

QUANTUM TOPOLOGY AND CATEGORIFICATION SEMINAR, SPRING 2017

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Part 1. Quantum topology: Chern-Simons theory and the Jones polynomial

1. THE JONES POLYNOMIAL: 1/24/17

Today, Hannah talked about the Jones polynomial, including how she sees it and why she cares about it as a topologist.

1.1. Introduction to knot theory.

Definition 1.1. A **knot** is a smooth embedding $S^1 \hookrightarrow S^3$. We can also talk about **links**, which are embeddings of finite disjoint unions of copies of S^1 into S^3 .

One of the major goals of 20th-century knot theory was to classify knots up to isotopy.

Typically, a knot is presented as a **knot diagram**, a projection of $K \subset S^3$ onto a plane with “crossing information,” indicating whether the knot crosses over or under itself at each crossing. Figure 1 contains an example of a knot diagram.



FIGURE 1. A knot diagram for the left-handed trefoil knot. Source: [Wikipedia](#).

Given a knot in S^3 , there's a theorem that a generic projection onto \mathbb{R}^2 is a knot diagram (i.e. all intersections are of only two pieces of the knot).

Link diagrams are defined identically to knot diagrams, but for links.

Theorem 1.2. Any two link diagrams for the same link can be related by planar isotopy and a finite sequence of Reidemeister moves.

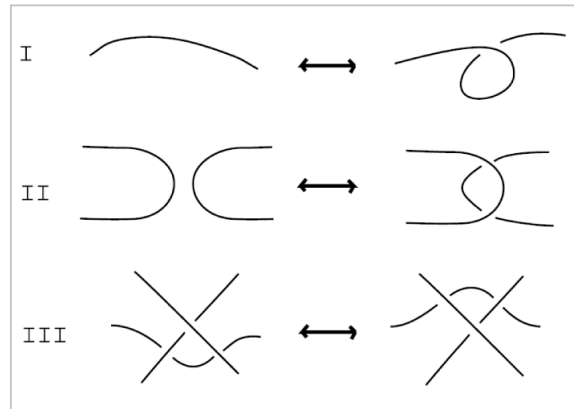


FIGURE 2. The three Reidemeister moves. Source: <https://www.computer.org/csdl/trans/tg/2012/12/ttg2012122051.html>.

1.2. Polynomials before Jones. The first knot polynomial to be defined was the Alexander polynomial $\Delta_K(x)$, a Laurent polynomial with integer coefficients that is a knot invariant, defined in the 1920s.

Here are some properties of the Alexander polynomial:

- It's symmetric, i.e. $\Delta_K(x) = \Delta_K(x^{-1})$.
- It cannot distinguish handedness. That is, if K is a knot, its **mirror** \bar{K} is the knot obtained by switching all crossings in a knot diagram, and $\Delta_K(x) = \Delta_{\bar{K}}(x)$.¹
- The Alexander polynomial doesn't detect the unknot (which is no fun): there are explicit examples of knots 11_{34} and 11_{42} whose Alexander polynomials agree with that of the unknot.²

So maybe it's not so great an invariant, but it's somewhat useful.

1.3. The Jones polynomial. The Jones polynomial was defined much later, in the 1980s. The definition we give, in terms of skein relations, was not the original definition. There are three local models of crossings, as in Figure 3.



FIGURE 3. The three local possibilities for a crossing in a knot diagram (technically, L_0 isn't a crossing). Source: https://en.wikipedia.org/wiki/Skein_relation.

The idea is that, given a knot K , you could try to calculate a knot polynomial for K in terms of knot polynomials on links where one of the crossings in K has been changed from L_- to L_+ (or vice versa), or **resolved** by replacing it with an L_0 . A relationship between the knot polynomials of these three links is a **skein relation**. This is a sort of inductive calculation, and the base case is the unknot. In particular, you can use the value on the unknot and the skein relations for a knot polynomial to describe the knot polynomial!

Example 1.3. The Alexander polynomial is determined by the following data.

- On the unknot, $\Delta(U) = 1$.
- The skein relation is $\Delta(L_+) - \Delta(L_-) = t\Delta(L_0)$.

Definition 1.4. The **Jones polynomial** is the knot polynomial v determined by the following data.

- For the unknot, $v(U) = 1$.

¹The mirror of the left-handed trefoil is the right-handed trefoil, for example.

²The notation for these knots follows Rolfsen.

- The skein relation is

$$(t^{1/2} - t^{-1/2})v(L_0) = t^{-1}v(L_+) - tv(L_-).$$

Example 1.5. Let's calculate the Jones polynomial on a Hopf link H , two circles linked together once. The standard link diagram for it has two crossings, as in Figure 4.

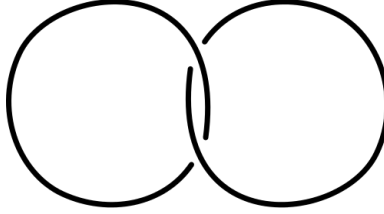


FIGURE 4. A Hopf link. Source: https://en.wikipedia.org/wiki/Link_group.

- Resolving one of the crossings produces an unknot: $v(L_0) = 1$.
- Replacing the L_- with an L_- produces two unlinked circles. One more skein relation produces the unknot, so $v(L_+) = -(t^{1/2} - t^{-1/2})$.

Putting these together, one has

$$(t^{1/2} - t^{-1/2}) \cdot 1 = t^{-1}(-(t^{1/2} + t^{-1/2}) - tv(H)),$$

$$\text{so } v(H) = -t^{-1/2} - t^{-5/2}.$$

◀

There are many different definitions of the Jones polynomial; one of the others that we'll meet later in this seminar is via the Kauffman bracket.

Definition 1.6. The **bracket polynomial** of an unoriented link L , denoted $\langle L \rangle$, is a polynomial in a variable A defined by the skein relations

- On the unknot: $\langle O \rangle = 1$.
- There are two ways to resolve a crossing C : as two vertical lines V or two horizontal lines H . We impose the skein relation

$$\langle C \rangle = A\langle V \rangle + A^{-1}\langle H \rangle.$$

- Finally, suppose the link L is a union of one unlinked unknot and some other link L' (sometimes called the **distant union**). Then,

$$\langle L \rangle = (-A^2 - A^{-2})\langle L' \rangle.$$

Example 1.7. Once again, we'll compute the Kauffman bracket for the Hopf link. (TODO: add picture).

The result is $\langle H \rangle = -A^4 - A^{-4}$.

◀

You can show that this bracket polynomial is invariant under type II and III Reidemeister moves, but not type I. We obviously need to fix this.

Definition 1.8. Let D be an *oriented* link, and $|D|$ denote the link without an orientation. The **normalized bracket polynomial** is defined by

$$X(D) := (-A^3)^{-\omega(D)}\langle |D| \rangle.$$

Here, $\omega(D)$ is the writhe of D , an invariant defined based on a diagram. At each crossing, imagine holding your hands out in the shape of the crossing, where (shoulder \rightarrow finger) is the positively oriented direction along the knot. If you hold your left hand over your right hand, the crossing is a **positive crossing**; if you hold your right hand over your left, it's a **negative crossing**.

Let ω_+ denote the number of positive crossings and ω_- denote the number of negative crossings. Then, the **writhe** of D is $\omega(D) := \omega_+ - \omega_-$. For example, the writhe of the Hopf link (with the standard orientation) is 2, and $X(H) = -A^{10} - A^2$.

Thankfully, this is invariant under all types of Reidemeister moves. The proof is somewhat annoying, however.

Theorem 1.9. By substituting $A = t^{-1/4}$, the normalized bracket polynomial produces the Jones polynomial.

So these two invariants are actually the same.
Here are some properties of the Jones polynomial.

- $v_{\bar{K}}(t) = v_K(t^{-1})$. Since the Jones polynomial is not symmetric, it can sometimes distinguish handedness, e.g. it can tell apart the left- and right-handed trefoils.
- It fails to distinguish all knots: once again, 11_{34} and 11_{42} have the same Jones polynomial.³
- It's unknown whether the Jones polynomial detects the unknot: there are no known nontrivial knots with trivial Jones polynomial.
- Computing the Jones polynomial is **P**-hard: there's no polynomial-time algorithm to compute it. (Conversely, the Alexander polynomial is one of very few knot invariants with a polynomial-time algorithm.)

If a knot does have trivial Jones polynomial, we know:

- it isn't an **alternating knot** (i.e. one where the crossings alternate between positive and negative).
- It has crossing number at least 18 (which is *big*).

One interesting application of what we'll learn in this seminar is that there are knots (9_{42} and 10_{11}) that can't be distinguished by the Jones or Alexander (or HOMFLY, or ...) but *are* distinguished by $SU(2)$ -Chern-Simons invariants.

2. INTRODUCTION TO QUANTUM FIELD THEORY: 1/31/17

Today, Ivan talked about quantum field theory (QFT), including what QFT is, why one might want to study it, how it relates to other physical theories, classical field theories, and quantum mechanics, and how to use canonical quantization to produce a QFT.

So, why should we study QFT? One good reason is that its study encompasses a specific example, the **Standard model**, the “theory of almost everything.” This is a theory that makes predictions about three of the four fundamental forces of physical reality (electromagnetism, the weak force, and the strong force), leaving out gravity. These predictions have been experimentally verified, e.g. by the Large Hadron Collider.

Unfortunately, the mathematical theory of QFTs is not well formulated; **free theories** are well understood, but if you can rigorously formulate the mathematical theory of **interacting QFTs**, you'll win a million-dollar prize! Perhaps that's a good reason to study QFT.

There's also the notion of a topological quantum field theory (TQFT), which has been rigorously formulated as mathematics, but many of the most important QFTs, including the Standard Model, do not fit into this framework.

QFTs fit into a table with other physical theories: the theory you want to use depends on how fast your particles move and how big they are.

- If your particles are larger than atomic scale and moving considerably slower than the speed of light c , you use *classical mechanics*.
- If your particles are atomic-scale, but moving much slower than c , you use *quantum mechanics*.
- If your particles are larger than atomic-scale, but moving close to the speed of light, you use *special relativity* or *general relativity*: the latter if you need to account for gravity, and the former if you don't.
- If your particles are at atomic-scale and moving close to the speed of light, but you don't need to take gravity into account, you use *quantum field theory*. In this sense, QFT is the marriage of special relativity and quantum mechanics.
- If your particles are small, but moving at about c , *and* you need to consider gravity, you end up in the domain of *string theory*. Here be dragons, of course: string theory hasn't been experimentally verified yet...

With the big picture in place, let's talk a little about classical field theory.

Let $\mathbb{R}^{1,3}$ denote **Minkowski spacetime**, \mathbb{R}^4 with the normal **Minkowski metric**

$$g_{\mu\nu} = \begin{pmatrix} 1 & & & \\ & -1 & & \\ & & -1 & \\ & & & -1 \end{pmatrix}.$$

Definition 2.1. A **field** is a section of a vector bundle over $\mathbb{R}^{1,3}$, or a connection on a principal G -bundle over $\mathbb{R}^{1,3}$. In the latter case, it's also called a **gauge field**.

³This is ultimately for the same reason as for the Alexander polynomial: there's a technical sense in which they're **mutant knots** of each other. It's notoriously hard to write down knot polynomials that detect mutations, and the Jones polynomial cannot detect them.

In this context, we care the most about trivial vector bundles and principal bundles!

Definition 2.2. A **classical field theory** is a collection of PDEs that specify the time evolution of a collection of fields.

Example 2.3. Electromagnetism is a famous example of a classical field theory: there are electric and magnetic fields \vec{E} and \vec{B} , respectively, and the **Maxwell equations** govern how they evolve in time:

$$\begin{aligned}\nabla \cdot \vec{E} &= \rho \\ \nabla \cdot \vec{B} &= 0 \\ \nabla \times \vec{E} + \frac{\partial \vec{B}}{\partial t} &= 0 \\ \nabla \times \vec{B} - \frac{\partial \vec{E}}{\partial t} &= J.\end{aligned}$$

Here, J is the **electric current** and ρ is the **charge density**. There may be some constants missing here. \blacktriangleleft

Usually (always?), you can present the evolution of the classical field theory as the “critical points” of a functional of the form⁴

$$S(\varphi_1, \dots, \varphi_n) = \int_{\mathbb{R}^4} d^4x \mathcal{L}(\varphi_1, \dots, \varphi_n, \partial_\mu \varphi_1, \dots, \partial_\mu \varphi_n).$$

The functional S is called the **action functional**, and the function \mathcal{L} is called the **Lagrangian**. Using calculus of variations, this notion of critical points is placed on sound footing. Physicists sometimes call these critical points **minimizers**, but sometimes we want to maximize S , not minimize it.

In this context, one can show that the critical points of S are the solutions to the **Euler-Lagrange equations**. In the case of a single field φ , these equations take the form

$$\frac{\partial \mathcal{L}}{\partial \varphi} - \partial_\mu \left(\frac{\partial \mathcal{L}}{\partial (\partial_\mu \varphi)} \right) = 0.$$

(We are using and will continue to use Einstein notation: any index μ that’s both an upper and lower index has been implicitly summed over.) So the Lagrangian contains all the information about the dynamics of the system.

Example 2.4. Let’s look at electromagnetism again: if $\rho = 0$ and $J = 0$, then let

$$A := A_\mu dx^\mu$$

be the **electromagnetic potential**. If $F = dA$, then

$$F = F_{\mu\nu} dx^\mu \wedge dx^\nu.$$

Then, the Lagrangian is

$$\mathcal{L}_{\text{Maxwell}} := -\frac{1}{4} F_{\mu\nu} F^{\mu\nu}, \quad (2.5)$$

where $F^{\mu\nu} = g^{\mu\alpha} g^{\nu\beta} F_{\alpha\beta}$ and $g^{\mu\nu}$ denotes the coefficients of the standard Minkowski metric. Then, the Euler-Lagrange equations and the fact that $dF = 0$ (since F is already exact) directly imply the Maxwell equations, where

$$F_{\mu\nu} = \begin{pmatrix} 0 & E_1 & E_2 & E_3 \\ -E_1 & 0 & B_3 & B_2 \\ -E_2 & B_3 & 0 & -B_1 \\ -E_3 & -B_2 & B_1 & 0 \end{pmatrix}.$$

\blacktriangleleft

Definition 2.6. A **free field theory** is one whose Lagrangian is quadratic in the fields and their partial derivatives. A field theory which is not free is called **interacting**.

In a free field theory, the Euler-Lagrange equations become linear, making them much easier to solve.

⁴To be precise, we should say what space of functions this takes place on. The right way to do this is to consider distributions, but we’re not going to delve into detail about this.

Example 2.7. One example of a free field theory uses the **Dirac Lagrangian**

$$\mathcal{L}_{\text{Dirac}} := \bar{\psi}(i\gamma^\mu \partial_\mu - m)\psi,$$

where $\mu: U \subset \mathbb{R}^{1,3} \rightarrow \mathbb{C}^4$, γ^μ are **Dirac matrices**, and m is a **mass parameter**, and $\bar{\psi} = \psi^\dagger \gamma^0$. This is used to describe the behavior of a free fermion (e.g. an electron). You can explicitly check this is quadratic in ψ and γ .

The Maxwell Lagrangian (2.5) also defines a free field theory. \blacktriangleleft

Example 2.8. Here's an example of an interacting field theory; its classical solutions don't represent anything physical, but we'll see it again.

$$\mathcal{L}_{\text{QED}} := \mathcal{L}_{\text{Dirac}} + \mathcal{L}_{\text{Maxwell}} + ie\bar{\psi}\gamma^\mu\psi A_\mu. \quad (2.9)$$

Here e is the charge of an electron, not ≈ 2.78 . The first two terms are free, but then it's coupled to an interacting term. \blacktriangleleft

From classical to quantum. To understand how we move from classical field theory to quantum field theory, we'll learn about quantum mechanics, albeit very quickly. This formalism extracts three aspects of a physical system.

- The **states** are the configurations that the system can be in.
- The **observables** are things which we can measure/observe about a system.
- **Time evolution** describes how observables or states evolve with time.

In quantum mechanics:

- The states are unit vectors in some (complex) Hilbert space \mathcal{H} .
- The observables are self-adjoint operators $A: \mathcal{H} \rightarrow \mathcal{H}$. They are not necessarily bounded. The things you can measure for A are in its spectrum $\text{Spec} A \subset \mathbb{R}$ (since A is self-adjoint). For example, if A represents the position in a coordinate you chose, the spectrum denotes the set of allowed positions in that coordinate.
- Time evolution has two equivalent formulations.
 - The **Schrödinger picture** describes time evolution of the states. There's a distinguished observable, usually representing the energy of the system, called the **Hamiltonian** $H: \mathcal{H} \rightarrow \mathcal{H}$. Then, a state $\psi \in \mathcal{H}$ in this system evolves as

$$i\hbar \frac{\partial}{\partial t} \psi(t) = H\psi(t).$$

- The **Heisenberg picture** describes time evolution of observables as satisfying the equation

$$\frac{d}{dt} A(t) = i\hbar [H, A(t)].$$

These two perspectives predict the same physics.

Generally, quantum field theories are obtained by taking a classical field theory and quantizing it. This is a process creating a dictionary based on the one between classical mechanics and quantum mechanics:

- The states in classical mechanics are points in T^*M , where M is a smooth manifold; quantum mechanics uses a Hilbert space.
- The observables in classical mechanics are smooth functions $T^*M \rightarrow \mathbb{R}$. In coordinates $(q^1, \dots, q^n, p_1, \dots, p_n)$, we have relations

$$\{q^i, q^j\} = 0 \quad \{p_i, p_j\} = 0 \quad \{q^i, p_j\} = \delta_j^i,$$

where $\{-, -\}$ is the Poisson bracket coming from the symplectic structure on T^*M . Quantum mechanics replaces functions with self-adjoint operators. In quantum mechanics, if X_i and P_i are the position and momentum operators in coordinate i , they satisfy the relations

$$[X_i, X_j] = 0 \quad [P_i, P_j] = 0 \quad [X_i, P_j] = i\hbar \delta_{ij}.$$

- Time evolution in classical mechanics satisfies $\frac{d}{dt} \gamma(t) = \{H, \gamma(t)\}$; quantum mechanics assigns the Schrödinger or Heisenberg pictures as above.

So for a classical field theory, we want a way to get a Hilbert space, a Hamiltonian, and position and momentum operators.

Example 2.10 (One-dimensional harmonic oscillator). The harmonic oscillator in one dimension satisfies the equation

$$H(x, p) = \frac{p^2}{2m} + n\omega^2 x^2.$$

Then, the Hilbert space is $\mathcal{H} = L^2(\mathbb{R})$, $Xf = x \cdot f$, and $Pf = -i\hbar \frac{\partial f}{\partial x}$.⁵ These automatically satisfy the relations $[X, P] = 1$, $[X, X] = 0$, and $[P, P] = 0$. The Hamiltonian is

$$H(X, P) = \frac{P^2}{2m} + m\omega^2 X^2.$$

◀

It's worth noting that quantization is not a deterministic process, more of an art: choosing the right position and momentum operators and showing why they satisfy the relations doesn't follow automatically from some general theory. But if you can get the commutation relations to work and it describes a physical system, congratulations! You've done quantization.

Next time, we'll discuss quantum field theory, where the fields are replaced with quantum fields.

3. CANONICAL QUANTIZATION AND CHERN-SIMONS THEORY: 2/7/17

Today, Jay will say some more things about quantum field theory. First, he'll discuss a setup for QFT, including some handwaving about canonical quantization, and some path integrals. Then, there will be some discussion of gauge theory and connections on G -bundles, and a little bit about Chern-Simons theory.

3.1. Quantum field theory and path integrals. Just like for quantum mechanics, fix a Hilbert space \mathcal{H} of states. Quantum fields are the things that we use to take measurements, more or less: operator-valued distributions over spacetime. It's helpful to think of them as simply operators on \mathcal{H} .

The general way canonical quantization works is that one desires quantum fields $\phi(x), \pi(x)$ which satisfy relations similar to the ones that positions and momenta do in classical mechanics.⁶ For example, their commutator $[\phi(x), \pi(y)]$ is another operator-valued distribution, and the position-momentum constraint is that $[\phi(x), \pi(y)] = \delta(x - y) \cdot 1$, analogous to the Poisson bracket for classical mechanics.

QFT actually computes things called **scattering amplitudes**: if you throw a bunch of particles at a bunch of other particles, sometimes new particles come out. Scattering amplitudes encode the probability of getting a particular new particle from a particular collision of old particles. These are computed with **propagators**, distributions of the form $D(x - y) = \langle 0 | \Phi(x)\Phi(y) | 0 \rangle$, where $|0\rangle$ is the vacuum state (the lowest-energy state), not the zero vector of \mathcal{H} . Here $x, y \in \mathbb{R}^{1,3}$, and we assume x is “after” y , in that $x_0 > y_0$. Physicists think of this as the creation of a particle at y , followed by its annihilation at x , and from this other things can be built, so if you want to compute anything, this is a good place to start.

The propagator can be computed in terms of a path integral, which has the advantage that the quantum fields are replaced with classical fields inside the integral:

$$\langle 0 | \phi(x)\phi(y) | 0 \rangle = \frac{\int D\phi \phi(x)\phi(y) e^{iS(\phi)/\hbar}}{\int D\phi e^{iS(\phi)/\hbar}},$$

where

$$S(\phi) = \int \mathcal{L}(\phi, \partial_\mu \phi) d^4x$$

and $D\phi$ is a “measure” on the space of fields, which famously still hasn't been made rigorous. Nonetheless, there is a theory for calculating path integrals, which boils down to Feynman diagrams and Wick contractions.

Unfortunately, approaching this systematically is generally done heuristically. The quantum fields are supposed to encode the amplitudes of the Fourier transforms of the classical fields, but making this precise doesn't come easily.

⁵Well, not all L^2 functions are differentiable, but there are ways of working around this, especially since differentiable functions are dense in L^2 .

⁶Some of what follows requires additional indices unless these are scalars (functions on the space), but thinking in terms of scalars is helpful for now.

The path integral is so nice because it allows you to avoid an explicit quantization — you can compute things such as expectation values in terms of the classical field theory. For example, if \mathcal{M} is a space of (classical) fields and F is a functional on those fields, the path integral allows you to compute the vacuum expectation value:

$$\langle 0 | F | 0 \rangle = \frac{\int D\phi e^{-iS(\phi)} F(\phi)}{\int D\phi e^{-iS(\phi)}},$$

where S is as before, written in terms of the Lagrangian.

Intuition for computing the path integral: for a zero-dimensional field theory, fields are numbers, and the path integral reduces to ordinary Gaussian integrals. The tricks that we use to compute these generalize to higher-dimensional theories, in some vague sense.

Another trick is to argue as to why the majority of the measure is concentrated at extremal points of the Lagrangian.

A third trick is to discretize the quantum field theory into a lattice model, which approximates the path integral by an ordinary integral which can be computed rigorously. The hope is to take a limit as the lattice approximation gets finer and finer, but this is mysterious in general.

3.2. Gauge theory. Gauge theories are those in which the fields are connections on principal G -bundles, where G is a compact Lie group.

That was a lot of words. Here's what some of them mean.

Definition 3.1. Let G be a compact Lie group.

- A **G -torsor** is a space with a simply transitive left action on G , necessarily a manifold diffeomorphic to G and with an isomorphic left action, but without an origin.
- A **principal G -bundle** is a fiber bundle $P \rightarrow M$, where the fibers are G -torsors in a smoothly varying way.
- A **connection** on a principal bundle $P \rightarrow M$ is a \mathfrak{g} -valued 1-form on the total space of P , where \mathfrak{g} is the Lie algebra of G .

The definition of a connection isn't super intuitive, but it's a way of defining parallel transport between the fibers of P . The tangent space to a point in P splits as the direct sum of the tangent space to M and the tangent space of G , which is \mathfrak{g} , so the 1-form is a way of projecting down to \mathfrak{g} , which is a local parallel transport.

If $P \rightarrow M$ is trivial (meaning it's just the projection $G \times M \rightarrow M$), the zero section defines a map $t : M \rightarrow G \times M$, and pulling back the connection along t allows one to think of it as a \mathfrak{g} -valued 1-form on M . In physics terminology, fixing this trivialization is called **fixing a gauge**.

Example 3.2. Electromagnetism is a gauge theory, with $G = U(1)$ and $\mathfrak{g} = \mathbb{R}$ with trivial Lie bracket. Physicists think of the vector potential as a covector field on spacetime, rather than on $U(1) \times \mathbb{R}^{1,3}$; this means that the gauge has already been implicitly fixed. \blacktriangleleft

There's a natural way to associate a \mathfrak{g} -valued 2-form of a connection called its **curvature**. The curvature of A is defined to be

$$F_A := dA + \frac{1}{2}[A, A].$$

In electromagnetism, this associates the electromagnetic 2-tensor F_{ij} to the vector potential.

The Lagrangian for a gauge theory usually depends only on the curvature of the connection, e.g. in Chern-Simons theory, which we'll discuss below.

Chern-Simons theory is motivated by a few theorems in pure mathematics. The **adjoint action** $\text{Ad } G$ of G on \mathfrak{g} is the derivative of the action of G on itself by conjugation.

Theorem 3.3 (Chern-Weil). *Given an $\text{Ad } G$ -invariant polynomial h on \mathfrak{g} of degree k , one can build a $2k$ -form $T(h, A) = h(F_A, \dots, F_A)$, and this 2-form is always exact: there's a $(2k-1)$ -form $\text{CS}(A)$ such that $d\text{CS}(A) = T(h, A)$, defined as*

$$\text{CS}(h, A) = \int_0^1 h(A, F_A, F_A, \dots, F_A) dt.$$

*This form is called the **Chern-Simons form**.*

You can do this basis-independently: the $\text{Ad } G$ -invariant polynomials on \mathfrak{g} of degree k are the space $\text{Sym}^k(\mathfrak{g}^*)$. If h is the quadratic polynomial associated to the Killing form, then the Chern-Simons form is

$$\text{CS}(h, A) = \text{Tr} \left(A \wedge dA + \frac{1}{3} A \wedge [A, A] \right).$$

It's possible to show that $\int_M \text{CS}(h, A) dM$ is a topological invariant, which suggests that the field theory with **Chern-Simons action**

$$S(A) := \frac{\ell}{4\pi} \int_M \text{CS}(h, A) dM$$

should have interesting topological properties. The Euler-Lagrange equations for this classical field theory boil down to $F_A = 0$ (these connections are called **flat connections**); in other words, the classical fields are flat connections. This is very useful: the moduli space of connections on G -bundles over a manifold M is really messy, but restricting to flat connections, you obtain something finite-dimensional.

The last thing we'll talk about today are examples of observables called Wilson loop operators. Let K be a functional on the space of fields, and $L \subset M$ be a link with components L_i . To each L_i , choose a finite-dimensional real representation R_i of G . Let

$$W_{R_i}(K_i) := \text{Tr} \exp \oint_{K_i} A.$$

Its expectation, called the **Wilson loop operator**, is a path integral

$$\langle K \rangle = \int DA \exp \left(\frac{1}{k} S(A) \right) \prod_{i=1}^r W_{R_i}(K_i).$$

The Jones polynomial will be one of these Wilson loop operators.

4. CHERN-SIMONS THEORY AND THE WESS-ZUMINO-WITTEN MODEL: 2/14/17

These are Arun's notes for his talk about Chern-Simons theory.

4.1. Functorial TQFT and CFT. Mired as we were in physics, let's zoom out a little bit. Chern-Simons theory ought to be a topological quantum field theory, meaning it should be possible to understand it with pure mathematics. Though Witten does not do this, Gill and Adrian will be speaking about a paper which adopts a more mathematical approach, and this mathematical notion of TQFT will also come up in the second half of the seminar.

Informally, a TQFT is the categorified notion of a bordism invariant. That is, equivalence classes closed n -manifolds up to bordism form an abelian group Ω_n under disjoint union, and a bordism invariant is a group homomorphism from Ω_n into some other abelian group. For example, the signature of the intersection pairing is a bordism invariant $\Omega_4^{\text{SO}} \rightarrow \mathbb{Z}$.

The categorified notion of an abelian group⁷ is a **symmetric monoidal category**, a category \mathcal{C} with a functor $\otimes : \mathcal{C} \times \mathcal{C} \rightarrow \mathcal{C}$ that is (up to natural isomorphism) associative, commutative, and unital. A **symmetric monoidal functor** is a functor between symmetric monoidal categories that preserves the product.

Example 4.1.

- (1) Complex vector spaces are a symmetric monoidal category under tensor product; the unit is \mathbb{C} .
- (2) For any n , there is a **bordism category** Bord_n whose objects are closed n -manifolds and whose morphisms are bordisms between them (so a bordism X with incoming boundary M and outgoing boundary N defines a morphism $M \rightarrow N$). Its monoidal product is disjoint union, and the unit is the empty set, regarded as an n -manifold. ◀

There are several related bordism categories, where the manifolds are oriented, spin, etc. Another example is manifolds and bordisms with **conformal structure**, the data of a Riemannian metric up to scaling, which form a bordism category we'll call $\text{Bord}_n^{\text{conf}}$.

Definition 4.2. An n -dimensional **topological quantum field theory** (TQFT) is a symmetric monoidal functor $Z : \text{Bord}_n \rightarrow (\text{Vect}_{\mathbb{C}}, \otimes)$.

⁷Technically of a commutative monoid.

So for every n -manifold M , there's a complex vector space $Z(M)$; for every bordism there's a linear map of vector spaces; and disjoint unions are mapped to tensor products.

This concept is called **functorial TQFT**. The idea is it should encompass the properties of a QFT that only depend on topological information, e.g. $Z(M)$ is the space of states of the theory on the manifold M .

Definition 4.3. An n -dimensional **conformal field theory** (CFT) is a symmetric monoidal functor $Z : \text{Bord}_n^{\text{conf}} \rightarrow (\text{Vect}_{\mathbb{C}}, \otimes)$.

This is basically the same; for CFT, though, the theory is allowed to depend on geometric information, as long as such information is scale-invariant.

Describing Chern-Simons theory as a functorial TQFT would be pretty cool, making all the calculations for the Jones polynomial much easier, and that's what Adrian and Gill are going to tell us about. But there's a wrinkle: Chern-Simons theory is **anomalous**, in that it depends on slightly more than a topological structure. There are a couple of approaches you can take here.

- One way to deal with this is to add structure that makes the anomaly vanish, so as a theory of manifolds with that structure, Chern-Simons theory is a TQFT. For example, the anomaly is trivial on framed manifolds, which is why Witten's calculations require playing with framings of knots. The anomaly vanishes on weaker structures, e.g. a signature structure or a trivialization of p_1 (the latter is the approach BHMV use), and these structures are easier to deal with.
- There's a sense in which Chern-Simons theory is "topological in one direction and conformal in the other two," and the anomaly arises from the conformal part. So you could choose a 3-manifold $M^3 = \Sigma^2 \times C^1$ such that the theory is topological in C and conformal in Σ and study each part separately. Witten uses this approach to compute some state spaces, which I'll tell you about in a little bit.

Today, though, we're going to continue the physics-based approach.

4.2. Connections and the Chern-Simons form. A principal G -bundle is a generalization of a covering space, together with its covering group. Let G be a finite group and $p : \tilde{X} \rightarrow X$ be a covering space with deck transformation group G . Then, every $x \in X$ has a neighborhood U such that $p^{-1}(U) \cong U \times G$ in a way that commutes with projection to U and the G -action. If you replace "finite group" with "Lie group" in the previous sentence, you obtain the definition of a principal G -bundle.

We also defined a connection on a principal G -bundle $P \rightarrow X$ to be a \mathfrak{g} -valued one-form on the total space P . These can be used to define parallel transport, just like connections on vector bundles. Given a connection A , $A \wedge dA + (1/3)A \wedge [A, A] \in \Omega_X^3(\mathfrak{g})$ (after pulling back along the zero section $X \rightarrow P$), so its trace is a real-valued 3-form that can be integrated. This is what the Chern-Simons action does:

$$S(A) := \frac{k}{2\pi} \int_M \text{Tr} \left(A \wedge dA + \frac{1}{3} A \wedge [A, A] \right), \quad (4.4)$$

where $k \in \mathbb{Z}$ is called the **level** of the theory. This defines a classical field theory as we discussed, and quantum Chern-Simons theory is its quantization. Working the quantization out is extremely difficult!

Anyways, assuming we can do that, we'd like to get some useful invariant out of the theory. If A is a connection on $P \rightarrow X$, it defines parallel transport. Just like parallel transport on a Riemannian manifold, this is locally unique but not globally unique, and given a curve $K \subset X$, we can ask what happens when we take something at a point, parallel-transport it around K , and compare the final result with the initial result.

Another analogy is to think about regular covering spaces (so for finite G): if you start with an x in the fiber $p^{-1}(y)$ and wind around K back to $p^{-1}(y)$, you might find yourself at $x' \neq x$. Since the action of G is transitive, $x' = gx$ for some $g \in G$, and g turns out to depend only on the conjugacy class of x . This g is called the **holonomy** of x around K , denoted $\text{hol}_K(x)$. Exactly the same definition applies to principal G -bundles in general (except that the connection is needed to define parallel transport).

It's easier to deal with numbers than with elements of G , so we'll take a trace: let $\rho : G \rightarrow \text{GL}(V)$ be a representation of G . Let

$$W_V(K) := \text{Tr}(\rho(\text{hol}_K(\text{id}_G))).$$

That is: parallel-transport the identity element around G , take its action on V , and take the trace.

Let $K = \coprod_i K_i$ be a link, and $\vec{V} = \{\rho_i : G \rightarrow \mathrm{GL}(V_i)\}_i$ be a collection of (finite-dimensional, real) G -representations. Then, we define the **Wilson loop operator** associated to K and \vec{V} to be

$$\langle K \rangle := \int DA e^{ikS(A)/4\pi} \prod_i W_{V_i}(K_i),$$

where $S(A)$ is the Chern-Simons action. This sketchy path integral integrates out the connection, so this is now an invariant of G , K , \vec{V} , and k . For judicious choices of \vec{V} , these will contain the data of the Jones polynomial.

From a physics perspective, Wilson loops encode generalized charge. If this were QED, a Wilson loop would measure the charge picked up by traveling around the loop in question, in the context of electromagnetic fields. It will sometimes be helpful to think of the representations as generalized charges.

4.3. Connections to Wess-Zumino-Witten CFT. Great, so how do you calculate anything? Witten's idea is to cut S^3 with an embedded link into pieces to derive the skein relation, meaning you just have to understand Chern-Simons theory on boundaries and calculate some state spaces for $\Sigma \times \mathbb{R}$ (a collar neighborhood for the Riemann surface Σ).

Herein lies a problem: the Chern-Simons action (4.4) is not gauge-invariant on a manifold with boundary. Oops! This is a manifestation of the anomaly mentioned earlier.

To fix this, you have to add a boundary term. This boundary term describes a 2D conformal field theory called the Wess-Zumino-Witten (WZW) model,⁸ and the state spaces of Chern-Simons theory are identified with the conformal blocks of the WZW theory. The conformal blocks can then be explicitly calculated. This has actually been proven mathematically, though we won't go into the proof.

The Wess-Zumino-Witten model is a σ -model (meaning the fields are maps to some space) associated to a Lie group G (which we assume to be compact, simply connected, and simple). Let Σ be a Riemann surface and $\gamma : \Sigma \rightarrow G$ be smooth and $k \in \mathbb{Z}$.

Let $B : \mathfrak{g} \times \mathfrak{g} \rightarrow \mathbb{R}$ be the Killing form, a symmetric bilinear form that is $x, y \mapsto 4\mathrm{Tr}(xy)$ for \mathfrak{su}_2 . The WZW action is a sum of two terms: the first is the **kinetic term**

$$S^{\mathrm{kin}}(\gamma) := -\frac{k}{8\pi} \int_{\Sigma} B(\gamma^{-1} \partial^\mu \gamma, \gamma^{-1} \partial_\mu \gamma) dA.$$

To define the second term, called the **Wess-Zumino term**, let X be a 3-manifold which Σ bounds and $\tilde{\gamma} : X \rightarrow G$ be an extension of γ .⁹ Let $\{e_i\}$ be a basis for \mathfrak{g} ; then, the expression $B(e_i, [e_j, e_k])$ is alternating, so

$$\omega^{\mathrm{ZW}} := B(e_i, [e_j, e_k]) dx^i \wedge dx^j \wedge dx^k$$

is an invariant 3-form on G . The Wess-Zumino term is then

$$S^{\mathrm{WZ}}(\gamma) := \int_X \tilde{\gamma}^* \omega^{\mathrm{WZ}}.$$

This depends on the choice of $\tilde{\gamma}$, but is well-defined in \mathbb{R}/\mathbb{Z} . Exponentiating the action functional fixes this. Anyways, the WZW action is

$$S(\gamma) := S^{\mathrm{kin}}(\gamma) + 2\pi k S^{\mathrm{WZ}}(\gamma).$$

The relationship between Chern-Simons theory and the WZW model is an example of the **holographic principle**: that in many contexts, information about an n -dimensional QFT (the “boundary”) determines and is determined by an $(n+1)$ -dimensional TQFT (the “bulk”). In general, this is an *ansatz* rather than a theorem or even a physics-motivated argument, but in some cases there's good evidence for it. In our case, it's a theorem!

Theorem 4.5 (CS-WZW correspondence). *There is an isomorphism between the state space of Chern-Simons theory for G on a surface Σ and the space of conformal blocks of the WZW model on Σ for G .*

This is understood, but the proof is complicated, and we won't go into it.

⁸There is also something called the Wess-Zumino model, and it is different!

⁹The obstruction to such an extension existing is $\pi_2(G)$, which vanishes because G is simply connected.

4.4. Some calculations. In this section, we assume the level k is large. The trivial representation is denoted $\underline{\mathbb{C}}$.

Witten calculates the Wilson loop operator for a link K in S^3 by breaking S^3 and K into pieces. This requires understanding the boundary, the state space for a punctured Riemann sphere S^2 , with the punctures given by $S^2 \cap K$. Each puncture is labeled by the representation V that was associated to its loop, where we replace V with V^* if the orientation of the link disagrees with the direction of time.

The holographic principle identifies the state space of Chern-Simons theory on the punctured sphere with the space of conformal blocks of the WZW model, which is the space of possible charges assigned at each puncture. That is, at each puncture p_i , we want an element R_i , and the global state is the tensor product of the local states. However, the total charge of the space must be zero (just as the total electric charge of our universe is predicted to be 0), which means restricting to elements that are G -invariant. If S_m^2 be the Riemann sphere with m punctures, this means that we want to calculate

$$\mathcal{H}_{S_m^2, \vec{V}} = \left(\bigotimes_i V_i \right)^G.$$

It's necessary to restrict to representations which are **positive-energy** considered as representations of the loop group LG , but all of the ones we use satisfy this property.

Lemma 4.6. *A G -representation $\underline{\mathbb{C}} \oplus V$, with V nonzero, is not positive-energy.*

Proposition 4.7. *For $m = 0$, $\mathcal{H}_{S^2, \emptyset} \cong \underline{\mathbb{C}}$.*

Proof. We want to compute the G -invariants of a tensor product over \emptyset , i.e. $\underline{\mathbb{C}}^G = \underline{\mathbb{C}}$. □

Proposition 4.8. *For $m = 1$, $\mathcal{H}_{S^2, V}$ is 1-dimensional if V is trivial, and is otherwise 0.*

Proof. The proposition is proven when V is trivial, so assume V is nontrivial. By Maschke's theorem, V^G is a subrepresentation of G and a direct sum of copies of the trivial representation, so by Lemma 4.6, we must have $V^G = 0$. □

Proposition 4.9. *For $m = 2$, $\mathcal{H}_{S^2, \{V_1, V_2\}}$ is 1-dimensional if $V_1^* \cong V_2$, and is otherwise 0.*

Proof sketch. If $V_1^* \cong V_2$, $V_1 \otimes V_2 = \text{End}(V_2, V_2)$ as G -representations, and the diagonal matrices are an invariant subspace. Since V_2 is positive-energy, there can't be any more invariants.

If otherwise, the direct-sum decomposition of $V_1 \otimes V_2$ into irreducibles cannot have any copies of the trivial representation (as otherwise they would appear in either V_1 or V_2). □

The three-point functions are tied to an interesting theory, but we won't need them to compute the Jones polynomial. The dimension of the state space attached to three *irreducible* representations V_i , V_j , and V_k is a coefficient N_{ij}^k expressed in terms of the S -**matrix** of the theory. This can be understood mathematically using modular tensor categories.

In principle, the 3-point functions determine all higher correlators in the theory, but we just need a specific 4-point function. Witten claims that if V is a G -representation and $V \otimes V$ decomposes into s distinct irreducible representations, then the Hilbert space associated to 4-punctured S^2 with points labeled by V , V , V^* , and V^* is s -dimensional. We will only need this for V the defining representation for SU_2 , where we can prove it directly.

Proposition 4.10. *Let V denote the defining representation of SU_2 . Then, $\mathcal{H}_{S^2, V, V, V^*, V^*}$ is one-dimensional.*

Proof. This is a fun calculation with the irreducible representations of SU_2 . Recall that there's an irreducible $(n+1)$ -dimensional representation P_n whose character on $\begin{pmatrix} z & 0 \\ 0 & z \end{pmatrix}$ is

$$\chi_n(z) = z^n + z^{n-2} + \dots + z^{-n+2} + z^{-n}.$$

The defining representation V and its dual are isomorphic to P_1 , so

$$\begin{aligned} \chi_{V \otimes V \otimes V^* \otimes V^*}(z) &= \chi_1(z)^2 \chi_1(z)^2 = (z + z^{-1})^4 \\ &= z^4 + 4z^2 + 6 + 4z^{-2} + z^{-4} \\ &= \chi_4(z) + 3\chi_2(z) + 2\chi_0(z). \end{aligned}$$

Since P_0 is the trivial representation, then as SU_2 -representations,

$$V \otimes V \otimes V^* \otimes V^* \cong P_4 \oplus (P_2)^{\oplus 3} \oplus \underline{\mathbb{C}}^{\oplus 2}.$$

The Hilbert space is the space of G -invariants factors through the direct sum. For an irreducible representation, $P_n^G = 0$ unless P_n is the trivial representation, so the G -invariants vanish on the copies of P_4 and P_2 , leaving a 2-dimensional space. \square

Part 2. Categorification: Khovanov homology