manifolds M, N in this subsection are not necessarily the same as ones appeared in previous subsections.

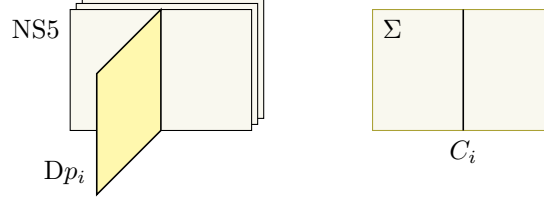


Figure 5: The Dp -brane Dp_i ending on the NS5-branes creates a defect \mathcal{E}_{Dp_i} along C_i .

fig:Dp_on_N
S5

is again regarded as a two-dimensional TQFT defined on Σ , the correlation function of the operators \mathcal{E}_i in the full theory T still should coincide with the partition function of an integrable lattice model $\mathsf{L}(\mathsf{T}[S \times M])$. The model is now defined on the lattice formed by the line operators $\mathcal{E}_i(u_i; N_i)[C_i]$ winding along one-cycles C_i on Σ , which are the image of \mathcal{E}_i in the two-dimensional theory $\mathsf{T}[S \times M]$.

As well as we can view the higher-dimensional theory as a two-dimensional TQFT, we may also view it as a theory $\mathsf{T}[\Sigma]$, which is a QFT on $S \times M$ specified by the surface Σ . In this theory, the operators \mathcal{E}_i are seen as extended operators $\mathcal{E}_i(C_i)$ supported on $\{u_i\} \times N_i$. Then, we have another relation

$$\left\langle \prod_{i=1}^l \mathcal{E}_i(C_i) [\{u_i\} \times N_i] \right\rangle_{\mathsf{T}[\Sigma], S \times M} = Z_{\mathsf{L}(\mathsf{T}[S \times M])}(\{\mathcal{E}_i(u_i; N_i)[C_i]\}). \quad (1.71)$$

eq:correspo
ndence_0

Thus, we finally arrive at a correspondence between a QFT on $S \times M$ equipped with extended operators and an integrable lattice model on Σ .

1.3.2 Brane construction and correspondence

We have seen above that a lattice model is realized by a lattice of line operators in a two-dimensional TQFT, and it is integrable if the TQFT is embedded in higher dimensions and the line operators come from extended operators localized in some directions of the extra dimensions. Now we are ready to explain how to get such structures of correspondence (1.71) using branes in string theory.

Consider a type II string theory in a ten-dimensional spacetime

$$\mathbb{R}^4 \times T^*\Sigma \times \mathbb{R}^2, \quad (1.72)$$

where Σ is a two-dimensional surface embedded in $T^*\Sigma$ as a zero section. Introduce a stack of N NS5-branes supported on $\mathbb{R}^4 \times \Sigma \times \{0\}$ in this spacetime, and Dp -branes Dp_i on $\mathbb{R}^{p-1} \times \Sigma_i \times \{0\}$ ending on the NS5-branes, where \mathbb{R}^{p-1} is a subspace of \mathbb{R}^4 (assuming $p \leq 5$) and Σ_i are surfaces in $T^*\Sigma$ such that $\Sigma_i \cap \Sigma = C_i$; see figure 5. Provided that Σ_i are suitably chosen, this brane system preserves four supercharges.

The low-energy dynamics of the NS5-branes is governed by a six-dimensional theory T_{NS5} on $\mathbb{R}^4 \times \Sigma$. The theory T_{NS5} is depending on whether IIA or IIB string theory we are

considering [30]:

IIA : $\mathcal{N} = (2, 0)$ superconformal QFT of type A_{N-1} ,

IIB : $\mathcal{N} = (1, 1)$ super Yang-Mills theory with gauge group $SU(N)$.

The theory T_{NS5} is formulated on $\mathbb{R}^4 \times \Sigma$, with topological twist along Σ which breaks half of sixteen supercharges. In this twisted theory, Dp_i create p -dimensional defects \mathcal{E}_{Dp_i} on $\mathbb{R}^{p-1} \times C_i$, reducing the number of unbroken supercharges to four. From the point of view of a four-dimensional observer, this brane configuration gives half-BPS defects $\mathcal{E}_{Dp_i}(C_i)$ in an $\mathcal{N} = 2$ theory $T_{NS5}[\Sigma]$. The total system is invariant under a $U(1)$ R-symmetry originating from the rotational symmetry on the \mathbb{R}^2 factor of the ten-dimensional spacetime.

Let us take a three-manifold M and $(p-2)$ -submanifolds N_i of M , and modify the above construction so that the world-volumes of the NS5-branes and the Dp -branes become $S^1 \times M \times \Sigma$ and $S^1 \times N_i \times \Sigma_i$, respectively. At low energies, we get the same theory T_{NS5} formulated on $S^1 \times M \times \Sigma$ and defects \mathcal{E}_{Dp_i} located on $S^1 \times N_i \times C_i$. In general, this modification completely breaks supersymmetry. For certain choices of M and N_i , however, there is a string background in which a fraction of supersymmetry is still preserved. In such a background, the path integral computes the twisted partition function, or *supersymmetric index* of T_{NS5} , defined by a trace with respect to the Hilbert space on $M \times \Sigma$ in the presence of defects \mathcal{E}_{Dp_i} inserted on $N_i \times C_i$. (Notice that for a closed manifold with thermal circle, QFT produces a number represented as a trace, such as (1.3).)

A salient feature of supersymmetric indices is that they are protected against continuous changes of various parameters of the theory. This means that the index of our theory is invariant under deformations of the geometric data of Σ and C_i , such as the metric on Σ and the shapes of C_i . In other words, the theory T_{NS5} on $S^1 \times M \times \Sigma$ is topological on Σ , as far as the computation of the index is concerned.

To relate the present setup to the previous situation (1.71), we apply T-duality along the circle S^1 . It turns Dp_i into $D(p-1)$ -branes $D(p-1)_i$, localized at points u_i along the dual circle \tilde{S}^1 , while sending the NS5-branes to those in the other type II string theory. The new NS5-branes produce the dual six-dimensional theory \tilde{T}_{NS5} on $\tilde{S}^1 \times M \times \Sigma$, and in this theory $D(p-1)_i$ create $(p-1)$ -dimensional defects $\mathcal{E}_{D(p-1)_i}$ on $\{u_i\} \times N_i \times C_i$. Thus we are in the situation studied before, and the correlation function of $\mathcal{E}_{D(p-1)_i}$ for this configuration coincides with the partition function of an integrable lattice model:

$$\left\langle \prod_{i=1}^l \mathcal{E}_{Dp_i}(C_i)[S^1 \times N_i] \right\rangle_{T_{NS5}[\Sigma], S^1 \times M} = Z_L(\tilde{T}_{NS5}[\tilde{S}^1 \times M])\left(\{\mathcal{E}_{D(p-1)_i}(u_i; N_i)[C_i]\}\right). \quad (1.73)$$

Here the left-hand side is expressed in the original frame; it implicitly depends on each spectral parameter u_i through the holonomy $\exp(2\pi i u_i)$ around S^1 of the gauge field for the flavor symmetry $U(1)_i$ supported on Dp_i . The holonomy appears in the index as a refinement parameter, or called *fugacity*, associated with $U(1)_i$.

1.3.3 Defects as transfer matrices

Finally, we apply the construction developed so far to the main theme of this thesis: integrable lattice models and defects as transfer matrices. Let us consider the brane construction for $p = 5$ case. To conform with the standard convention, take S-duality first and we still have D5- and NS5-branes

$$\begin{aligned} \text{ND5} & S^1 \times M \times \Sigma, \\ \text{NS5}_i & S^1 \times N_i \times \Sigma_i, \end{aligned}$$

where NS5_{*i*} create defects $\mathcal{E}_{\text{NS5}_i}$ on $S^1 \times N_i \times C_i$.

One should notice that in this setup we necessarily have $N_i = M$ and thus the defects $\mathcal{E}_{\text{NS5}_i}$ fill the whole $S^1 \times M$, which produces a four-dimensional theory

$$\mathsf{T}_{\text{D5NS5}}[\Sigma]$$

with $\mathcal{N} = 1$ supersymmetry. Then we now have

$$\langle 1 \rangle_{\mathsf{T}_{\text{D5NS5}}[\Sigma], S^1 \times M} = Z_{\mathsf{L}(\mathsf{T}_{\text{D5NS5}}[S^1 \times M])}(\{\mathcal{E}_i(S^1 \times M)[C_i]\}). \quad (1.74)$$

eq:correspondence1

For example, when $M = S^3$, the left-hand side is given by the supersymmetric index for $\mathcal{N} = 1$ theory and the right-hand side corresponds to the partition function of Bazhanov-Sergeev integrable lattice model [31–34]. When $M = L(p, 1)$, lens space, the left-hand side is computed in [35], which defines a new integrable lattice model through this correspondence.

The brane tiling construction of integrable lattice models can be enriched by introducing additional defects. In addition to the previously defined D5NS5-brane system, let us consider a D3-brane such as

$$\text{D3} \quad S^1 \times N \times C \times \mathbb{R}^{\geq}, \quad (1.75)$$

eq:D3

$$\text{D3}' \quad S^1 \times \{0\} \times C' \times D, \quad (1.76)$$

eq:D3'

where N is a curve in M , C and C' are one-cycles on Σ , and D is some two-manifold perpendicular to 5-branes. In field theory side, these D3-branes create new extended defect operators elongated in $S^1 \times N$ and S^1 , respectively. A single D3-brane insertion corresponds to a new oriented line C , C' in the integrable lattice model, which we represent by a dashed line. Now that we have two kinds of lines originating from 5-branes and a D3-brane, we can define three kinds of R-matrices:

$$R = \begin{array}{c} \uparrow \\ \text{---} \text{---} \text{---} \\ \downarrow \end{array}, \quad L = \begin{array}{c} \uparrow \\ \text{---} \text{---} \text{---} \\ \downarrow \end{array}, \quad \mathcal{R} = \begin{array}{c} \uparrow \\ \text{---} \text{---} \text{---} \\ \downarrow \end{array}. \quad (1.77)$$

The middle one is usually called *L-operator*. Correspondingly, we have four Yang-Baxter equations, involving zero to three dashed lines. Those that involving one or two dashed line,

$$\begin{array}{c} \uparrow \\ \text{---} \text{---} \text{---} \\ \downarrow \end{array} = \begin{array}{c} \uparrow \\ \text{---} \text{---} \text{---} \\ \downarrow \end{array} \quad \text{and} \quad \begin{array}{c} \uparrow \\ \text{---} \text{---} \text{---} \\ \downarrow \end{array} = \begin{array}{c} \uparrow \\ \text{---} \text{---} \text{---} \\ \downarrow \end{array} \quad (1.78)$$

are called *RLI relations*. The effect of the insertion of such an additional defect on the lattice model is seen in terms of the L-operator. The neighborhood of the dashed line looks like

$$\hookrightarrow \begin{array}{c} \uparrow \\ | \\ \rightarrow \end{array} \begin{array}{c} \uparrow \\ | \\ \rightarrow \end{array} \cdots \cdots \begin{array}{c} \uparrow \\ | \\ \rightarrow \end{array} \hookrightarrow \quad . \quad (1.79)$$

This diagram shows that the defect acts on the lattice model by a transfer matrix constructed from L-operators. Thus, the insertion of a defect operator in four-dimensional theory represented by (1.75) or (1.76) is mapped into lattice model side as the action of a transfer matrix constructed from L-operators.

What we will discuss in detail in the subsequent sections are the introduction of a single D3-brane (1.75) and (1.76) to the D5NS5-brane system, and investigate the correspondence between supersymmetric gauge theories and integrable lattice models with the additional defects. In the next two sections, we will be studying the followings:

1. For the case of single D3 (1.75), let $M = S^3$, $\Sigma = T^2$, and $N = S^1$. Then the correspondence (1.74) is explicitly written by the equivalence between the supersymmetric index and the partition function of the Bazhanov-Sergeev integrable lattice model [31–34]:

$$\langle 1 \rangle_{\mathcal{T}_{\text{D5NS5}[\Sigma], S^1 \times S^3}} = \mathcal{I}_{S^1 \times S^3}(p, q, t), \quad (1.80)$$

$$Z_{\mathcal{L}(\mathcal{T}_{\text{D5NS5}[S^1 \times S^3])}(\{\mathcal{E}_i(S^1 \times S^3)[C_i]\}) = Z_{\text{Bazhanov-Sergeev}}. \quad (1.81)$$

In this system, the D3 creates a surface defect on $S^1 \times S^1 \subset S^1 \times S^3$ and it acts on the supersymmetric index as a difference operator [36, 37],

$$\langle S_{(r,s)} \rangle_{\mathcal{T}_{\text{D5NS5}[\Sigma], S^1 \times S^3}} = \mathfrak{S}_{(r,s)} \mathcal{I}_{S^1 \times S^3}(p, q, t). \quad (1.82)$$

In the next section we will show that this difference operator are indeed identified with the transfer matrix constructed from L-operators of elliptic type:

$$\mathfrak{S}_{(r,s)} = \text{elliptic transfer matrix}. \quad (1.83)$$

2. For the case of single D3' (1.76), let $M = \mathbb{R}^3$ and $\Sigma = T^2$. D3' creates a line defect wrapping the thermal circle S^1 and it acts on the quantum Hilbert space of a spin chain as a transfer matrix, since now we have non-compact three-manifold \mathbb{R}^3 . This spin chain has an equivalent lattice model description, and it is really an integrable model of trigonometric type. The line defect in the four-dimensional theory is realized as a Wilson-'t Hooft line operator T , whose magnetic charge and electric charge (\mathbf{m}, \mathbf{e}) also have the counterpart in the lattice model side. The vacuum expectation values (vevs) of the Wilson-'t Hooft lines realize the deformation quantization of the Hitchin moduli space [38, 39], and are naturally quantized by the Weyl quantization:

$$\text{Weyl quantization of } \langle T_{(\mathbf{m}, \mathbf{e})} \rangle = \text{trigonometric transfer matrix}. \quad (1.84)$$

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