

M392C NOTES: APPLICATIONS OF QUANTUM FIELD THEORY TO GEOMETRY

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These notes were taken in UT Austin's M392C (Applications of Quantum Field Theory to Geometry) class in Fall 2017, taught by Andy Neitzke. I live-TeXed them using vim, so there may be typos; please send questions, comments, complaints, and corrections to a.debray@math.utexas.edu. Any mistakes in the notes are my own. Thanks to Michael Hott and Andy Neitzke for a few corrections.

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Lecture 1.

Donaldson invariants and supersymmetric Yang-Mills theory: 8/31/17

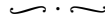
“The wind blowing on it, well, that’s not the worst thing that could happen to a pond! Now imagine you have a laser...”

The course website is <https://www.ma.utexas.edu/users/neitzke/teaching/392C-qft-geometry/>. There are also lecture notes which are hosted at <https://github.com/neitzke/qft-geometry>, and are

currently a work in progress; if you have contributions or improvements, feel free to contribute them, as a pull request or otherwise. (I'm also taking notes, of course, and if you find problems or typos in my notes, feel free to let me know.) There's also a Slack channel for course-related discussions, which may be easier to use than office hours.

There will be exercises in this course, and you should do at least one-fourth of them for the best grade. Of course, you also want to do them in order to gain understanding. Some worked-out computations could be useful for submitting to the professor's lecture notes.

This course will be relatively wide-ranging; today's prerequisites involve some gauge theory, but the next few lectures won't as much.



Suppose you want to study the topology of smooth manifolds X . Surprisingly, it's really effective to introduce a geometrical gadget, e.g. a Riemannian metric g . Using it, we can define the *Laplace operator* on differential forms $\Delta: \Omega^k(X) \rightarrow \Omega^k(X)$, which has the formula

$$\Delta := dd^* + d^*d,$$

where $d: \Omega^k(X) \rightarrow \Omega^{k+1}(X)$ is the de Rham differential, and $d^*: \Omega^{k+1}(X) \rightarrow \Omega^k(X)$ is its adjoint in the L^2 -inner product on differential forms induced by the metric. Thus d is canonical, but d^* depends on the choice of metric.

Next we consider the equation

$$(1.1) \quad \Delta \omega = 0.$$

This is a linear equation, so its space of solutions $\mathcal{H}_{k,g} := \ker(\Delta: \Omega^k \rightarrow \Omega^k)$, called the *space of harmonic k -forms*, is a vector space. If X is compact, it's even a finite-dimensional vector space, which is a consequence of the ellipticity of the Laplace operator.¹ Hence we can define a nonnegative integer

$$b_k(X) := \dim \mathcal{H}_{k,g},$$

called the k^{th} *Betti number* of X . It's a fact that $b_k(X)$ does not depend on the choice of the metric! Thus they are invariants of the smooth manifold X .

In fact, there's even a categorified version of this. This reflects a recent (last decade or so) trend of replacing numbers with vector spaces, sets with categories, etc.

Theorem 1.2. *If X is compact,² there is a canonical isomorphism $\mathcal{H}_{k,g} \cong H^k(X; \mathbb{R})$, where the latter is the singular cohomology of X with coefficients in \mathbb{R} .*

This shows $b_k(X)$ doesn't depend on the smooth structure of X , and is even a homotopy invariant. This will not be true for the Donaldson invariants that we'll discuss later.

Exercise 1.3. Work out some of these spaces of harmonic forms for a metric on S^1 and S^2 .

You have to choose a metric, and there are more or less convenient ones to pick. But no matter how you change the metric, there will be a canonical way to identify them.³

If X is oriented and $4n$ -dimensional, there's a small refinement of the middle Betti number b_{2n} and space of harmonic forms \mathcal{H}_{2n} . The *Hodge star operator*

$$\star: \Omega^p(X) \longrightarrow \Omega^{\dim X - p}(X)$$

is an involution on $\Omega^{2n}(X)$.

Remark 1.4. Let's recall the Hodge star operator. This is an operator on differential forms defined using the Riemannian metric satisfying $\star^2 = 1$ in even dimension, and $[\star, \Delta] = 0$. Hence it acts on harmonic forms. On \mathbb{R}^2 with the usual metric, $\star(1) = dx \wedge dy$, and $\star(f dx) = f dy$. ◀

Hence we can decompose $\Omega^{2n}(X)$ into the (± 1) -eigenspaces of \star : let $\Omega^{2n,\pm}(X)$ denote the ± 1 -eigenspace for \star . Similarly, $\mathcal{H}_{2n}(X)$ splits into $\mathcal{H}_{2n}^\pm(X)$. Thus b_{2n} also splits:

$$b_{2n}(X) = b_{2n}^+(X) + b_{2n}^-(X).$$

These spaces and numbers are also topological invariants, and can be understood in that way.

¹For a general differential operator on differential forms, nothing like this is true.

²Compactness is really necessary for this.

³Interesting question: if you change the metric infinitesimally, how does \mathcal{H}_k change?

Exercise 1.5. In dimension $4n + 2$, the Hodge star squares to -1 . You can still extract topological information from this; what do you get?

Linear equations seem to behave more or less the same in all dimensions. But nonlinear equations behave very differently in different dimensions. In the 1980s, Donaldson [13] used nonlinear equations to produce new and interesting invariants of 4-manifolds. Let X be a connected, oriented 4-manifold with a Riemannian metric g .

Fix a compact Lie group G . For Donaldson, $G = \mathrm{SU}(2)$, and it's probably fine to assume that for much of this class. Fix a principal G -bundle $P \rightarrow X$. We'll consider connections on P .

Remark 1.6. If you don't know what a connection is, that's OK. Locally, a connection on P is represented by a Lie algebra-valued 1-form $A \in \Omega_X^1(\mathfrak{g})$, and has a *curvature* 2-form $F \in \Omega_X^2(\mathfrak{g}_P)$, which locally is written

$$F = dA + A \wedge A.$$

Because $\mathrm{SU}(2)$ is nonabelian, $A \wedge A$ isn't automatically zero. ◀

Since F is a 2-form and $\dim X = 4$, we can decompose F into its *self-dual* part F^+ and its *anti-self-dual* part F^- , defined by the splitting of Ω^2 by the Hodge star.

Exercise 1.7. Show that if you reverse the orientation of X , F^+ and F^- switch.

Donaldson studied the *anti-self-dual Yang-Mills equation* (ASD YM):

$$(1.8) \quad F^+ = 0.$$

By Exercise 1.7, this is not really different than studying the self-dual Yang-Mills equation; the reason one prefers the ASD version is that it occurs more naturally on certain complex manifolds which were test cases for Donaldson theory.

If G is abelian, e.g. $\mathrm{U}(1)$, (1.8) is linear. But if G is nonabelian, e.g. $\mathrm{SU}(2)$, then (1.8) is nonlinear.

Definition 1.9. The *instanton moduli space* is the space \mathcal{M} of equations on P obeying (1.8), modulo the action of the *gauge group* \mathcal{G} , the bundle automorphisms of P .⁴

Exercise 1.10. Show that if $G = \mathrm{U}(1)$, then \mathcal{M} is only governed by linear algebra in that

$$\mathcal{M} \cong H^1(M; \mathbb{R}) / H^1(X; \mathbb{Z}).$$

So in this case we don't find anything new, though the way we found it is still interesting.

When G is nonabelian, this is not a vector space. It still has some reasonable structure. We now fix $G = \mathrm{SU}(2)$. In this case, (topological) isomorphism classes of principal $\mathrm{SU}(2)$ -bundles are classified by the integers, given by the formula

$$k := \int_X c_2(P) \in \mathbb{Z},$$

where c_2 denotes the second Chern class.

This means the moduli of instantons is a disjoint union over \mathbb{Z} of spaces \mathcal{M}_k .

Theorem 1.11. If $k > 0$ and g is chosen generically, \mathcal{M}_k is a finite-dimensional manifold.

Hence one could learn topological information about X by studying topological properties of \mathcal{M}_k . The first idea would be the Betti numbers, but these turn out not to depend on the smooth structure.

Proposition 1.12. Assuming $k > 0$ and g is generic,

$$\dim \mathcal{M}_k = 8k - 3(1 - b_1(X) + b_2^+(X)).$$

But there's more to \mathcal{M}_k than the dimension. Donaldson introduced an orientation on \mathcal{M}_k , which is canonically defined (and a lot of hard work!), and one can produce classes $\tau_\alpha \in \Omega^*(\mathcal{M}_k)$ labeled by classes $\alpha \in H_*(X)$. Using these, the *Donaldson invariants* are the real numbers

$$(1.13) \quad \langle \mathcal{O}_{\alpha_1} \cdots \mathcal{O}_{\alpha_t} \rangle := \int_{\mathcal{M}} \tau_{\alpha_1} \wedge \cdots \wedge \tau_{\alpha_t} \in \mathbb{R}.$$

Theorem 1.14. If $b_2^+(X) > 1$, the Donaldson invariants are independent of g .

⁴**TODO:** not sure if I got this right.

Moreover, they really depend on smooth information: it's not possible to reconstruct them out of algebraic or differential topology, unlike the Betti numbers. So these are very powerful. Their study is called *Donaldson theory*. One good reference is Donaldson and Kronheimer's book [14].

Unfortunately, Donaldson theory is technically very hard: the ASD YM equation is hard to study: \mathcal{M}_k is usually noncompact, and (1.13) is an integral over a noncompact space, which is no fun.

What does this have to do with quantum field theory? In 1988, Witten [28], following a suggestion of Atiyah, found an interpretation of the Donaldson invariants in terms of quantum field theory (hence the suggestive notation in (1.13)).

There are many different quantum field theories: the Standard Model describes three of the four fundamental forces of the universe; quantum electrodynamics describes electromagnetism. Witten interpreted the Donaldson invariants in terms of a specific QFT, called “(a topological twist of) $\mathcal{N} = 2$ supersymmetric Yang-Mills theory (SYM) with gauge group $SU(2)$.”

One imagines X to be a “spacetime” or “universe” whose laws of physics are governed by $\mathcal{N} = 2$ supersymmetric Yang-Mills theory, and to compute the Donaldson invariants, one conducts “experimental measurements” (correlation functions). According to the rules of Lagrangian quantum field theory, this means computing an integral over an infinite-dimensional space (which is alarming, but so it goes):

$$\langle \mathcal{O}_\alpha \rangle = \int_{\mathcal{C}} d\mu \Phi_\alpha e^{-S},$$

where

- \mathcal{C} is the *space of fields*, some sort of infinite-dimensional space akin to the space of functions on X or forms on X ,
- $S: \mathcal{C} \rightarrow \mathbb{R}$ is a functional called the *action*,
- $\Phi_\alpha: \mathcal{C} \rightarrow \mathbb{R}$ is a (set of) *observables*,
- and $d\mu$ is some measure on \mathcal{C} .

In general, computing these correlation functions are very hard,⁵ but in $\mathcal{N} = 2$ SYM, Witten found localization, a way to reduce it to Donaldson's integrals over finite-dimensional spaces.

This is undoubtedly cool, and brings geometric topology into quantum field theory, but it does not make it much easier to actually compute Donaldson invariants.

The next step was taken in 1995, by Seiberg and Witten [25, 26], who were interested in a different but related physics problem. They answered a fundamental question about SYM: how it behaves at low energies.

To make an analogy, suppose you have a pond, and you're pond-ering what happens when wind goes across the surface. You're good at physics, so you model the pond as a system of 10^{30} molecules of water and other things, then rent some time on a supercomputer where you model the action on the wind and... somehow this seems wrong. Instead, you model the water and the wind using things like the Navier-Stokes equations. This is not easy, but it's much, much easier.

The idea is there's a “high-energy” description, in terms of 10^{30} particles, but the “low-energy” description⁶ involves things like temperature, pressure, liquid, and other things that are hard to define from the high-energy approach. The low-energy picture is very useful for calculations, though if you fire a laser into your pond it wouldn't suffice. Obtaining the description of the low-energy physics from the high-energy physics is typically very hard; in this case, one would have to define temperature and pressure and a lot of things starting from fundamentals. But you just have to do it once, then can apply it to all bodies of water, etc.

Seiberg and Witten applied this to $\mathcal{N} = 2$ SYM with gauge group $SU(2)$, and showed that its low-energy description is (roughly) $\mathcal{N} = 2$ SYM with gauge group $U(1)$, coupled to matter (sometimes called monopoles). Since the gauge group is abelian, this is much easier. Now, one can imagine that there's an easier description of the Donaldson invariants in terms of the low-energy theory (though, again, this was not the original intent of Seiberg and Witten), and this is given by the *Seiberg-Witten equations*. They look more complicated but are actually vastly simpler.⁷

In the Seiberg-Witten equations, the fields are

- a connection Θ in a $U(1)$ -bundle \mathcal{E} , or equivalently a determinant line of a Spin^c -structure, and

⁵Unless $\dim X = 0$, where \mathcal{C} is finite-dimensional. We'll talk about this in the next few lectures.

⁶The term “low-energy,” despite sounding pejorative, is actually a very useful thing to have.

⁷For a reference, see Morgan [23].

- a section ψ of S^+ , a spinor bundle associated to a Spin^c -structure.

In this case, there's a *Dirac operator* \not{D} and a pairing

$$q: S^+ \otimes S^+ \longrightarrow \Lambda_+^2 T^*X.$$

Then, the Seiberg-Witten equations are

$$(1.15a) \quad F^+ = q(\psi, \bar{\psi})$$

$$(1.15b) \quad \not{D}\psi = 0.$$

Let $\widetilde{\mathcal{M}}$ denote the moduli space of pairs (Θ, ψ) satisfying (1.15) modulo the action of some group. For generic g , this is a compact manifold, so understanding its topology is much easier, and the correlation functions for the low-energy theory can be written as integrals over $\widetilde{\mathcal{M}}$, and there's a simple formula relating these to the correlation functions for the high-energy theory. Once this was realized, there was very rapid progress of its use in applications, though understanding precisely why it's the same came more slowly, beginning from a physical argument by Moore and Witten [21] and proceeding to a very different-looking mathematical proof much more recently.

This is an application of QFT to geometry, as we will study in this course. Somehow the most powerful applications involve taking a low-energy limit, and many of them also involve localization in supersymmetric QFT (from an infinite-dimensional integral to a finite-dimensional one).

We will start more slowly: first considering QFT where $\dim X = 0$, then $\dim X = 1$ (which is quantum mechanics); in these cases, the physics can be made completely rigorous (though it's not necessarily easy). We'll briefly talk about $\dim X = 2$, then jump into $\dim X = 4$.

Lecture 2.

Zero-dimensional QFT and Feynman diagrams: 9/5/17

Last time, we talked about two perspectives on physics, high-energy (or *fundamental*) and low-energy (or *effective*). For example, the high-energy description of a pond is the physics of the 10^{30} or so particles in it, and the low-energy description is the Navier-Stokes equations. We're interested in the relationship between Donaldson theory in the high-energy perspective and Seiberg-Witten theory in the low-energy perspective, which is a story about four-dimensional QFT. But over the next few lectures, we're going to learn about this passage from fundamental to effective in 0-dimensional QFT, one of the few cases where it's known how to make everything rigorous. Nonetheless, it's still an interesting theory, e.g. it has Feynman diagrams.

We also discussed that in the Lagrangian formalism to QFT on a spacetime X , one evaluates integrals over a space $\mathcal{C}(X)$, which is some kind of function space. Hence, it's usually infinite-dimensional, unless $\dim X = 0$. Hence, let's assume $X = \text{pt}$, so $\mathcal{C}(X) = \{X \rightarrow \mathbb{R}\} = \mathbb{R}$. There are many choices for $S: \mathcal{C} \rightarrow \mathbb{R}$,⁸ such as

$$S(x) = \frac{m}{2}x^2 + \frac{\lambda}{4!}x^4,$$

where $m, \lambda > 0$. Here m might mean some kind of mass, and λ measures the interaction in the system.

Now we can define something important and fundamental: the *partition function*

$$Z := \int_{-\infty}^{\infty} dx e^{-S(x)}.$$

The *observables* are polynomial functions $f: \mathcal{C} \rightarrow \mathbb{R}$, and their (*unnormalized*) *expectation values* are

$$\langle f \rangle := \int_{-\infty}^{\infty} dx f(x) e^{-S(x)}.$$

We require f to be polynomial so that this integral converges. All of these are functions in m and λ . Also, notice that all of these are completely well-defined; maybe this is a trivial observation, but it won't be true when we ascend to higher dimensions.

⁸One can also use \mathbb{C} -valued actions.

Computing these quantities is less trivial. Let's start with Z , or even $Z_0 := Z(m, \lambda = 0)$. This is a Gaussian:

$$Z_0 = \int_{-\infty}^{\infty} dx e^{-mx^2/2} = \sqrt{\frac{2\pi}{m}}.$$

In order for this to be well-defined, we need $m = 0$ of course, but there's a physical reason to throw out this case, as it corresponds to a system with more than one vacuum state and a degenerate critical point of the action.

To compute the partition function for $\lambda > 0$, we're not sure how to directly evaluate the integral, but we can try to expand it out as a Taylor series in λ around 0. This will allow us to understand the system in the presence of weak interactions, which is often exactly what physicists want to know. We'll leave $e^{-mx^2/2}$ alone, since we know how to integrate it exactly. The $\lambda x^4/4!$ term expands to

$$Z(m, \lambda) = \int_{-\infty}^{\infty} dx \sum_{n=0}^{\infty} \left(-\frac{\lambda}{4!}\right)^n \frac{x^{4n}}{n!} e^{-mx^2/2}.$$

We'd like to switch the sum and integral to obtain

$$(2.1) \quad = \sum_{n=0}^{\infty} \left(-\frac{\lambda}{4!}\right)^n \underbrace{\int_{-\infty}^{\infty} \frac{x^{4n}}{n!} e^{-mx^2/2}}_I,$$

but we have to be careful about convergence. If this works, though, the integral I is tractable.

Exercise 2.2. Show that

$$\int_{-\infty}^{\infty} dx x^{2k} e^{-mx^2/2} = \sqrt{\frac{2\pi}{m}} \frac{1}{m^k} \frac{(2k)!}{k! 2^k}.$$

Hence, modulo the assumption we made before, if $\tilde{\lambda} := \lambda/m^2$,

$$(2.3) \quad \begin{aligned} Z(m, \lambda) &= \sqrt{\frac{2\pi}{m}} \sum_{n=0}^{\infty} \left(-\frac{1}{96}\right)^n \frac{(4n)!}{n!(2n)!} \tilde{\lambda}^n \\ &= \sqrt{\frac{2\pi}{m}} \left(1 - \frac{1}{8}\tilde{\lambda} + \frac{35}{384}\tilde{\lambda}^2 + \cdots + (1390.1)\tilde{\lambda}^{10} + \cdots\right). \end{aligned}$$

This is called the *perturbation series* for this partition function. Though this partition function is a scalar multiple of a Bessel function, often these series are actually divergent for any $\tilde{\lambda} > 0$. This means the assumption we made in (2.1) was wrong. There's various ways to think about this — if this function did converge to its Taylor series, it would do so in a neighborhood of 0 in \mathbb{C} , hence for negative λ . Physically, this doesn't make sense.

Nonetheless, the perturbation series is still useful in those cases.

Definition 2.4. Let $f : \mathbb{R}_+ \rightarrow \mathbb{C}$ be a function and $s := \sum_{n=0}^{\infty} c_n t^n$ be a formal series. We say that s is an *asymptotic series* for f as $t \rightarrow 0^+$ if for all $N \geq 0$,

$$\lim_{t \rightarrow 0^+} t^{-N} \left| f(t) - \left(\sum_{n=0}^N c_n t^n \right) \right| = 0.$$

In this case, we write

$$f(t) \underset{t \rightarrow 0^+}{\sim} \sum_{n=0}^{\infty} c_n t^n.$$

In particular, this means that

$$\begin{aligned} \lim_{t \rightarrow 0^+} |f(t) - c_0| &= 0 \\ \lim_{t \rightarrow 0^+} \frac{1}{t} |f(t) - c_0 t + c_1| &= 0, \end{aligned}$$

and so on. So even if s doesn't converge, it's still useful, capturing the limits, linear behavior, quadratic behavior, etc., of f . You have encountered other asymptotic series in your life: Stirling's formula for the factorial is an asymptotic series for the gamma function at ∞ : it doesn't actually converge in a sensible way, but it captures a lot of useful information.

Proposition 2.5. *The series (2.3) is an asymptotic series for the partition function $Z(m, \lambda)$ as $\lambda \rightarrow 0^+$.*

So it's not equality, but it's a useful and interesting approximation.

You might wonder whether there's some better series approximating $Z(m, \lambda)$ that actually converges, but this is not true.

Proposition 2.6. *If f has a convergent Taylor series at x_0 , then its Taylor series is an asymptotic series for f at x_0 .*

Proposition 2.7. *Every smooth function f can have at most one perturbation series as $x \rightarrow x_0$.*

Sometimes none exists.

We will interpret (2.3) in terms of Feynman diagrams. The basic object is a vertex with four half-edges attached:



A Feynman diagram for (2.3) is a placement of some of these vertices and a way of connecting the half-edges. (Feynman diagrams for other systems may look different.)

placeholder

FIGURE 1. Some Feynman diagrams with one or two vertices.

Let D_n denote the set of diagrams with n vertices.

Proposition 2.8. *The number of ways to pair up $2k$ objects is $(2k)!/k!$.*

Corollary 2.9.

$$|D_n| = \frac{(4n)!}{(2n)!2^{2n}}.$$

There's also a group action of a group $G_n := (S_4)^n \rtimes S_n$ on D_n , where the i^{th} copy of S_4 permutes the half-edges for the i^{th} vertex, and S_n shuffles the n vertices. In other words, we can restate the asymptotic series for the partition function (2.3) in a more combinatorial manner: since $Z_0 = \sqrt{2\pi/m}$.

$$\frac{Z(m, \lambda)}{Z_0} \sim \sum_{n=0}^{\infty} (-\tilde{\lambda})^n \frac{|D_n|}{|G_n|}.$$

We want to describe $|D_n|/|G_n|$ as the cardinality of some kind of quotient set, but this is only literally true if the G_n -action on D_n is free. The proper thing to do, as suggested by the orbit-stabilizer theorem, is to sum over orbits, weighted by the order of their stabilizers.⁹ Thus

$$\frac{Z(m, \lambda)}{Z_0} \sim \sum_{n=0}^{\infty} (-\tilde{\lambda})^n \sum_{[\Gamma] \in D_n/G_n} \frac{1}{|\text{Aut } \Gamma|}.$$

Since $\tilde{\lambda} = \lambda/m^2$ and a Feynman diagram in D_n has n^2 edges, we can rewrite (2.3) in a way that is completely a combinatorial sum over Feynman diagrams:

$$\frac{Z(m, \lambda)}{Z_0} \sim \sum_{n \geq 0} \sum_{[\Gamma] \in D_n/G_n} \frac{(-\lambda)^{|V(\Gamma)|}}{m^{|E(\Gamma)|}} \cdot \frac{1}{|\text{Aut}(\Gamma)|}.$$

Here, $V(\Gamma)$ is the set of vertices of Γ , and $E(\Gamma)$ is the set of edges. This leads to the *Feynman rules* for summing over the Feynman diagrams for this theory:

- Draw one representative Γ for each orbit in D_n/G_n .
- Define its weight w_Γ as the product of factors $-\lambda$ for each vertex and $1/m$ for each edge, weighted by $1/|\text{Aut}(\Gamma)|$.

Then,

$$\frac{Z}{Z_0} \sim \sum_{[\Gamma]} w_\Gamma.$$

⁹Another way to think about this is to consider the quotient groupoid D_n/G_n , and sum over it in the groupoid measure, which amounts to the same thing.

Example 2.10. Let's calculate some low-order terms.

- The empty Feynman diagram has the weight 1.
- The action of $G_1 \cong S_4$ on D_1 is transitive, so we only need a single representative, such as the “figure-8 diagram.” Its stabilizer group has order 8, so there's a contributing factor of $(-\lambda)/8m^2$.
- There are three orbits in D_2/G_2 , represented by a graph with zero self-loops, which contributes a term of $\lambda^2/48m^4$, one with one self-loop on each vertex, which contributes $\lambda^2/16m^4$, and one with two self-loops on each vertex, which contributes $\lambda^2/128m^4$.

Thus, the perturbative expansion is

$$\begin{aligned} \frac{Z}{Z_0} &\sim 1 - \frac{\lambda}{8m^2} + \frac{\lambda^2}{48m^4} + \frac{\lambda^2}{16m^4} + \frac{\lambda^2}{128m^4} + O(\lambda^3) \\ &= 1 - \frac{\lambda}{8m^2} + \frac{35}{384} \frac{\lambda^2}{m^4} + O(\lambda^3). \end{aligned}$$

The higher-order terms correspond to diagrams with 3 or more vertices. ◀

If you know the automorphism group of a diagram Γ , then the automorphism group of $\Gamma \amalg \Gamma$ is very similar: a copy of $\text{Aut}(\Gamma)$ for each component, plus the S_2 switching them. If you follow your nose in this line of thought, you can determine the sum in terms of only nonempty, connected diagrams.

Proposition 2.11.

$$\sum_{\Gamma} w_{\Gamma} = \exp \left(\sum_{\Gamma \text{ connected, nonempty}} w_{\Gamma} \right).$$

This suggests that $\log(Z/Z_0)$ is an important physical quantity, and indeed, it's called the *free energy* of the system, as in statistical mechanics. We'd like to say that

$$\log \left(\frac{Z(m, \lambda)}{Z_0} \right) \sim \sum_{\Gamma \text{ connected, nonempty}} w_{\Gamma},$$

though there's an analysis argument to check here.

Now we want to compute expectation values. Let's start with

$$\langle x^k \rangle := \int_{-\infty}^{\infty} x^n e^{-S} dx.$$

If k is odd this is 0, but for k even, we can compute an asymptotic series for this function with a similar sum over Feynman diagrams, but with different rules:

- In addition to the 4-valent vertices from before, each diagram must have exactly k univalent vertices.
- We only consider automorphisms which fix these vertices.

You can work this out with a similar argument as for Z/Z_0 .

To compute the *normalized expectation values* $\langle x^k \rangle / Z$, use the same diagrams, but with the rule that every connected component of Γ must have at least one univalent vertex. You can then draw out the first few diagrams and conclude things such as

$$\frac{\langle x^2 \rangle}{2} \sim \frac{1}{m} - \frac{\lambda}{2m^3} + O(\lambda^2).$$

More generally, there's no need to constrain ourselves to a quartic interaction: we can instead consider the action

$$(2.12) \quad S = \frac{m}{2} x^2 + \sum_{k=3}^{\infty} \frac{\lambda_k x^k}{k!}.$$

In this case, we consider Feynman diagrams with vertices of arbitrary valence ≥ 3 , and sum with the rules that an edge contributes $-1/m$ and an n -valent vertex contributes $-\lambda_n$. We can actually carry out the analysis even if (2.12) doesn't converge (in which case we don't get an asymptotic series for a function, but that's OK). Anyways, tabulating the Feynman diagrams we get the beginning of the normalized perturbative expansion

$$\frac{Z}{Z_0} \sim 1 - \frac{\lambda_4}{8m^2} + \frac{\lambda_3^2}{12m^3} + \dots$$

Yet another generalization is to consider actions on $\mathcal{C} = \mathbb{R}^N$, rather than \mathbb{R} , corresponding to considering the theory on N points, rather than one point. Now, the quartic term is some 4-tensor, so (using the Einstein summation convention) the most general action is

$$S = \frac{1}{2} x^i M_{ij} x^j + \frac{1}{4!} C_{ijkl} x^i x^j x^k x^\ell,$$

and Z_0 is again a Gaussian:

$$Z_0 = \int_{\mathbb{R}^n} e^{-x^i M_{ij} x^j / 2} = \frac{(2\pi)^{N/2}}{\sqrt{\det M}}.$$

In this case, one can compute with Feynman diagrams again, but this time labeling the edges with labels $1, \dots, N$.

Lecture 3.

A Little Effective Field Theory: 9/7/17

Today, we're going to illustrate the passage from the fundamental to the effective using zero-dimensional QFT: the fundamental theory will be an action $S(x, y)$ in two variables, and its effective theory S_{eff} will be a simpler theory in a single variable.

Last time, we discussed the fields $\mathcal{C} = \mathbb{R}^N$ in a zero-dimensional QFT with an action

$$S := \frac{1}{2} x^i M_{ij} x^j + \frac{1}{4!} C_{ijkl} x^i x^j x^k x^\ell.$$

As $C \rightarrow 0$, one wants to compute the asymptotic series, which amounts to a sum over Feynman diagrams. In this context, one can sum over unlabeled diagrams Γ , but with the weight incorporating the labels of the half-edges in $\{1, \dots, N\}$. Explicitly, the weight of an edge i to j should be $(M^{-1})^{ij}$, and that of a vertex with half-edges i, j, k , and ℓ is C_{ijkl} .

More abstractly, if V is a finite-dimensional vector space with a measure μ , you can choose an $M \in \text{Sym}^2 V^*$ and a $C \in \text{Sym}^4 V^*$, and define the action

$$S(x) := \frac{1}{2} M(x, x) + \frac{1}{4!} C(x, x, x, x).$$

Then, one would compute the partition function

$$\int d\mu e^{-S(x)}.$$

Now let's focus on a specific example. We can start with fields $\mathcal{C} = \mathbb{R}^2$ with coordinates x, y and an action

$$(3.1) \quad S(x, y) := \frac{m}{2} x^2 + \frac{M}{2} y^2,$$

which is two uncoupled systems. So let's turn on coupling in (3.1):

$$(3.2) \quad S(x, y) := \frac{m}{2} x^2 + \frac{M}{2} y^2 + \frac{\mu}{4} x^2 y^2.$$

Say that we're actually interested in x : we want to compute Z and $\langle x^n \rangle$, but *not* $\langle y \rangle$ or $\langle f(x, y) \rangle$ that depends on y . This might happen in a system which naturally comes with both x and y , but y is some extra degrees of freedom. We'll see this is natural when $M \gg m$.

There are only a few kinds of labels in the Feynman diagram, because M and C in (3.2) have a lot of zeroes: we'll use a solid line for $1/m$ (corresponding to x^2) and a dashed line for $1/M$ (for y^2); all vertices must have two solid half-edges and two dashed half-edges, weighted by $-\mu$.

Let's compute $\log(Z/Z_0)$; by Proposition 2.11, this allows us to only sum over connected diagrams. There is only one diagram with a single vertex (order μ), and three with two vertices (order μ^2). Their respective computations are

$$\log\left(\frac{Z}{Z_0}\right) \sim -\frac{\mu}{4mM} + \frac{\mu^2}{16m^2M^2} + \frac{\mu^2}{16m^2M^2} + \frac{\mu^2}{8m^2M^2} + O(\mu^3).$$

For correlation functions, we must add n univalent vertices for x^n . The μ^0 -term (the “tree level”) calculates exactly the noninteracting theory. When we enumerate the diagrams for $\langle x^2 \rangle$, there’s one with zero 4-valent vertices, one for a single 4-valent vertex, and three with two 4-valent vertices, and the sum is

$$\frac{\langle x^2 \rangle}{Z} \sim \frac{1}{m} - \frac{\mu}{2m^2M} + \frac{\mu^2}{4m^3M^2} + \frac{\mu^2}{2m^3M^2} + \frac{\mu^2}{4m^3M^2} + O(\mu^3).$$

This is not the logarithm: since we’ve normalized this calculation, it’s a sum over Feynman diagrams for which every connected component contains a univalent vertex.

This explodes more quickly than other ones we considered: to compute $\langle x^4 \rangle$, there are a lot of diagrams to sum over, even just at the μ^2 . The answer will be

$$\frac{\langle x^4 \rangle}{Z} \sim \frac{3}{m^2} - \frac{3\mu}{m^3M} + \frac{33\mu^2}{4m^4M^2} + O(\mu^3).$$

And since we only care about x , there should be some way to simplify this and get all of the dashed lines out of the way first. One idea is: if we only want

$$\langle x^n \rangle = \int_{\mathbb{R}^2} dx dy x^n e^{-S(x,y)},$$

then by Fubini’s theorem, we can integrate out the dependence on y , defining S_{eff} such that

$$e^{-S_{\text{eff}}(x)} := \int_{\mathbb{R}} dy e^{-S(x,y)}.$$

Then

$$\langle x^n \rangle = \int_{\mathbb{R}} dx x^n e^{-S_{\text{eff}}(x)}.$$

In this particular example, we can compute S_{eff} , or at least its asymptotic series (which suffices if we want to do the asymptotic series for $\langle x^n \rangle$ in the original theory). The answer for the asymptotic series for $\mu \rightarrow 0$ is

$$(3.3) \quad S_{\text{eff}}(x) \sim \frac{m_{\text{eff}}}{2} x^2 + \sum_{k \geq 3} \lambda_k x^k,$$

where m_{eff} is some effective mass. The interacting term is interesting — there are interactions between multiple x s (vertices with four solid edges). These arise because of Feynman diagrams such as the one in Figure 2, where by “ignoring y ” we close the gap between these two vertices and obtain an interaction between two copies of x .

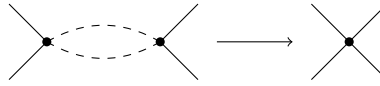


FIGURE 2. Left: a Feynman diagram for the action (3.2). In the effective field theory (3.3), the dashed lines correspond to terms which are integrated out, so this diagram becomes a quartic x - x interaction (on the right).

Specifically, in (3.3), the terms are

$$m_{\text{eff}} = m + \frac{\mu}{2M}$$

$$\lambda_k = \begin{cases} 0, & k \text{ odd} \\ -\left(-\frac{\mu}{M}\right)^{k/2} \frac{1}{2^{k/2+2k}}, & k \text{ even.} \end{cases}$$

Thus, as $M \rightarrow \infty$, $m_{\text{eff}} \rightarrow m$: when $M \gg m$, this is a more reasonable approximation.

This is our first baby example of an effective field theory. The fact that we integrated out the degrees of freedom we didn’t care about is a useful heuristic to have around.

Symmetries. Let's go back to $\mathcal{C} = \mathbb{R}$ and

$$S = \frac{m}{2}x^2 + \frac{\lambda}{4!}x^4.$$

This is in a sense the simplest nontrivial example: if you had a cubic term instead of a quartic term, $\int e^{-S}$ wouldn't be well-defined (it goes to ∞ as $x \rightarrow \pm\infty$).

Proposition 3.4. $\langle x^n \rangle = 0$ when n is odd.

Proof.

$$\begin{aligned} \langle x^n \rangle &= \int_{-\infty}^{\infty} dx x^n e^{-S(x)} \\ &= \int_{-\infty}^{\infty} d(-x) (-x)^n e^{-S(-x)} \\ &= (-1)^n \int_{-\infty}^{\infty} dx x^n e^{-S(x)} \\ &= (-1)^n \langle x^n \rangle. \end{aligned}$$

□

One takeaway is that this theory is symmetric under the group $\mathbb{Z}/2$ acting on \mathcal{C} as multiplication by $\{\pm 1\}$. This leads to a very general principle.

Proposition 3.5. Let $S: \mathcal{C} \rightarrow \mathbb{R}$ and the measure on \mathcal{C} are both G -invariant for a group G , then $\langle \mathcal{O} \rangle = \langle \mathcal{O}^g \rangle$ for any observable $\mathcal{O}: \mathcal{C} \rightarrow \mathbb{R}$, where $\mathcal{O}^g = g^* \mathcal{O}$.

If G is a Lie group, we can differentiate this equation: take $g = \exp(tX)$ for some $X \in \mathfrak{g}$: taking

$$\left. \frac{d}{dt} \right|_{t=0} (\langle \mathcal{O} \rangle = \langle \mathcal{O}^{tX} \rangle),$$

we conclude that $\langle X(\mathcal{O}) \rangle = 0$.

In general, symmetries are an extremely important ingredient in QFT.

Fermions and super-vector spaces. You might remember that we wanted to do something topological, but our computations, as functions in the parameters (m, λ) , were not deformation-invariant (you could think of them as nonconstant functions on a moduli space of QFTs). To get things that are, we need one more ingredient: fermions.

The way to do this, which will return again and again in this course, is to replace the manifold \mathcal{C} by a supermanifold! Since we've so far only considered vector spaces, we'll get a slightly gentler introduction in the form of super-vector spaces.

For a reference on this material, check out Etingof's course notes for a class on the mathematics of QFT.¹⁰

Definition 3.6. A super-vector space is a $\mathbb{Z}/2$ -graded vector space $V = V^0 \oplus V^1$.

For example, if $V^0 = \mathbb{R}^p$ and $V^1 = \mathbb{R}^q$, V is denoted $\mathbb{R}^{p|q}$. This can be done over any field, but we're only going to consider \mathbb{R} or \mathbb{C} .

These are not so terrible. But how we do algebra with them is also different: if you are taking tensor products, super-vector spaces are not the same as $\mathbb{Z}/2$ -graded vector spaces!¹¹

Definition 3.7. The symmetric monoidal category of super-vector spaces $(\text{sVec}, \otimes, s_{-, -})$ is the same as that for ordinary $\mathbb{Z}/2$ -graded vector spaces $\text{Vect}^{\mathbb{Z}/2}$, except for the symmetry

$$s_{V,W}: V \otimes W \rightarrow W \otimes V.$$

For $\text{Vect}^{\mathbb{Z}/2}$, this is the map $v \otimes w \mapsto w \otimes v$, but in sVec , it's defined on homogeneous v, w by

$$v \otimes w \mapsto (-1)^{|v||w|} w \otimes v,$$

where $v \in V^{|v|}$ and $w \in W^{|w|}$; non-homogeneous elements are sums of homogeneous ones, so this determines $s_{V,W}$.

¹⁰For supermanifolds specifically, see <https://ocw.mit.edu/courses/mathematics/18-238-geometry-and-quantum-field-theory-fall-2006/lecture-notes/sec9.pdf>.

¹¹If the base field has characteristic 2, these two notions are actually the same, which quickly follows from Definition 3.7. But this will not be important to us.

So the point is if v or w is in V^1 , we multiply by -1 :

$$s(v \otimes w) = \begin{cases} -w \otimes v & v \text{ or } w \text{ is in } V^1 \\ w \otimes v, & v, w \in V^0. \end{cases}$$

This category is considerably more useful than it looks. There's a sense in which $s\text{Vec}$ and $\text{Vect}^{\mathbb{Z}/2}$ are the only two symmetric monoidal structures that can be placed on the monoidal category $(\text{Vect}^{\mathbb{Z}/2}, \otimes)$.

Other algebraic constructions are also different.

Definition 3.8. The *symmetric algebra* on a super-vector space V is the superalgebra $(\mathbb{Z}/2\text{-graded algebra})$

$$\text{Sym}^*(V) := T^*V / \langle v \otimes w - s(w \otimes v) \rangle.$$

Thus, if $V = V^0$, $\text{Sym}^*(V)$ is the usual symmetric algebra, but if $V = V^1$, $\text{Sym}^*(V) = \Lambda^*(V^1)$, the exterior algebra! In general, it'll be a mix of these two things.

We can use this to define polynomial functions: in ordinary algebra, there's a canonical isomorphism between the algebra of polynomials on a vector space V and $\text{Sym}^*(V^*)$.

Definition 3.9. Motivated by this, if $V \in s\text{Vec}$, we define its *algebra of polynomial functions* $\mathcal{O}(V)$ to be

$$\mathcal{O}(V) := \text{Sym}^*(V^*).$$

Here $V^* := \text{Hom}_{s\text{Vec}}(V, \mathbb{R}^{1|0}) = (V^0)^* \oplus (V^1)^*$. $\mathcal{O}(V)$ is itself a super-vector space, in fact a (super)commutative algebra! That is, $p \cdot q = (-1)^{|p||q|} q \cdot p$.

In physics, the even direction corresponds to bosonic stuff, and the odd direction to fermionic stuff. So \mathcal{C} may be a super-vector space, and we can take the action function $S \in \mathcal{O}^0(\mathcal{C})$.

Example 3.10. Let's consider a purely fermionic theory, such as $\mathcal{C} = \mathbb{R}^{0|2}$. Then, \mathcal{C} has coordinate functions $\psi^1, \psi^2 \in \mathcal{O}^1(\mathcal{C})$, which have odd statistics in the sense that

$$\begin{aligned} \psi^1 \psi^2 &= -\psi^2 \psi^1 \\ (\psi^1)^2 &= 0 \\ (\psi^2)^2 &= 0. \end{aligned}$$

This, $\mathcal{O}^0(\mathcal{C})$ has basis $\{1, \psi^1 \psi^2\}$ and $\mathcal{O}^1(\mathcal{C})$ has basis $\{\psi^1, \psi^2\}$. Thus $\text{Sym}^* \mathcal{C}$ is four-dimensional, which is as expected, since it should be $\Lambda^* \mathbb{R}^2$.

Since there's no quartic terms in ψ^1 and ψ^2 , we actually can't introduce interactions, so our action functional is

$$(3.11) \quad S := \frac{1}{2} M \psi^1 \psi^2.$$

This is somewhat like a function, but it behaves very weirdly: $S^2 = 0$!

We'd like to make sense of the partition function in this setting. In order to do this, we need rules for integrating over odd variables. To integrate over $\mathbb{R}^{0|1}$ with odd coordinate ψ , the most general function is $a\psi + b$, so we can stipulate that its integral is

$$\int_{\mathbb{R}^{0|1}} d\psi (a\psi + b) := a.$$

We'll define the exponential via its power series, which means it's much simpler than for bosons!

Now, on $\mathbb{R}^{0|k}$, we have to specify order of integration: to compute

$$\int_{\mathbb{R}^{0|k}} d\psi^1 d\psi^2 \cdots d\psi^k F = \int_{\mathbb{R}^{0|1}} d\psi^1 \left(\int_{\mathbb{R}^{0|1}} d\psi^2 \left(\cdots \int_{\mathbb{R}^{0|1}} F \right) \cdots \right),$$

first evaluate the innermost integral, then the next innermost, and so on, ending at the outermost ($d\psi^1$ in the above equation).

Hence the partition function is

$$\begin{aligned}
 Z &= \int_{\mathbb{R}^{0|2}} d\psi^1 d\psi^2 e^{-S(\psi^1, \psi^2)} \\
 &= \int_{\mathbb{R}^{0|2}} d\psi^1 d\psi^2 \left(1 - \frac{1}{2} M \psi^1 \psi^2\right) \\
 &= -\frac{1}{2} M \int_{\mathbb{R}^{0|2}} d\psi^1 d\psi^2 \psi^1 \psi^2 \\
 &= \frac{1}{2} M \int_{\mathbb{R}^{0|2}} d\psi^1 d\psi^2 \psi^2 \psi^1 \\
 &= \frac{1}{2} M.
 \end{aligned}$$

For bosons (i.e. even fields), we had a Gaussian

$$\int_{-\infty}^{\infty} e^{-Mx^2/2} dx = \frac{\sqrt{2\pi}}{\sqrt{M}}.$$

This is suggestive: if you arrange the masses of bosons and fermions right, things might cancel out to produce a theory whose dependence on the mass cancels out and is deformation-invariant.

Lecture 4.

Supersymmetry in zero dimensions: 9/12/17

We've been doing zero-dimensional quantum field theory, and we will continue to do so today. Last time, we introduced supersymmetry, so \mathcal{C} is a super-vector space. We looked at a particular specific example where \mathcal{C} is odd, e.g. $\mathbb{R}^{0|2}$, which has two odd coordinate functions $\psi^1, \psi^2 \in \mathcal{O}^1(\mathcal{C})$. The total coordinate algebra is $\mathcal{O}(\mathcal{C}) = \Lambda^*(\mathbb{R}^2)$, and the even functions are spanned by $1, \psi^1 \psi^2 \in \mathcal{O}^0(\mathcal{C})$. Since ψ^1 and ψ^2 are odd, $\psi^1 \psi^2 = -\psi^2 \psi^1$.

Let's introduce the action (3.11): since there are only odd terms, there can be no interacting terms, because higher-order powers of ψ^1 and ψ^2 vanish! The partition function is

$$(4.1) \quad Z = \int_{\mathcal{C}} d\mu e^{-S} = \int_{\mathcal{C}} d\mu \left(1 - \frac{1}{2} M \psi^1 \psi^2\right).$$

Last time, we discussed a heuristic way to understand the measure $d\mu$; today we'll be more explicit.

Definition 4.2. The *parity change operator* $\Pi: \text{sVec} \rightarrow \text{sVec}$ sends a super-vector space $V = V^0 \oplus V^1$ to the super-vector space with even part V^1 and odd part V^0 .

That is, Π just switches the odd and even parts of a super-vector space.

Definition 4.3. A *translation-invariant measure* on an odd super-vector space $V = V^1$ is a $d\mu \in \Lambda^{\text{top}}(\Pi V)$. For an $f \in \mathcal{O}(V) \cong \Lambda^*((\Pi V)^*)$, let $f^{\text{top}} \in \Lambda^{\text{top}}((\Pi V)^*)$ be its top-degree component; then, the integral of f with respect to $d\mu$ is

$$\int d\mu f := (d\mu) \cdot (f^{\text{top}}).$$

There's a one-dimensional space of measures, determined up to a scalar.

Exercise 4.4. For $V = \mathbb{R}^{0|1}$ with odd coordinate ψ , show there's a measure $d\mu$ on V such that for all $a, b \in \mathbb{R}$,

$$\int d\mu (a\psi + b) = a.$$

This measure is called $d\psi$. Notice that

$$\int d\psi \psi = 1 \quad \text{and} \quad \int d\psi = 0.$$

The fact that the integral of a constant in an odd direction is 0 is one of the striking features of this “Grassmann integration.”

We can also give names to some more measures:

Exercise 4.5. For $V = \mathbb{R}^{0|1}$ and $c \in \mathbb{R}$, show there are measures $c d\psi$ and $d(c\psi)$ on V such that

$$\begin{aligned} \int_V (c d\psi) f(\psi) &= c \int_V d\psi f(\psi) \\ \int_V d(c\psi) f(c\psi) &= \int_V d\psi f(\psi). \end{aligned}$$

Then prove the Grassmann change-of-variables formula

$$d(c\psi) = \frac{1}{c} d\psi,$$

or equivalently that

$$\int_V d(c\psi) c\psi = 1.$$

Similarly, on $\mathbb{R}^{0|q}$, define $d\psi = d\psi^1 d\psi^2 \cdots d\psi^q$ to be the unique measure such that

$$\int_{\mathbb{R}^{0|q}} d\psi \psi^q \psi^{q-1} \cdots \psi^1 = 1.$$

This definitely depends on how the ψ^i are ordered; we'll stick with this convention, which is common in physics.

These behave more like measures than top-degree forms: you need no choice of orientation to integrate. These definitions might be strange, but they're forced on you if you want a good change-of-variables formula.

Now, we know how to calculate the partition function (4.1):

$$Z = \int_{\mathbb{R}^{0|2}} d\psi \left(1 - \frac{1}{2} M \psi^1 \psi^2 \right) = \frac{1}{2} M.$$

More fermions. If we add more fermions, we can turn on interactions: if \mathcal{C} is any even-dimensional¹² odd super-vector space with translation-invariant measure $d\mu$, let

$$M \in \text{Sym}^2(V^*) = \Lambda^2((\Pi V^1)^*)$$

$$C \in \text{Sym}^4(V^*) = \Lambda^4((\Pi V^1)^*).$$

If V is at least four-dimensional, C can be nonzero. In that case we can change the action (3.11) to one with interactions:

$$S := \frac{1}{2} M + \frac{1}{4!} C \in \mathcal{O}(\mathcal{C}).$$

After choosing a basis for ΠV^1 , equivalently an isomorphism $V \cong \mathbb{R}^{0|q}$, we can rewrite this in coordinates:

$$(4.6) \quad S = \frac{1}{2} M_{IJ} \psi^I \psi^J + \frac{1}{4!} C_{IJKL} \psi^I \psi^J \psi^K \psi^L,$$

where M_{IJ} is an antisymmetric matrix with real entries, and C_{IJKL} is a totally antisymmetric tensor. Again the partition function is $Z = \int d\mu e^{-S}$, but this time it's possible to evaluate it algebraically.

Exercise 4.7. If $\mathcal{C} = \mathbb{R}^{0|4}$ and

$$(4.8) \quad S = m \psi^1 \psi^2 + m \psi^3 \psi^4 + \lambda \psi^1 \psi^2 \psi^3 \psi^4,$$

show that

$$Z = m^2 - \lambda.$$

Remark 4.9. This is much easier than the bosonic case, where calculations like this flowed through asymptotic series, Feynman diagrams, etc. There is a perturbation-theoretic description of the fermionic case as a Feynman diagram expansion; the rules are quite similar to those for bosons, but with some extra signs. ◀

In the theory (4.8), $Z_0 = m^2$.

¹²If \mathcal{C} is odd-dimensional, (4.6) still makes sense, but skew-symmetry forces us to leave out one fermion, so the partition function is 0. However, some correlation functions will be nonzero.

Remark 4.10. For the more general theory of the form (4.6), $Z_0 = \text{Pf}(M)$, the *Pfaffian* of the antisymmetric matrix M ; this is a number which squares to the determinant. ◀

To compute Z/Z_0 in (4.8), you can again sum over Feynman diagrams with four-valent vertices, but skew-symmetry introduces a sign rule which forces all Feynman diagrams with more than one vertex to have weight 0.

Bosons and fermions together. Now we consider $\mathcal{C} = V = V^0 \oplus V^1$ with both odd and even parts. We need a theory of integration for such spaces, but that won't be so hard: we'll first integrate over the odd part, then over the even part.

We also want some functions to integrate; polynomials don't have finite integrals on V^0 .

Definition 4.11. Let $C^\infty(V) := C^\infty(V^0) \otimes \mathcal{O}(V^1)$.

We also need a measure to integrate with.

Definition 4.12. Let V be a super-vector space.

- The *Berezinian line* of V is

$$\text{Ber}(V) := \Lambda^{\text{top}} V^0 \otimes (\Lambda^{\text{top}}(\Pi V^1))^*.$$

- An *integration measure* is an element of $\text{Ber}(V^*)$.¹³

If V is purely odd, this reduces to the above definition of the space of measures.

Since V now has an even subspace, integration will depend on orientation again.

Definition 4.13. Let V be an oriented super-vector space and $d\mu = \omega^0 \otimes \omega^1 \in \text{Ber}(V^*)$ be an integration measure. For any $f = f^0 \otimes f^1 \in C^\infty(V)$, its *integral* is

$$\int_V d\mu f := \int_{V^0} \omega^0 f^0 \left(\int_{V^1} \omega^1 f^1 \right).$$

That is: integrate the odd part, then the even part.

On $\mathbb{R}^{p|q}$ there's a canonical measure

$$d\mu = dx d\psi := (dx^1 \wedge \cdots \wedge dx^p) \otimes (d\psi^1 \cdots d\psi^q).$$

Example 4.14. Take $\mathcal{C} = \mathbb{R}^{1|2}$ with action

$$(4.15) \quad S(x, \psi^1, \psi^2) := S_1(x) + S_2(x) \psi^1 \psi^2.$$

Then, the partition function is

$$Z = \int dx d\psi e^{-S} = \int dx S_2(x) e^{-S_1(x)}.$$

In other words, in the purely bosonic theory with action S_1 , this is just the correlation function $\langle S_2(x) \rangle$. This can be a helpful perspective, but it also obscures why this is happening. ◀

For general S_1 and S_2 , these are not super interesting.¹⁴ But there is a special case that is much better. Fix an $h: \mathbb{R} \rightarrow \mathbb{R}$ such that as $|x| \rightarrow \infty$, $h(x) \rightarrow \infty$. Then, set

$$\begin{aligned} S_1(x) &:= \frac{1}{2} h(x)^2 \\ S_2(x) &:= h'(x). \end{aligned}$$

Hence

$$(4.16) \quad S = \frac{1}{2} h(x)^2 + h'(x) \psi^1 \psi^2,$$

using the action (4.15). This action is invariant under a certain odd vector field on \mathcal{C} ; we're going to explain what this means.

¹³To be completely precise, this would be a measure twisted by the orientation bundle, since measures don't require orientation to integrate.

¹⁴No pun intended.

Definition 4.17. Let A be a super-commutative super-algebra, a *derivation* D on A with degree $|D|$ is a function $D: A \rightarrow A$ such that $D(a + a') = D(a) + D(a')$ and

$$D(aa') = (Da)a' + (-1)^{|a||D|}a(Da').$$

The set of all derivations of $\mathcal{O}(V)$ is a super-vector space, which we'll denote $\text{Vect}(V)$.

Exercise 4.18. Show that $\text{Vect}(V)$ is a *super-Lie algebra*, in the same way that vector fields on a vector space are a Lie algebra.¹⁵

On $\mathbb{R}^{p|q}$, we have the usual derivations/vector fields $\partial_{x_i} \in \text{Vect}^0(V)$, but now also some odd vector fields $\partial_{\psi^I} \in \text{Vect}^1(V)$, defined to satisfy

$$\begin{aligned}\partial_{\psi^I}(x^i) &= 0 \\ \partial_{\psi^I}(\psi^J) &= \delta_I^J.\end{aligned}$$

Hence

$$\partial_{\psi^1}(\psi^1\psi^2) = \psi^2 \quad \text{and} \quad \partial_{\psi^1}(\psi^2\psi^1) = -\psi^2.$$

Now we'll discuss the symmetry in the action (4.16). Let

$$\begin{aligned}(4.19) \quad Q_1 &:= \psi^1\partial_x + h(x)\partial_{\psi^2} \\ Q_2 &:= \psi^2\partial_x - h(x)\partial_{\psi^1}.\end{aligned}$$

Then,

$$Q_1 S = \psi^1 h'(x)h(x) + h(x)h'(x)\partial_{\psi^2}(\psi^1\psi^2) = 0,$$

and similarly for $Q_2 S$.

Exercise 4.20. This means that if $X = [Q_1, Q_2]$, then X is an even vector field and $XS = 0$. Find X and show this explicitly.

There's also a sense in which Q_1 and Q_2 are divergence-free.

Definition 4.21. Let

$$X := h^i \partial_{x^i} + g^I \partial_{\psi^I}.$$

Then, the *Lie derivative* along X of a section of $\text{Ber}(V^*)$ is

$$\mathcal{L}_X(dx d\psi) := (\partial_{x^i} h^i + \partial_{\psi^I} g^I) dx d\psi.$$

If $\mathcal{L}_X(d\mu) = 0$, we say X is *divergence-free*.

There is a coordinate-free definition of this, which can be found in [31]. Other references on the general theory:

- Deligne-Morgan, "Notes on supersymmetry (following Joseph Bernstein)" [12].
- Witten recently wrote some notes on integration on supermanifolds in [30], which are pretty down-to-Earth.

Lemma 4.22. Let Q be a divergence-free vector field on an oriented super-vector space V with measure $d\mu$. For any $f \in C^\infty(V)$,

$$\int_V d\mu Qf = 0.$$

Proof. We'll compute in coordinates: suppose $Q = h^i \partial_{x^i} + g^I \partial_{\psi^I}$. Then,

$$\begin{aligned}\int_V d\mu Qf &= \int_{V^0} dx (Qf)^{\text{top}} \\ &= \int_{V^0} dx (h^i \partial_{x^i} f + g^I \partial_{\psi^I} f)^{\text{top}} \\ &= \int_{V^0} dx \left(-(\partial_{x^i} h^i) f + (-1)^{|g^I|} (\partial_{\psi^I} g^I) f \right)^{\text{top}} \\ &= 0,\end{aligned}$$

¹⁵There's a whole theory of super-manifolds and Lie super-groups and more. But it's possible to go a long way before needing to understand the whole package.

because Q is divergence-free. ⊠

Using this, we can show that a certain deformation of these theories is actually constant.

Proposition 4.23. *Let V be an oriented super-vector space with measure $d\mu$, $S \in C^\infty(V)$, and Q be a divergence-free odd vector field on V with $[Q, Q] = 0$ and $QS = 0$. For any smooth family of odd elements $\{\Psi_t\} \in C_c^\infty(V)$ with $\Psi_0 = 0$, let*

$$S_t := S + Q\Psi_t,$$

which is called a Q -exact deformation of S . Then, Z_t is independent of t .

Proof. Let $\Psi'_t := \partial_t \Psi_t$. Since

$$Z_t = \int_{\mathcal{C}} d\mu e^{-(S+Q\Psi_t)},$$

then

$$\begin{aligned} \partial_t Z_t &= - \int_{\mathcal{C}} d\mu (Q\Psi'_t) e^{-(S+Q\Psi_t)} \\ &= - \int_{\mathcal{C}} d\mu Q(\Psi'_t e^{-(S+Q\Psi_t)}) \\ &= 0 \end{aligned}$$

by Lemma 4.22. ⊠

Lecture 5.

Localization in supersymmetry: 9/14/17

Today, we're going to use the $(0+1)$ -dimensional field theory that we've been developing to do something actually topological. Recall that our state space is $\mathcal{C} = \mathbb{R}^{1|2}$, and given a smooth $h: \mathbb{R} \rightarrow \mathbb{R}$ such that $|h(x)| \rightarrow \infty$ as $|x| \rightarrow \infty$, we defined the action (4.16), and we'd like to compute its partition function $Z = \int_{\mathcal{C}} e^{-S}$.

Rather than boldly going forward as in previous lectures, we first observed that the partition function is invariant under two symmetries Q_1 and Q_2 (4.19). If

$$Q := Q_1 + Q_2 = (\psi^1 + \psi^2)\partial_x + h(x)(\partial_{\psi^2} - \partial_{\psi^1}),$$

then Q acts on $C^\infty(\mathcal{C})$ and $[Q, Q] = 0$. By Proposition 4.23, for any deformation $\psi_t \in C_c^\infty(\mathcal{C})$. That is, if $S_t := S + Q(\psi_t)$ and $Z_t := \int e^{-S_t}$, then $\partial_t Z_t = 0$. One way to think of this is to take Q as a differential operator and consider “ Q -cohomology” — then, Proposition 4.23 tells us that Z_t only depends on the cohomology class of S .

Consider deforming $h(x)$ to a family $h_t(x)$ in a compactly supported manner, which defines a variation S_t of S . Using dots to denote $\frac{d}{dt}$,

$$\dot{S}(x) = h(x)\dot{h}(x) + \dot{h}'(x)\psi^1\psi^2.$$

Since $\dot{S}(x) = Q\Psi$ with $\Psi = -\dot{h}(x)\psi^1$, Proposition 4.23 tells us that Z does not depend on $h(x)$, as long as you only take compactly supported deformations.

Exercise 5.1. Bootstrap this to show that Z only depends on the behavior at infinity: it's only a function of ε_\pm , where $\lim_{h \rightarrow \pm\infty} \varepsilon_\pm = \varepsilon_\pm \infty$.

This is in a sense topological; certainly, there's no dependence on the metric.

One way to think of this which will come up again and again is that the action makes the configuration space only care about compact things. If you switch the behavior of h at $\pm\infty$, which requires doing something noncompact, it will change the invariants. Donaldson theory has the same behavior, with chambers in which the invariants do not change (where $b_2^+(X) > 1$), plus “wall-crossing phenomena” on their boundaries (where $b_2(X) = 1$).

Now let's compute Z , using topological invariance and a trick called localization. Since Z doesn't depend on our choice of h , let's do something nice: Z does not depend on λ in the variation $h(x) \rightarrow \lambda h(x)$ for $\lambda > 0$, so let's compute the limit as $\lambda \rightarrow \infty$. That is, we need to understand the asymptotics of

$$(5.2) \quad \lim_{\lambda \rightarrow \infty} \int_{-\infty}^{\infty} dx e^{-\lambda F(x)},$$

where $F(x) \rightarrow \infty$ as $x \rightarrow \infty$. The asymptotics are controlled by something called the *method of steepest descent*, which may be surprising at first.

Proposition 5.3. Assume F has a unique global minimum at x_c .¹⁶ Then, as $\lambda \rightarrow \infty$,

$$(5.4) \quad \int_{-\infty}^{\infty} dx e^{-\lambda F(x)} \sim \sqrt{\frac{2\pi}{\lambda F''(x_c)}} e^{-\lambda F(x_c)}.$$

That is, neither side of (5.4) has a limit at $\lambda \rightarrow \infty$, but their ratio does, and its limit is 1.

The proof is a little bizarre: first you only look at a tiny neighborhood of x_c , and then expand that neighborhood to the whole real line, and each of these contributes an exponentially small amount to the integral. For a full proof, check out [6]; it has a big list of cool tricks for computing asymptotic expansions like this one.

Another way to interpret (5.3) is that it gives us permission to truncate $F(x)$ to quadratic order around F_c . Thus, let's reshape h such that all of its global minima are 0, and make a quadratic approximation of S by summing over all of these local minima. The fermionic part is already quadratic, so we just have to look at the bosonic part. As $\lambda \rightarrow \infty$, we get that

$$Z(\lambda) \simeq \sum_{x_c: h(x_c)=0} \int dx d\psi \exp\left(-\frac{1}{2}\lambda^2 h'(x_c)^2 (x - x_c)^2 - \lambda h'(x_c) \psi^1 \psi^2\right).$$

This is a Gaussian in the bosonic and in the fermionic parts:

$$\begin{aligned} &= \sum_{x_c} \left(\sqrt{\frac{2\pi}{\lambda^2 h'(x_c)^2}} (\lambda h'(x_c)) \right) \\ &= \sqrt{2\pi} \sum_{h_c} \frac{h'(x_c)}{|h'(x_c)|} \\ &= \sqrt{2\pi} \sum_{h_c} \text{sign}(h'(x_c)). \end{aligned}$$

This is actually not so hard to evaluate directly, first integrating out the fermionic part then the bosonic part, but it's a useful example nonetheless.

Now we can look at one function in each deformation class.

- If $\lim_{x \rightarrow \pm\infty} h(x) = \pm\infty$, then there are an odd number of points x_c with sign +1 and an even number with sign -1, so we get $Z/\sqrt{2\pi} = -1$.
- If $\lim_{x \rightarrow \pm\infty} h(x) = \mp\infty$, this reverses: there are an even number with sign +1 and an odd number with sign -1, so $Z/\sqrt{2\pi} = -1$ again.
- If $\lim_{x \rightarrow \pm\infty} h(x) = \infty$ (or if it goes to $-\infty$), the number of critical points with positive and negative signs are the same, so $Z = 0$.

Hence you can express this in terms of ε_{\pm} :

$$Z = \sqrt{2\pi} \frac{|\varepsilon_+ - \varepsilon_-|}{2}.$$

Localization in a 0-dimensional σ -model. Let (M, ω) be a compact, $2n$ -dimensional symplectic manifold: this means ω is a differential 2-form on M with $d\omega = 0$ and $\omega^n \neq 0$, and assume there is $U(1)$ -action on M generated by the vector field $Y := \omega^{-1}(dH)$, where $H: M \rightarrow \mathbb{R}$ is some function.

Let's assume all the fixed points of Y are isolated,¹⁷ and pick an $\alpha \in \mathbb{R}$. We're going to use all this stuff to prove something cool, a formula for

$$(5.5) \quad \int_M \frac{\omega^n}{n!} e^{i\alpha H}.$$

¹⁶Otherwise, there would be a sum over local minima.

¹⁷It is possible to excise this assumption, but it's helpful for now.

Example 5.6. In examples, this integral is something people actually care about. Let $M = S^2$ with $\omega := \sin \theta d\theta \wedge d\varphi$ and $H := z = \cos \theta$, so $Y = \partial_\varphi$. Then,

$$\int_M \frac{\omega^n}{n!} e^{i\alpha H} = \int_{S^2} e^{i\alpha \cos \theta} \sin \theta d\theta \wedge d\varphi = 2\pi \int_0^\pi e^{i\alpha \cos \theta} \sin \theta d\theta.$$

This is a Bessel function, and it's also funny to notice it is a great example of the kinds of integrals you teach for u -substitution and never expect to see anywhere else. It is hence easy to solve:

$$\begin{aligned} 2\pi \int_0^\pi e^{i\alpha \cos \theta} \sin \theta d\theta &= -2\pi \int_1^{-1} e^{i\alpha z} dz \\ &= \frac{2\pi}{i\alpha} (e^{i\alpha} - e^{-i\alpha}) \\ &= 4\pi \frac{\sin \alpha}{\alpha}. \end{aligned}$$

◀

This answer demonstrates a localization phenomenon: it's a sum of contributions only from the north and south poles. In general, the integral (5.5) is a sum of contributions

$$(\pm) \frac{2\pi}{i\alpha} e^{i\alpha H(x_c)},$$

summed over the fixed points x_c of the $U(1)$ -action. This is an instance of the Duistermaat-Heckman theorem [15], and we're going to prove it using localization in supersymmetry.

To do this, we're going to need a supermanifold that's not a super-vector space, but it's not so bad.

Definition 5.7. Let $E \rightarrow M$ be a vector bundle. Its *parity change* ΠE is a supermanifold whose algebra of functions is $C^\infty(\Pi E) := C^\infty(M, \Lambda^*(E))$.

We won't go into the general theory of supermanifolds here. Concretely, for $E = TM$, in local coordinates on M , we have even coordinates x^i , $i = 1, \dots, 2n$, and odd coordinates ψ^i for $i = 1, \dots, 2n$, and we can translate between functions on $\Pi(TM)$ and differential forms on M by exchanging ψ^i and dx^i .

If we want to write down a zero-dimensional quantum field theory, we ought to have an action. Let $\mathcal{C} := \Pi(TM)$ and take

$$S := -i\alpha(H + \omega),$$

or in coordinates,

$$= -i\alpha(H + \omega_{ij}\psi^i\psi^j).$$

There's a canonical measure (up to scaling) $dx d\psi$ on $\Pi(TM)$, which in local coordinates is exactly $dx d\psi$ from before, and is invariant under change-of-charts. This might be surprising. A more abstract way to think of this is that the super-tangent bundle $T\mathcal{C}$ to \mathcal{C} factors into a short exact sequence

$$0 \longrightarrow \Pi(\pi^* TM) \longrightarrow T\mathcal{C} \longrightarrow TM \longrightarrow 0,$$

so $\text{Ber}(T\mathcal{C}) = \text{Ber}(TM) \otimes \text{Ber}(\Pi(TM))$, hence must be trivial. Hence the partition function is

$$Z := \int_{\mathcal{C}} dx d\psi e^{-S}.$$

If we integrate over fermions first, we get

$$Z = (i\alpha)^n \int_M \frac{\omega^n}{n!} e^{i\alpha H}.$$

We want to compute this by localization. This means first writing down a vector field under which S is invariant. We'll take

$$Q := d + \iota_Y = \psi^i \partial_{x^i} + Y^j \partial_{\psi^j},$$

where ι_Y is contraction,¹⁸ then, Q is odd and S is invariant under it:

$$\begin{aligned} QS &= (d + \iota_Y)(H + \omega) \\ &= dH + \iota_Y \omega \\ &= 0, \end{aligned}$$

Moreover,

$$\begin{aligned} \frac{1}{2}[Q, Q] &= [d, \iota_Y] = \mathcal{L}_Y \\ &= \psi^i (\partial_{x^i} Y^j) \partial_{\psi^j} + Y^i \partial_{x^i}. \end{aligned}$$

We want to use localization to obtain the fixed points of Y through a perturbation $S \rightarrow S + \lambda Q\Psi$ as $\lambda \rightarrow \infty$ for some Ψ . To do this, we need to choose a $U(1)$ -invariant metric g on M , which we can always do, and the answer will turn out to not depend on it. Then, let

$$\Psi := g(Y) = g_{ij} \psi^i Y^j = \psi^i Y_i.$$

Exercise 5.8. Check that

$$(5.9) \quad Q\Psi = g(Y, Y) - d(gY)$$

$$(5.10) \quad Q^2\Psi = 0,$$

e.g. by computing in coordinates.¹⁹

Thus, we know the perturbation (5.9). As before, Z is independent of λ , which uses (5.10).

Now, you can take $\lambda \rightarrow \infty$ and use the method of steepest descent to conclude that

$$(5.11) \quad Z \sim \sum_{x_c \in M: Y(x_c)=0} e^{iaH(x_c)} (2\pi)^n \frac{(d(gY)(x_c))^n / n!}{\sqrt{\det(D^2(g(Y, Y)))(x_c)}}$$

This looks complicated, but just like before, the fermionic and bosonic pieces almost cancel each other out, leaving behind a topological contribution. Here, $g(Y, Y)$ is a real-valued function on M , so we can take its Hessian $D^2(g(Y, Y))$ and evaluate it at the critical point x_c .

Both the numerator and the denominator of the fraction in (5.11) are naturally valued in $\Lambda^{\text{top}} T_{x_c}^* M$. We'll exploit this to calculate their ratio in a local model by diagonalizing the $U(1)$ -action on $T_{x_c} M$. That is, we choose an isomorphism

$$T_{x_c} M \cong \bigoplus_{i=1}^n \mathbb{R}_i^2,$$

where $U(1)$ acts on \mathbb{R}_i^2 by

$$\theta \mapsto \begin{pmatrix} \cos k_i \theta & \sin k_i \theta \\ -\sin k_i \theta & \cos k_i \theta \end{pmatrix}$$

for some weights $k_1, \dots, k_n \in \mathbb{R}$.

Example 5.12. The 2-dimensional case is simplest: take $T_{x_c} M = \mathbb{R}^2$ with weight k . Let $\omega = r dr \wedge d\theta$ be the symplectic form and the standard metric $g := dr^2 + r^2 d\theta^2$ is $U(1)$ -invariant. Then, $Y = k\partial_\theta$, $H = (1/2)kr^2$, and $g(Y, Y) = k^2 r^2$, then

$$(5.13) \quad d(gY) = 2kr dr \wedge d\theta \sqrt{\det(D^2(g(Y, Y)))} = 2k^2 r dr \wedge d\theta.$$

Again, almost everything cancels out, so we get

$$Z = (2\pi)^n \sum_{x_c} \frac{e^{iaH(x_c)}}{\prod_{i=1}^n k_i(x_c)},$$

i.e.

$$(5.14) \quad \int_M \frac{\omega^n}{n!} e^{iaH} = \left(\frac{2\pi}{ia} \right)^n \sum_{x_c} \frac{e^{iaH(x_c)}}{\prod_{i=1}^n k_i(x_c)}.$$

¹⁸This is the differential in the Cartan model for the $U(1)$ -equivariant cohomology for M , but that's not important right now.

¹⁹Even though the derivations only form a (super)-Lie algebra, so Q^2 doesn't make sense on that level, it's still acting on vector fields, and we can iterate its action. This differs from $[Q, Q]$ by $1/2$, so it doesn't make a difference.

(5.14) is known as the *Duistermaat-Heckman formula*. We’ve just given a completely rigorous proof of it, which probably differs greatly from their original proof in [15].

Next time, we’ll wrap up this story and begin thinking about higher dimensions.

Lecture 6.

One-dimensional QFT: 9/19/17

Reminder: there are exercises in the professor’s notes, and you should try them!

We’ve been talking about localization in the past week, and we’re on the way to thinking about effective field theory. Last time, we discussed the Duistermaat-Heckman formula (5.14) for a compact symplectic manifold (M, ω) (i.e. M is a compact $2n$ -dimensional manifold, and $\omega \in \Omega^2(M)$ is a closed form with $\omega^n \neq 0$). We assumed we had a vector field Y which generates a $U(1)$ -action,²⁰ This Y is generated by a Hamiltonian $H: M \rightarrow \mathbb{R}$, in the sense that $Y = \omega^{-1}(dH)$; we assume H has isolated fixed points.²¹

Then, we showed that the integral

$$\int_M \frac{\omega^n}{n!} e^{i\alpha H}$$

depends only on the fixed points of H , and the precise formula is (5.14). The equation uses the infinitesimal $U(1)$ -action on the tangent space of a fixed point.

Exercise 6.1. Suppose V is a $2n$ -dimensional vector space with an orientation and an inner product, i.e. a reduction of the structure group from $GL_{2n}(\mathbb{R})$ to $SO(2n)$. Then, define a natural line $\text{Pf}(V)$ with $\text{Pf}(V)^{\otimes 2} = \text{Det}(V) := \Lambda^{2n}(V)$.

Exercise 6.2. We saw what happens for $M = S^2$ last time, in Example 5.6. Try it with \mathbb{CP}^2 .²²

Remark 6.3. Another way to interpret (5.14) is that “the stationary phase approximation to

$$\int \frac{\omega^n}{n!} e^{i\alpha H}$$

is exact.” This is an asymptotic analysis as $\alpha \rightarrow \infty$; we’ve already done this for things like $\int e^{-tF}$, but in this case there’s something weirder going on: as $\alpha \rightarrow \infty$, the function is oscillating more and more rapidly. The fact that it only depends on the critical points in the end is a manifestation of the fact that these oscillations cancel each other out.

This stationary phase analysis is much like the method of steepest descent that we’ve been doing: approximate the integrand by its quadratic Taylor expansion around each critical point. There are some tricky technicalities, and you have to make rigorous the idea that you’re integrating something only conditionally convergent

The point is, if you hear someone saying the stationary phase approximation is exact, that’s a different statement with a different proof than the approach we used. There’s a really great exposition of this approach in [5]. ◀

In our proof of the Duistermaat-Heckman formula, we used localization for

$$\int_{\mathcal{C}} d\mu e^{-S},$$

where $\mathcal{C} = \Pi TM$, the parity change of the tangent bundle, and

$$S = -i\alpha(\omega + H) \in C^\infty(\mathcal{C}) = \Omega^*(M).$$

If $Q := \iota_Y + d$, then $QS = 0$.

There’s an interpretation of this in terms of $U(1)$ -equivariant cohomology which allows for a more general formula than (5.14). Namely, we think of Q as an “equivariant differential,” and we can generalize to any $S \in C^\infty(\mathcal{C})$ with $QS = 0$, i.e. any *equivariantly closed* form $\alpha \in \Omega^*(M)$ on any compact manifold M with a $U(1)$ -action.

²⁰This is a strong assumption: a general smooth function $H: M \rightarrow \mathbb{R}$ can be taken for a Hamiltonian, and we can let $Y = \omega^{-1}(dH)$, which generally does not generate a $U(1)$ -action, as $e^{2\pi Y} \neq \text{id}$. So as an equation, we’re assuming $e^{2\pi Y} = \text{id}$.

²¹Unlike the previous assumption, this is generically true. We also assumed H is Morse in the final step; this assumption probably can be removed, but the argument will be nontrivial.

²²**TODO:** which circle action?

Theorem 6.4 (Atiyah-Bott-Berline-Vergne [5, 8]). *Let M be a compact manifold with a $U(1)$ -action with isolated fixed points $\{x_c\}$, and let $\beta \in \Omega^*(M)$ be an equivariantly closed form. Then,*

$$\int_M e^\beta = (-2\pi i)^n \sum_{x_c} \frac{e^{\beta^{\text{bot}}(x_c)}}{\prod_{i=1}^n k_i(x_c)}.$$

Here, β^{bot} denotes the piece of β in Ω^0 ; equivariantly closed forms are generally non-homogeneous.

The equivariant folks also call this theorem “the localization theorem in equivariant cohomology,” and like this formulation of it better.

We can also generalize to non-isolated fixed points, and we will need to use this later. In this case, the steepest descent analysis of

$$\int d\mu e^{-S+\lambda Q\Psi}$$

as $\lambda \rightarrow \infty$ is now localized on the fixed set P , and the integrand is determined by the local structure around P .²³

Let NP denote the normal bundle, and recall that in the steepest descent analysis, we introduced a $U(1)$ -invariant metric g on M . Since the volume form ω^n on a symplectic manifold defines an orientation, NP is also oriented, and the orientation and the metric g define an $SO(2n)$ -structure.

Definition 6.5. Let X be a manifold with a $U(1)$ -equivariant vector bundle $E \rightarrow X$ together with a reduction of its structure group to $SO(2n)$ compatible with the $U(1)$ -action. Let $Y \in \Omega^0(\mathfrak{so}(E))$ denote the action of $U(1)$, and choose a $U(1)$ -invariant metric g on E and let $F \in \Omega^2(\mathfrak{so}(E))$ denote its curvature form. Then, the *equivariant Euler form* of E is

$$\text{Eul}(E) := \text{Pf}\left(\frac{1}{2\pi}(Y + F)\right).$$

In general, $\text{Eul}(E)$ is concentrated in even degrees of $\Omega^*(X)$.²⁴ If $n = 1$, the Euler form has a simpler formula:

$$\text{Eul}(E) = \frac{1}{2\pi}(ik + F),$$

with k as in Example 5.12.

More generally, the bottom piece of the Euler form is $\prod ik_i/2\pi \neq 0$, so if all $k_i \neq 0$, there’s an inverse to the Euler form $1/\text{Eul}(E) \in \Omega^*(X)$, using the fact that

$$\frac{1}{a+x} = \frac{1}{a} \left(\frac{1}{1+a^{-1}x} \right) = \frac{1}{a} (1 - a^{-1}x + a^{-2}x^2 - \dots).$$

This leads to the most general version of the ABBV formula, which is one of the coolest things you can do with 0-dimensional supersymmetric quantum field theory.

Theorem 6.6 (Atiyah-Bott-Berline-Vergne [5, 8]). *With M , β , and P as above,*

$$\int_M e^\beta = \int_P \frac{e^\beta}{\text{Eul}(NP)}.$$

Quantum field theory in one dimension. Now we’ll move on to the one-dimensional case, which specializes to undergraduate quantum mechanics. Choose a compact Riemannian 1-manifold (X, η) : either $X = [0, T]$, or $X \cong S^1$ with circumference T . We’ll parametrize X by t , which you can think of as time. Now, the space \mathcal{C}_X of (some kind of generalized) functions on X will be infinite-dimensional.

Let’s define a theory. Fix a Riemannian manifold (Y, g) , which we’ll call the *target*, and $V: Y \rightarrow \mathbb{R}$, called the *potential*.

- For $X \cong S^1$, let $\mathcal{C}_{S^1} := \{\phi: S^1 \rightarrow Y\}$, the C^∞ maps from S^1 to Y , and
- for $X = [0, T]$, fix $y_0, y_1 \in Y$ and let $\mathcal{C}_{[0,T]_{y_0}^{y_1}} := \{\phi: [0, T] \rightarrow Y \mid \phi(0) = y_0, \phi(T) = y_1\}$.

So for S^1 we get loops, and for $[0, T]$ we get paths with chosen endpoints. Let dV_X denote the volume form on X and R denote the scalar curvature of Y ; then, we define the action $S: \mathcal{C}_X \rightarrow \mathbb{R}$ to be

$$(6.7) \quad S(\phi) := \int_X dV_X \frac{1}{2} \left(g(\dot{\phi}, \dot{\phi}) + V(\phi) - \frac{1}{3} R(\phi) \right).$$

²³We may need to make a transversality assumption on P , but it’s OK.

²⁴**TODO:** Does the cohomology class of the Euler form depend on the metric?

Or in coordinates,

$$(6.8) \quad = \int_X dt \sqrt{\eta_{tt}} \left(\frac{1}{2} g_{ij}(\phi(t)) \dot{\phi}^i(t) \dot{\phi}^j(t) \eta^{tt}(t) + V(\phi(t)) - \frac{1}{3} R(\phi(t)) \right).$$

If you parametrize X by arc length, $\eta_{tt} = 1$ and this simplifies:

$$(6.9) \quad = \int_X \frac{1}{2} g_{ij} \dot{\phi}^i \dot{\phi}^j + V(\phi(t)) - \frac{1}{3} R(\phi(t)).$$

We would like to define the partition function

$$(6.10) \quad Z_X \text{ “=” } \int_{\mathcal{C}_X} d\phi e^{-S(\phi)},$$

but here we run aground: $d\phi$ is now a measure on an infinite-dimensional Banach space. There’s no analogue of the Lebesgue measure here: a unit ball contains infinitely many balls of radius $1/4$, so there’s no consistent way to define the volume of anything to be nonzero and finite. Nonetheless, in a sense $d\phi$ doesn’t exist, but the whole expression (6.10) will exist; it’s something that statistical mechanics researchers call the Wiener measure.

Physicists make sense of (6.10) by discretization. For concreteness, set $X = [0, T]$ and fix $y_0, y_1 \in Y$. We’ll replace X by a *lattice*: for some $N > 0$, let $t_0, t_1, \dots, t_N \in X$ such that $\Delta_t = t_j - t_{j-1} = T/N$. The discretized field space $\mathcal{C}_{X;N}$ is the space of piecewise geodesic paths $\phi : X \rightarrow Y$ that are smooth on (t_{j-1}, t_j) and such that the path from $\phi(t_{j-1})$ to $\phi(t_j)$ is the unique minimal geodesic between them.²⁵ The map $\phi \mapsto (\phi(t_0), \dots, \phi(t_N))$ defines an embedding $\mathcal{C}_{X;N} \subset Y^{N+1}$, and this is a finite-dimensional manifold, so we can use the *product measure*

$$d\mu_N := \frac{1}{(4\pi\Delta t)^{N \dim Y/2}} \prod_{n=1}^{N-1} du dy(\phi(t_n)).$$

Then, we can define the discretized partition function

$$Z_{X;N} = \int_{\mathcal{C}_{X;N}} e^{-S} d\mu_N,$$

and try to take the limit as $N \rightarrow \infty$. This does exist!

Theorem 6.11. *The limit $\lim_{N \rightarrow \infty} Z_{X;N}$ exists, and is the heat kernel $k_T(y_0, y_1)$.*

Interestingly, it only depends on the endpoints y_0, y_1 and the total length.

Definition 6.12. Fix Y and V as above. For $t \in \mathbb{R}_+$, the *heat kernel (deformed by V)* is a smooth function $k_t : Y \times Y \rightarrow \mathbb{R}$ satisfying the *heat equation*

$$(6.13) \quad \partial_t k_t(x, y) + (-\Delta_X + V(x))k_t(x, y) = 0,$$

and as a distribution,

$$(6.14) \quad \lim_{t \rightarrow 0} k_t(x, y) = \delta(x, y).$$

You can also characterize the heat kernel as the fundamental solution to the heat equation (6.13).

Exercise 6.15. Show that when $Y = \mathbb{R}^n$ and $V = 0$, the heat kernel is

$$(6.16) \quad k_t(x, y) = \left(\frac{1}{4\pi t} \right)^{n/2} \exp\left(-\frac{1}{4t} \|x - y\|^2 \right).$$

The heat kernel is the kernel of an integral operator, the operator U_t of heat evolution for time t . This is the operator evolving solutions to (6.13) forward in time. As an integral kernel, this has the formula

$$(U_t f)(x) = \int_M du dy k_t(x, y) f(y) dy.$$

U_t is a *smoothing* operator: it maps distributions to smooth functions. It also defines a linear operator on $L^2(M)$ which has the formula

$$U_t = e^{-t(-\Delta + V)}.$$

²⁵One might be surprised to learn this stuff was formalized and written down surprisingly recently, in the mid-2000s.

Heuristic proof of Theorem 6.11 when $V = 0$. Let's discretize the heat operator: let $U_T = (U_{\Delta t})^N$ and

$$(6.17) \quad K_T(y_N, y_0) := \int_{Y^{N-1}} \prod_{n=1}^{N-1} \text{dvol} \prod_{n=0}^{N-1} k_{\Delta t}(y_{n-1}, y_n).$$

When Δt is sufficiently small (N is sufficiently large), we have short-time asymptotics of $k_{\Delta t}$:

$$k_{\Delta t}(x, y) \sim \left(\frac{1}{4\pi\Delta t} \right)^{\dim Y/2} \exp\left(-\frac{1}{4\Delta t} d(x, y)^2 \right).$$

This is the piece that we're not making precise. If you substitute this into (6.17), you get

$$\begin{aligned} k_T(y_N, y_0) &\sim \int_{Y^{N-1}} \prod_{n=1}^{N-1} \text{dvol} \prod_{n=0}^{N-1} \left(\frac{1}{4\pi\Delta t} \right)^{\dim Y/2} \exp\left(-\frac{1}{4\pi\Delta t} d(y_{n+1}, y_n)^2 \right) \\ &= \int \text{d}\mu_N \exp\left(-\frac{\Delta t}{4} \left(\frac{d(y_{n+1}, y_n)}{\Delta t} \right)^2 \right) \\ &= Z_{X;N}. \end{aligned}$$

⊠

Though this was not a proof, this proof can be made rigorous; see [11]. Exactly where the scalar curvature goes is somewhat of a mystery, though some more careful analysis of the asymptotics above can be found in [7].

Lecture 7.

Local observables: 9/21/17

One thing that came up a few times in the past few lectures about localization is the question of if M is a compact manifold with a $U(1)$ -action with isolated fixed points, why is the infinitesimal action nontrivial? The idea is that in a local model, i.e. \mathbb{R}^n with a single fixed point, the $U(1)$ -action is diffeomorphic to the standard rotation action with the origin as its fixed point, whose infinitesimal action is nontrivial.

More rigorously, one can fix a $U(1)$ -invariant Riemannian metric on M , which means the exponential map is a $U(1)$ -invariant map between a tubular neighborhood of the fixed-point set and its normal bundle, which implies the action must be nontrivial. Thus we do not need to worry about transversality, etc. Choosing such a metric requires averaging over $U(1)$, and therefore crucially requires $U(1)$ to be compact.

Last time, we also talked about the heat kernel: on a compact Riemannian manifold Y and for $t \in (0, \infty)$, given a potential function $V : Y \rightarrow \mathbb{R}$, we obtained a heat kernel function $k_t : Y \times Y \rightarrow \mathbb{R}$ defined to satisfy (6.13) and (6.14), which uniquely characterizes it. If $V = 0$, this is the usual heat equation, and in general it's a perturbation. The idea (when $V = 0$) is that if there's a point source of heat at y at $t = 0$, $k_t(x, y)$ calculates the amount of heat at x at time t , so for t small, it looks like an approximation to a δ -function, and when t is large, heat is spread evenly (on a compact manifold).

We then considered a 1-dimensional quantum field theory whose space of fields $\mathcal{C}_{[0,T]_{y_0}^{y_1}}$ on the interval $[0, T]$ with the usual metric is the space of functions $\phi : [0, T] \rightarrow Y$ with $\phi(0) = y_0$ and $\phi(T) = y_1$. The action is

$$S = \int_0^T \frac{1}{4} g(\dot{\phi}, \dot{\phi}) + V(\phi) - \frac{1}{3} R(\phi),$$

where R is the scalar curvature on Y . Then, in Theorem 6.11, we showed that the partition function

$$Z_{[0,T]_{y_0}^{y_1}} = k_T(y_0, y_1),$$

though defining this as an integral on $\mathcal{C}_{[0,T]_{y_0}^{y_1}}$ doesn't quite make sense; instead, we had to discretize $[0, T]$ and the action and the path integral and take a continuum limit. In each case there is a finite-dimensional space of paths and the integral does make sense, and what you get is kind of a Riemann sum for the path integral.

The proof is trickier than one thinks: if you just use the leading term, you don't get the scalar curvature, and the next term is where the scalar curvature comes from, but there's an extra factor of 2 to account for.

Precisely, the leading term in the asymptotic is

$$k_{\Delta t}(x, y) \underset{\substack{\Delta t \rightarrow 0 \\ x \rightarrow y}}{\sim} \left(\frac{1}{4\pi\Delta t} \right)^{\dim Y/2} \exp\left(-\frac{1}{4\Delta t} d(x, y)^2 \right).$$

The scalar curvature comes up in the next term, but there are corrections in Δt and $x - y$, and these account for the spurious factor of 2.

Remark 7.1. People care about the heat kernel for a lot of different reasons, but this is a good one: the simplest version of quantum mechanics (one-dimensional QFT) has the heat kernel as its partition function, so this is really something fundamental. ◀

One new feature of 1-dimensional QFT is that you can glue intervals together. If you trace this through the argument, this turns into the semigroup law for the heat kernel: the heat equation describes time evolution of something, and composing the intervals $[0, T_1]$ and $[T_1, T_2]$ corresponds to first evolving the system for time T_1 , then using that as the initial condition and evolving for $T_2 - T_1$. In higher dimensions there are more and more ways to do this, and therefore there are more and more interesting structures.

Remark 7.2. We only gave the proof for $V = 0$; in the case $V \neq 0$ there's an analogous proof using *Trotter's formula*

$$e^{A+B} = \lim_{N \rightarrow \infty} (e^{A/N} e^{B/N})^N. \quad \blacktriangleleft$$

You can also formulate this QFT on a sphere. In this case the partition function is a trace:

$$Z_{S^1(T)} = \int_Y k_T(y, y) dy.$$

If $U_T: L^2(Y) \rightarrow L^2(Y)$ is the time evolution operator by T , with formula

$$U_T = e^{-(\Delta+V)T},$$

which is unitary, then U_T is a trace-class operator (meaning that, though it's infinite-dimensional, its trace can still be rigorously defined).

Local observables. There's yet another structure present in the 1-dimensional case that we didn't have in dimension zero. Let $\mathcal{O}: \mathcal{C}_X \rightarrow \mathbb{R}$ be a function which only depends on the k -jet of ϕ at a $t \in X$, i.e. only depends on the first k derivatives of ϕ at t (including $\phi(t)$).

Example 7.3. One quick example of a local observable would be to choose a function $F: Y \rightarrow \mathbb{R}$ and define

$$(\mathcal{O}_F(t))(\phi) := F(\phi(t)). \quad \blacktriangleleft$$

We can then define its *expectation* using a path integral:

$$(7.4) \quad \langle \mathcal{O}_F(t) \rangle := \int_{\mathcal{C}_X} d\phi (\mathcal{O}_F(t))(\phi) e^{-S} = \int_{\mathcal{C}_X} d\phi F(\phi(t)) e^{-S(\phi)}.$$

Once again, to make this rigorous one must discretize and show that the continuum limit exists, but when one does, there's a nice answer.

Let $\widehat{F}: L^2(Y) \rightarrow L^2(Y)$ denote the operator sending $\psi \mapsto F \cdot \psi$. Also let $H := -\Delta + V$.

Theorem 7.5. $\langle \mathcal{O}_F(t) \rangle_{[0,T]_{y_0}^{y_1}}$ is the kernel representing

$$e^{-H(T-t)} \widehat{F} e^{-Ht}.$$

Exercise 7.6. Give a heuristic proof for this, similar to the one we gave for Theorem 6.11. (Or a rigorous proof, if you want; it should look very similar to the one in [11].)

So the idea is that we've stuck the operator into the path integral, and this computes a modified heat flow, where we've stuck in a related operator into the time evolution.

If one has multiple local observables $\mathcal{O}_{F_1}(t_1), \dots, \mathcal{O}_{F_k}(t_k)$, there's a similar definition for the expectation $\langle \mathcal{O}_{F_1}(t_1) \cdots \mathcal{O}_{F_k}(t_k) \rangle$, where you stick them all into the path integral. The answer is very similar to the one in Theorem 7.5: if we do this on a circle and assume that $t_1 < \dots < t_k$, then

$$\langle \mathcal{O}_{F_1}(t_1) \cdots \mathcal{O}_{F_k}(t_k) \rangle = \text{tr}(e^{-(T-t_k)H} \widehat{F}_k \cdots \widehat{F}_2 e^{-(t_2-t_1)H} \widehat{F}_1 e^{-t_1 H}).$$

There's a similar answer on the interval (though without a trace).

The general slogan is the *path integral converts local observables into operators*, and this process is called *path-integral quantization*.

So far, everything has been commutative. But for more general observables, path-integral quantization can turn commutative operators into noncommutative ones!

So let's define a local observable depending on the 1-jet, i.e. a function $F: TY \rightarrow \mathbb{R}$. Now we let

$$(\mathcal{O}_F(t))(\phi) := F(\phi'(t)),$$

which defines a function $\mathcal{O}_F: \mathcal{C}_X \rightarrow \mathbb{R}$, and define its expectation as in (7.4), heuristically as a path integral and rigorously as a limit over discretizations. Path integral quantization will turn \mathcal{O}_F into an operator $\hat{F}: L^2(Y) \rightarrow L^2(Y)$ in that, for example,

$$\langle \mathcal{O}_F(t) \rangle_{S^1(T)} = \text{tr}(e^{-HT} \cdot \hat{F}).$$

Exercise 7.7. Let $Y = \mathbb{R}$ with the standard metric and introduce coordinates (x, p) on TY ($x \in \mathbb{R}, p \in T_x\mathbb{R}$). We know the operator x will quantize to the operator \hat{x} , multiplication by x on $L^2(\mathbb{R})$,²⁶ and show that p quantizes to $2\partial_x: L^2(\mathbb{R}) \rightarrow L^2(\mathbb{R})$. You can also use the flat metric on S^1 instead of \mathbb{R} if you'd like, though you have to replace x by e^{ix} .

This is very weird: commutative things became noncommutative. Where did this come from?

We'll compute the commutator, which is somehow the most fundamental object to come out of this question. Let $Y = \mathbb{R}$ or S^1 . Naïvely,

$$(7.8a) \quad \text{Kernel}(e^{-t_3 H} \hat{p} e^{-t_2 H} \hat{x} e^{-t_1 H}) = \int_{\mathcal{C}} d\phi \phi'(t_1 + t_2) \phi(t_1) e^{-S(\phi)}$$

$$(7.8b) \quad \text{Kernel}(e^{-t_3 H} \hat{x} e^{-t_2 H} \hat{p} e^{-t_1 H}) = \int_{\mathcal{C}} d\phi \phi(t_1 + t_2) \phi'(t_1) e^{-S(\phi)}.$$

With the path integral defined by discretization, these are actually both literally true. Moreover, as $t_2 \rightarrow 0$, they look equal. But what happens when we discretize? Let $y_1 := t_1 - \Delta t$, $y_2 = t_1$, and $y_3 = t_1 + \Delta t$. We know $\phi(t_1) = y_2$, but what about the derivative? We have the two choices

$$\frac{1}{\Delta t}(y_3 - y_2) \quad \text{or} \quad \frac{1}{\Delta t}(y_2 - y_1).$$

If you take the continuum limits of (7.8a) and (7.8b), you'll wind up with terms like these.²⁷ So in the path integral, there are again two possibilities:

$$(7.9a) \quad \frac{1}{\Delta t} \int dy_2 y_2 (y_3 - y_2) e^{-F}$$

$$(7.9b) \quad \frac{1}{\Delta t} \int dy_2 y_2 (y_2 - y_1) e^{-F},$$

where

$$F = \frac{(y_1 - y_2)^2 + (y_2 - y_3)^2}{(\Delta t)^2}.$$

You can compare (7.9a) and (7.9b) directly by integration by parts:

$$\begin{aligned} \frac{d}{dy_2}(y_2 e^{-F}) &= e^{-F} - y_2 \frac{dF}{dy_2} e^{-F} \\ &= e^{-F} - \frac{y_2(-2(y_1 - y_2) + 2(y_2 - y_3))}{(\Delta t)^2} e^{-F}, \end{aligned}$$

so the difference between (7.9a) and (7.9b) is

$$2 \int dy_2 e^{-F},$$

which was the value of the path integral with the operator 2 inserted. Therefore we conclude

$$\lim_{t_2 \rightarrow 0} \langle \cdots \hat{p} e^{-t_2 H} \hat{x} \cdots \rangle - \langle \cdots \hat{x} e^{-t_2 H} \hat{p} \cdots \rangle = 2 \langle \cdots 1 \cdots \rangle,$$

²⁶This is not technically an operator $L^2(\mathbb{R}) \rightarrow L^2(\mathbb{R})$, but there are ways to work around this: as soon as you apply e^{-HT} it does make sense. You can also work with a compactly supported version if you'd like.

²⁷You can also average, in which case you'll average (7.8a) and (7.8b).

or

$$[\hat{p}, \hat{x}] = 2.$$

This is how we got noncommutativity: it's important to be careful when you're doing two things at the same point.

What's next? We've seen that the one-dimensional QFT of maps into Y has to do with heat flow on Y , e.g.

$$Z_{S^1(T)} = \text{tr}(e^{-TH}).$$

This knows all of the eigenvalues of the Laplacian on Y , and is in particular very far from being topological. Next time, we'll cure this just as we did in dimension 0: we'll add fermions to make a supersymmetric quantum field theory. For example, instead of considering maps $S^1 \rightarrow Y$, we'll consider maps from the supermanifold ΠTS^1 into Y (so a bosonic part that looks like the circle and additional fermionic directions). This again has to do with heat flow, but for spinors, and the answer will be a super-trace of e^{-TH} , and this is the index of a Dirac operator — all but the zero eigenspace cancel out, and we'll obtain something topological.

Lecture 8.

The harmonic oscillator and some partition functions: 9/26/17

“So far, we don't look directly at the sun.”

We've been talking about one-dimensional QFT, with data a Riemannian manifold (Y, g) and a potential $V : Y \rightarrow \mathbb{R}$. The space of fields is $\mathcal{C}_X = \text{Map}(X, Y)$, where X is an interval, though we only approach it through its discretization.

We're going to modify the action slightly, in order to match the standard quantum-mechanical convention for the Hamiltonian in quantum mechanics,

$$H = -\frac{1}{2}\Delta + V.$$

Consequently the action should be

$$(8.1) \quad S(\phi) = \int_X \left(\frac{1}{2}g(\dot{\phi}, \dot{\phi}) + V(\phi) - \frac{1}{3}R(\phi) \right) dx.$$

This changes the conventions with the heat kernel, but nothing substantive is different.

We also saw that the interval of length T is associated with the operator $e^{-HT} : L^2(Y) \rightarrow L^2(Y)$, which sends

$$\psi(x) \mapsto \tilde{\psi}(x) := \int_Y k_T(x, y) \psi(y) d\text{vol}_Y,$$

where k_T is the heat kernel. The reason it has the name e^{-HT} is that if you differentiate it, you get the heat equation:

$$\underbrace{\frac{\partial}{\partial T}(e^{-HT}\psi)}_{\Psi(T,x)} = -He^{-HT}\psi,$$

and $\psi(T, x)$ satisfies the heat equation, somewhat tautologically. In functional analysis, this is sometimes called *functional calculus*.

In a similar way, the circle with circumference T is associated with the operator $\text{tr}_{L^2(Y)} e^{-HT}$.

Example 8.2 (Harmonic oscillator). The harmonic oscillator is familiar to students of quantum mechanics. In this model we let $Y = \mathbb{R}$ and

$$V(x) = \frac{1}{2}\omega^2 x^2,$$

for some $\omega \in \mathbb{R}$. The idea is that a particle can move around on a line, but is constrained by the potential to stay close to the origin, in that paths far from the origin aren't weighted very heavily.

The Hamiltonian is the operator

$$(8.3) \quad H = -\frac{1}{2} \frac{\partial^2}{\partial x^2} + \frac{1}{2} \omega^2 x^2,$$

which isn't *a priori* an operator $L^2(\mathbb{R}) \rightarrow L^2(\mathbb{R})$, but it's defined densely enough that the theory still makes sense.

The eigenvalues are $\{1/2 + n \mid n = 0, 1, 2, \dots\}$, and the eigenfunctions look like

$$\begin{aligned}\psi_0 &= e^{-\omega x^2/2} \\ \psi_1 &= x e^{-\omega x^2/2} \\ \psi_2 &= \left(x^2 - \frac{1}{2\omega}\right) e^{-\omega x^2/2},\end{aligned}$$

and in general, these are built out of *Hermite polynomials* H_n : up to some fixed constant $C(\omega)$,

$$\psi_n(x) = C(\omega) H_n(x\sqrt{\omega}) e^{-\omega x^2/2}.$$

The partition function on S^1 is nice:

$$\begin{aligned}Z_{S^1(T)} &= \text{tr}(e^{-TH}) \\ &= \sum_{n=0}^{\infty} \exp\left(-\omega\left(n + \frac{1}{2}\right)T\right),\end{aligned}$$

which is a geometric series, so its sum is

$$= \frac{1}{2 \sinh(\omega T/2)}.$$

Example 8.4. There's another example in which one can obtain things pretty explicitly is where Y is compact, and a simple choice is S^1 . When $V = 0$, this is called a σ -model on S^1 , or the particle on a circle.

So let $Y = S^1(R)$ (i.e. the circle of circumference R , or $\mathbb{R}/R\mathbb{Z}$) and $V = 0$, so

$$H = -\frac{1}{2} \frac{\partial^2}{\partial x^2}$$

acting on $\mathcal{H} = L^2(S^1(R))$. The eigenfunctions are a natural Fourier basis for $L^2(S^1)$:

$$\begin{aligned}\psi_0(x) &= 1 \\ \psi_{2n-1}(x) &= \sin\left(\frac{2\pi n x}{R}\right) \\ \psi_{2n}(x) &= \cos\left(\frac{2\pi n x}{R}\right).\end{aligned}$$

You can check these are periodic with period R . The eigenvalues scale with n^2 (unlike in Example 8.2, where they're evenly spaced): we have 0 with multiplicity 1, and $2\pi^2 n^2/R^2$ with multiplicity 2 for $n \geq 1$ (coming from ψ_{2n-1} and ψ_{2n}).

Thus the partition function on $S^1(T)$ is

$$\begin{aligned}Z_{S^1(T)} &= \text{tr}(e^{-TH}) \\ &= 1 + 2 \sum_{n=1}^{\infty} \exp\left(-\frac{2\pi^2 n^2 T}{R^2}\right) \\ &= \sum_{n=-\infty}^{\infty} \exp\left(-\frac{2\pi^2 n^2 T}{R^2}\right) \\ &= \vartheta\left(\tau = \frac{2\pi i T}{R^2}, z = 0\right).\end{aligned}$$

This $\vartheta(\tau, z)$ is called the *Jacobi ϑ -function*, where $\text{Im}(\tau) > 0$ and $z \in \mathbb{C}$ (regarded as a point on an elliptic curve). This is a little bit of a mystery: why does this function appear in one of the simplest quantum-mechanical model? Is it possible to get the ϑ -function for nonzero z ?

To answer that question, we'll use another nice thing about the system: a symmetry. There's an action of $U(1)$ on this theory: for an $\alpha \in U(1)$, we have an operator $S_\alpha: \mathcal{H} \rightarrow \mathcal{H}$ such that

$$(S_\alpha \psi)(x) := \psi(x + \alpha).$$

That this is a symmetry of the system means it commutes with the Hamiltonian:

$$[S_\alpha, H] = 0.$$

Remark 8.5. For those who are interested in bridges between topology and quantum field theory, it's worth mentioning that this QFT has a “global $U(1)$ -symmetry.” One thing you can do to these kinds of theories is formulate them on Riemannian manifolds X (i.e. the spacetime) equipped with a principal $U(1)$ -bundle with connection. This is a general principle, but in this case means to take $X = S^1$ with a flat principal $U(1)$ -bundle, which is characterized by its holonomy $\alpha \in U(1)$. In this case, instead of integrating over loops, you're integrating over “twisted loops” that don't completely close up, but have a “twisted boundary condition”

$$\phi(x + T) = \phi(x) + \alpha.$$

This intertwines the $U(1)$ -action on the principal bundle and on the target. \blacktriangleleft

Exercise 8.6. Compute $\text{tr}_{\mathcal{H}} e^{-T^H S_\alpha}$, and show that it's $\vartheta(\tau, z)$ for some nonzero z depending on α .

These days, people such as Dan Freed have been thinking of this perspective, especially in higher dimensions, as spreading the theory out over the moduli space of principal $U(1)$ -bundles. This is not the historical way to think about symmetries, but is interesting and fruitful. \blacktriangleleft

These examples were very nice, in that the Hamiltonian only has a point spectrum. Examples where the Hamiltonian has a continuous spectrum exist and are physically relevant, e.g. scattering phenomena. Often these also have discrete spectra.

Determinants. Recall that if V is a finite-dimensional vector space, $M : V \otimes V \rightarrow \mathbb{R}$ is a positive-definite bilinear form, $d\mu$ is a translation-invariant measure on V , and $c \in \mathbb{R}$, then we had a formula

$$(8.7) \quad (2\pi c)^{-\dim V/2} \int_V d\mu e^{-M(x,x)/2} = \frac{d\mu}{\sqrt{\det cM}}.$$

Something tricky is going on: $\sqrt{\det cM}$ is a density, hence defines a measure, so the right-hand side is a ratio of two measures, hence a number! This is because of how it transforms under change-of-coordinates: if $M \mapsto A^T M A$, then

$$\det M \mapsto (\det A)^2 \det M,$$

so $\sqrt{(\det M)} \mapsto |\det A| \sqrt{\det M}$, which is why it's a density.

Remark 8.8. Recall that the space of densities on V is $\Lambda^{\text{top}}(V^*) \otimes_{\mathbb{Z}/2} u(V)$, where $u(V)$ is the orientation bundle. That is, a density is a pair (ω, u) , where u is an orientation of V and ω is a volume form, such that $(\omega, u) \sim (-\omega, -u)$. And this is fine because if you use u and ω to integrate a function, using $-u$ and $-\omega$ gives you the same answer. Densities form a one-dimensional vector space, and unlike for volume forms, there is a canonical notion of a positive density. \blacktriangleleft

We're going to try to understand this for V infinite-dimensional: because the action (8.1) is quadratic in ϕ , our discretized path integral looks like the left-hand side of (8.7), where $c = \Delta t = T/N$. When we let $N \rightarrow \infty$ (so $\dim V \rightarrow \infty$), the left-hand side exists, but the right-hand side doesn't make sense, since we can't choose nontrivial translation-invariant measures on an infinite-dimensional vector space.

In the finite-dimensional case, we can avoid talking about the measure by choosing an inner product on V such that $\|d\mu\| = 1$ (that is, a density specified by a top form with norm 1). This lets us identify $V = V^*$. This simplifies the linear algebra a bit: we can use this identification to replace $M\delta t : V \otimes V \rightarrow \mathbb{R}$ with an $A : V \rightarrow V$, and

$$\frac{d\mu}{\sqrt{\det(M\Delta t)}} = \frac{1}{\sqrt{\det A}},$$

where this time, it's the ordinary determinant of a matrix.

Example 8.9. For the harmonic oscillator, the matrix is

$$A = \begin{pmatrix} 2 + \omega^2 T^2 / N^2 & -1 & 0 & \cdots & -1 \\ -1 & 2 + \omega^2 T^2 / N^2 & -1 & & \\ 0 & -1 & \ddots & & \\ \vdots & & & \ddots & \\ -1 & & & & \ddots \end{pmatrix}.$$

If A is an $N \times N$ matrix, then pleasantly,

$$\lim_{N \rightarrow \infty} \frac{1}{\sqrt{\det A}} = \frac{1}{2 \sinh(\omega T/2)},$$

which is what we got for the partition function in Example 8.2 in a totally different way! But it's less obvious how to generalize A itself to infinite dimensions. This was worked out in a recent paper of Ludewig [20].

In the infinite-dimensional case, we'll choose the space V of functions $\phi : S^1 \rightarrow \mathbb{R}$ such that a certain norm is finite. That is, using the eigenbasis we found above, we can write

$$(8.10) \quad \phi(t) = c\sqrt{T} + \sum_{n=1}^{\infty} \frac{\sqrt{2T}}{2\pi n} \left(a_n \sin\left(\frac{2\pi n}{T}t\right) + b_n \cos\left(\frac{2\pi n}{T}t\right) \right).$$

Then, we let the norm be

$$\|\phi\|^2 := c^2 + \sum_{n=1}^{\infty} a_n^2 + b_n^2.$$

This is not the usual L^2 -norm, and looks more like a Sobolev norm (and in fact is equivalent to one). Anyways, we take V to be the space of functions for which this is finite.

Using (8.10), the action

$$S = \frac{1}{2} \int dt \left(g(\dot{\phi}(t), \dot{\phi}(t)) + \omega^2 \phi(t)^2 \right)$$

becomes

$$S(\phi) = \frac{1}{2} \left(\omega^2 T^2 c^2 + \sum_{n=1}^{\infty} \left(1 + \frac{\omega^2 T^2}{4\pi^2 n^2} \right) (a_n^2 + b_n^2) \right).$$

The eigenvalues of this operator are

$$\lambda = \omega^2 T^2, 1 + \frac{\omega^2 T^2}{4\pi^2 n^2},$$

where the latter has multiplicity 2 for each $n > 0$, and therefore one can show that

$$\sqrt{\det A} = 2 \sinh\left(\frac{1}{2} \omega T\right),$$

as desired. Usually in physics this is heuristically done with some sort of ζ -regularization, but in the one-dimensional case everything can be made rigorous! \blacktriangleleft

Lecture 9.

Symmetries and effective field theory in 1D QFT: 9/28/17

In the last few lectures, we've been learning about one-dimensional QFT, though in a specific class of examples: Lagrangian quantum field theories (so the partition function is an integral) and specifically, a σ -model with action

$$S = \int_X dt \left(\frac{1}{2} g(\dot{\phi}, \dot{\phi}) + V(\phi) - \frac{1}{6} R \right),$$

and we showed that the partition function is about heat flow (on an interval of length T , it's heat flow for time T ,²⁸ and on a circle of circumference T , it's the trace of that operator). There are two ways to think about this (for concreteness, let $X = S^1(T)$):

- (1) an integral over loops in the target Y , or
- (2) as the trace of e^{-HT} in $L^2(Y)$,

and it's possible to rigorously show these are equal. These are generic in one-dimensional QFT: the Hilbert space might not always be $L^2(Y)$, but the fact that we recover traces of operators on S^1 is a recurring theme.

In Example 8.2 (where $Y = \mathbb{R}$ and $V = \omega^2 x^2/2$), these interpretations turn into an infinite-dimensional determinant of $g(\dot{\phi}, \dot{\phi})$ (coming from what is, more or less, an infinite-dimensional Gaussian integral) for (1) and a sum

$$\sum_n e^{-T\lambda_n},$$

²⁸The thing that allows you to get the whole operator on the interval is that it has a boundary, and so you're freely able to choose boundary conditions and therefore understand what the heat kernel does to them.

where λ_n is the n^{th} eigenvalue of H for (2).

These two perspectives have established names: (1) is called the *Lagrangian formulation*, and (2) is called the *Hamiltonian formulation*. They're supposed to be formally equivalent, though showing this is difficult.

Remark 9.1. These two perspectives also exist in classical mechanics, and can be recovered from these by taking a classical limit. Classically, one restricts to the extrema of the action S , and there the proof that the Lagrangian and Hamiltonian formulations are equivalent is easier. \blacktriangleleft

Exercise 9.2. Figure out these two interpretations for the other example we considered, Example 8.4, where the Hamiltonian interpretation produces

$$Z_{S^1(T)} = \vartheta\left(\tau = \frac{2\pi iT}{R^2}, z = 0\right).$$

Show that for the Lagrangian formulation (discretize the path integral), you get

$$Z_{S^1(T)} = \vartheta\left(\tau = \frac{iR^2}{2\pi^2 T}, z = 0\right).$$

These are indeed equal, thanks to the modularity of the ϑ -function, or the Poisson summation formula. So in this case you again recover something mathematically interesting.

We can also add symmetry to the picture. Recall that in the zero-dimensional case, we found a Lie algebra action of vector fields on \mathcal{C} which annihilate S (or, exponentiated, a Lie group action of G on \mathcal{C} preserving S). This produced constraints on the correlation functions: $\langle \mathcal{O} \rangle = \langle \mathcal{O}^g \rangle$ and $\langle X \mathcal{O} \rangle = 0$ if $X \in \mathfrak{g}$.

In the model we've been considering, we can choose two kinds of symmetries:

- isometries of X , or
- isometries of Y that preserve V .

The first exists for any choice of parameters (though there's not much to say for an interval), but the second might be the trivial group for some choices of Y and V .

For $X = S^1(T)$, the isometry group is $U(1)$, acting by $t \mapsto t + c$. This produces a constraint on the correlation functions:

$$\langle \mathcal{O}_1(t_1) \mathcal{O}_2(t_2) \cdots \mathcal{O}_n(t_n) \rangle = \langle \mathcal{O}_1(t_1 + c) \mathcal{O}_2(t_2 + c) \cdots \mathcal{O}_n(t_n + c) \rangle.$$

Exercise 9.3. Show this from the Hamiltonian perspective (it comes from the cyclicity of the trace).

Any isometry of Y preserving V produces a similar formula.

We'll use this as an opportunity to introduce some useful notation describing how symmetries act on the fields.²⁹ For an action of $U(1)$, which is connected, it suffices to use the Lie algebra $\mathfrak{u}_1 \cong \mathbb{R}$, so we'll describe the action of the shift $t \mapsto t + \varepsilon$. To first-order in ε , this is

$$\phi(t + \varepsilon) = \phi(t) + \varepsilon \dot{\phi}(t).$$

We represent this by writing

$$(9.4) \quad \delta \phi = \varepsilon \dot{\phi}.$$

It feels like this should be more complicated, but this notation encapsulates the fact that this symmetry is completely local, and therefore actually very simple.

Effective field theory in 1 dimension. Consider a system where $Y = \mathbb{R}^2$ with coordinates (x, y) and potential

$$V = \frac{1}{2}x^2 + \frac{1}{2}\omega^2 y^2 + \frac{\mu}{4}x^2 y^2.$$

Quantum-mechanically, this is telling us about a system with two kinds of particles with a slight coupling between them.

Suppose $\omega \gg 1$, which means that y is oscillating very rapidly. Therefore, if you care about x , you should be able to eliminate y by integrating out over all of those oscillations and replacing them with their average values.³⁰

²⁹This notation will not work for discrete symmetries; for example, $t \mapsto -t$ on the interval or the circle is an isometry preserving the action, called *time-reversal symmetry*. This is a $\mathbb{Z}/2$ -symmetry having to do with orientation-reversal, and many theories do not have it. Nonetheless, we're not going to worry about discrete symmetries for now, and for supersymmetry, which we're going to switch to soon, this notation will be extremely convenient.

³⁰In atomic physics, this is called the *Born-Oppenheimer approximation*.

Said plainly, this is somewhat crazy: we're going to integrate over one infinite-dimensional space to end up in another. But this will be very helpful.

Suppose we're interested in normalized expectation values involving only x , such as

$$\frac{\langle x(0)x(t) \rangle_{S^1(T)}}{Z_{S^1(T)}} = \frac{1}{Z_{S^1(T)}} \int_{\mathcal{C}_{S^1(T)}} dx dy x(0)x(t) e^{-S(x,y)},$$

where by pullback along the map $S^1(T) \rightarrow Y$, $x, y: S^1 \rightarrow \mathbb{R}$. This is fine because the infinite-dimensional determinants present for the expectation value and for $Z_{S^1(T)}$ cancel each other out. Since this is quadratic, it's possible to attach it explicitly with a change of variables, proceeding somewhat similarly to before, but there's a more general method.

Before we go to the effective field theory, we're going to remember how we used Feynman diagrams for zero-dimensional QFT to calculate asymptotic series. In this case we summed over four-valent vertices with labels i, j, k, ℓ , weighted by λ_{ijkl} , and edges i to j weighted by M_{ij}^{-1} (where M is the quadratic piece of the action).

The basic new ingredient in 1-dimensional QFT is the **Green's functions** D_x and D_y on S^1 , defined to (distributionally) satisfy

$$\begin{aligned} (\partial_t^2 - 1)D_x(t) &= \delta(t) \\ (\partial_t^2 - \omega^2)D_y(t) &= \delta(t). \end{aligned}$$

Explicitly,

$$\begin{aligned} D_x(t) &= \frac{1}{2} \sum_{n \in \mathbb{Z}} e^{-|t+nT|} \\ D_y(t) &= \frac{1}{2\omega} \sum_{n \in \mathbb{Z}} e^{-\omega|t+nT|}. \end{aligned}$$

The Feynman diagram expansion is a little more complicated.

- At zeroth-order, we have a segment with endpoints labeled $0, t$, which contributes a factor $D_x(t)$.
- At first order, we put a dashed loop (for y) at $t' \in (0, t)$, for every value of t' , and attach a Green's function for y there. This involves an integral (thankfully, over a one-dimensional space):

$$\frac{\mu}{2} \int_{S^1(T)} dt' D_x(t') D_x(t-t') D_y(0).$$

- In higher dimensions you'll have more diagrams and more integrals over finite-dimensional spaces.

This is known as *perturbation theory* in quantum mechanics. Each of these things gives you a number, and so while this is complicated, you can calculate it in principle.

The effective field theory description will be easier. We decide to rid ourselves of these dashed lines (corresponding to terms in y). Formally,

$$S_{\text{eff}}(x) = -\log \int dy e^{-S(x,y)}.$$

In this case, at second order we get a Feynman diagram that looks like Figure 2, with edges labeled by t and t (one side) and t' and t' (the other side). We still have integrals, but they're easier.

- The terms that are second-order in x are

$$\int dt \frac{1}{2} \dot{x}(t)^2 + \frac{1}{2} x(t)^2 + \frac{\mu}{2} \int dt x(t)^2 D_y(0),$$

where the second term comes from a dashed loop. We can rewrite this as

$$\int dt \frac{1}{2} \dot{x}(t)^2 \left(\frac{1}{2} + \frac{\mu}{2} D_y(0) \right) x(t)^2.$$

- The term that's fourth-order in x (from the diagram like in Figure 2) is

$$\frac{\mu^2}{2} \int dt dt' x(t)^2 x(t')^2 D_y(t-t')^2.$$

One particularly weird consequence of the fourth-order term (and appearing more strongly in higher-order terms) is the presence of nonlocal phenomena, coming from $dt dt'$. The idea is that x -particles may be connected by y -fields, coupled by $D_y(t - t')$. So

$$S_{\text{eff}} \neq \int dt \dot{x}(t)^2 + V(x(t))$$

and therefore the effective theory is more complicated.

But not all hope is lost: $D_y(t - t')$ decays exponentially away from $t - t' = 0$. So we can expand the non-local interaction in powers of $t - t'$:

$$\begin{aligned} & \int dt dt' x(t)^2 x(t')^2 D_t(t - t')^2 \\ &= \int dt dt' \left(x(t)^4 + 2x(t)^3 \dot{x}(t)(t - t') + \left(x(t)^2 \dot{x}(t)^2 + \frac{1}{2} x(t)^3 \ddot{x}(t) \right) (t - t')^2 + O((t - t')^3) \right) D_y(t - t')^2. \end{aligned}$$

The idea is, we can replace the nonlocal term with something local, as long as we're willing to take derivatives of the fields.

Without evaluating in detail, we can learn something about the shape of the answer by integrating over t' . The action you get is

$$(9.5) \quad S = \int dt c_1 x(t)^4 + \frac{c_2}{\omega} \left(x^2 \dot{x}^2 + \frac{1}{2} x^3 \ddot{x} \right) + \dots$$

That is, you have an infinite-series of “higher-derivative interactions” suppressed by larger and larger negative powers of ω . Some correlation functions will be dominated by paths where only the first few terms are large enough to really contribute (which is what we'll mean by a low-energy limit), and in this case we have a systematic expansion in terms of powers of ω^{-1} , and the theory looks local after all!

This is interesting: when you integrate out a field, there's no reason to expect that the theory you get is local. Even in this case, where the heuristics suggested a low-energy limit was possible, we had to make some estimates to recover locality, some of which were physically rather than mathematically justified.³¹

Introducing supersymmetry. We'll now begin studying supersymmetry quantum mechanics. There was nothing topological about the model we've been studying: indeed, it computed the trace of the heat kernel. To make it topological, we'll do the same thing that we did in zero dimensions: making a new quantum field theory whose configuration space is

$$\mathcal{C} = \Pi(T \text{Map}(X, Y)),$$

the parity change of the tangent space of the infinite-dimensional space of maps from X to Y , an infinite-dimensional supermanifold.

Let's say what this actually means. A point of $\text{Map}(X, Y)$ is a map $\phi : X \rightarrow Y$ (again, X is an interval with boundary conditions $[0, T]_{y_0}^{y_1}$ or a circle $S^1(T)$). Taking the tangent bundle (and parity change) means also specifying a fermionic term $\psi \in \Pi T(\phi^*TY)$: pull back the tangent bundle by ϕ , then take a section of it.

Remark 9.6. If M is oriented, ΠTM has a canonical measure. This reassures us that we're on relatively safe ground. ◀

Now let's write the action, which depends on these bosonic and fermionic directions (ϕ, ψ) . Let ∇_t be the pullback of the Levi-Civita connection to ϕ^*TY by ϕ ; then, the action is

$$(9.7) \quad S(\phi, \psi) = \int dt \frac{1}{2} (g(\dot{\phi}, \dot{\phi}) + g(\psi, \nabla_t \psi)).$$

These two terms are qualitatively different: the first is second-order in time derivatives, but the second is only first-order: it looks like $\int dt \psi \dot{\psi}$. This reflects another difference between bosons and fermions: using integrating by parts, for a boson ϕ ,

$$\int dt \phi \dot{\phi} = - \int dt \dot{\phi} \phi = - \int dt \phi \dot{\phi},$$

³¹Mathematical justifications should be possible, just require more thought.

so this term, which might look meaningful locally, is zero. But fermions anticommute and hence pick up an extra sign:

$$\int dt \psi \dot{\psi} = - \int dt \dot{\psi} \psi = \int dt \psi \dot{\psi},$$

So we could have inserted a bosonic term like this one, but it would vanish. Similarly, a term like $\int \dot{\psi} \dot{\psi}$ would have vanished as well.

The action S is invariant under time translation again, and we write

$$(9.8) \quad \begin{aligned} \delta \phi &= \varepsilon \dot{\phi} \\ \delta \psi &= \varepsilon \dot{\psi}. \end{aligned}$$

There's also an additional odd symmetry (i.e. an odd vector field on the supermanifold),

$$(9.9) \quad \begin{aligned} \delta \phi &= \varepsilon \psi \\ \delta \psi &= -\varepsilon \dot{\phi}. \end{aligned}$$

To understand what this means, you have to think of ε as having odd parity. More explicitly, one has an odd vector field

$$Q = \int dt \psi(t) \frac{\delta}{\delta \phi(t)} - \dot{\phi} \frac{\delta}{\delta \psi(t)},$$

and $QS = 0$.

Lecture 10.

Perturbation theory in quantum mechanics and spin structures: 10/3/17

The last computation we did might not be terribly easy to follow, so today we're going to start with something different, but in the same spirit, and that should be a little clearer.

Recall that if you're doing an integral over paths with a Gaussian action, you're going to get an infinite-dimensional determinant. For non-Gaussian actions, you can make an asymptotic expansion in the "non-Gaussianity."

Example 10.1 (Quartic oscillator). As usual, this will be a one-dimensional σ -model, whose fields are maps $X \rightarrow Y$. The target Y is \mathbb{R} with the usual metric, and the action is

$$(10.2) \quad V = \frac{1}{2} \omega^2 x^2 + \frac{\lambda}{4!} x^4.$$

The Hamiltonian is therefore

$$\begin{aligned} H &:= -\frac{1}{2} \frac{\partial^2}{\partial x^2} + V(x) \\ &= -\frac{1}{2} \frac{\partial^2}{\partial x^2} + \frac{1}{2} \omega^2 x^2 + \frac{\lambda}{4!} x^4. \end{aligned}$$

As usual in physics, we'd like to compute the eigenvalues of H acting on $L^2(\mathbb{R})$. In quantum mechanics, you solve this by directly looking at the Hamiltonian, but we're trying to use this as a toy example of a QFT, so we're going to get them out of the partition function

$$Z_{S^1(T)} = \sum_n e^{-E_n T},$$

where

$$H\psi_n = E_n \psi_n,$$

i.e. E_n is the n^{th} eigenvalue.

Let Z_0 denote the partition function for the $\lambda = 0$ theory; then, we can compute $Z_{S^1(T)}/Z_0$ as a sum over Feynman diagrams.

Recall that if W is a finite-dimensional state space and we have an action

$$S = \frac{1}{2} M(x, x) + \frac{1}{4!} C(x, x, x, x),$$

where $M : W^{\otimes 2} \rightarrow \mathbb{R}$ and $C : W^{\otimes 4} \rightarrow \mathbb{R}$, we summed over Feynman diagrams with 4-valent vertices (corresponding to $C \in (V^*)^{\otimes 4}$) and with edges weighted by M^{-1} .

Now, though, the state space is infinite-dimensional, the functions $x : S^1 \rightarrow \mathbb{R}$, and the second-order term is

$$M(x, x) = \int_0^T dt \frac{1}{2} \|\dot{x}(t)\|^2 + \frac{1}{2} \omega^2 x^2.$$

Its “inverse”³² is the Green’s function for the Laplace operator:

$$(10.3) \quad G(t, t') = G(t - t') = \frac{1}{2\omega} \sum_{n \in \mathbb{Z}} e^{-\omega|t-t'+nT|}.$$

This inverts M in the sense that

$$M(G(t, t'), f(t)) = f(t').$$

In finite dimensions, recall that

$$M(M^{-1}(\eta, \cdot), \nu) = \eta(\nu),$$

justifying our choice to call it an inverse.

Now, weighting by M^{-1} means weighting by an integral of a Green’s function.

- The empty diagram contributes a factor of 1.
- In first-order, we have a single vertex and the “figure-8 diagram.” All half-edges are labeled with time t , and the contribution is

$$-\frac{\lambda}{8} \int_0^T dt G(t, t) G(t, t).$$

- At second-order, two vertices have two different times t and t' associated with their half-edges. For example, the diagram with four edges between the vertices v and v' contributes³³

$$\frac{\lambda^2}{48} \int_0^T \int_0^T dt dt' G(t - t')^4.$$

Our explicit formula (10.3) for $G(t, t')$ means these can be concretely evaluated, and the answer is, to first order,

$$\log\left(\frac{Z}{Z_0}\right) \sim -\frac{\lambda T}{32\omega^2} \left(\coth \frac{\omega T}{2}\right)^2 + O(\lambda^2).$$

We want to use this to calculate eigenvalues, or at least the zeroth eigenvalue. We know

$$\log Z = \log\left(\sum_n e^{-TE_n}\right),$$

and as $T \rightarrow \infty$,

$$\begin{aligned} \log Z &\sim -TE_0(\lambda) \\ \log Z_0 &\sim -TE_0(\lambda = 0). \end{aligned}$$

Therefore

$$\log\left(\frac{Z}{Z_0}\right) \sim -T(E_0(\lambda) - E_0),$$

and a similar method gives you the corrections for higher eigenvalues. Concretely,

$$\begin{aligned} E_0(\lambda) - E_0 &\sim \lambda \left(\lim_{T \rightarrow \infty} \frac{1}{32\omega^2} \left(\coth \frac{\omega T}{2}\right)^2 \right) \\ &= -\frac{\lambda}{32\omega^2}, \end{aligned}$$

and this is precisely the result you obtain by more conventional quantum-mechanical methods. ◀

³²In finite dimensions, this is an inverse matrix, but not in infinite dimensions.

³³In Minkowski signature, which is common for relativistic QFT, there are some additional complications due to boundary terms, etc.

Let's return to supersymmetric quantum mechanics. We want to write down a 1-dimensional QFT whose space of fields $\mathcal{C} = \Pi T \text{Map}(X, Y)$. Formally, one defines \mathcal{C} as Spec of some \mathbb{C} -algebra in supergeometry, so it's difficult to speak precisely of its points, but they should be maps $\phi : X \rightarrow Y$ (bosonic) and $\psi \in \Pi \Gamma(\phi^* TY)$ (fermionic).

The action $S \in C^\infty(\mathcal{C})$ is given in (9.5). There's a time-translation symmetry (9.8); concretely, this means that in local coordinates on Y ,

$$\delta \phi^I = \varepsilon \dot{\phi}^I$$

$$\delta \psi^I = \varepsilon \dot{\psi}^I.$$

Here δ is the shift or translation. This symmetry was present in ordinary quantum mechanics.

But in the supersymmetric case, there's an additional odd symmetry (9.9), which defines a derivation $Q \in \text{Vect}^1(\mathcal{C}_X)$. This will be the engine that makes supersymmetry behave so differently.

Proposition 10.4. $QS = 0$.

Proof (when $Y = \mathbb{R}^n$). We just have to compute (though there are a few steps left implicit): in local coordinates,

$$S = \frac{1}{2}(\dot{\phi}^I \dot{\phi}^I + \psi^I \dot{\psi}^I),$$

and therefore

$$\begin{aligned} \delta S &= \frac{1}{2}(2\varepsilon \dot{\psi}^I \dot{\phi}^I + (-\varepsilon \dot{\phi}^I \dot{\psi}^I - \psi^I \ddot{\phi}^I)) \\ &= 0, \end{aligned}$$

since we can commute $\dot{\psi}^I$ and $\dot{\phi}^I$. \(\square\)

Exercise 10.5. Show that Q is “a square-root of time-translations,” in that

$$\frac{1}{2}[Q, Q] = H.$$

We'd like to obtain something topological out of this, but there's a metric around. It turns out the variation is Q -exact, so the partition function is invariant.³⁴

Exercise 10.6. Show that if we vary the metric g on Y under the variation

$$g_{IJ} \mapsto g_{IJ} + \delta g_{IJ},$$

then the action varies by

$$S \mapsto S + Q\Psi,$$

where

$$\Psi := \int dt \left(\frac{1}{2}(\delta g)_{IJ} \dot{\phi}^I \dot{\phi}^J \right).$$

Thus we expect $Z_{S^1(T)}$ to be independent of the metric on Y , and indeed this is true, but it *does* depend on a spin structure on Y !

Spin structures and spinors. We now need to discuss spin structures and spinors. For this section, Morgan [23] is a good reference.

Definition 10.7. Let V be a finite-dimensional vector space. Its *tensor algebra* is the free (noncommutative) algebra on V ; explicitly, this is the \mathbb{Z} -graded algebra

$$T(V) := \bigoplus_{n \geq 0} V^{\otimes n},$$

where $V^{\otimes 0} = \mathbb{C}$, and the multiplication rule is

$$(v_1 \otimes \cdots \otimes v_k) \times (v'_1 \otimes \cdots \otimes v'_{k'}) = v_1 \otimes \cdots \otimes v_k \otimes v'_1 \otimes \cdots \otimes v'_{k'}.$$

(Often, this multiplication rule is also denoted \otimes .)

This is the progenitor of all sorts of useful algebraic structures, such as symmetric algebras, exterior algebras, and the one we need, Clifford algebras.

³⁴There are some infinite-dimensional subtleties, but it can be shown rigorously.

Definition 10.8. Let V be a finite-dimensional real (or complex) vector space together with a symmetric, positive-definite quadratic form $\langle \cdot, \cdot \rangle$. Let \mathcal{I} denote the two-sided ideal of $T(V)$ generated by $\{v \otimes v - \langle v, v \rangle\} \mid v \in V\}$. Then, the *Clifford algebra* of V is

$$\text{Cliff}(V) := T(V)/\mathcal{I}.$$

The idea is to impose the relation $v^2 = -\langle v, v \rangle$ with minimal other choices. This does not respect the \mathbb{Z} -grading on $T(V)$, but it does respect it mod 2, so the Clifford algebra is $\mathbb{Z}/2$ -graded.

In the Clifford algebra, we have a relation

$$\frac{1}{2}(vw + wv) = -\langle v, w \rangle,$$

and therefore if e_1, \dots, e_n is a basis of V , then $\text{Cliff}(V)$ has a basis consisting of $1, e_1, \dots, e_n, e_i e_j$ for $i < j$, $e_i e_j e_k$ for $i < j < k$, and so on. Thus it's 2^n -dimensional.

Remark 10.9. The exterior algebra $\Lambda^*(V)$ is also 2^n -dimensional, leading one to suspect it's related to the Clifford algebra. Indeed, the Clifford algebra can be interpreted as a deformation of $\Lambda^*(V)$. ◀

Definition 10.10. The *pin group* $\text{Pin}(V)$ is the group of all elements $v_1 v_2 \dots v_m \in \text{Cliff}(V)$ for which each $v_i \in V$ and $\langle v_i, v_i \rangle = 1$.³⁵ The *spin group* $\text{Spin}(V)$ is $\text{Pin}(V) \cap \text{Cliff}^0(V)$.

The pin and spin groups are Lie groups, and in fact compact. There's a canonical action of $\text{Spin}(V)$ on V given by

$$g \cdot v = g v g^{-1} \in V \subset \text{Cliff}(V),$$

where we interpret $g v g^{-1}$ as multiplication in $\text{Cliff}(V)$. This action preserves the metric on V , so we obtain a map $\text{Spin}(V) \rightarrow \text{SO}(V)$. This map is in fact a double cover, and if $\dim V \geq 3$, this is the universal cover of $\text{SO}(V)$.

Notationally, we will let $\text{Cliff}(n)$ denote the Clifford algebra of \mathbb{R}^n with the usual inner product. It is generated by e_1, \dots, e_n , with a relation

$$\frac{1}{2}[e_i, e_j] = -\delta_{ij}.$$

Example 10.11. In low dimensions, these are familiar objects.

- $\text{Cliff}(1) \cong \mathbb{C}$, $\text{Cliff}^0(1) \cong \mathbb{R}$, and $\text{Spin}(1) \cong \mathbb{Z}/2$.
- $\text{Cliff}(2) \cong \mathbb{H}$, $\text{Cliff}^0(2) \cong \mathbb{C}$, and $\text{Spin}(2) \cong \text{U}(1)$.
- $\text{Cliff}(3) \cong \mathbb{H} \oplus \mathbb{H}$, $\text{Cliff}^0(3) \cong \mathbb{H}$, and $\text{Spin}(3) \cong \text{SU}(2)$.

In higher dimensions, though, they're harder to explicitly identify with familiar objects. ◀

We'll need some topology of the spin group and some representation theory.

Definition 10.12. Fix an oriented Riemannian manifold Y . The *bundle of orthonormal frames*³⁶ is the principal $\text{SO}(n)$ -bundle $P \rightarrow Y$ whose fiber at $p \in Y$ is the space of oriented, orthonormal bases of $T_p Y$. A *spin structure* is a lift of P to a principal $\text{Spin}(n)$ -bundle (across the covering map $\text{Spin}(n) \rightarrow \text{SO}(n)$).

Not every oriented Riemannian manifold admits a spin structure, and there may be multiple spin structures (isomorphism classes of lifts of P), but we'll say more about that later.

Example 10.13. Let $Y = S^1$ with the standard metric and orientation. Then, $\text{SO}(1)$ is trivial, so $P = Y$ again. Thus a spin structure is a double cover of Y ; there are two of these up to isomorphism, hence two spin structures on S^1 . ◀

Lecture 11.

Clifford algebras and spin structures: 10/5/17

Last time, we defined and discussed Clifford algebras: given a real, finite-dimensional vector space V with a positive-definite inner product, one can construct its Clifford algebra $\text{Cliff}(V)$, a $\mathbb{Z}/2$ -graded associative algebra. Inside this, we constructed a Lie group called $\text{Pin}(V)$, and the intersection of $\text{Pin}(V)$ and $\text{Cliff}^0(V)$ we called the spin group $\text{Spin}(V)$. There's a double cover $\text{Spin}(V) \rightarrow \text{SO}(V)$.

³⁵For $V = \mathbb{R}^n$ and the usual inner product, this group is usually denoted $\text{Pin}^+(n)$; if we started with a negative definite form, we'd obtain its sibling $\text{Pin}^-(n)$. For $V = \mathbb{C}^n$, there's only one kind, $\text{Pin}^c(n)$.

³⁶This definition sounds much scarier than it actually is!

Remark 11.1.

$$\pi_1(\mathrm{SO}(n)) = \begin{cases} 1, & n = 1 \\ \mathbb{Z}, & n = 2 \\ \mathbb{Z}/2, & n \geq 3. \end{cases}$$

It turns out that $\mathrm{Spin}(n)$ is connected for $n \geq 2$, and therefore is the universal cover for $\mathrm{SO}(n)$ when $n \geq 3$. \blacktriangleleft

For $V = \mathbb{R}^n$, we denoted these $\mathrm{Spin}(n)$, $\mathrm{Cliff}(n)$, etc.

If (M, g) is an oriented Riemannian manifold, it has a bundle of orthonormal frames $P \rightarrow M$, which is a principal $\mathrm{SO}(n)$ -bundle; we then defined a spin structure to be a lift of P to a principal $\mathrm{Spin}(n)$ -bundle \tilde{P} . Two spin structures \tilde{P} and \tilde{P}' are equivalent if there's an isomorphism of principal $\mathrm{Spin}(n)$ -bundles $\tilde{P} \cong \tilde{P}'$ that commutes with the projections back to P . In Example 10.13 we showed that the circle has two spin structures, which relates to its double covers. There's a more general fact.

Exercise 11.2. Let M be a manifold which admits a spin structure (sometimes called *spinnable*). Show that the set of spin structures on M up to equivalence is a torsor for $H^1(M; \mathbb{Z}/2)$.³⁷ Idea: given a spin structure Q and a double cover C , one can “twist Q by C ” to obtain another spin structure $Q \otimes_{\mathbb{Z}/2} C$, and the abelian group of isomorphism classes of double covers of M is canonically identified with $H^1(M; \mathbb{Z}/2)$.

Not all manifolds are spinnable; in general, this is a codimension 2 phenomenon.

Example 11.3. Let $M = \mathbb{CP}^2$, with the Fubini-Study metric and the usual complex orientation; this manifold does not admit a spin structure!³⁸

Let H denote a hyperplane in \mathbb{CP}^2 , i.e. an embedded \mathbb{CP}^1 . Then, $T\mathbb{CP}^2$, restricted to H , is $\mathcal{O}(1) \oplus \mathcal{O}(2)$: $\mathcal{O}(2)$ is $T\mathbb{CP}^1$ and $\mathcal{O}(1)$ is the normal bundle.³⁹

\mathbb{CP}^1 has two charts, so we can explicitly write down transition functions for $\mathcal{O}(1) \oplus \mathcal{O}(2)$, which are maps $S^1 \rightarrow \mathrm{U}(1) \times \mathrm{U}(1) \subset \mathrm{SO}(4)$. Explicitly, let $R_\theta \in \mathrm{SO}(2)$ denote the matrix which acts through rotation by θ ; then one of the transition functions is

$$(11.4) \quad \theta \mapsto \begin{pmatrix} R_\theta & \\ & R_{2\theta} \end{pmatrix}.$$

A spin structure is a lift of this map to $\mathrm{Spin}(4)$, the universal cover of $\mathrm{SO}(4)$. But the loop defined by (11.4) is the nontrivial element of $\pi_1(\mathrm{SO}(4))$, and therefore this map cannot lift to $\mathrm{Spin}(4)$. \blacktriangleleft

So we see that not every oriented 4-manifold is spinnable. This is minimal.

- Every compact, oriented 2-manifold Σ is spinnable. This ultimately is because $\chi(\Sigma)$ is always even.
- Every compact, oriented 3-manifold M is spinnable, and this is for a more striking reason: the tangent bundle of M is trivial!
- For 4-manifolds, we got stuck on a codimension-2 submanifold with odd intersection number. It turns out that a compact, oriented 4-manifold is spinnable iff the self-intersection number of all embedded surfaces is even.

We're stressing the embedded-loops perspective, rather than a more abstract one, because it is the way spin phenomena will appear in this course.

Proposition 11.5. *There's an irreducible, complex, $\mathbb{Z}/2$ -graded representation $S = S^0 \oplus S^1$ of $\mathrm{Cliff}(2n)$, with $\dim S^0 = \dim S^1 = 2^{n-1}$. Up to isomorphism and shifting the grading, this is the unique irreducible representation of $\mathrm{Cliff}(2n)$.*

This in particular means there's an action of $\mathrm{Spin}(2n)$ on S^0 and S^1 . For example, when $n = 1$, one can explicitly write down the matrices

$$e_1 \mapsto \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}, \quad e_2 \mapsto \begin{pmatrix} 0 & i \\ i & 0 \end{pmatrix},$$

and check that they satisfy the Clifford relations $e_i^2 = -1$ and $e_1 e_2 = -e_2 e_1$. Under the explicit identification $\mathrm{Cliff}(2) \cong \mathbb{H}$, this is the defining representation of \mathbb{H} , i.e. acting on itself by left multiplication.

We'd like to bring this theory to vector bundles over a manifold.

³⁷This means that a choice of a spin structure defines an isomorphism of abelian groups from the set of spin structures to $H^1(M; \mathbb{Z}/2)$.

³⁸This fact does not depend on the choice of metric or orientation.

³⁹A generic section of the normal bundle intersects itself at one point, which is the reason why the normal bundle is $\mathcal{O}(1)$; a similar argument gets you $\mathcal{O}(2)$ for the tangent bundle.

Definition 11.6. Let M be a $2n$ -dimensional spin manifold, and let $Q \rightarrow M$ be its spin structure. Let

$$SM := Q \times_{\text{Spin}(2n)} S,$$

which is called the *spinor bundle*.

This bundle has some additional structure.

- The Levi-Civita connection induces a connection ∇ on SM .
- There's an action of the tangent bundle, i.e. a map of vector bundles $\rho: TM \rightarrow \text{End}(SM)$, induced from the action of $\text{Cliff}(n)$ on \mathbb{R}^n . This requires checking that this action is equivariant for the action of $\text{Spin}(n)$, which follows because we more or less defined the $\text{Spin}(n)$ -action on \mathbb{R}^n using the Clifford algebra action.

The key player in our story will be a canonical first-order differential operator. On the tangent bundle, there's no canonical first-order operator (though the Laplacian is a canonical second-order operator), so this is something new and cool.

Definition 11.7. The *Dirac operator* $\not{D}: C^\infty(SM) \rightarrow C^\infty(SM)$ is defined locally by the formula

$$\not{D} = \sum_{i=1}^{2n} \rho(e_i) \circ \nabla_{e_i},$$

where $\{e_1, \dots, e_n\}$ is a local orthonormal basis for TM .

This is an odd operator, interchanging S^0 and S^1 .

Example 11.8. Suppose $M = \mathbb{R}^2$, so there's a unique spin structure, and SM is the trivial bundle $\mathbb{C} \rightarrow \mathbb{R}$. Then, the Dirac operator is

$$\not{D} = e_1 \partial_1 + e_2 \partial_2 = \begin{pmatrix} 0 & \partial_1 + i \partial_2 \\ -\partial_1 + i \partial_2 & 0 \end{pmatrix}.$$

If you square it, you get

$$\not{D}^2 = \begin{pmatrix} -\partial_1^2 - \partial_2^2 & 0 \\ 0 & -\partial_1^2 - \partial_2^2 \end{pmatrix},$$

which looks a lot like the Laplacian, but acts on the spinor bundle. ◀

Definition 11.9. Let M be an even-dimensional spin manifold.⁴⁰ Define the *spinor Laplacian* $\Delta := -\partial^2: C^\infty(SM) \rightarrow C^\infty(SM)$.

This is an even operator, preserving S^0 and S^1 . In general, the Laplacian is a self-adjoint operator, so a square root abstractly exists, but having a concrete description is part of the magic of spinor bundles.

Back to supersymmetric quantum mechanics. We've been considering a 1-dimensional supersymmetric QFT, whose space of fields is $\mathcal{C}_X = \Pi T \text{Map}(X, Y)$, with bosonic fields $\phi: X \rightarrow Y$ and fermionic fields $\psi \in \Pi \Gamma(\phi^* TY)$; the action is (9.7).

We showed that there's an odd vector field Q on \mathcal{C}_X and an even vector field H on \mathcal{C}_X with $[Q, Q] = H$ and $QS = 0$. Using the theory of spinors we've just developed, we'll associate Q to the Dirac operator and H to the Laplacian.

Let $X = S^1(T)$. We'll try to calculate

$$Z_{S^1(T)} = \int_{\mathcal{C}_X} d\phi d\psi e^{-S(\phi, \psi)}$$

by discretization. The supergeometry adds some nuance, but the general story still works.

Define $\mathcal{C}_{X;N}$ be the space of piecewise geodesic paths $S^1 \rightarrow Y$, e.g. ϕ changing direction at t_1, \dots, t_n , along with odd elements $\psi_i \in T_{\phi(t_i)} Y$. We'll then discretize the action, and define

$$Z_{X;N} := \int_{\mathcal{C}_{X;N}} d\psi d\phi e^{-S_{\text{disc}}}.$$

⁴⁰There is an analogous story in odd dimensions, which we will not need; there are slightly different statements, though.

The first apparent obstacle to writing down S_{disc} is: at a turning point, which direction do we apply the Levi-Civita connection in? There's no standard answer in the literature; what we're going to do is use the formula

$$\nabla_t = \frac{\partial}{\partial t} + A_t.$$

This relies on a choice of a local frame, but we have one around: choose the trivialization F_ϕ of ϕ^*TY coming from a fixed trivialization of TS^1 . It will turn out that the limit of the fermionic integral as $N \rightarrow \infty$ exists, but depends on F_ϕ .

On \mathbb{R}^n , the fermionic piece of the discretized action $S_{\text{disc}} \in C^\infty(\mathcal{C}_{X;N})$ is

$$S_{\text{disc}} = \sum_i \psi_i^I (\psi_{i+1}^I - \psi_i^I),$$

where $I = 1, \dots, 2n$. On a general Riemannian manifold, if $A_t = \alpha dt$, we would instead have

$$S_{\text{disc}} = \sum_i \psi_i^I (\psi_{i+1}^I - \psi_i^I + \alpha_j^I \psi_i^J).$$

This looks coordinate-dependent, but will turn out to be okay.

The space of trivializations F_ϕ of ϕ^*TY is a torsor for $\mathcal{L}SO(2n)$: any two choices differ by a loop. This loop group is disconnected, and its connected components are canonically $\pi_1(SO(2n))$; let τ be the generator of this group. One can show that the integral over fermions is

$$\omega(F_\phi) = -\omega(\tau F_\phi),$$

which is the precise sense in which it depends on the trivialization, which is topological.

This means we have a problem in defining Z_X : we need an extra structure on Y which picks out a “good” class of trivializations F_ϕ . This is exactly where we'll use a spin structure! Choose a spin structure $Q \rightarrow Y$; then, we can restrict to the F_ϕ which lift to the spin structure (i.e. F_ϕ maps into $SO(2n)$, and we want it to lift across the map $\text{Spin}(2n) \rightarrow SO(2n)$). This cures the sign problem.

Example 11.10. Let's look at a toy model: let Y be a spun surface and $\phi: S^1(T) \rightarrow Y$. Let $P \rightarrow Y$ be the bundle of oriented frames, a principal $SO(2)$ -bundle with the Levi-Civita connection; then, trivializing $E := \phi^*P$ produces a trivial principal $SO(2)$ -bundle over S^1 with a (perhaps nontrivial) connection

$$\nabla_t = \partial_t + \alpha R,$$

where $\alpha \in \mathbb{R}$ and

$$R = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}.$$

Then, we can make an educated guess for the (fermionic part of the) discretized action,

$$(11.11) \quad S_{\text{disc}} := \frac{1}{2} \sum_{i=1}^N \psi_i^1 \psi_{i+1}^1 + \psi_i^2 \psi_{i+1}^2 + \frac{\alpha T}{N} (\psi_i^1 \psi_{i+1}^2 - \psi_i^2 \psi_{i+1}^1).$$

Let $W = \int_{\mathcal{C}_X} e^{-S}$.

Proposition 11.12.

$$\lim_{N \rightarrow \infty} W = \frac{1}{2} \sin\left(\frac{1}{2} \alpha T\right).$$

This is a fermionic integral, so lots of stuff is nilpotent, and this quickly reduces to an algebraic, then a combinatorial problem. It's probably true that

$$W = \text{Re} \left(1 + \frac{i \alpha T}{2N} \right)^N,$$

and as $N \rightarrow \infty$, we get $e^{i \alpha T}$. (There may be details wrong, but we'll sort them out.)

Notice that this changes by a sign under $\alpha T \mapsto \alpha T + 2\pi$, and this is the sign change that we've been concerned by. ◀

Lecture 12.

The index of the Dirac operator: 10/10/17*“Not many people know this, but you have to put your pants on one leg at a time!”*

We’ve been studying supersymmetric quantum mechanics, a theory of “super maps” to a target Y , which is a Riemannian manifold with a spin structure. The space of fields \mathcal{C}_X consists of ordinary maps $\phi : X \rightarrow Y$ as well as a fermionic part, $\psi \in \Gamma(\Pi\phi^*TY)$. To compute partition functions, we’ll first integrate over fermions, then over bosons.

Remark 12.1. There’s a slight extension of the story of bosonic (i.e. non-supersymmetric) quantum mechanics that we discussed earlier: in addition to the input data of a compact Riemannian manifold Y and a potential $V : Y \rightarrow \mathbb{R}$, suppose that one also has a vector bundle $E \rightarrow Y$ with a metric (equivalently, one could take its principal $O(n)$ -bundle of frames) together with a compatible connection ∇ .

In this situation, one can formulate the heat equation “compatible with E ,” i.e. a similar-looking equation to (6.13), but for sections of E . Let $f_t \in \Gamma(Y, E)$ vary with C^2 regularity in t ; then, the heat equation is

$$(12.2) \quad \frac{\partial}{\partial t} f_t + H f_t = 0,$$

where

$$\begin{aligned} \Delta &= \nabla^* \nabla \\ H &= -\frac{1}{2} \Delta + V. \end{aligned}$$

Previously, the heat kernel (6.16) was a function on $Y \times Y$; in this situation, the heat kernel coupled to E is a section $k_t \in \Gamma(E^* \boxtimes E)$. The bundle $E^* \boxtimes E \rightarrow Y \times Y$ is the vector bundle whose fiber over (y_0, y_1) is

$$E_{y_0}^* \otimes E_{y_1} = \text{Hom}(E_{y_0}, E_{y_1}).$$

This is precisely what we need, for integrating with respect to this kernel maps sections to sections, as time evolution ought to.

Just as we did before, you can get the heat kernel coupled to E out of a one-dimensional quantum field theory: the fields are the same, and the action is the same. But we add something to the partition function

$$(12.3) \quad Z_{S^1(T)}(E) = \int_{\mathcal{C}_{S^1(T)}} d\phi \, \text{tr}(\text{Hol}_\phi) e^{-S(\phi)},$$

(and similarly for open boundary conditions). Here, we use the connection on E to define the holonomy Hol_ϕ generated by parallel transport for any loop $\phi : S^1 \rightarrow Y$.⁴¹ Coupling a theory to a bundle is a common technique in physics; if the vector bundle has rank n , this represents having n flavors of particles instead of one, and the nontriviality of the vector bundle encodes particle-particle relations.

One can rigorously show, with a proof whose sketch looks similar to how we got the heat kernel out of the untwisted case, that (12.3) is $\text{tr}(e^{-HT})$, where the trace is now in the space of L^2 sections of E . ◀

We’re talking about this today because we’ll need it in supersymmetric quantum mechanics. When we integrate over fermions, we’ll get an effective field theory of maps $X \rightarrow Y$ coupled to the spinor bundle $S \rightarrow Y$. More precisely, $S = S^0 \oplus S^1$, and we’ll find that

$$\begin{aligned} Z_{S^1(T), \text{eff}} &= Z(S^0) - Z(S^1) \\ &= \text{tr}_{L^2(S^0)}(e^{-HT}) - \text{tr}_{L^2(S^1)}(e^{-HT}). \end{aligned}$$

This quantity is also called the *super-trace* in super-algebra: if V is a super-vector space and $A \in \text{End } V$, the super-trace of A is

$$\text{Str}_V A := \text{tr}_{V^0} A - \text{tr}_{V^1} A.$$

Thus $Z_{S^1(T), \text{eff}} = \text{Str}_{L^2(S)} e^{-HT}$.

⁴¹This formula (12.3) looks a lot like the expectation for the observable tr Hol_ϕ , and indeed you can think of it as an expectation in the original theory. But this other perspective is also useful.

The Dirac operator \not{D} on S has a block form on S^0 and S^1 :

$$\not{D} = \begin{pmatrix} 0 & \not{D}^1 \\ \not{D}^0 & 0 \end{pmatrix}.$$

Precisely, $\not{D}^i: S^i \rightarrow S^{1-i}$ for $i = 0, 1$.

Definition 12.4. The *index* of an operator $A: V \rightarrow V$ on a (possibly infinite-dimensional) vector space V is $\text{ind}(V) := \dim \ker(V) - \dim \text{coker}(V)$.

This is only interesting on infinite-dimensional vector spaces — and even there, many bounded operators do not have finite indices.

Remark 12.5. The Dirac operator \not{D} is a first-order *elliptic* differential operator, meaning its principal symbol is invertible. For a more concrete example, in local coordinates the Laplacian on functions has the form

$$\Delta = -\sum_i \frac{\partial^2}{\partial x_i^2} + \text{lower-order terms}.$$

Its symbol $\sigma_\Delta: M \rightarrow \mathbb{R}$,⁴² where $\pi: T^*X \rightarrow X$ is projection is obtained by replacing $\frac{\partial}{\partial x_i}$ with k_i .

Ellipticity means that σ is invertible off of the zero section. On a compact manifold, the theory of elliptic operators is awesome: the kernel and cokernel of the operator are finite-dimensional, and, no matter what kind of sections you start with, the kernel consists of C^∞ sections! So the index of the Dirac operator is something interesting. \blacktriangleleft

Since the Dirac operator is self-adjoint, then $\dim \text{coker } \not{D}^0 = \dim \ker \not{D}^1$.

Proposition 12.6. *The partition function*

$$Z_{S^1(T), \text{eff}} = \text{Str } e^{-HT} = \text{ind}(\not{D}^0) = \dim \ker(\not{D}^0) - \dim \ker(\not{D}^1).$$

It's pretty cool that you get this quantity out of supersymmetric quantum mechanics.

Proof. Since $[\not{D}, \Delta] = -\Delta$ (again, this is a supercommutator), we can consider $L^2(S)$ as a unitary representation of the super-Lie algebra generated by \not{D} and Δ with that relation, which will have a nice decomposition into irreducibles.

Since Δ is central, we can diagonalize it, so it acts as a scalar E (called the energy) in each $\mathbb{Z}/2$ -graded irreducible V . There are three cases.

- (1) $\dim V = 1|1$, and $E > 0$. These are the states with $\psi \in V^0$ and $\not{D}\psi \in V^1$.
- (2) $\dim V = 1|0$ and $E = 0$, for a state ψ with $\not{D}\psi = 0$.
- (3) $\dim V = 0|1$ and $E = 0$, for a state ψ with $\not{D}\psi = 0$.

All irreducibles of the first kind will cancel out in the supertrace, since \not{D} switches V^0 and V^1 .

If $\Delta\psi = -E\psi$, then $E > 0$ and $\ker \Delta = \ker \not{D}$. This is because

$$\begin{aligned} -E\langle \psi, \psi \rangle &= \langle \psi, \Delta\psi \rangle \\ &= \langle \psi, -\Delta^2\psi \rangle \\ &= -\langle \Delta\psi, \Delta\psi \rangle \leq 0. \end{aligned}$$

Now let's see what happens to the supertrace in each of the three cases.

- For representations of type (1),

$$\text{Str}_V e^{T\Delta} = e^{TE} - E^{TE} = 0.$$

- For representations of type (2),

$$\text{Str}_V e^{T\Delta} = \text{Str}_V 1 = 1.$$

- For representations of type (3),

$$\text{Str}_V e^{T\Delta} = -\text{Str}_V 1 = -1.$$

Hence we get a factor of 1 for each piece of $\ker \not{D}^0$ and a -1 for each piece of $\ker \not{D}^1$, as desired. \boxtimes

⁴²More generally, for a differential operator $E \rightarrow F$, its symbol lives in $\Gamma(\pi^* \text{Hom}(E, F))$.

This proof was a toy example of something common in supersymmetry: there's an action of a super-Lie algebra on some super-vector space, and it decomposes into irreducibles, most of which cancel out.

Remark 12.7. The super-trace, and hence the partition function, doesn't depend on T ! The limit $T \rightarrow \infty$ makes it clearer why Proposition 12.6 is true; the limit $T \rightarrow 0$ makes heat flow local, which is how we'll approach the index theorem. \blacktriangleleft

To use this to prove interesting mathematics, we'll provide a formula for $\text{ind } \not{D}$, called the index theorem, which relates it to topology, specifically with characteristic classes.

Let C be a symmetric function on countably many variables $\{y_i\}$. Using C , one can define a characteristic class of principal $\text{SO}(2n)$ -bundles $E \rightarrow X$.⁴³ Choose a metric on X and a compatible connection on E with curvature form $F \in \Omega^2(\mathfrak{so}_{2n})$. In local coordinates, we can block-diagonalize

$$F = \bigoplus_{i=1}^n \begin{pmatrix} 0 & F_i \\ -F_i & 0 \end{pmatrix},$$

where $F_i \in \Omega^2(X)$.

Now, consider the form $C(\{F_i\}/2\pi i) \in \Omega^*(X)$.

Proposition 12.8. *The cohomology class of $C(\{F_i\}/2\pi i)$ does not depend on the choice of metric or connection.*

There is something to prove here.

Example 12.9 (Pontrjagin classes). Let X be a compact Riemannian manifold and $E = TX$, and let

$$C = \prod_i (1 + y_i^2).$$

The characteristic class associated to C is called the *(total) Pontrjagin class*

$$p(X) \in \Omega^*(X) = 1 + p_1(X) + p_2(X) + \cdots.$$

p_k is called the k^{th} Pontrjagin class. \blacktriangleleft

Since C is a function in y_i^2 , the Pontrjagin class $p_k \in \Omega^{4k}(X)$. For example,

$$p_1(X) = -\frac{1}{4\pi^2} \text{tr}(F \wedge F).$$

Remark 12.10. Strictly speaking, we've defined a representative for the Pontrjagin cohomology class in de Rham cohomology. If one works more topologically, one can define Pontrjagin classes for principal $\text{O}(n)$ -bundles in integer cohomology $p_k \in H^{4k}(X; \mathbb{Z})$. \blacktriangleleft

Example 12.11. The lowest-dimensional interesting examples are four-manifolds, where we have

$$\begin{aligned} \int_{S^4} p_1(S^4) &= 0 \\ \int_{\mathbb{CP}^2} p_1(\mathbb{CP}^2) &= 3 \\ \int_{K3} p_1(K3) &= -48. \end{aligned}$$

This already implies some interesting topology: suppose that \mathbb{CP}^2 had an orientation-reversing diffeomorphism. This would multiply $\int p_1(\mathbb{CP}^2)$ by -1 , but since the (cohomology class of the) Pontrjagin classes of a vector bundle don't depend on the orientation, this would force $3 = -3$, which is probably not true. So we deduce that \mathbb{CP}^2 and $K3$ admit no orientation-reversing diffeomorphisms! That is, \mathbb{CP}^2 is not diffeomorphic to $\overline{\mathbb{CP}^2}$, which is kind of strange: someone living in \mathbb{CP}^2 would observe some inherent chirality in their universe. \blacktriangleleft

Remark 12.12. The Hirzebruch signature theorem in four dimensions implies that if $\dim X = 4$, $p_1(X) = 3\sigma(X)$, where $\sigma(X)$ is the signature of the intersection pairing in middle cohomology. See, for example, the computations in Example 12.11. \blacktriangleleft

⁴³This construction is in fact universal, and works for n odd, but the wording is a little different.

Definition 12.13. Let X be a compact Riemannian manifold. Its \widehat{A} -genus⁴⁴ $\widehat{A}(X) \in \Omega^*(X)$ is the characteristic class of TX associated to the symmetric function

$$C = \prod_i \frac{y_i/2}{\sinh(y_i/2)} = \prod_i \left(1 - \frac{y_i^2}{24} + \frac{7y_i^4}{5760} + \cdots \right).$$

In low dimensions, the \widehat{A} -genus is

$$\widehat{A}(X) = 1 - \frac{1}{24}p_1(X) + \frac{7p_1(X)^2 - 4p_2(X)}{5760} + \cdots.$$

We'll use supersymmetric quantum mechanics to derive a formula for the Dirac operator in terms of characteristic classes:

Theorem 12.14 (Atiyah-Singer). *Let X be a closed spin manifold. Then,*

$$\text{ind } \not{D}^0 = \int_X \widehat{A}(X).$$

This will explain some surprising divisibility results in the indices of Dirac operators of spin manifolds.

Lecture 13.

Index theory and supersymmetric quantum mechanics, I: 10/12/17

Though we have been studying and will continue to study supersymmetric quantum mechanics, there are many things that are called that, and ours isn't necessarily the most commonly studied one. To clarify, physicists use a number \mathcal{N} , which keeps track of the size of the odd part of the super-Lie algebra of symmetries. In our case, this is the super-Lie algebra generated by Q and H with $[Q, Q] = H$, so there's just one odd piece. Thus, the theory we've been studying would be called $\mathcal{N} = 1$ supersymmetric QFT.⁴⁵

Though there are different flavors of quantum mechanics, which are all theories of maps $X \rightarrow Y$, where Y is a Riemannian manifold of any dimension and X is 1-dimensional. All of them have the same formal structure of a one-dimensional quantum field theory: there is a Hilbert space \mathcal{H} associated to a point, and a Hamiltonian $H: \mathcal{H} \rightarrow \mathcal{H}$; time evolution on $[0, T]$ acts by the operator $e^{-TH}: \mathcal{H} \rightarrow \mathcal{H}$. In this way, the formal structure of a one-dimensional quantum field theory: a functor from the category of 0-dimensional Riemannian manifolds and one-dimensional Riemannian cobordisms to the category of topological vector spaces.

\mathcal{N}	\mathcal{H}	H	Q
0	$L^2(Y)$	$-(1/2)\Delta$ (Laplace-Beltrami)	n/a
1	$\Gamma_{L^2}(Y, SY)$	$-(1/2)\Delta$ (spinor Laplacian)	\not{D}
2	$\Omega_{L^2}(Y) := \Gamma_{L^2}(\Lambda^* TY)$	$-(1/2)\Delta, \Delta = [d, d^*]$	d, d^*
4	$\Omega_{L^2}(Y)$	$-(1/2)\Delta, \Delta = [d, d^*]$	$\partial, \bar{\partial}, \partial^*, \bar{\partial}^*$

TABLE 1. Various kinds of supersymmetric quantum mechanics. The $\mathcal{N} = 4$ case requires a Kähler structure on Y ; there's a similar $\mathcal{N} = 8$ case when Y is hyperkähler (i.e. Kähler with respect to two different anticommuting complex structures).

See Table 1 for some common examples of supersymmetric quantum mechanics. There are others not in the table, e.g. some for which \mathcal{N} isn't a power of 2. Just as how the $\mathcal{N} = 4$ and $\mathcal{N} = 8$ cases underline interesting geometry on Y , these other models lead to new and interesting geometric structures (e.g. hyperkähler with torsion, etc.). However, the $\mathcal{N} = 2^k$ theories are usually studied the most frequently because they are the dimensional reductions of higher-dimensional theories that we want to understand.

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⁴⁴This is pronounced "A-hat."

⁴⁵Some people will call this $\mathcal{N} = 1/2$ supersymmetric QFT: be wary!

We now return to the story in progress: we're studying $\mathcal{N} = 1$ supersymmetric quantum mechanics, whose space of fields \mathcal{C}_X has bosonic terms $\phi : X \rightarrow Y$ and fermionic terms $\psi \in \Pi\Gamma(\phi^*TY)$, and a partition function

$$Z_{S^1(T)} = \int_{\mathcal{C}_{S^1(T)}} d\phi d\psi e^{-S(\phi, \psi)}.$$

We'd like to solve this by integrating out the fermions first, and the claim is that what you get is

$$W(\phi) = e^{-S(\phi)} \text{Str}_S \text{Hol}(\phi^*(TY)),$$

as we mentioned last time.

Remark 13.1. The usual method of proving this would be discretizing the path integral, but that doesn't work, because of something called the *fermion doubling problem*. This is a standard gotcha in lattice field theory: you want to take a continuum limit, but when you do, you get a theory with twice as many fermions as you started with. It should be possible to overcome this, discretize the path integral, and show that what you get agrees with the infinite-dimensional determinant \blacktriangleleft

So instead we're going to use the other method we've discussed: computing an infinite-dimensional determinant. We want to compute

$$\int d\psi e^{-(1/2) \int g(\psi, \nabla_t \psi)},$$

so we want the Pfaffian for the skew pairing on $\Gamma(\phi^*TY)$ defined by

$$(13.2) \quad (\psi_1, \psi_2) \mapsto \int dt \frac{1}{2} \langle \psi_1, \nabla_t \psi_2 \rangle.$$

As before, the determinant of a bilinear form isn't quite a number, until we choose a metric on $\Gamma(\phi^*TY)$; you'll see where we need it in the proof.

Let's start by looking at ∇_t . Its eigenvalues are

$$\lambda_{k, \pm i} = \frac{2\pi i}{T} \left(k \pm \frac{\alpha_i}{2\pi} \right),$$

where $k \in \mathbb{Z}$ and $e^{\pm i\alpha_i}$ are the eigenvalues of the holonomy operator on $\phi^*(TY)$. You can see this by choosing a local trivialization in which

$$\nabla_t = \partial_t + \begin{pmatrix} i\alpha_1/T & & & & \\ & -i\alpha_1/T & & & \\ & & \ddots & & \\ & & & i\alpha_N/T & \\ & & & & -i\alpha_N/T \end{pmatrix}.$$

Now to get a number for the determinant, we need a metric. Let's try the one that makes Fourier modes an orthonormal basis (the naïve L^2 norm). If we do this, then

$$\det (13.2) = \prod_{k \in \mathbb{Z}} \prod_{i=1}^n \frac{2\pi i}{T} \left(k + \frac{\alpha_i}{2\pi} \right) \left(k - \frac{\alpha_i}{2\pi} \right),$$

which diverges. Oops.

Instead, let's take the Sobolev norm

$$\|\psi\|^2 = \int dt g(\psi, \psi) + g(\nabla_t \psi, \nabla_t \psi).$$

Then the norm of the k^{th} Fourier mode is asymptotically about k^2 , and these functions are still orthogonal. For large k , the determinant relative to this norm is

$$\det (13.2) \sim \prod_{|k| > M} \prod_{i=1}^n \frac{2\pi i}{T} \left(1 + \frac{\alpha_i}{2\pi k} \right) \left(1 - \frac{\alpha_i}{2\pi k} \right),$$

and this converges.

Therefore the determinant is an honest function of the $\alpha_i \in \mathbb{C}$ with zeros of multiplicity 2 at $\alpha_i = 2\pi k$; moreover, it's periodic under $\alpha_i \mapsto \alpha_i + 2\pi$, and it's real for $\alpha_i \in \mathbb{R}$. This pins down the determinant up to a constant multiple: it must be

$$\det (13.2) \propto \prod_{i=1}^n \left(2 \sin \left(\frac{\alpha_i}{2} \right) \right)^2.$$

This already comes to us as the square of something, so we have an obvious candidate,

$$\prod_{i=1}^n 2 \sin \left(\frac{\alpha_i}{2} \right).$$

Unfortunately, this is not 2π -periodic: it depends on $\alpha \bmod 4\pi$. But this is exactly the sign ambiguity that the spin structure resolves: it allows the holonomy to be lifted to the double cover, and unambiguously get something 4π -periodic.

Remark 13.3. There are some great references for this, including papers by Witten [27] and Atiyah [2]. ◀

Let's see more carefully what's going on with this spin structure.

Lemma 13.4. *If $A \in \mathrm{SO}(n)$ has eigenvalues $e^{i\pm\alpha_i}$, then*

$$\det(1 - A) = \prod_{i=1}^n \left(2 \sin \frac{\alpha_i}{2} \right)^2.$$

Proposition 13.5. *If $A \in \mathrm{Spin}(2n)$ and $\rho : \mathrm{Spin}(2n) \rightarrow \mathrm{SO}(2n)$ is the double cover map, then*

$$\det(1 - \rho(A)) = (-1)^n (\mathrm{Str}_S A)^2,$$

where S is the spinor representation.

If you think it's weird that the determinant turns into a trace, you're not wrong. This is particular to $\mathrm{SO}(2n)$, and is completely false for other groups.

Proof when $n = 1$. Since $\mathrm{SO}(2)$ is rotation matrices, then there's some $\alpha \in \mathbb{R}$ such that

$$\rho(A) = \begin{pmatrix} \cos \alpha & \sin \alpha \\ -\sin \alpha & \cos \alpha \end{pmatrix},$$

and therefore

$$\det(1 - \rho(A)) = \left(2 \sin \left(\frac{\alpha}{2} \right) \right)^2.$$

The spinor representation of $\mathrm{Spin}(2)$ has complex dimension $1|1$, so

$$\mathrm{Str}_S(A) = e^{i\alpha/2} - e^{-i\alpha/2} = 2i \sin \left(\frac{\alpha}{2} \right). \quad \boxtimes$$

The proof for $n > 1$ follows the same strategy.

Thus

$$W(\phi) = \mathrm{Str}_S \mathrm{Hol} \phi^*(TY)$$

and

$$Z_{S^1(T)} = \mathrm{ind} \not{D}^0.$$

We can use this to compute $\mathrm{ind} \not{D}^0$ using localization, like we did for the Duistermaat-Heckman formula. We will introduce a perturbation $S \mapsto S + tQ\Psi$ and let $t \rightarrow \infty$.

Before, we did this for $\mathcal{C} = \Pi T M$ and $S = e^{-(H+\omega)}$, where M is a finite-dimensional manifold, H is a function on M generating a $U(1)$ -action on M , and ω is a symplectic form on M . This led to the ABBV formula

$$Z = \int_F \frac{e^{H+\omega}}{\mathrm{Eul}(NF)},$$

where F is the $U(1)$ -fixed locus of M .

Today, though, we're looking at $M = \mathcal{L}Y$ (since $\mathcal{C} = \Pi T(\mathcal{L}Y)$), which has a $U(1)$ -action by rotating the loops, and we take

$$H(\phi) = \frac{1}{2} \int g(\dot{\psi}, \dot{\psi})$$

$$\omega(\psi) = \frac{1}{2} \int g(\psi, \nabla_t \psi).$$

Again, let F denote the $U(1)$ -fixed locus of M , which is the space of constant loops. Thus $F \cong Y \subset \mathcal{L}Y$. On constant loops, H and ω are both 0, so if we can prove an analogue of the ABBV formula, we would conclude that

$$(13.6) \quad \text{ind } \not{D}^0 = Z = \int_Y \frac{1}{\text{Eul}(NF)}.$$

Our next goal is to understand NF , i.e. loops that are small deformations of constant loops. We have

$$T_{y_0} \mathcal{L}Y = \text{Map}(S^1, T_{y_0} Y) \cong T_{y_0} Y \oplus \bigoplus_{k \geq 1} (T_{y_0} Y \oplus T_{y_0} Y),$$

so

$$NF \cong \bigoplus_{k \geq 1} (TY \oplus TY),$$

where $U(1)$ acts on the k^{th} summand with weight k . The curvature of the connection is induced from the curvature on TY .

Next time, we'll compute the equivariant Euler class, and again get something in terms of an infinite-dimensional Pfaffian, which will turn out to be the \widehat{A} -genus.

Lecture 14.

Index theory and supersymmetric quantum mechanics, II: 10/17/17

We've been studying supersymmetric quantum mechanics (which will always mean a 1D supersymmetric QFT, and in our case an $\mathcal{N} = 1$ supersymmetric QFT), the super-theory of bosonic fields $\phi : X \rightarrow Y$ and fermionic fields $\psi \in \Gamma(\Pi \phi^* TY)$. As with the purely bosonic case, we related this to a heat kernel, and computed that $Z_{S^1(T)}$ is the supertrace of e^{-TH} , the index of \not{D}^0 , the even-graded part of the Dirac operator $\not{D} : \Gamma_{L^2}(Y, S) \rightarrow \Gamma_{L^2}(Y, S)$. Today we're going to use this to present a proof sketch of the Atiyah-Singer index theorem for \not{D} (which can be made precise, though we will not do so).

In the language of path integrals,

$$Z_{S^1(T)} = \int_{\mathcal{C}_{S^1(T)}} dx \, d\psi \, e^{-S},$$

and by a localization formula akin to the ABBV formula,

$$\int_Y \frac{1}{\text{Eul}(NF)}.$$

Here $U(1)$ acts on $\mathcal{L}Y$ by rotating loops, and the fixed locus F is the space of constant loops, which is homeomorphic to Y . The normal space $N_y F$ for some $y \in Y$ picks up a $U(1)$ -action, and as a representation decomposes as

$$(14.1) \quad N_y F \cong \bigoplus_{k > 0} T_y Y \oplus \mathbb{R}^2,$$

where U_1 acts in \mathbb{R}^2 with weight k .⁴⁶ This is because a deformation of the constant loop y is an infinitesimal loop, hence a map $S^1 \rightarrow T_y Y$.

Now we need to understand what the equivariant Euler class is. Previously, given a vector bundle E with a $U(1)$ -action, we defined its equivariant Euler class to be the Pfaffian $\text{Pf}(Z + R)$, where $Z \in \Omega^0(\mathfrak{so}(E))$ is an infinitesimal generator for the $U(1)$ -action and $R \in \Omega^2(\mathfrak{so}(E))$ is the curvature.

⁴⁶TODO: is this right?

In our case, $E = NF$, so using (14.1), the curvature is

$$(14.2a) \quad R \otimes 1 = \bigoplus_{i=1}^n R_i \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix} \otimes 1$$

and the infinitesimal generator for its $U(1)$ -action is

$$(14.2b) \quad Z = 1 \otimes \bigoplus_{k>0} k \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}.$$

Computing the Pfaffian is a finite-dimensional question about linear algebra.

Exercise 14.3. On $\mathbb{R}^2 \otimes \mathbb{R}^2$, show that

$$\text{Pf}\left(a \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix} \otimes 1 + b \left(1 \otimes \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}\right)\right) = a^2 - b^2.$$

Formally, we can conclude that

$$\text{Eul}(NF) = \prod_{i=1}^n \prod_{k>0} (k + R_i)(k - R_i).$$

If we try regularizing it as we did before, when the R_i were numbers instead of forms, we get

$$\text{Eul}(NF) = \prod_{i=1}^n \frac{\sin \pi R_i}{\pi R_i}.$$

Here, by $\sin(\pi R_i)$ we mean the formal power series, where multiplication is wedge product of forms. Since the zeroth-order term is πR_i , the quotient makes sense.

Remark 14.4. This is one of the places where a rigorous proof would need to be careful — why is this the right regularization for this infinite-dimensional determinant? ◀

Anyways, $\text{Eul}(NF)$ is almost $1/\widehat{A}(Y)$, though it's off by a factor of $(4\pi^2)^k$. There are two ways to fix this.

- (1) A more precise treatment would get rid of this discrepancy, and product the \widehat{A} -genus on the nose.
- (2) Alternatively, compute one concrete example in each dimension to fix the discrepancy.

Thus we can (mostly) conclude

$$Z_{S^1(T)} = \int_Y \widehat{A}(Y),$$

implying the index theorem for \not{D}^0 .

Remark 14.5. This is only one piece of the general Atiyah-Singer index theorem and its many extensions and variants. One simple one is to fix another vector bundle E with a metric and connection over Y . Then, there's a twisted Dirac operator $\not{D}_E: S \otimes E \rightarrow S \otimes E$. The Atiyah-Singer index theorem for this Dirac operator says that

$$\text{ind } \not{D}_E^0 = \int_Y \widehat{A}(Y) \cdot \text{ch}(E).$$

Here, $\text{ch}(E)$ is the *Chern character*, a characteristic class of principal $U(n)$ -bundles associated to the symmetric function

$$P(z_i) = \sum e^{z_i}.$$

It has an expansion in terms of the usual Chern classes, which begins

$$\text{ch}(E) = \text{rank}(E) + c_1(E) + \frac{1}{2}(c_1(E)^2 - 2c_2(E)) + \cdots.$$

Unlike the Pontrjagin classes and \widehat{A} -genus, these characteristic classes are even-dimensional, i.e. $c_i(E) \in \Omega^{2i}$, and therefore are interesting sooner. For example, this variant implies the existence of anomalous theories on surfaces, where the anomaly comes entirely from the Chern character.

This variant can again be proved (more or less rigorously) using supersymmetric quantum mechanics coupled to E : the partition function is now

$$Z_{S^1(T)} = \int d\phi d\psi \, \text{tr} \widehat{\text{Hol}}(\phi) e^{-S(\phi, \psi)},$$

where we need to replace $\text{trHol}(\phi)$ with something supersymmetric, namely

$$\overline{\text{tr}(\text{Hol } \phi)} := \exp\left(i \oint A + \psi \psi F\right).$$

For a reference, see Friedan-Winney [16].

In supersymmetric localization for this supersymmetric quantum mechanics, one will get

$$\int_Y \frac{e^{H+\omega}}{\text{Eul}(NF)},$$

and now $\omega = F$, so we recover

$$\int_Y \frac{e^F}{\widehat{A}(Y)^{-1}},$$

suggesting how the Chern character appears.⁴⁷

Coupling the theory to E will allow you to recover many of the classical corollaries of the Atiyah-Singer index theorem: e.g. for the Gauss-Bonnet theorem, one uses $E = S$ (since $S \otimes S = \Omega^*$). ◀

Remark 14.6. Another important extension is the Atiyah-Singer index theorem for families [3, 4]: given a family of twisted Dirac operators smoothly varying over a parameter space B , one can define a super-vector bundle $I \rightarrow B$, the *index bundle*, which carries a natural connection, and you can ask about its holonomy, curvature, and so on.⁴⁸ Originally, Atiyah and Singer computed the K -theory class of this bundle, but more can be understood, and a lot of it was worked out by Bismut-Freed [9, 10].

In quantum mechanics, this means we have a family of one-dimensional supersymmetric QFTs, or in other words a family of supersymmetric QM theories. Such a family always gives a vector bundle of ground states over its parameter space, with a connection (called the *Berry phase*, which is well-known in physics). Alvarez-Winney [1] give a supersymmetric explanation of the index theorem for families using the Berry connection. ◀

Of course, there's more to the index theorem than the Atiyah-Singer theorem for families, but it's not clear what the limit is. How much of index theory can be absorbed into supersymmetric quantum mechanics?

This wraps up our discussion of 1D QFT; we now move on to the two-dimensional case. One special case will be $X = T^2$, so the space of fields will be (some supersymmetric version of) $\mathcal{L}^2 Y$. One might then expect an index theorem for Dirac operators on $\mathcal{L}Y$, which will require a spin structure on $\mathcal{L}Y$ (and hence some higher structure on Y). This will compute something called the elliptic genus, and here something interesting happens – topological properties of the elliptic genus were first suggested by the physics, then later proven mathematically!

Physically, this means spacetime is two-dimensional: one can imagine a condensed-matter system on a wire, or a field theory whose fields are worldlines of strings.

The free boson in two dimensions. We're going to define a two-dimensional analogue of the σ -model we've been studying. Let X be a 2-dimensional Riemannian manifold and Y , the target, be a Riemannian manifold,⁴⁹ and $V : Y \rightarrow \mathbb{R}$ be a function. Let $\mathcal{C} := \text{Map}(X, Y)$. Similarly to before, we'll let the action be

$$(14.7) \quad S(\varphi) := \int_X \left(\frac{1}{2} \|\text{d}\varphi(x)\|^2 + V(\varphi(x)) \right) \text{dvol}_X,$$

where dvol_X is the volume *density* on X (since we didn't require X to be oriented). The term $\|\text{d}\varphi\|$ may look confusing, but the idea is that $\text{d}\varphi : T_x X \rightarrow T_{\varphi(x)} Y$ is a bounded linear map between normed vector spaces, and therefore has the usual norm in that sense. We'll try to proceed and see where new difficulties arise.

Exercise 14.8. Suppose $X = S^1(L) \times S^1(T)$, $Y = \mathbb{R}$, and $V = 0$. Then, we can Fourier-expand φ in modes:

$$\varphi(x, t) = \sum_{n \in \mathbb{Z}} a_n(t) e^{2\pi i n x / L}.$$

⁴⁷Here and in some other places, we've implicitly assumed the dimension of Y is even; in odd dimensions, these things tend to vanish, or behave differently.

⁴⁸The dimension of this may jump, but in ways that can be accounted for and dealt with.

⁴⁹We will not impose any orientation, spin, etc. until we introduce supersymmetry. This is generally how the story goes.

This allow us to simplify (14.7):

$$(14.9) \quad S(\varphi) = L \int dt \frac{1}{2} \dot{a}_0(t)^2 + \sum_{n>0} \left(|\dot{a}_n(t)|^2 + \frac{4\pi^2 n^2}{L^2} |a_n(t)|^2 \right).$$

We've gone from one function of two variables to infinitely many functions $a_n(t)$ of one variable $t \in S^1(T)$,⁵⁰ which doesn't sound simpler, but each of $\text{Re}(a_n)$ and $\text{Im}(a_n)$ is something we've seen before: the total action is

$$S(\mathbf{a}(t)) = \sum_{n \geq 0} S_n(a_n(t)),$$

and each S_n is separately the action of a harmonic oscillator with potential

$$V(a) = \omega_n^2 a_n^2,$$

where $\omega_n := 2\pi n/L$.

Now we want to compute the partition function. To be fully rigorous, one should discretize the model and put it on a lattice. This is harder than before but probably tractable, and has probably been done. But this is hard, so we're going to try to carefully guess the right answer. Naïvely, we would like

$$Z_{S^1(L) \times S^1(T)} = \text{tr}_{\mathcal{H}} e^{-TH},$$

for some Hilbert space \mathcal{H} of the theory. We know the Hilbert space for the harmonic oscillator is $L^2(\mathbb{R})$, so we're going to guess that we get a copy of $L^2(\mathbb{R})$ for each $\text{Re}(a_n)$ and $\text{Im}(a_n)$, i.e.

$$\mathcal{H} = \bigotimes_{n \in \mathbb{Z}} L^2(\mathbb{R}).$$

The Hamiltonian should also decompose in this way, because the action did:

$$H = \sum_{n \in \mathbb{Z}} H_n,$$

where $H_n: \mathcal{H} \rightarrow \mathcal{H}$ is the identity on all but the n^{th} copy of $L^2(\mathbb{R})$.

This seems reasonable enough, so let's talk about the eigenvalues of H . The smallest one, the ground state, should be the sum of the ground state energies of for each \mathcal{H}_n . This is

$$(14.10) \quad E = \sum_n E_n = 2 \sum_{n>0} \frac{\omega_n}{2} = \frac{2\pi}{L} \sum_{n>0} n.$$

This is a problem, and says that we need to think more carefully; it's the first divergence people usually encounter in quantum mechanics.

Next time, we'll discuss how to resolve that infinite sum. The answer, of course, is $-1/12$, and we'll discuss why this is physically meaningful.

Lecture 15.

Free bosons and modular forms: 10/19/17

Last time, we briefly looked at a 2D σ -model for a "free bosonic" theory, of maps $\phi: X \rightarrow \mathbb{R}$, where X is a surface with a Riemannian metric. The action is

$$(15.1) \quad S(\phi) = \int_X \text{dvol}_X \frac{1}{2} \|\text{d}\phi\|^2.$$

Let's specifically let X be the flat torus with length L and width T ; we want to compute the partition function for X . There are two routes: since (15.1) is a quadratic form on some huge vector space, we could evaluate the partition function as the infinite-dimensional determinant of a form

$$\phi_1, \phi_2 \mapsto \int \langle \text{d}\phi_1, \text{d}\phi_2 \rangle.$$

Alternatively, as we did last time, we could write it as

$$Z_X = \text{tr}_{\mathcal{H}(L)} e^{-TH},$$

⁵⁰We're going to ignore the zero mode, which is a map to \mathbb{R} with no potential, for the time being.

by Fourier-expanding

$$\phi(t, x) = \sum_n a_n(t) e^{inx/L}$$

and correspondingly decomposing the Hilbert space as

$$\mathcal{H}(L) = L^2(\mathbb{R}) \otimes \bigotimes_{n>0} (L^2(\mathbb{R})^{\otimes 2}).$$

The Hamiltonian also decomposes: the first copy has Hamiltonian

$$H_0 := -\frac{1}{2} \frac{\partial^2}{\partial a^2},$$

and the n^{th} piece has Hamiltonian

$$H_n := -\frac{1}{2} \frac{\partial^2}{\partial a^2} + \frac{1}{2} \omega_n^2 a^2,$$

where $\omega := 2\pi n/L$.

Remark 15.2. We'd like to imitate what we did last time, but have a major stumbling block: $\mathcal{H}(L)$ has no countable basis. A naïve choice of basis is

$$\psi_{k_0} \otimes \psi_{k_1} \otimes \cdots, \quad k_i \in \mathbb{N},$$

but this is uncountable. This relates to trying to define ϕ in terms of initial conditions of its Fourier modes; now we need countably many numbers rather than just one. \blacktriangleleft

More generally, given an n -dimensional quantum field theory formulated with a path integral and an $(n-1)$ -manifold M , you want to define a Hilbert space that the theory canonically associates to M . We haven't been trying to do this in general, and have been explicitly checking that the Hamiltonian and Lagrangian (path integral) points of view agree. The heuristic, which is still far from proven in general, is that one should restrict attention to the extrema of the action S when the theory lives on $M \times \mathbb{R}$. Let \mathcal{S} denote the space of extrema (in our examples so far, these have been spaces of harmonic maps, and in 1D QFT, we specifically had parameterized geodesics). Then, \mathcal{S} canonically carries a symplectic structure, which is a formal construction. This approach is called *path integral quantization*.

Example 15.3. In quantum mechanics, the space of geodesics is $\mathcal{S} = TY$. We can use the Riemannian metric to identify $TY \cong T^*Y$; the canonical symplectic structure is the Liouville form coming from T^*Y . Physically, for harmonic maps $\phi_1, \phi_2: M \times \mathbb{R} \rightrightarrows \mathbb{R}$, the symplectic form is

$$\omega(\delta\phi_1, \delta\phi_2) = \int_M (\langle \delta\phi_1(x), \delta\dot{\phi}_2(x) \rangle - \langle \delta\phi_2(x), \delta\dot{\phi}_1(x) \rangle). \quad \blacktriangleleft$$

This is all classical field theory — to obtain the Hilbert space for M , take the geometric quantization of (\mathcal{S}, ω) . This is very difficult in general, but in special cases useful things are known. While we won't study this in detail, it's a general context for this story that may be useful to know.

In top dimension, one gets a number associated to every closed manifold, which is its partition function. In codimension 1, one should get the Hilbert space obtained by this geometric quantization. If X is an n -dimensional manifold with boundary M , the partition function should be a number, provided initial conditions on M . That is, X is an element of $\text{Hom}(\mathcal{H}_M, \mathbb{C})$ (or, if one has multiple boundary components M_1 and M_2 , $\text{Hom}(\mathcal{H}_{M_1}, \mathcal{H}_{M_2})$); we saw this for the heat kernel as time evolution on $\mathcal{H} = L^2(\mathbb{R})$. This uses the fact that a neighborhood of ∂X in X looks like $X \times \mathbb{R}_{\geq 0}$, and we got the Hilbert space from $M \times \mathbb{R}$.

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Returning to our specific example, we're going to try to compute the smallest eigenvalue of H . We did this in (14.10), and quickly ran into a divergent series. We're going to try to work around this by doing some sort of cutoff.

Suppose we discretized the theory. This roughly means cutting off the higher Fourier modes and leaving the smaller ones unchanged. The answer will scale E_n by some cutoff function $f(E_n)$, which is close to 1 below about $1/\varepsilon$, close to 0 above about $1/\varepsilon$, and analytic. For example, $f_\varepsilon(E) = e^{-\varepsilon E}$ works.

This makes the divergence in (14.10) better behaved:

$$\begin{aligned} E &= \frac{2\pi}{L} \sum_{n>0} n e^{-\varepsilon n/L} \\ &= \frac{2\pi}{L} (-L) \frac{d\varepsilon}{d \sum_{n>0} e^{-n\varepsilon/L}} \\ &= -2\pi \frac{d}{d\varepsilon} \left(\frac{e^{-\varepsilon/L}}{1 - e^{-\varepsilon/L}} \right). \end{aligned}$$

This converges, and we conclude

$$E(\varepsilon) = \frac{2\pi}{L} \frac{e^{\varepsilon/L}}{(e^{\varepsilon/L} - 1)^2}.$$

This has a singularity as $\varepsilon \rightarrow 0$, corresponding to getting rid of the cutoff function. If you expand around $\varepsilon = 0$, you'll find

$$E(\varepsilon) = \frac{2\pi L}{\varepsilon^2} - \frac{\pi}{6L} + \dots$$

Thus, the only piece of this that diverges is proportional to L , so we could absorb it by adding a local term to the action

$$\delta S = \frac{2\pi}{\varepsilon^2}.$$

This doesn't change much about the physics,⁵¹ since it's a purely constant local term, and after making this change, you obtain

$$E = -\frac{\pi}{6L},$$

which is finite!⁵²

That was just the ground state energy... but fortunately, the rest of the spectrum is easier: the difference between

$$\psi_0 \otimes \psi_0 \otimes \psi_0 \otimes \dots$$

and

$$\psi_{n_1} \otimes \psi_{n_2} \otimes \psi_{n_3} \otimes \dots$$

(as long as only finitely many $n_i \neq 0$) is

$$E = -\frac{\pi}{6L} + \sum_i n_i \omega_i,$$

the sum of the energy shifts for ψ_0 to ψ_{n_i} in the one-dimensional theory.

If you try to excite infinitely many states, you'll get something with infinite energy, but that's okay, because those states do not contribute to the partition function. We can use this to compute the partition function to be

$$Z_X = ?? \cdot \eta(q)^{-2},$$

where $q := \exp(-2\pi T/L)$ and

$$(15.4) \quad \eta(q) := q^{1/24} \prod_{n=1}^{\infty} (1 - q^n).$$

Each term in the product tracks one of the components of the Hilbert space in the decomposition we made. The remaining term is the contribution from the zero mode, a harmonic-oscillator-like system but with zero potential. This produces an infinite factor in the path integral just by itself, for a simple reason: the configuration space is $\mathcal{C} = \text{Map}(S^1 \times S^1, \mathbb{R})$, and \mathcal{C} and the action (15.1) share a translation symmetry $y \mapsto y + c$ (on the target) under which the integrand is invariant. Thus the path integral cannot be invariant, because it must contain a factor of $\int_{-\infty}^{\infty} dt$: if you integrate something translation-invariant and nonzero on \mathbb{R} , the result must be infinity. This just shifts the zero mode, which is why we didn't see it until now.

⁵¹Specifically, it changes the vacuum energy, but none of the normalized correlation functions.

⁵²In physics, this is thought of as getting rid of the cosmological constant using renormalization.

We're going to deal with \mathcal{H}_0 by regulating: replace $L^2(\mathbb{R})$ with $L^2(S^1(V))$, where V is some large length. So just for regulating this one piece, we pretend the target is a circle. This produces a finite answer (though it does depend on V):

$$Z_X = \frac{V}{2\pi\sqrt{T/L}} \eta(q)^{-2}.$$

The η -term is the most interesting piece, and it has a cool modular property encoding the fact that Z is invariant under the symmetry exchanging L and T : if $\tau := T/L$ so $q = e^{2\pi i\tau}$, then there's a modular transformation sending $\tau \mapsto \tau' = iL/T = -1/\tau$. The modular property says that Z_X is invariant under $L \leftrightarrow T$. Realistically, we've done a bunch of heuristics, so this modular property is justification that the heuristics were right. Even the $q^{1/24}$ is important: without it, modularity is lost.

In order to think of η as a modular form, we should also say something about translation-symmetry on the upper half-plane (in τ), and this does happen. The physics can still be formulated on $X_\tau := \mathbb{C}/(\mathbb{Z} \oplus \tau\mathbb{Z})$ for some $\tau \in \mathbb{H}$, a "tilted torus" (from a parallelogram lattice, not a square lattice), and again you can compute the partition function, which looks very similar:

$$(15.5) \quad Z_{X_\tau} = \frac{V}{2\pi\sqrt{\text{Im } \tau}} |\eta(q)|^{-2},$$

where, as before, $q = e^{2\pi i\tau}$. This η -function is, somewhat mysteriously, one of the more fundamental modular forms, appearing in plenty of other contexts, including purely algebro-geometric questions about K3 surfaces.

We've learned something very interesting here: 2D QFT is a factory of modular forms: on the torus, if we choose a theory which is invariant under symmetries of the torus, the partition function ought to be a modular form! We chose a simple target, so we're going to get a simple modular form. But for fancier targets, we'll get cooler and crazier stuff, including something called the Witten genus once we turn on supersymmetry: it will come from one of the simpler supersymmetric theories, an $\mathcal{N} = (0, 1)$ -supersymmetric 2D σ -model. But all in due time.

Lecture 16.

Free bosons with target $S^1(R)$: 10/26/17

Last time, we considered the free bosonic QFT in 2 dimensions, the theory of maps $\phi : X \rightarrow \mathbb{R}$, where X is a surface with a Riemannian metric and the action is

$$S(\phi) = \int_X \text{dvol}_X \|\text{d}\phi\|^2.$$

We then provided some heuristics for why the partition functions make sense, even though they produce infinities if evaluated naïvely.⁵³

Though you can do this for any surface, and it depends on the genus in a nice way, we focused on the torus $X := S^1(T) \times S^1(L)$, for which the partition function is

$$Z_X = \frac{V}{2\pi\sqrt{T/L}} \eta(q)^{-2},$$

where $q = e^{-2\pi T/L}$ and $\eta(q)$ is the Dedekind η -function (15.4).

V came from the need to regulate an *IR divergence*, which ultimately arises from the noncompactness of the field space. It came from a translation action $\phi \mapsto \phi + c$. It's like adding a small mass to the physics, and fixes the infinity. Instead of computing Z_X , take instead a point $p \in X$ and compute $\langle e^{-\varepsilon(\phi(p))^2} \rangle$ and take the limit as $\varepsilon \rightarrow 0$. Other ways of fixing this produce the same result, suggesting it's actually something meaningful. You can also think of this fix as replacing \mathbb{R} with an S^1 , but only for the zero mode, foreshadowing that you might be interested in compactifying the target for the entire theory.

This destroys the translation symmetry, which might be sad, but it preserves a much more interesting symmetry, modular symmetry on the space of tori. This began from the observation that the answer is invariant under switching T and L . Then, for a tilted torus $X_\tau := \mathbb{C}/\mathbb{Z} \oplus \tau\mathbb{Z}$, the partition function is

$$Z = \frac{V}{2\pi\sqrt{\text{Im } \tau}} |\eta(\tau)|^{-2},$$

⁵³One can formulate the same action with target Y , where Y is any Riemannian manifold, but the argument that the partition functions are finite would be different, and in fact might not hold.

where $q := e^{2\pi i \tau}$. This is invariant under the actions $\tau \mapsto \tau + 1$ and $\tau \mapsto -1/\tau$, hence under the modular group.

There's yet more symmetry: this partition function only depends on the conformal class of the metric. That is, if you rescale the metric, the partition function is the same (since it only uses the ratio T/L); alternatively, it makes sense as a theory of surfaces solely with a conformal structure. One says that the free boson in 2 dimensions is an example of a *conformal field theory*.

This is an important fact, and can be seen before calculating the partition function: how does the action change when you rescale the metric g_{ij} ? In coordinates, the action is

$$S(\phi) = \int_X \sqrt{\det g} \partial_i \phi \partial_j \phi (g^{-1})^{ij},$$

so if we rescale $g \mapsto \lambda g$ for some $\lambda > 0$, $\sqrt{\det g} \mapsto \lambda \sqrt{\det g}$, and $(g^{-1})^{ij} \mapsto \lambda^{-1} (g^{-1})^{ij}$, so the action is invariant!⁵⁴ Hence everything about the theory – the partition function, correlation functions – is scale-invariant.⁵⁵

Remark 16.1. A similar argument holds in other dimensions, though one also has to rescale ϕ to cancel out the effect of λ . ◀

This argument hasn't interacted with the way we regulated the IR divergence, but this is fine as long as the way we regular does not destroy this symmetry.

Alternatively, we could compactify the target space: let's consider the theory of maps $X \rightarrow Y$, where X is still the torus $S^1(L) \times S^1(T)$, but $Y := S^1(2\pi R)$. In the language of path integrals, the space of fields $\mathcal{C} = \text{Map}(X, Y)$ is disconnected! It breaks up as a disjoint union indexed by winding numbers:

$$\mathcal{C} = \bigsqcup_{n_1, n_2 \in \mathbb{Z}} \mathcal{C}_{n_1, n_2},$$

where $\phi : X \rightarrow Y$ is in \mathcal{C}_{n_1, n_2} iff

$$\begin{aligned} \phi(x + L, t) &= \phi(x, t) + 2\pi R n_1 \\ \phi(x, t + T) &= \phi(x, t) + 2\pi R n_2. \end{aligned}$$

So the path integral would be a sum, indexed by $n_1, n_2 \in \mathbb{Z}$, of partition functions Z^{n_1, n_2} . You could compute this using Gaussians like we've done before.

Dually, in the Hamiltonian formalism, the space \mathcal{S} of classical solutions on $S^1(L) \times \mathbb{R}$ (i.e. the spatial piece of X times \mathbb{R}) doesn't immediately see both winding numbers, as the time-direction winding number is global. Thus in this case,

$$\mathcal{S} = \bigsqcup_{n_1 \in \mathbb{Z}} \mathcal{S}_{n_1},$$

again sorted by winding number.

When you do path-integral quantization of \mathcal{S} , you'll get a direct sum of spaces, indexed by n_1 . Where does n_2 come from? Each \mathcal{S}_{n_1} is acted on by the group of isometries of Y , which is an S^1 . Hence $\mathcal{S}_{n_1} = S^1 \times \mathcal{S}_{n_1, \text{reduced}}$, which means after quantizing, you'll get a copy of $L^2(S^1)$, and that decomposes via Fourier series, producing the second \mathbb{Z} -valued index:

$$\mathcal{H}^{n_1} \cong \bigoplus_{m_2 \in \mathbb{Z}} \mathcal{H}^{n_1, m_2}.$$

Our notation is deliberate: m_2 and n_2 are not the same. The interpretation of m_2 is the center-of-mass momentum, but n_2 is the winding number for time. The contributions from \mathcal{C}_{n_1, n_2} are not the same as the contributions from \mathcal{H}^{n_1, m_2} when $n_2 = m_2$; they're related by Poisson resummation.

When one does all of the calculations, the lowest eigenvalue of H in \mathcal{H}^{n_1, m_2} is

$$(16.2) \quad E_{n_1, m_2} = \frac{1}{L} (n_1 R)^2 + \left(\frac{m_2}{R} \right)^2.$$

The first term comes from the winding number in the space direction, and the second from Fourier analysis on $S^1(R)$: the eigenfunctions are $\psi_{m_2} = e^{2\pi i m_2 / R}$, and H restricts to $\frac{d^2}{dx^2}$. This is not complicated – though there's *a priori* a string propagating around a circle, we're able to focus just on its center of mass, which is a particle propagating around a circle, and that's the same as in quantum mechanics.

⁵⁴More than this is true: even if λ depends on x , the action is still invariant.

⁵⁵This is not completely immediate: potentially, there could be an issue with the conformal anomaly, though in this case it doesn't arise.

The first term in (16.2), however, really does depend on the string. Consider $\phi(x, t) := x \cdot n_1 2\pi R/L$. The action is

$$S(\phi) = \int_{S^1(L)} \|d\phi\|^2 = L \frac{(2\pi n_1 R)^2}{L^2}.$$

TODO: there's a spurious factor of $1/L$ somewhere.

Anyways, the derivation will produce a partition function which is a sum over sectors:

$$Z_{T^2} = \frac{1}{|\eta(\tau)|^2} \sum_{n_1, m_2 \in \mathbb{Z}} q^{(m_2/R - n_1 R)^2/4} \bar{q}^{(m_2/R + n_1 R)^2/4},$$

where again $q := e^{-2\pi T/L}$.

This has a modular symmetry as before, which isn't too much of a surprise. But another symmetry pops up: the partition function is invariant under $R \leftrightarrow 1/R$! This is something different about strings compared to point products: this theory cannot tell the difference between a circle of radius R and a circle of radius $1/R$. This kind of unexpected duality happens a lot in higher-dimensional QFT: we have a single QFT with two different classical descriptions, even though the theory itself does not contain this symmetry when you fix an R .

This duality extends to correlation functions. For example, for any $x \in X$, the map $\phi \mapsto d\phi(x)$ is a map $\mathcal{C} \rightarrow T_x^*X$, giving a T_x^*X -valued observable. (In coordinates, you'd look at $\partial_t \varphi$ and $\partial_x \varphi$.) Similarly, one has $\star d\phi(x): \mathcal{C} \rightarrow T_x^*X$, where one applies the Hodge star operator. Then,

$$\langle d\phi(x_1) d\phi(x_2) \rangle_{Y=S^1(R)} = \langle \star d\phi(x_1) \star d\phi(x_2) \rangle_{Y=S^1(1/R)}.$$

This is a little spooky: these are two completely different integrals, which the notation belies: these are maps into two seemingly unrelated circles! And there's an analogue of this for any observable.

This is a little strange, but there is motivation for it. If $\mathcal{O} = d\phi$, then there's a trivial kinematic relation $d\mathcal{O} = 0$, i.e. any correlation function $\langle d\mathcal{O}(x) \cdots \rangle = 0$. But $\mathcal{O}_D := \star d\phi$ also satisfies $d\mathcal{O}_D = 0$, for the fancier reason that $d \star d\phi = 0$.

This duality is an instance of *T-duality*.

Remark 16.3. When $R = 1$, this is still an interesting self-duality: sure, it's the same theory, but you get identities for partition functions, which make for an enhanced symmetry. This is very hard to see classically. ◀

Supersymmetry. A supersymmetric analogue of the theory of maps $T^2 \rightarrow Y$ will again produce interesting topology, this time mixed with the modular symmetry.

Recall that in $\mathcal{N} = 1$ supersymmetric quantum mechanics, $Z_{S^1(T)}$ recovered for us the \hat{A} -genus, a characteristic class associated to the symmetric function

$$\prod_i \frac{z_i/2}{\sinh(z_i/2)},$$

and the Atiyah-Singer index theorem equates this with the index of a Dirac operator.

The answer in 2D is the *Witten genus*, a characteristic class associated to the symmetric function

$$\widetilde{W}(q, z_1, \dots) = \prod_i \frac{z_i/2}{\sinh(z_i/2)} \left(\prod_{n \geq 1} \frac{(1 - q^n)^2}{(1 - q^n e^{2i})(1 - q^n e^{-2i})} \right).$$

The z_i get eaten up in defining the characteristic class, but if you expand this in powers of q , you obtain a series of characteristic classes

$$W(q) = \sum_{n \geq 0} q^n W_n(q),$$

where each $W_n(q)$ is some explicit function of the curvature.

For any particular manifold Y , you get some q -series $W_Y(q)$. Recall that the \hat{A} -genus is an integer on a spin manifold, but the analogue of integrality for the Witten genus is crazier: if Y is spin, all coefficients $W_{n,Y} \in \mathbb{Z}$, and if furthermore $p_1(Y) = 0$, then $W_Y(q)$ is a modular form! This is ultimately because it comes from the partition function of a QFT with modular symmetry like we've seen today.

Lecture 17.

 T -duality for free bosons in 2D: 10/31/17

Last time, we studied the 2D free boson theory with target S^1 , the theory of maps $X \rightarrow S^1(R)$. We previously studied the partition function on $X = S^1(T) \times S^1(L)$ with target \mathbb{R} , so we expanded

$$\phi(x, t) = \frac{1}{\sqrt{2L}} \sum_{n \in \mathbb{Z}} a_n(t) e^{2\pi i n x / L},$$

where $1/\sqrt{2L}$ is a normalization: you can pick a different one, but this one is convenient.

For maps to $S^1(R)$, we modified this approach in two ways.

(1) First we added a term for the winding number of ϕ :

$$\phi(x, t) = \frac{1}{\sqrt{2L}} \sum_{n \in \mathbb{Z}} a_n(t) e^{2\pi i n x / L} + w x \frac{R}{L},$$

where $w \in \mathbb{Z}$ is such that

$$\phi(x + L, t) = wR + \phi(x, t).$$

(2) Now, a_0 is a map $\mathbb{R} \rightarrow S^1(R)$ (where \mathbb{R} is the time coordinate).

But after these modifications, the action $S(\phi)$ in terms of $a_n(t)$ is exactly the same, except with an additional term

$$\int_{S^1(L)} \left(\omega \frac{R}{L} \right)^2 = \frac{\omega^2 R^2}{L}.$$

Therefore for the partition function $Z_{S^1 \times S^1}$, the contributions from all fields $a_n(t)$, $n \neq 0$, are the same as before, $\eta(q)^{-2}$. The winding number produces a contribution of

$$\sum_{w \in \mathbb{Z}} e^{-w^2 R^2 T / L},$$

and $a_0(t)$ has action just like in our study of quantum mechanics:

$$\int dt \frac{1}{2} \dot{a}_0(t)^2,$$

with periodicity $a_0 \sim a_0 + R\sqrt{L}$. Thus the contribution of a_0 to Z_X can be understood from the Hamiltonian perspective: $\mathcal{H} = L^2(S^1(R\sqrt{L}))$, and the n^{th} eigenvalue is

$$E_n = \left(\frac{n}{R\sqrt{L}} \right)^2 = \frac{l^2}{R^2 L},$$

and, as we computed before, the contribution of a_0 to Z_X is

$$\sum_n e^{-n^2 T^2 / R^2 L}.$$

Here n is the momentum of the center of mass around S^1 . A couple things stand out:

- The partition function, but not the action, is invariant under switching $R \leftrightarrow 1/R$. Physically, this means exchanging n (momentum) and w (winding number). In string theory, this is an instance of T -duality.
- The answer only depends on T/L , which says that this theory is actually a conformal field theory: it only depends on the conformal class of the metric on X . The action is also conformally invariant (before all of this Fourier expansion). This is related to the fact that the energies E_n of the states are proportional to $1/L$: the partition function is

$$Z = \sum_n e^{-T E_n},$$

where $E_n = c_n/L$. Therefore

$$Z = \sum e^{-c_n T / L},$$

which is invariant under rescalings $T \mapsto \lambda T$ and $L \mapsto \lambda L$ (in fact any conformal change of the metric). But be careful: if you take this to its logical extreme, you get something untrue: if you compute Z_X for a non-flat metric on the torus, you'll discover that it's not quite conformally invariant. For example, computing $1/\sqrt{\det \Delta}$ with ζ -regularization as usual produces something that transforms according to the

beautiful *Polyakov formula* under conformal transformations. This is an instance of something called the *conformal anomaly*.

Returning to T -duality, you might want to know why it's true. It looks like something that just happens to pop out of the partition function, but we can understand it in a way that makes it look less like a coincidence, and which shows that correlation functions are also exchanged under the duality. We will do this by producing some larger theory with more fields, such that when you integrate out some fields, you get the theory of maps $X \rightarrow S^1(R)$ and when you integrate out some other fields, you get the theory of maps $X \rightarrow S^1(1/R)$. The integration will be so simple to make the duality manifest.

Fix a Riemann surface X and let the space of fields \mathcal{C}_{big} be the space of pairs $\varphi : X \rightarrow S^1(2\pi)$ and $B \in \Omega^1(X)$. The action is

$$(17.1) \quad S_{\text{big}} = \frac{1}{2\pi} \int_X \frac{1}{2R^2} \|B\|^2 + \frac{i}{2\pi} \int B \wedge d\varphi.$$

Observe that the action doesn't depend on any derivatives of B , which says that B has no dynamics. It behaves like a Lagrange multiplier. The partition function is, as usual,

$$Z = \int_{\mathcal{C}_{\text{big}}} D\varphi DB e^{-S_{\text{big}}}.$$

We said we wanted to integrate something out, and there's an obvious choice: B . Completing the square in (17.1) produces

$$(17.2) \quad S_{\text{big}} = \frac{1}{2\pi} \int_X \frac{1}{2R^2} \|B - iR^2 \star d\varphi\|^2 + \frac{R^2}{4\pi} \int \|d\varphi\|^2.$$

If $\tilde{B} := B - iR^2 \star d\varphi$, then the first term in (17.2) is a Gaussian integral in \tilde{B} , and contributes some infinite-dimensional determinant of φ . So we can drop it and consider the second term, which has an effective action of

$$S_{\text{eff}} = \frac{R^2}{4\pi} \|d\varphi\|^2 = \frac{1}{4\pi} \int \|R\varphi\|^2,$$

which is (proportional to) the action for the theory with target $S^1(R)$.

Bereft of other options, we'll integrate out φ , and then something cool will happen. In (17.1), φ appears only linearly. Recall that, in a suitably distributional sense,

$$\int_{-\infty}^{\infty} dx e^{ixy} = \delta(y).$$

We'll need to use the discrete analogue

$$(17.3) \quad \sum_{n \in \mathbb{Z}} e^{2\pi i n a} = \sum_{m \in \mathbb{Z}} \delta(a - m).$$

That is, when $a \in \mathbb{Z}$, we pick up exactly one δ -function, and if $a \notin \mathbb{Z}$, we get 0, and this is what happens on the left: if $a \notin \mathbb{Z}$, it oscillates itself to death, and otherwise it accumulates.

We'd like to use Stokes' theorem to say that

$$\int D\varphi e^{(i/2\pi) \int_X B \wedge d\varphi} = \int D\varphi e^{(i/2\pi) \int_X dB \wedge \varphi},$$

but as φ is a map to S^1 , we can't literally do this: there are issues with exactness of forms. So we'll expand:

$$d\varphi = d\varphi_0 + \sum_{i=1}^{2g} 2\pi n_i \omega^i,$$

where $\varphi_0 : X \rightarrow \mathbb{R}$, g is the genus of X , $\{n_i\} \in \mathbb{Z}^{2g}$, and ω^i are representatives for a basis of $H^1(X; \mathbb{Z})$. In this case, to integrate over φ means to integrate over φ_0 and sum over $\{n_i\}$.

We can now say that

$$\frac{i}{2\pi} \int B \wedge d\varphi_0 = \frac{i}{2\pi} dB \wedge \varphi_0,$$

so integrating over φ_0 produces a δ -function

$$\int D\varphi_0 e^{(i/2\pi) \int dB \wedge \varphi_0} = \delta(dB),$$

which imposes the constraint that $dB = 0$.

Now, write

$$B = d\vartheta_0 + 2\pi \sum_{i=1}^{2g} a^i \omega_i,$$

where $\vartheta_0: X \rightarrow \mathbb{R}$, $a_i \in \mathbb{R}$, and the ω_i are differential 1-forms such that

$$\int \omega^i \wedge \omega_j = \delta_j^i$$

and $d\omega^i = 0$. (That is, ω^i is dual to ω_i). We plug this in to find the effective action is

$$\begin{aligned} S(\vartheta_0, a^i, n_j) &= \frac{1}{4\pi R^2} \int \|B\|^2 + i \int_X \left(d\vartheta_0 + \sum_i a^i \omega_i \right) \wedge \sum_j 2\pi n_j \omega^j \\ &= \frac{1}{4\pi R^2} \int \|B\|^2 + 2\pi i \sum a^i n_j \int_X \omega_j \omega^i \\ &= \frac{1}{4\pi R^2} \int \|B\|^2 + 2\pi i \sum a^i n_i. \end{aligned}$$

Now we must sum over $\{n_i\} \in \mathbb{Z}^{2g}$: using (17.3),

$$\sum_{\{n_i\}} e^{2\pi i \sum a^i n_i} = \sum_{k^i \in \mathbb{Z}} \delta(a^i - k^i).$$

This imposes the constraint that B has integer periods, hence that $B = d\tilde{\varphi}$ for some $\tilde{\varphi}: X \rightarrow S^1$. Hence the effective action is

$$S_{\text{eff}} = \frac{1}{4\pi R^2} \int \|d\tilde{\varphi}\|^2 = \frac{1}{4\pi} \int \|d(\tilde{\varphi}/R)\|^2,$$

which is (proportional to) the action for the $S^1(1/R)$ theory.

Why does this imply that the theories are dual? We found two different ways to compute the path integral of this theory: one as the path integral of the $S^1(R)$ theory, and the other as the path integral of the $S^1(1/R)$ theory. The same is true if you put operators into the path integral, providing a way of identifying correlation functions for one theory with different correlation functions for the other.

This was a formal argument, but the nonrigorous argument involved an infinite-dimensional Gaussian, which is something we've discussed before, and so it's probably possible to make this rigorous in a way that's not too terrible.

This instance of T -duality is a simple prototype for a very broad phenomenon of dualities between 2D σ -models, called *mirror symmetry*. This is a huge active area of research, and is based on a similar equivalence between two different σ -models with targets X and Y , where X and Y are “mirror” manifolds.

This T -duality is also a prototype of a phenomenon in 4D gauge theory, which is what we're interested in for this class: electric-magnetic duality. Though it exists in a more general setting, we're going to set it up in the Gaussian world, and begin discussing abelian gauge theory in 4D.

Fix a compact oriented 4-manifold X , and let the space of fields \mathcal{C} be the space of principal $U(1)$ -bundles $P \rightarrow X$ and connections ∇ on P . Let $F \in \Omega^2(X)$ denote the curvature of ∇ , and fix $g, \vartheta \in \mathbb{R}$. Then, the action is

$$S := \frac{1}{2g^2} \int F \wedge \star F + \frac{i\vartheta}{4\pi^2} \int F \wedge F.$$

Just to get oriented, let's figure out the classical equations of motion, which are the critical points of the action. Imagine a small deformation $\nabla \mapsto \nabla + \delta\alpha$, where $\delta\alpha \in \Omega^1(X)$. (This uses the fact that connections are an affine space modeled on $\Omega^1(X)$, so we can add a 1-form to a connection.)

The second term is topological, so its variation will be zero, and we get

$$\delta S = \frac{1}{2g^2} \int_X d(\delta\alpha) \wedge \star F + F \wedge \star d(\delta\alpha) + \frac{2i\vartheta}{4\pi^2} \int d(\delta\alpha) \wedge F.$$

$dF = 0$ by the Bianchi identity and so the second term vanishes. Then, we integrate by parts to simplify the first term:

$$= \frac{1}{g^2} \int_X (\delta\alpha) \wedge d(\star F),$$

so this vanishes if $d \star F = 0$, and this is precisely the equations of motion.

Remark 17.4. This is the succinct version of Maxwell's equations! $dF = 0$ and $d \star F = 0$ are (the source-free versions of) Maxwell's equations, in that you can expand them out in terms of components \mathbf{E} and \mathbf{B} of F and recover the familiar-looking equations for electromagnetism that you learned in high school.⁵⁶ ◀

Decomposing F into its self-dual and anti-self-dual components

$$F_{\pm} := \frac{1}{2}(F \pm \star F)$$

and letting

$$\tau := \frac{\vartheta}{\pi} + i \frac{2\pi}{g^2},$$

the action is

$$S = \frac{i\bar{\tau}}{4\pi} \int_X \|F_+\|^2 - \frac{i\tau}{4\pi} \int_X \|F_-\|^2.$$

This makes the duality phenomenon more manifest. In 2D, we had the duality $R \leftrightarrow 1/R$ switching $d\varphi$ and $\star d\varphi$, and in 4D, we have $T \leftrightarrow -1/\tau$, which exchanges F and $\star F$. In a similar way to our discussion of T -duality, one can compute correlation functions of curvatures of connections and duality implies

$$\langle F_{12}(x_1) F_{34}(x_2) \cdots F_{13}(x_n) \rangle = \langle (\star F)_{12}(x_1) (\star F)(x_2) \cdots \rangle.$$

Remark 17.5. You might wonder why we're using $U(1)$ -connections instead of just 1-forms (or \mathbb{R} -connections). This is because of the way this theory will arise for us: we will see Seiberg-Witten theory as a low-energy theory for an $SU(2)$ -theory, and this theory has no noncompact analogue. We will also want magnetic monopoles in our theory, which again comes from the $SU(2)$ -theory, and these are present in the $U(1)$ -theory but not the \mathbb{R} -theory. ◀

There's another symmetry, sending $\tau \mapsto \tau + 2$, and so the symmetry group is the infinite subgroup of the unitary group generated by $(+2)$ and $\tau \mapsto -1/\tau$: this theory has infinitely many descriptions!

Lecture 18.

U(1) gauge theory in 4D: 11/2/17

We've just started discussing gauge theory in 4D and how it relates to electromagnetism. The action of *pure* $U(1)$ gauge theory or *quantum Maxwell theory*⁵⁷ is

$$\begin{aligned} S &= \frac{1}{2g^2} \int_X F \wedge \star F + \frac{i\vartheta}{4\pi^2} \int F \wedge F \\ &= \frac{i\bar{\tau}}{4\pi} \int_X \|F_+\|^2 - \frac{i\tau}{4\pi} \int_X \|F_-\|^2, \end{aligned}$$

where $g, \vartheta \in \mathbb{R}$ and $\tau := \vartheta/\pi + 2\pi i/g^2$. The space of fields \mathcal{C} is the space of $U(1)$ -bundles on X with connection, and $F \in \Omega^2(X)$ is the curvature of the connection.

The first term in the action can be formulated in any dimension, but the second part really requires $\dim X = 4$.

⁵⁶That said, I did not learn these in high school.

⁵⁷The quantum theory of electromagnetism, quantum electrodynamics, is this theory plus a term corresponding to an electron.

If you solve for the classical equations of motion (the extremizers of the action), you get the equations for a harmonic form, which were first mathematically studied by Hodge:

$$(18.1) \quad \begin{aligned} dF &= 0 \\ d\star F &= 0. \end{aligned}$$

Since F is a 2-form, we can (locally) write it as a skew-symmetric matrix in coordinates (t, x^1, x^2, x^3) :

$$(18.2) \quad F = \begin{bmatrix} 0 & iE_1 & iE_2 & iE_3 \\ -iE_1 & 0 & B_3 & -B_2 \\ -iE_2 & -B_3 & 0 & B_1 \\ -iE_3 & B_2 & -B_1 & 0 \end{bmatrix},$$

where E_i is the i^{th} component of the electric field and B_i is the i^{th} component of the magnetic field. With this notation, (18.1) reduces to Maxwell's equations (specifically as formulated by Heaviside):

$$(18.3a) \quad \operatorname{div} \mathbf{E} = 0$$

$$(18.3b) \quad \operatorname{div} \mathbf{B} = 0$$

$$(18.3c) \quad i \operatorname{curl} \mathbf{E} = -\frac{d\mathbf{B}}{dt}$$

$$(18.3d) \quad \operatorname{curl} \mathbf{B} = i \frac{d\mathbf{E}}{dt}.$$

Hodge theory is interesting for entirely mathematical reasons, and it's cool that electromagnetism comes out of it.

Remark 18.4.

- (1) The theory with coupling τ is equivalent to the theory with coupling $\tau + 2$, i.e. $\vartheta \mapsto \vartheta + 2\pi$. This is relatively straightforward: e^{-S} is invariant under this shift. This is because

$$\int_X F \wedge F \in 4\pi^2 \mathbb{Z},$$

because $[F/2\pi] = c_1(X) \in H^2(X; \mathbb{Z})$.

- (2) There's a harder-to-show equivalence $\tau \leftrightarrow -1/\tau$, whose derivation resembles the path-integral argument we gave for T -duality of the free boson theory in 2D. This equivalence appears at the quantum level, meaning e^{-S} is not invariant.

◀

These symmetries generate the group $\Gamma_0(2) := \langle \tau \mapsto \tau+2, \tau \mapsto -1/\tau \rangle \subset \operatorname{SL}_2(\mathbb{Z})$,⁵⁸ so this theory is parameterized by $\tau \in \mathbb{H}/\Gamma_0(2)$. This space is the moduli space of these QFTs (abelian quantum Maxwell theories in dimension 4). Mathematically, $\mathcal{H}/\Gamma_0(2)$ is a torus $T_\tau^2 := \mathbb{C}/\mathbb{Z} \oplus \tau\mathbb{Z}$, so we're led to wonder whether there is a torus around.

Remark 18.5. It turns out the answer is yes! There is a six-dimensional QFT \mathcal{X} such that the low-energy limit of \mathcal{X} on $T_\tau^2 \times X$ is (a supersymmetric extension of) the Maxwell theory on X with coupling τ .⁵⁹ Theory \mathcal{X} comes in abelian and nonabelian versions — this one is cool and relatively accessible, and the nonabelian one is cooler. ◀

Coupling to matter. In classical electromagnetism, an electric charge/current is a 1-form

$$j = \begin{bmatrix} i\rho \\ j_1 \\ j_2 \\ j_3 \end{bmatrix} \in \Omega^1(\mathbb{R}^4)$$

such that $d\star j = 0$ (which is called *conservation of current*). Here ρ is the charge density and \mathbf{j} is the current. This modifies the Maxwell equations (18.1) to

$$(18.6) \quad \begin{aligned} dF &= 0 \\ d\star F &= ig^2 \star j. \end{aligned}$$

⁵⁸TODO: check notation.

⁵⁹There's a little more structure needed to define theory \mathcal{X} than a Riemannian metric, but the dimensional reduction story ends up working out.

We'd like to generalize our QFT to a theory of electromagnetic fields interacting with this fixed j . This entails adding a term to the action which looks locally like

$$(18.7) \quad S_{\text{curr}} := i \int A \wedge \star j.$$

Here, $A \in \Omega^1(X)$ should be a local representative of the connection ∇ , i.e. in coordinates, $\nabla = d + A$.

Exercise 18.8. Show that if P is trivial, this is possible, and S_{curr} is independent of the choice of global trivialization of P . (Hint: this uses conservation of current.)

If P is not trivial, then it's more elaborate. We'll return to this point.

Gauge invariance. The space of fields $\mathcal{C} = \{(P, \nabla)\}$, where P is a principal $U(1)$ -bundle and ∇ is a connection on P , is way too big to make sense of! We're going to attempt to quotient by some equivalence relation to make this more concrete.

First, fix one $U(1)$ -bundle P_i in each isomorphism class of $U(1)$ -bundles; these are a discrete set. Then, we'll integrate over all connections on P_i . But there are still automorphisms of P , and we want two connections on P_i to be equivalent if there's an automorphism carrying one to the other. Locally, the automorphism group is $U(1)$, so the group of automorphisms is $\mathcal{G} = \text{Map}(X, S^1)$: for any $\chi : X \rightarrow \mathbb{R}/\mathbb{Z}$, we produce the automorphism

$$\nabla \mapsto \nabla + d\chi.$$

We need to check that the action is invariant under this automorphism, but since the action was built only from intrinsic data, the original action is fine, and we can integrate over the smaller quotient space. However, (18.7) adds a term depending on choices, so we really have to check whether the action is gauge-invariant after adding S_{curr} .

Suppose that χ lifts to a map $X \rightarrow \mathbb{R}$. Then

$$\begin{aligned} \int_X A \wedge \star j &= \int_X (A + d\chi) \wedge \star j \\ &= \int_X A \wedge \star j + \int_X \chi \wedge d\star j \\ &= \int_X A \wedge \star j, \end{aligned}$$

so we're fine in this case.

Exercise 18.9. For a more general $\chi : X \rightarrow \mathbb{R}/\mathbb{Z}$, show that gauge invariance holds iff $\star j \in H^3(X; \mathbb{Z})$, i.e. it's *integrally quantized*.

This is called the *quantization of charge*, and physically means that the current has to track charged particles. For example, pick a 1-manifold $\gamma \subset X$, representing the worldline of a charged particle. Then, its dual current δ_γ is a distribution-valued 3-form, in that if C is a 3-chain,

$$\int_C \delta_\gamma = \gamma \cap C.$$

If j is integrally quantized, then $\star j$ must be a finite \mathbb{Z} -linear combination of these δ_γ , hence coming from a finite set of particles with integer charges.

Remark 18.10. If P is not trivial, writing the coupling to j is considerably harder; when $\star j$ is exact, there's a proof using partitions of unity, but in the general case it's not clear what happens. ◀

Let $\gamma \subset X$ be a loop. If $\star j = k\delta_\gamma$ ($k \in \mathbb{Z}$), we can think of adding S_{curr} as multiplying e^{-S} by

$$\exp\left(ik \int_X A \wedge \star j\right) = \exp\left(ik \oint_\gamma A\right) = \text{Hol}_{\gamma, V_k}(\nabla).$$

That is, if V_k is the irreducible representation with weight k (which is 1-dimensional), this last term is the holonomy of the connection around γ , computed in the representation V_k .

This is the same thing as inserting the observable $\text{Hol}_{\gamma, V_k}(\nabla)$ into the path integral, which is called a *Wilson line* in the representation V_k , which represents the track of a charged particle with charge k . This is a little elementary for an abelian gauge group, but is extremely useful in the nonabelian case, where the intuition of Wilson lines as paths of charged particles “charged” by representations is useful.

Coupling to dynamical fields. Let’s expand the already-huge space of fields to $\mathcal{C} = \{(P, \nabla, \varphi)\}$, where P and ∇ are as before, and φ is a section of the *associated bundle* $E_k := P \times_{U(1)} V_k$. As before, V_k is the irreducible representation of $U(1)$ with weight k . Thus E_k is a complex vector bundle with a Hermitian metric induced from the one on V_k (which is unitary), and ∇ induces a connection D on E_k such that for any $f : X \rightarrow \mathbb{C}$,

$$Df = (df + ikAf) \in \Omega_X^1(\mathbb{C}).$$

The action has some new terms: if S_{gauge} represents the first action we wrote down,

$$\begin{aligned} S &= S_{\text{gauge}} + \frac{1}{2} \int_X \|D\varphi\|^2 \, \text{dvol}_X + \frac{1}{2} \int_X m^2 |\varphi|^2 \, \text{dvol}_X \\ &= S_{\text{gauge}} + \frac{1}{2} \int_X \|d\varphi + ikA\varphi\|^2 + m^2 \|\varphi\|^2 \\ &= S_{\text{gauge}} + \int_X \eta^{ij} (\partial_i \varphi \partial_j \bar{\varphi} + 2ik \partial_i \bar{\varphi} A_j \varphi - 2ik \partial_j \bar{\varphi} A_i \varphi + k^2 A_i A_j |\varphi|^2) + n^2 |\varphi|^2 \, \text{dvol}_X. \end{aligned}$$

Here η^{ij} is the (matrix inverse of the) metric on X . In this case, the equations of motion are now

$$\begin{aligned} d\star F &= g i^2 k \underbrace{(\bar{\varphi} D\varphi - \varphi D\bar{\varphi})}_{\star j} \\ D\star D\varphi &= m^2 \varphi. \end{aligned}$$

This theory is almost quantum electrodynamics; it’s just missing a fermion. Hence it’s sometimes called *scalar QED*, the theory of an electromagnetic field coupled to a massive, electrically charged field.

Remark 18.11. What if you want magnetically charged particles (magnetic monopoles)? There are some very interesting quantum field theories that have them, but nobody knows how to write down an action which also includes magnetically charged particles. This is special to dimension 4, for which both electric and magnetic charges can be carried by particles. ◀

Of course, we’d like to be able to actually calculate stuff. Let’s do some perturbation theory, expanding around $g = 0$. First note that

$$\begin{aligned} 0 &\leq \frac{1}{2} \int \|F \pm \star F\|^2 = \frac{1}{2} \int (F \pm \star F) \wedge \star (F \pm \star F) \\ &= \pm \int F \wedge F + \int F \wedge \star F, \end{aligned}$$

so

$$\int F \wedge \star F \geq \left| \int F \wedge F \right|.$$

Thus the path integral sector with $\int F \wedge F = 4\pi^2 n$ (n is called the *instanton number*, and these sectors are called *instanton sectors*) is bounded below by $2\pi^2 |n|/g$. In other words,

$$e^{-S} \leq e^{-2\pi^2 |n|/g^2}.$$

Something funny just happened: the right-hand side is a classic counterexample in analysis, as it’s smooth but not analytic. It was constructed solely for a counterexample in that setting, so it’s a surprise to see it in physics. And this fact has consequences for us: if you expand around $g = 0$, all terms in the Taylor series vanish. This means that the contributions from sectors with $n \neq 0$ are invisible in this series expansion around $g = 0$. Hence one says that instantons are a nonperturbative effect. They have physical relevance, and are crucial in Donaldson theory, but can’t be seen by perturbation theory.

Another interesting aspect of this expansion is that g behaves like a coupling, as

$$S = \frac{1}{g^2} \|F\|^2 + \cdots + \partial \bar{\varphi} A \varphi + \cdots,$$

and under $A \mapsto gA$,

$$S = \|F\|^2 + \cdots + g \partial \bar{\varphi} A \varphi + \cdots.$$

Next time, we'll talk about why interacting field theories exist, and when they appear, and so on.

Lecture 19.

UV cutoffs in a 4D σ -model: 11/7/17

Last time, we discussed U(1)-gauge theory in 4D with a charged scalar, whose action is

$$\begin{aligned} S(A, \phi) &:= \frac{1}{g^2} \int_X \|F\|^2 + \|D\phi\|^2 + \frac{i\theta}{2\pi} \int F \wedge F \\ &= (\text{quadratic}) + kA\phi \, d\phi + kA^2\phi^2. \end{aligned}$$

The statement that ϕ is charged means that it's transforming in a nontrivial representation of U(1); in this case, the representation is V_k , the one-dimensional unitary U(1)-representation given by $e^{i\theta} \mapsto e^{ik\theta}$.

Another way to think of the notion that ϕ is charged is that it interacts in a nontrivial way with the gauge field because of the $\|D\phi\|^2$ term in the action. Since $D\phi = d\phi + kA\phi$, this is the same thing as the other notion of charge.

The action is not just quadratic, which means that everything we compute will involve interactions — we can't just do a bunch of Gaussian integrals like we've been doing before. To warm up, we're going to look at a simple toy model of a 4D σ -model with interactions. Specifically, let

$$\mathcal{C}_X := \text{Map}(X, \mathbb{R})$$

with action

$$S(\phi) = \int_X \text{dvol}_X \left(\frac{1}{2} \|d\phi\|^2 + \frac{m^2}{2} \phi^2 + \frac{\lambda}{4!} \phi^4 \right).$$

Remark 19.1. If $\dim X = 1$, this is the *anharmonic oscillator*, a commonly studied variant of the harmonic oscillator. ◀

For $X = \mathbb{R}^4$ and $x_1, x_2 \in \mathbb{R}^4$, let's compute the normalized two-point correlation function $\langle \phi(x_1)\phi(x_2) \rangle / Z$. There are a few places that divergences could arise, including the infrared divergence we grappled with in dimension 2.

The rules of perturbation theory say that up to order λ (which will suffice to see the problem), there are only two Feynman diagrams to worry about: the line x_1 to x_2 , and that line with an x' loop in it.

To the first diagram, we associate the function $D(x, y)$, the Green's function for the operator $(-\Delta + m^2)$. That is, it's defined to satisfy

$$(19.2) \quad (-\Delta_x + m^2)D(x, y) = \delta(x, y).$$

Remark 19.3. For $X = \mathbb{R}$ and $m = 0$, $D(x, y) = (1/2)|x - y|$. So it's continuous, but not smooth. For $d \neq 2$,

$$D(x, y) = \frac{1}{\|x - y\|^2},^{60}$$

(and when $d = 2$ there's a log term). Differences in the behavior of Green's functions in different dimensions is a big reason quantum field theories behave so differently in different dimensions. ◀

Therefore the expansion of this two-point function, up to order λ , is

$$D(x_1, x_2) + \frac{\lambda}{2} \int_{\mathbb{R}^4} D(x_1, x') D(x_2, x') D(x', x') \, dx' + O(\lambda).$$

Here we have a problem: $D(x', x')$ is $1/0$, and there doesn't seem to be a good way to avoid this with this approach.

⁶⁰Possibly up to a constant.

Let's try something different: in dimension 1, we made an argument rigorous by putting it on a lattice, so let's try that.⁶¹ This changes the space of fields: if the lattice has edge length L and $\Lambda := 1/L$, then the Fourier transform of any field ϕ on the lattice has a cutoff: $\widehat{\phi}(p) = 0$ when $\|p\| > \Lambda$; this is called a *UV cutoff*. Heuristically, this is telling us that we're cutting out the high-momentum pieces.

This suggests defining a *cutoff theory* for a $\Lambda \in \mathbb{R}_+$, whose space of fields is

$$\mathcal{C}_\Lambda = \{\phi : \mathbb{R}^4 \rightarrow \mathbb{R} \mid \text{supp}(\widehat{\phi}) \subset B_\Lambda(0)\}.$$

This changes the Green's function.

Exercise 19.4. Show the Fourier-space formula for the original Green's function $D(x, y)$:

$$D(x, y) = \int_{(\mathbb{R}^4)^*} dp \frac{e^{ip(x-y)}}{\|p\|^2 + m^2}.^{62}$$

This makes it clear where the divergence in $D(x', x')$ is coming from: it is

$$D(x, x) = \int_{(\mathbb{R}^4)^*} dp \frac{1}{\|p\|^2 + m^2},$$

so letting $r = \|p\|$,

$$\sim \int dr \frac{r^3}{r^2 + m^2}.$$

So the divergence of $D(x', x')$ comes from the region of large p , suggesting that the UV cutoff might fix this. And indeed, the Green's function for the cutoff theory is

$$(19.5) \quad D_\Lambda(x, y) = \int_{\|p\| < \Lambda} \frac{e^{ip(x-y)}}{\|p\|^2 + m^2},$$

so $D_\Lambda(x', x')$ is finite.

Replacing D with D_Λ in our calculation, we find that the asymptotic series expansion is

$$\langle \phi(x_1)\phi(x_2) \rangle = D(x_1, x_2) + (\lambda \cdot \Lambda^2 + \Lambda^4 \|x_1 - x_2\|^2 + \Lambda^6 \|x_1 - x_2\|^4 + \dots),$$

(possibly with some other terms).

This is a strange perturbation series: if $\Lambda \gg 1/\|x_1 - x_2\|$ (so the lattice distance is small), the second term can be much larger than the first term, especially when x_1 and x_2 are far apart, even if λ is small! Therefore in this regime, the expansion is not well-behaved. This theory is not very useful for calculating.

What should we do about this? Let's try to formulate our computations with an effective action $S_{\text{eff}}(\Lambda')$, where the cutoff is $\Lambda' \approx E$, smaller than Λ , with the hope that it's more useful for calculations. To obtain this action from our original one, we need to integrate out all of the modes $\widehat{\phi}(p)$ for $\Lambda' < \|p\| < \Lambda$.

The general expectation is that $S_{\text{eff}}(\Lambda')$ will be non-local: we'll obtain infinitely many terms if we expand it in powers of fields and their derivatives. This is about the place in the movie where ominous music starts playing — it's not clear how to use this effective theory to compute meaningful information.

Remark 19.6. The fact that the action is invariant under $\phi \mapsto -\phi$ is nice: it means that, for example, we can ignore terms like $\phi^{10} \|d\phi\|^{10}$. This will be a slight boon to us. ◀

One convenient way to investigate this is to make use of the idea that taking the effective action should be a continuous process. We'll define an infinite-dimensional space \mathcal{A} of all possible (already cut off) actions $S(\phi)$ invariant under $\phi \mapsto -\phi$. We will then define a flow F on \mathcal{A} , where $F_t(S)$ is the effective action obtained by integrating out the modes for which $\Lambda e^{-t} < \|p\| < \Lambda$. As t varies, this defines trajectories in \mathcal{A} . This can be made rigorous; the canonical reference is Polchinski [24].

Theorem 19.7. As $t \rightarrow \infty$, this flow is driven to a three-dimensional space. For this specific theory, the parameters are ϕ^2 , $|d\phi|^2$, and ϕ^4 .

⁶¹Today's perspective on QFT regards the lattice as pretty natural: rather than a trick used to avoid infinities, the lattice is really informing you about short-range interactions in physics.

⁶²Something about this formula is related to the fact that the Laplacian is elliptic, but I didn't see why.

This is an amazing result: even though \mathcal{A} is infinite-dimensional, we only need to worry about three parameters.

The reason is not really about quantum field theory, but a more general fact about scaling phenomena. There's an action ρ of \mathbb{R}^\times on $\mathcal{C}_X = \{\phi : \mathbb{R}^4 \rightarrow \mathbb{R}\}$ sending

$$\rho_\varepsilon : \phi(x) \mapsto \frac{1}{\varepsilon} \phi(\varepsilon x),$$

and this leaves the term

$$(19.8) \quad \frac{1}{2} \int \|\mathrm{d}\phi\|^2$$

invariant.

Remark 19.9. This says that the *free theory*, whose action is (19.8) and no other terms, is (classically) conformally invariant. There is a question about anomalies for the quantum theory, though. ◀

The \mathbb{R}^\times -action transforms the other terms:

$$\begin{aligned} \rho_\varepsilon^* \left(\int \mathrm{d}x \phi^n \right) &= \varepsilon^{4-n} \int \phi^n \\ \rho_\varepsilon^* \left(\int \|\mathrm{d}\phi\|^m \phi^n \right) &= \varepsilon^{4-(2m+n)} \int \|\mathrm{d}\phi\|^m \phi^n. \end{aligned}$$

Definition 19.10. Motivated by this, define the *scaling dimension* of the coupling to be $\dim \phi^n := n$ and $\dim \|\mathrm{d}\phi\|^m \phi^n := 2m + n$.

Formally, in the theory without the cutoff, the correlation function

$$\frac{1}{\varepsilon^2} \langle \phi(0) \phi(\varepsilon x) \rangle$$

computed with S is equal to $\langle \phi(0) \phi(x) \rangle$ computed with $\rho_\varepsilon^*(S)$.

Hence, at least formally, for large ε (i.e. over long distances), the effect of the terms with scaling dimension greater than 4 is very small. We call such terms *irrelevant*.

Lecture 20.

$\mathcal{N} = 2$ supersymmetric Yang-Mills theory: 11/9/17

Last time, we studied a scalar field theory in dimension 4. To avoid divergences we imposed a cutoff Λ , only integrating over fields ϕ with $\widehat{\phi}(p) = 0$ for $\|p\| > \Lambda$. For such ϕ , the action is

$$S(\phi) := \int \mathrm{d}^4x \|\mathrm{d}\phi\|^2 + \frac{m^2 \phi^2}{2} + \frac{\lambda \phi^4}{4!}.$$

Among the simplest thing you can compute is a normalized two-point correlator $\langle \phi(x_1) \phi(x_2) \rangle / Z$. The first term contributes a factor of $D_\Lambda(x_1, x_2)$, where D_Λ is defined in (19.5), and the second term (a line from x_1 to x_2 plus a loop at x' in the middle) contributes a factor of

$$\int \mathrm{d}^4x' D_\Lambda(x', x_1) D_\Lambda(x', x_2) D_\Lambda(x', x').$$

Together, these are asymptotically approximately

$$\frac{\langle \phi(x_1) \phi(x_2) \rangle}{Z} \sim \left(\frac{1}{\|x_1 - x_2\|^2} + \dots \right) + \lambda(\Lambda^2 + \dots).$$

Now, if $\lambda^{1/2} \Lambda \gg E := \|x_1 - x_2\|^{-1}$, the subleading term is greater than the leading term. This is the issue mentioned last time: it suggests we might need to worry about higher-order terms in this regime.

Remark 20.1. One concern with this approach to the 4D scalar theory is that it makes it much harder to define on more general manifolds — it crucially uses Fourier theory on \mathbb{R}^4 . In principle, something should work: we're trying to cut off long-range interactions, and locally every 4-manifold is \mathbb{R}^4 , but it will require more setup than in this case. ◀

Today we're going to discuss the motivating ideas behind this course, enriched by all that we've learned so far.

Everything we've been learning is with the goal of studying a specific 4D QFT, $\mathcal{N} = 2$ supersymmetric Yang-Mills theory with gauge group $SU(2)$. Fix a 4-manifold X with a spin structure, and let S be the spinor bundle. The fields are:

- (P, ∇) , a principal $SU(2)$ -bundle on spacetime X together with a connection ∇ ,
- ϕ , a section of $P \times_G \mathfrak{g}_{\mathbb{C}} = (\text{ad } P)_{\mathbb{C}}$, and
- an odd field ψ , a section of $\Pi((P \times_G \mathfrak{g}_{\mathbb{C}}) \otimes S \otimes R)$, where R is a two-dimensional vector space.⁶³

Let \mathcal{C}_X denote the space of such fields.

The action contains a few terms. The first term is the usual action for Yang-Mills gauge theory, and the term corresponding to ϕ allows the theory to have charged particles, just as in the $U(1)$ -case we considered:

$$S = \underbrace{\frac{1}{g^2} \int_X \text{tr} \|F\|^2}_{YM} + \frac{i\theta}{2\pi} \int_X \text{tr} F \wedge F + \frac{1}{2} \int_X \|D\phi\|^2 + (\text{a term in } \psi).$$

Here YM is called the *Yang-Mills term* for this theory, and $D\phi$ is the covariant derivative of ϕ induced from ∇ . For the fermionic term, see the professor's notes.

This complicated-looking action is nonetheless one of the simplest $\mathcal{N} = 2$ supersymmetric Yang-Mills actions physicists consider in 4D. It also has a key property: when $X = \mathbb{R}^4$, \mathcal{C}_X admits odd vector fields Q_i with $Q_i S = 0$. In lower dimensions, this foreshadowed the theory's ability to recover topological information. It's hard to write down supersymmetric theories in higher dimensions, and this suggests where this action comes from: physicists use ordinary Yang-Mills theory to describe nature, and then one wants to add terms to make it supersymmetric.⁶⁴

There are eight such odd vector fields Q_i .⁶⁵ So why is this not an $\mathcal{N} = 8$ theory? The action is clearly invariant under the *Poincaré algebra* $\mathfrak{iso}(4)$, but these symmetries get extended to a super-Lie algebra

$$\tilde{\mathfrak{g}} = \underbrace{\mathfrak{iso}(4)}_{\tilde{\mathfrak{g}}_0} \oplus \tilde{\mathfrak{g}}_1,$$

and $\tilde{\mathfrak{g}}_1$ is an 8-dimensional vector space. But through the adjoint action, it acquires an action of $\tilde{\mathfrak{g}}_0$, and (at least after complexification), $\tilde{\mathfrak{g}}_1$ is a direct sum of two copies of the spinor representation $S^+ \oplus S^-$.

It's *believed* that there exists a quantum field theory defined by this action, meaning that if you place it on a lattice and let the lattice length tend to zero, there's a sensible limit. Equivalently, we would like to define the theory with a cutoff, and let the cutoff go to infinity. This is *not* true for every action, e.g. it's false for the action

$$S = \frac{1}{2} \int \|\text{d}\phi\|^2 + m^2 \phi^2 + \lambda \phi^4$$

in this dimension. This theory isn't chiral, and therefore fermion doubling isn't an issue.

Witten had the idea to study this theory on compact Riemannian manifolds X with a spin structure. This kills the $\mathfrak{iso}(4)$ - and $\tilde{\mathfrak{g}}_1$ -actions, but Witten found a way to modify the theory in such a way that, even on an arbitrary manifold, there's a single odd vector field Q on \mathcal{C}_X with $QS = 0$.

This is called a *twist* of the theory. Recall that ψ is a section of $\Pi((\text{ad } P) \otimes S \otimes R)$; a twisting of the theory replaces R with a vector bundle over X . Specifically, we choose $R = S^+$, so ψ is a section of $\Pi((\text{ad } P) \otimes S \otimes S^+)$. The kernel of $\text{Spin}(n) \rightarrow \text{SO}(n)$ acts nontrivially on S , but acts trivially on $S \otimes S^+$, which means this theory does not require a spin structure!

Since $S \otimes S^+$ is canonically trivial on \mathbb{R}^4 , this is the same theory on \mathbb{R}^4 . But on a compact 4-manifold, it has a single odd symmetry!

Subsequently, there's a similar story as in dimensions 1 and 2: the odd symmetry implies that Z and the correlation functions (of Q -closed operators) are invariant under changes $S \mapsto S + Q(\psi)$. Hence, they must be invariant of the metric on X : the computation looks at the variation $\frac{\delta S}{\delta g_{\mu\nu}}$ and see that it's of the form $Q(\psi)$. This is

⁶³This copy of R is often thought of as the defining representation of another copy of $SU(2)$, denoted $SU(2)_R$, and this $SU(2)_R$ -action is called *R-symmetry*.

⁶⁴It would be very interesting if $\mathcal{N} = 2$ supersymmetric Yang-Mills were to describe our universe, but at low energies it looks completely different, so it does not seem to. So our justification for it is not phenomenological. There are other justifications, e.g. wanting to write down the most general action possible reflecting all symmetries present. It's also believed this is the unique supersymmetric theory with these terms in this dimension.

⁶⁵Producing four is not so hard, but eight is difficult.

a totally local calculation, though it's important that X is compact: the independence is up to bounding terms for the initial boundary problem on \mathcal{C}_X .

Thus, you should get a *topological* invariant Z of compact 4-manifolds X , and for every list of Q -closed operators $\mathcal{O}_1, \dots, \mathcal{O}_n$, the correlation function $\langle \mathcal{O}_1 \cdots \mathcal{O}_n \rangle$.⁶⁶

Which \mathcal{O}_i are allowed? If we look at local operators $\mathcal{O}^{(0)}(x)$, then we impose $\mathcal{O}^{(0)}(x) = \text{tr}(\phi(x))^2$. This theory also has Q -invariant nonlocal operators. There are also operators $\mathcal{O}^{(1)}(\gamma)$ associated to a loop σ , $\mathcal{O}^{(2)}(\Sigma)$ associated to a surface Σ , and so on, and these have explicit formulas.⁶⁷

Thus our topological invariant is actually for 4-manifolds together with inserted closed 0-manifolds, 1-manifolds, 2-manifolds, etc. Witten then claims these are equal to the invariants recently introduced by Donaldson. This involves computing via localization, so the path integral is reduced to an integral over a much smaller space, the fixed locus of Q .

This is a fancier version of what we did in lower dimension: using the fact that the path integral is invariant under Q -exact deformations, cleverly choose a deformation $S \mapsto S + tQ\psi$, and take the limit as $t \rightarrow \infty$. This replaces the integral with a finite-dimensional integral. In supersymmetric quantum mechanics, this was an integral over the space of constant loops, and in this case, it's an integral over the *instanton moduli space*, the space of $SU(2)$ -connections ∇ with zero self-dual part.

This is cool, but not useful for computations — that requires the low-energy description, which we'll talk about next time.

Lecture 21.

Twists of $\mathcal{N} = 2$ supersymmetric Yang-Mills theory: 11/14/17

Lecture 22.

Localization to the moduli space of instantons: 11/21/17

A lot of people (myself included) weren't here last time, when we discussed twisted $\mathcal{N} = 2$ supersymmetric Yang-Mills theory, and we'll briefly review it.

We're going to fix the gauge group $G = SU(2)$; some of what we say today, including the action, makes sense more generally, but other things do not. We formulate the theory on a Riemannian 4-manifold X , with a parameter $g \in \mathbb{R}$.

The space of fields \mathcal{C}_X is the space of tuples of the following data: the bosonic fields

- (P, ∇) , where $P \rightarrow X$ is a principal $SU(2)$ -bundle and ∇ is a connection for it,
- a $\phi \in \Gamma((\text{ad } P)_{\mathbb{C}})$, and
- $D \in \Gamma((\text{ad } P)_{\mathbb{C}} \otimes \Lambda_+^2 T^*X)$,

and the fermionic fields

- $\psi \in \Pi\Gamma((\text{ad } P)_{\mathbb{C}} \otimes T^*X)$,
- $\eta \in \Pi\Gamma((\text{ad } P)_{\mathbb{C}})$, and
- $\chi \in \Pi\Gamma((\text{ad } P)_{\mathbb{C}} \otimes \Lambda_+^2 T^*X)$.

The action is a complicated-looking integral, and there is an odd vector field Q on \mathcal{C}_X such that $QS = 0$. It'll be important later that S is not positive definite.

There are various invariant local operators associated to points, lines, surfaces, etc: for example, if $x \in X$, $\mathcal{O}^{(0)}(x) = \text{tr}(\phi(x)^2)$. These satisfy a key relation: if C is a k -chain,

$$(22.1) \quad Q(\mathcal{O}^{(k)}(C)) = \mathcal{O}^{(k-1)}(\partial C).$$

For example, if γ is a path from x to y , then $\phi(x) - \phi(y)$ is Q of something, or, in Q -homology, $[\phi(x)] = [\phi(y)]$. This means that supersymmetric correlation functions for these operators don't depend on their position. More generally, if C is a k -chain, $[\mathcal{O}^{(k)}(C)]$ only depends on the homology class of C .

⁶⁶So is this a topological field theory? It depends what you mean exactly by that: this was the first QFT considered where the correlation functions were distance-independent, and indeed Witten's paper [28] is called "Topological Quantum Field Theory." But in the modern, mathematical perspective on TQFT, this does not precisely satisfy the axioms, though certainly it's a lot closer than anything else we've seen so far in this class.

⁶⁷More generally, there are operators associated to any k -chain, but Q -invariance forces us to consider only those arising from cycles.

In order to compute partition and correlation functions, we're going to perturb the action by adding Q of something. In this case, the particular perturbation is very simple — $S = Q\Psi$ for some Ψ already, and therefore we can shift S by a multiple of itself! Or, said differently, we can rescale S by taking $g \rightarrow 0$. This says in particular that the correlation functions are independent of g .

Remark 22.2. Is this a topological field theory? Different people in physics mean different things by topological field theory, but in this case it's independent of the metric — and this follows formally because the action is Q -exact for any metric, and therefore Q -exact deformations of the metric do not change it at all. \blacktriangleleft

As before, we're going to compute the partition functions and correlation functions by saddle point integration around the minima of S . We got confused in class about where the fixed points are (**TODO**), but at these fixed points, $F^+ = D$, $\nabla\phi = 0$, and $[\phi, \bar{\phi}] = 0$. The equations of motion $\delta S = 0$ further implies that $D = 0$.

We'd like to reduce these equations to $F^+ = 0$ and $\phi = 0$, because then integrating over the minima would be an integral over the instanton moduli space $\{(P, \nabla) \mid F^+ = 0\} / \sim$, where \sim denotes gauge equivalence.

Proposition 22.3. *If $b_2^+ \neq 0$, $F^+ = 0$, $\nabla\phi = 0$, and $[\phi, \bar{\phi}] = 0$, then $\phi = 0$ for a generic Riemannian metric.*

Proof. Suppose for the sake of contradiction there exists some $\phi \neq 0$, but such that $\nabla\phi = 0$ and $[\phi, \bar{\phi}] = 0$. Since ϕ is an infinitesimal automorphism of P , it generates a one-parameter group of automorphisms of (P, ∇) . If V denotes the defining representation of $SU(2)$, then we can decompose the associated vector bundle $E := P \times_{SU(2)} V$ under the action of this group.

Because $[\phi, \bar{\phi}] = 0$, then (the matrix of) ϕ is actually diagonalizable, and therefore E decomposes into two line bundles; since ϕ is covariantly constant, the connection also splits:

$$(E, \nabla) = (\mathcal{L}', \nabla') \oplus (\mathcal{L}'', \nabla'').$$

Assume $F \neq 0$, so this isn't a flat connection; thus at least one of these bundles (without loss of generality, \mathcal{L}'). The property $F^+ = 0$ is also inherited by these line bundles, so (\mathcal{L}', ∇') becomes a $U(1)$ instanton. In particular, if F' is the curvature 2-form of ∇' , then $F' \in \Omega^{2,-}(X)$ and $dF' = 0$. Chern-Weil theory says that

$$\frac{1}{2\pi}[F'] \in H^2(X; \mathbb{Z}) \subset H^2_{\text{dR}}(X).$$

The existence of a nonzero element of $H^{2,-}(X) \cap H^2(X; \mathbb{Z})$ is a condition on (X, g) (where g denotes the Riemannian metric on X). If $b_2^+ := \dim H^{2,+}(X; \mathbb{R}) \neq 0$, then $H^{2,-}(X)$ has nonzero codimension, so generically does not intersect the lattice $H^2(X; \mathbb{Z})$, so there can be no such ϕ (for a generic metric). \boxtimes

In Donaldson theory, it's standard to assume $b_2^+ \neq 0$, again to avoid these *reducible instantons*. In fact, one often takes $b_2^+ > 1$ so there are no reducible instantons even in a one-parameter family of metrics. The space of metrics therefore has a singular piece of the metrics with reducible instantons: if $b_2^+ > 1$, it's at least codimension 2, so its complement is path-connected. This is convenient for proving that things are independent of metrics, because any two metrics may be joined by a path.

Let \mathcal{M} denote the moduli space of $SU(2)$ instantons modulo gauge equivalence. If F is the curvature of a $(P, \nabla) \in \mathcal{M}$, then

$$\frac{1}{4\pi^2} \int F \wedge F \in \mathbb{Z},$$

so let \mathcal{M}_k denote the component with that integral equal to k . Pick a point $(P, \nabla) \in \mathcal{M}_k$ (for $k \neq 0$; the $k = 0$ case is a little more complicated).

We still have the fermionic fields $\phi \in \Omega^0((\text{ad } P)_{\mathbb{C}})$, $\eta \in (\text{ad } P)_{\mathbb{C}}$, and $\chi \in \Omega^{2,+}((\text{ad } P)_{\mathbb{C}})$.⁶⁸

There's a linear operator

$$L := \nabla + (d_{\nabla}^+)^* : \Omega^0((\text{ad } P)_{\mathbb{C}}) \oplus \Omega^{2,+}((\text{ad } P)_{\mathbb{C}}) \longrightarrow \Omega^1((\text{ad } P)_{\mathbb{C}}).$$

and there's a Laplacian $\Delta := LL^* + L^*L$.

If $f := \eta + \psi + \chi$, then the fermion kinetic term is $(1/g^2)\langle f, Lf \rangle$; similarly, if $b = \phi + \delta A + D$, the boson kinetic term is $(1/g^2)\langle b, \Delta b \rangle$ (here g is the initial parameter, not the metric). The nonzero eigenspaces of L and Δ contribute something depending on g , so they must cancel (which you can also see explicitly), but there's no such argument for the zero eigenspace, which we have to study more closely.

⁶⁸**TODO:** I was more confused by this than usual and it may be wrong.

More precisely, we're going to study the zero modes, the f with $\langle f, Lf \rangle = 0$ and the b with $\langle b, \Delta b \rangle = 0$. Inside $\ker L$, we have the η with $\nabla \eta = 0$, but since ∇ is irreducible this forces $\eta = 0$. For χ , we know $(d_\nabla^+)^* \chi = 0$, so $\chi = 0$ if ∇ is regular.

For the adjoint, $\ker(L^*) = \ker(\nabla^*) \cap \ker(d_\nabla^+)$. So the contribution of f is⁶⁹

$$\langle \nabla \eta, \psi \rangle + \langle (d_\nabla^+)^* \chi, \psi \rangle.$$

For the bosonic zero modes, one only needs $\delta A \in \ker(L^*) = \ker(\nabla^*) \cap \ker(d_\nabla^+)$.

We did this relative to a specific ∇ , so working over the entire moduli space, we have two bundles of fermionic and bosonic zero modes over \mathcal{M}_k . The bosonic bundle is easy: it's just $T\mathcal{M}_k$, because they have some first-order behavior. The fermionic part is a parity change of the same vector space over each point, hence is $\Pi T\mathcal{M}_k$.

This means that we can integrate over these finite-dimensional spaces, and will do so next time.

Lecture 23.

The Higgs mechanism: 11/28/17

Last time, we discussed topologically twisted $\mathcal{N} = 2$ supersymmetric Yang-Mills theory on a Riemannian 4-manifold (X, g) , with the goal of calculating correlation functions of Q -closed observables. We claimed this localizes to the moduli space \mathcal{M} of $SU(2)$ -instantons on X , i.e. the $SU(2)$ -bundles with connection (P, ∇) with curvature $F^+ = 0$, modulo gauge equivalence.

If $Q \in \text{Vect}^1(\mathcal{C}_X)$ denotes the odd vector field defining the supersymmetry, then $S = Q\Psi$ for some Ψ , so we can choose the deformation $S = tQ\Psi$ and let $t \rightarrow \infty$, which was the reason we can localize to \mathcal{M} .

Remark 23.1. Last time, we wrote that the action was

$$S = \int_X \|\nabla \phi\| - \|F\| + \|D\| + \dots$$

but it should be

$$S = \int_X \|\nabla \phi\| + \|F\| - \|D\| + \dots,$$

which solves the positive definiteness issue from before. In the literature, there's a typo in [19], which is for the most part an excellent reference; [21] writes the action correctly. \blacktriangleleft

There's a very geometric way to understand what's happening here: the zero modes form a super-vector bundle over \mathcal{M} , which is isomorphic to $T\mathcal{M} \oplus \Pi T\mathcal{M}$. Integration over all nonzero modes produces a function in the zero modes, and it ends up only depending on the odd part. In particular, this implies it's a differential form on \mathcal{M} . This is how Seiberg and Witten's approach connects to Donaldson's — the latter approach uses differential forms on \mathcal{M} , and here's how we got to it.

However, the fermion zero modes are only zero modes up to quadratic order,⁷⁰ which is why this integral is nonzero at all. However, they have *Yukawa couplings*, which are nonzero cubic couplings $\phi[\psi_\mu, \psi^\mu]$ and $\overline{\phi}[\psi_\mu, \psi^\mu]$.

To understand what's going on, let's turn to the bosonic case. We're more or less trying to compute a Gaussian

$$\int e^{-S(x,y)}$$

where

$$S(x, y) = ax^2 + 0y^2 + xy^2$$

(plus possibly some higher-order terms in y). Up to quadratic order, y is a zero mode, but there's a nonzero cubic coupling, and that's what happenong in our computations of super-Yang-Mills, except y is a fermion. The point is that integrating over x produces a function of the zero mode y .

⁶⁹Again, I'm more skeptical than usual that I wrote everything down correctly.

⁷⁰This is also true for the bosonic zero modes, but we're not going to use that.

So if we want to compute, e.g. $\langle \phi^{(0)} \rangle$ (recall $\phi^{(0)} = \text{tr } \phi^2$), we have to compute the Gaussian integral

$$(23.2) \quad \begin{aligned} I(\psi) &= \int \mathcal{D}\phi \text{tr}(\phi(x)^2) \exp\left(-\int_X \text{tr} \|\nabla \phi\|^2 - \frac{i}{\sqrt{2}} \bar{\phi}[\psi_\mu, \psi^\mu]\right) \\ &= \int \mathcal{D}\phi \text{tr}(\phi(x)^2) \left(-\frac{i}{\sqrt{2}} \int dy \text{tr} \bar{\phi}(y)[\psi_\mu(y), \psi^\mu(y)]\right)^2 \exp\left(-\int \text{tr} \|\nabla \phi\|^2\right). \end{aligned}$$

Though this integral has a lot of stuff in it, we can attack it with the standard approaches to integration theory. If

$$H(x) = -\frac{i}{\sqrt{2}} \int d^4y G(x, y)[\psi_\mu(y), \psi^\mu(y)],$$

where G is the Green's function of the operator $\nabla^* \nabla$, then (23.2) is equal to $\text{tr}(H(x)^2)$. The Green's function appears twice in a four-point Feynman diagram with vertices $\phi(x)$, $\phi(x)$, $\bar{\phi}(y)$, and $\bar{\phi}(y)$.

The next step is to expand ψ with respect to a basis of $\ker(L^*)$, where L is the fermion kinetic term:

$$\psi = \sum \varepsilon_i f_n,$$

where ε_n is odd, and $\{f_n\}$ is a basis for $\ker(L^*)$. This implies $I(\psi)$ is quartic in the ε_n , and therefore as a function on $\Pi T \mathcal{M}$, it's some differential 4-form $\mathcal{J}^{(0)} \in \Omega^4(M)$.

Witten here remarks that the physical arguments may lack the mathematical rigor we're used to, but they have nonetheless produced a concrete formula.

Claim 23.3. $\mathcal{J}^{(0)}$ is the same 4-form on \mathcal{M} that was described by Donaldson.

Donaldson described it in a very different way. Let $\mathcal{M}^{\text{conn}}$ denote the moduli space of all $\text{SU}(2)$ -connections on principal bundles on X , without requiring $F^+ = 0$. On $\mathcal{M}^{\text{conn}} \times X$, there's a universal principal $\text{SU}(2)$ -bundle $\pi^* P$ with connection ∇^{univ} . If F_{univ} denotes its curvature, then it has a differential form (sort of a universal Pontrjagin class) $\omega := \text{tr}(F_{\text{univ}}^2) \in \Omega^4(\mathcal{M}^{\text{conn}} \times X)$. This decomposes into pieces on $\mathcal{M}^{\text{conn}}$ and pieces on X : let $\omega^{(\ell)}$ denote the piece of ω in $\Omega^\ell(\mathcal{M}^{\text{conn}}) \otimes \Omega^{4-\ell}(X)$.

Donaldson's description of $\mathcal{J}^{(0)}$ was $\omega^{(4)}$ (and the other pieces of ω correspond to correlation functions of line, surface, etc. operators). Witten doesn't prove that these are the same, but it's probably been proven.

But the Seiberg-Witten solution to the low-energy physics of $\mathcal{N} = 2$ super-Yang-Mills theory for $G = \text{SU}(2)$ tells us something new and interesting.

Returning to $X = \mathbb{R}^4$, we'd like to understand the physics of this theory at low energies. We might make an unrefined guess that the physics can be modeled by an action which includes only massless fields. That is, if you expand around $A_\mu = 0$, $\psi = 0$, and $\phi = 0$, then all fields are massless (e.g. $m^2 |\phi|^2$). But there's an important and tricky point: the minima of the action are degenerate.

Example 23.4. For a more concrete example of this idea, consider a ball rolling in a potential $V(x, y) = xy^2$.
TODO: I missed what came next; sorry! ◀

The upshot is that this theory has degenerate vacua, which are the ϕ such that $[\phi, \bar{\phi}] = 0$, and we consider them up to gauge equivalence. Such a ϕ can be diagonalized: there's some $a \in \mathbb{C}$ such that in some basis,

$$\phi = \frac{1}{\sqrt{2}} \begin{pmatrix} a & 0 \\ 0 & -a \end{pmatrix},$$

where the fields given by a and by $-a$ are gauge equivalence. Thus, the moduli space is the quotient of \mathbb{C} by the antipodal action $a \mapsto -a$ of $\mathbb{Z}/2$.

Let $u = a^2$, so $u = \text{tr}(\phi^2) = \phi^{(0)}$. Expanding around some minimum (vacuum) with $u \neq 0$, some fields will acquire mass in a process called the *Higgs mechanism*.

Remark 23.5. The Higgs mechanism is a general process that occurs in gauge theories coupled to matter. That is, the fields are principal G -bundles with connection (P, ∇) together with a $\phi \in \Gamma(V_P)$, where V_P is the associated bundle for P and some fixed G -representation V . For example, we could let $H: V \rightarrow \mathbb{R}$ be G -invariant and introduce an action

$$(23.6) \quad S = \frac{1}{4g^2} \int F \wedge \star F + \frac{1}{2} \int \|D\phi\|^2 + \int_X H(\phi),$$

where g is a parameter. Suppose H has a minimum other than $\phi = 0$, and for simplicity assume the minima form a single G -orbit. For example, if $G = \mathrm{SU}(2)$ and $V = \mathbb{C}^2$ is the defining representation, we could let

$$H(\phi) = \|\phi\|^4 - 2m^2\|\phi\|^2.$$

The space of minima is an S^3 in \mathbb{R}^4 with radius m .

Now, though, we can expand around the orbit $\|\phi\| = m$. Fix P to be trivial, and choose $\phi := \begin{pmatrix} 0 \\ 1 \end{pmatrix} f$, where f is real. Then we're left with a space of fields which is the connections A on the trivial $\mathrm{SU}(2)$ -bundle, with ϕ fixed.

Write

$$\phi_0 := m \begin{pmatrix} 1 \\ 0 \end{pmatrix} \quad \text{and} \quad \phi := \phi_0 + \delta\phi.$$

Then,

$$\|D\phi\|^2 = \|D(\phi_0 + \delta\phi)\|^2 = \|A\phi_0\|^2 + (\text{terms involving } \delta\phi).$$

This quadratic term in A is new. Let's write A in the Pauli basis for \mathfrak{su}_2 :

$$(23.7) \quad A = A^1 \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} + A^2 \begin{pmatrix} 0 & i \\ i & 0 \end{pmatrix} + A^3 \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix}.$$

Then

$$\|A\phi_0\|^2 = m^2\|A^1\|^2 + \|A^2\|^2 + \|A^3\|^2.$$

Thus all components of A are massive: at energies $E \ll m$ they should be integrated out. $\delta\phi$ is also massive. The idea is that near the minimum of H , tangential directions don't contribute anything significant to the action, so we only care about normal directions, which acquire mass from the potential. This is the Higgs mechanism: there's no gauge symmetry and the gauge fields instead become massive. \blacktriangleleft

The Higgs mechanism will be crucial in supersymmetric Yang-Mills.

Lecture 24.

Running coupling: 11/30/17

Last time we talked about the Higgs mechanism within the context of a gauge theory with gauge group G and matter field ϕ in a representation V of G . For us, this will be a 4D gauge theory with $G = \mathrm{SU}(2)$ and V is the defining representation.

Suppose there exists a *classical vacuum* of the theory in which $\phi \neq 0$, i.e. a minimum of the action S in which ϕ is nonzero and constant. Let's also assume (P, ∇) is trivial.

In this case, the space of such vacua is acted on by G . Fix one such vacuum, labeled by a $\phi_0 \in V$. If the stabilizer of ϕ_0 inside G is trivial, then the low-energy physics has *no* gauge symmetry: the entire symmetry group has been broken. The low-energy approximation is just expanding the action around ϕ_0 (and the trivial connection).

More generally, the low-energy physics has gauge symmetry $H := \mathrm{stab}(\phi_0) \subset G$.

Example 24.1. Suppose $G = \mathrm{SU}(2)$, V is the adjoint representation, and

$$\phi_0 = \begin{pmatrix} a & 0 \\ 0 & -a \end{pmatrix}.$$

In the action we then have

$$\|\nabla\phi_0\|^2 = \|d\phi_0 + [A, \phi_0]\|^2 = \cdots + \|[A, \phi_0]\|^2.$$

Like last time, this decomposes in terms of the components of A in the Pauli basis (23.7): there's an $a^2(\|A_2\|^2 + \|A_3\|^2)$ term, so A_2 and A_3 acquire a mass. This says that if you turn on the gauge field a little bit, these have a little mass, and if you turn it on stronger, they have more mass.

The upshot is that in the low-energy theory, you integrate A_2 and A_3 out. This ought to produce something like

$$S_{\mathrm{eff}} \approx \frac{1}{4g^2} \int F^1 \wedge \star F^1,$$

which looks like the $\mathrm{U}(1)$ theory! \blacktriangleleft

We're studying quantum field theory, but this argument felt very classical. How can we make sure it's justified?

To get an accurate picture of the physics (i.e. all of the usual quantities one computes in QFT) from a Lagrangian and avoid divergences, one should use a Lagrangian with a cutoff close to the energy scale one wants to study.

In 4D gauge theory, the coupling g depends on the energy scale E ; this phenomenon is called a *running coupling*. The dependence isn't completely understood, and depends on the group you choose; for details, see [24].

- For $G = U(1)$ where the theory is not coupled to matter (so the pure Maxwell theory we studied, which is noninteracting), E doesn't affect g : there's different discrete values of g that can exist, and changing E doesn't change them.
- For $G = U(1)$ with matter charged in the fundamental representation, the coupling gets weaker at low energies and stronger at high energies. If you just calculate in one loop, there's a singularity for some large enough E , called a *Landau pole*. This is a problematic feature if you want this QFT to actually exist, suggesting that QED isn't "the" theory of the universe and needs a correction at high energies. It might be the case that two-loop or three-loop corrections fix this issue, but it's unclear. Nonetheless, the real calculations in quantum electrodynamics and effective field theory are not affected by high-energy phenomena.

Suppose the charged matter has mass m . Then it should be possible to integrate the particle out, so at energies lower than about m , the coupling shouldn't depend on E at all!

- Suppose G is a simple group, e.g. $SU(2)$, with no matter. This is an interacting theory, and there's a phenomenon called *anti-screening* or *asymptotic freedom*: the coupling decreases as the energy increases. This is good, because it means you can extrapolate it to high energies, so it makes sense. But because the coupling is large at low energies, computing in the low-energy theory is hard — and there's some energy value Λ at which the coupling blows up!

It's believed that in low energies, g isn't a parameter of the theory. Instead, Λ is, and it's dimensionful, unlike g .

- Suppose $G = SU(2)$ and the theory has matter in the adjoint representation, and let ϕ_0 be as in Example 24.1. At high energies, the energy-coupling graph looks like the pure $SU(2)$ -theory. Below $E \approx a$, it will look like the pure $U(1)$ -theory, meaning g doesn't depend on E . This prevents the blowup at low energies. As a gets larger and larger, the transition happens at lower energies, so there are more options for the coupling.
- If you couple the $SU(2)$ theory to a small amount of matter (a small representation), there's asymptotic freedom again, but with a lot of matter (e.g. 12 copies of the fundamental representation), the running turns around, and g increases with E . The person who first understood this won a Nobel for it.⁷¹ In between these two regimes, g is independent of E , and the theory is actually conformal! The fundamental example of this is $\mathcal{N} = 4$ supersymmetric Yang-Mills theory.⁷²

This means that we want $|a| \gg \Lambda$ in order for the naïve classical calculation of the Higgs mechanism in Example 24.1 to be right, because in this case, the coupling resembles the $U(1)$ -theory on a larger region of possible energies. When $|a| \gg \Lambda$, the whole picture is qualitatively different, and this argument doesn't work.

Now let's add supersymmetry. We're again considering $\mathcal{N} = 2$ super-Yang-Mills theory for $G = SU(2)$. The classical picture is that the moduli space of vacua has the fields

$$\phi \sim \begin{pmatrix} a & 0 \\ 0 & -a \end{pmatrix},$$

where we identify the ϕ produced by a and $-a$. Let $u = a^2$.

One classical approximation (which we don't expect to be particularly accurate) sets all the massive fields to zero. Then, we get a field theory whose fields are tuples $(a, p, \nabla, \lambda^\pm, \dots)$, where

- $a: \mathbb{R}^4 \rightarrow \mathbb{C}$,
- (p, ∇) is a principal $U(1)$ -bundle with connection,
- $\lambda^\pm \in \Pi\Gamma(S^\pm \otimes \mathbb{C}^2)$, where S^\pm is the spinor bundle, and

⁷¹The problem with g increasing with E and the absence of a high-energy description is that it means the theory probably doesn't exist in quite that way. This is one of the reasons it's hard to write down quantum field theories (and supersymmetric quantum field theories): you might expect to be able to write down anything and use it to compute topological invariants, but you have to be careful to get anything that works or makes sense at all.

⁷²This is not good for the Donaldson-theoretic perspective: the existence of a running coupling is good, because it makes the invariants easier to compute using effective field theory. For other applications, such as geometric Langlands, it's not an obstacle, though.

- there are a few more fermionic terms.

This theory is $\mathcal{N} = 2$ super-Yang-Mills theory for gauge group $U(1)$. (Part of) the action is

$$\begin{aligned} S &= \frac{1}{g^2} \int \|da\|^2 + i\vartheta \int F \wedge F + \dots \\ &= (\text{Im } I) \int \|da\|^2 + \frac{i\tau}{4\pi} \int \|F_+\|^2 + \frac{i\bar{\tau}}{4\pi} \int \|F_-\|^2 + \dots \end{aligned}$$

Here $\tau = i/g^2 + \vartheta/2\pi$. Recall that $u = a^2$, so we can also write this as

$$S = \frac{1}{g^2} \int \frac{\|du\|^2}{|u|^2} + \dots$$

This piece of the action describes a σ -model of maps $\mathbb{R}^4 \rightarrow \mathcal{M}^{\text{classical}} = \mathbb{C}_u$, carrying the Kähler metric $g_{u\bar{u}} = 1/|u|^2$. This is fine except when $u = 0$, where the effective physics looks singular — and that's reasonable, since we got rid of all of the mass terms.

So the naïve picture of the low-energy physics is:

- for $u \neq 0$, an $\mathcal{N} = 2$ super-Yang-Mills theory with gauge group $U(1)$, and
- for $u = 0$, an $\mathcal{N} = 2$ super-Yang-Mills theory with gauge group $SU(2)$.

This is how the story was understood prior to the work of Seiberg and Witten. But this naïve approach is wrong, and in fact wrong enough to change the answer! We neglected to include interactions between low-energy fields, which can arise when you integrate out high-energy fields. At the very least, we should try to write them down, which would be a more systematic approach. At large $|u|$, i.e. $|u| \gg \Lambda$, the Higgs picture should be accurate or close to accurate, and then we'll try to work down to $|u|$ small.

For example, we haven't yet incorporated the running coupling of the $SU(2)$ theory, so let's do that. After a calculation which we'll skip, let g_{eff}^2 denote the effective $SU(2)$ coupling at scale a . Perturbation theory implies

$$(24.2) \quad \frac{1}{g_{\text{eff}}^2} = \frac{1}{4\pi^2} \log \frac{|u|}{\Lambda}.$$

This comes entirely from the one-loop Feynman diagram.⁷³ This is a miracle of supersymmetric field theory: there are lots of higher-loop diagrams, and they all cancel each other out.

Now the effective theory has a coupling $\tau_{U(1)}$ which depends on a , and the action will have a $\int (\text{Im}(\tau_{U(1)}(a)) \|da\|^2$ term. You might wonder whether it exists; it's not a big problem if it doesn't, since it's an effective field theory. Next you'd wonder if it's supersymmetric, and the answer is yes,⁷⁴ but only if τ is holomorphic in a .

Thus we have two strong constraints on τ : it should satisfy (24.2), and it should be holomorphic. So let's write a holomorphic function whose imaginary part is (24.2), e.g.

$$(24.3) \quad \tau(a) = \frac{i}{\pi} \log\left(\frac{u}{\Lambda}\right).$$

This tells us the running, at least for $|u|$ large. Something weird happens, though: the complex logarithm has multiple branches, and this causes a shift $\tau \mapsto \tau + 2$. This is not a problem because this is a symmetry of the $U(1)$ theory, as we discussed, and tells us that ϑ was only defined mod 2π anyways.

Lecture 25.

The exact moduli space for the Seiberg-Witten solution: 12/5/17

Last time, we began discussing the Seiberg-Witten solution of the low-energy physics of the $\mathcal{N} = 2$ super-Yang-Mills theory with gauge group $G = SU(2)$. Classically, this theory has a moduli space $\mathcal{M}^{\text{cl}} := \mathbb{C}/(\mathbb{Z}/2)$, where the vacua are labeled by

$$\phi = \begin{pmatrix} a & 0 \\ 0 & -a \end{pmatrix},$$

⁷³Dimensional analysis of (24.2) is a little weird, which is part of the reason Λ is usually considered to be the square of the transition point in Seiberg-Witten theory.

⁷⁴This is not a theorem: there are no theorems in this world yet.

where a and $-a$ describe the same theory (i.e. the $\mathbb{Z}/2$ -action on \mathbb{C} sends $a \mapsto -a$). We then calculated the effective action expanded around any $a \neq 0$, and found that it is the action for $\mathcal{N} = 2$ super-Yang-Mills theory, but for gauge group $U(1)$, and because $U(1)$ is abelian, this makes life much simpler. More precisely, the coupling of the $U(1)$ is the same as the coupling in the original $SU(2)$ theory.

We next turn to the quantum picture. Define

$$u := \frac{1}{2} \langle \text{tr } \phi^2(x) \rangle,$$

with $x \in \mathbb{R}^4$. Classically, $u = a^2$, but u should be a function on the exact space of vacua \mathcal{M} . So the effective theory will be broadly similar: it will include a σ -model into \mathcal{M} , plus some additional fields. In the limit where $|u| \rightarrow \infty$, \mathcal{M} should look like \mathcal{M}^{cl} , and the effective action should be close to the classical effective action (this was the argument from renormalization group running).

Computing the 1-loop term in the perturbation theory for the original $SU(2)$ theory told us that the coupling in the effective quantum action is not just the same as the coupling in the $SU(2)$ theory, and is given in (24.3), where Λ is the dynamical scale of the original $SU(2)$ theory. This is just the one-loop term, so it's not completely accurate, but it's more accurate.

Thus the coupling $\tau_{U(1)}$ is a function on \mathcal{M} . As long as $|u| \gg \Lambda$, $\text{Re } \log(u/\Lambda) \gg 0$, so $\tau \in \mathbb{H}$ and so the coupling is in the upper half-plane. Another concern is that \log is multivalued, and therefore τ is: as $u \mapsto ue^{2\pi i}$, $\tau \mapsto \tau + 2$.

Remark 25.1. We'd like to say that u is a coordinate on \mathcal{M} , i.e. a function $\mathcal{M} \rightarrow \mathbb{C}$. More generally, any local operator \mathcal{O} should define a local coordinate on \mathcal{M} whose value at some point is the local operator for that theory. If A denotes the algebra of local operators, it's very often true that $\mathcal{M} \cong \text{Spec } A$, but it's not clear if there's a general reason for this to be true. \blacktriangleleft

Anyways, the $\tau \mapsto \tau + 2$ monodromy is not an issue, because $U(1)$ -super-Yang-Mills with coupling τ is isomorphic to $U(1)$ -super-Yang-Mills with coupling $\tau + 2$. So everything that we can get out of this effective theory is well-defined.

Remark 25.2. Electromagnetic duality also defines an isomorphism between the theory with coupling τ and the theory with coupling $-1/\tau$. \blacktriangleleft

Thus, to determine the exact picture, we need

- a complex curve \mathcal{M} ,
- a holomorphic function $u: \mathcal{M} \rightarrow \mathbb{C}$,
- an open subset $U \subset \mathcal{M}$ such that $u: U \rightarrow \{|u| > C\} \subset \mathbb{C}$ is an isomorphism, and
- a holomorphic map $\tau: \mathcal{M} \rightarrow \mathbb{H}/\text{SL}_2(\mathbb{Z})$ such that as $|u| \rightarrow \infty$ in U ,⁷⁵

$$\tau \rightarrow \frac{i}{\pi} \log\left(\frac{u}{\Lambda^2}\right).$$

Remark 25.3. This last bullet point is one of the more interesting pieces of the Seiberg-Witten solution: rather than a single effective Lagrangian to describe the whole theory, there's a family of Lagrangians on different coordinate charts of the space that agree on overlaps. Somehow this works, but it's somewhat mysterious what's going on, and the general idea has not been fully sussed out. \blacktriangleleft

We'll use this to conclude some facts about \mathcal{M} . First of all, τ produces a *special Kähler metric* on \mathcal{M} .

The coupling τ provides multiple different effective descriptions of the theory at a particular point. Once one is chosen, the field ϕ becomes a local coordinate. But these are particularly nice: these local coordinate systems are parameterized by an $\text{SL}_2(\mathbb{Z})$ -principal bundle $X \rightarrow \mathcal{M}$ (so an $\text{SL}_2(\mathbb{Z})$ torsor over each point, varying smoothly): given a $T \in X$, we get a local coordinate a , and $\text{SL}_2(\mathbb{Z})$ acts on coordinates in the obvious way. Another way to say this is: let $V := \mathbb{Z}^2$ denote the fundamental representation of $\text{SL}_2(\mathbb{Z})$. Then, there's an associated bundle

$$\Gamma := X \times_{\text{SL}_2(\mathbb{Z})} V,$$

which is a bundle of lattices. Physically, these are the lattices of electromagnetic charges. Then, there's a homomorphism $Z: \Gamma \rightarrow \mathbb{C}$, which in a local trivialization is just (a, a_D) (a and its dual), which are local coordinate functions $V \rightarrow \mathbb{C}$. a_D is related to a by electromagnetic duality. Both a and a_D are good local coordinates, and more generally any linear combination of them is a coordinate, and this is what Z encodes.

⁷⁵We really should be considering a subgroup of $\text{SL}_2(\mathbb{Z})$, because we only have $\tau \mapsto \tau + 2$ symmetry, rather than $\tau \mapsto \tau + 1$. However, this would pass to $\mathbb{H}/\text{SL}_2(\mathbb{Z})$ anyways, so it's all right for now.

So \mathcal{M} has a lot of extra structure, rather than just some complex curve. This implies that the theory is rigid: not all transformations preserve it, but the ones coming from $\mathrm{SL}_2(\mathbb{Z})$ do.

The QFT also provides a metric on \mathcal{M} , which is

$$g = (\mathrm{Im} \tau(a))^{-1} |da|^2$$

in any local coordinate chart, and this metric is (geodesically) complete. This ought to be true physically, as an effective field theory description with holes in the moduli space can't account for all of the physics, and there's also a mathematical argument for this.

Proposition 25.4. *The existence of a u which is one-to-one on U implies $\mathcal{M} \cong \mathbb{C}$.*

Proof. It suffices to show that \mathcal{M} can't contain any "holes" (i.e. it has no higher-genus pieces). If it did, then $\overline{\mathcal{M}}$ is a surface with genus greater than 0, and u extends to a function $\bar{u}: \overline{\mathcal{M}} \rightarrow \mathbb{CP}^1$ which has degree $d > 1$. This means at a generic point in U , u is d -to-one.

We also need \mathcal{M} to be complete, which was addressed above. \square

So the topology of the moduli space is simple.

But there's something strange going on: \mathcal{M} is simply connected, but these systems have monodromy. Because \mathcal{M} is simply connected, $\tau: \mathcal{M} \rightarrow \mathbb{H}/\mathrm{SL}_2(\mathbb{Z})$ lifts to a map $\tilde{\tau}: \mathcal{M} \rightarrow \mathbb{H}$, i.e. there is no monodromy!

So evidently this is not the right picture. What's the next thing you might think to try? Maybe there are some points of \mathcal{M} which aren't described by this effective theory. Concretely, this would mean there's a divisor $D \subset \mathcal{M}$ where the effective $\mathrm{U}(1)$ theory breaks down.

If D is a single point, then τ would lift to a map $\tilde{\tau}: \mathcal{M} \rightarrow \mathbb{H}/\mathbb{Z}$, corresponding to nontrivial monodromy around a point. In this case, $\mathrm{Im} \tau$ is a globally defined harmonic function on $\mathcal{M} \setminus D \cong \mathbb{C}^\times$.

Exercise 25.5. Show that there is no nonconstant harmonic map $f: \mathbb{R}^2 \setminus \{(0,0)\} \rightarrow \mathbb{R}_{>0}$.

One way to do this is to look at the equations that the average on a circle centered on the origin must satisfy.

Ok, so the next guess: what about two points? Let m_1 and m_2 be the monodromies around these two singular points x_1 and x_2 . Relative to some fixed choice of basis,

$$(25.6) \quad \begin{aligned} m_1 &= \begin{pmatrix} 3 & 2 \\ -2 & -1 \end{pmatrix} \\ m_2 &= \begin{pmatrix} 1 & 0 \\ -2 & 1 \end{pmatrix}. \end{aligned}$$

These matrices generate a subgroup of $\mathrm{SL}_2(\mathbb{Z})$ called $\Gamma(2)$, the subgroup of matrices equal to the identity matrix mod 2.

We make the somewhat optimistic assumption that $\mathcal{M} = \mathbb{C}$. We'll make τ using a weird trick from algebraic geometry: we'll make a genus-1 elliptic curve and let τ be the modular parameter of that curve. Specifically, consider the family of elliptic curves over $\mathcal{M} \setminus \{u = \pm\Lambda^2\}$ defined by

$$(25.7) \quad \Sigma_u := \{y^2 = (x-u)(x-\Lambda^2)(x+\Lambda^2)\},$$

a one-dimensional subvariety of \mathbb{C}^2 . It's a standard fact from algebraic geometry that a cubic curve inside \mathbb{C}^2 , once projectivized, is a genus 1 curve, hence a torus $\mathbb{C}/\mathbb{Z} \oplus \mathbb{Z}\tau$. This $\tau \in \mathbb{H}/\mathrm{SL}_2(\mathbb{Z})$ is called the *modular parameter*.

So for every $u \in \mathcal{M} \setminus \pm\Lambda^2$, Σ_u is a smooth torus, hence defined a modular parameter $\tau(u)$. The effective physics is described by this family of algebraic curves, and the family degenerates at two points, which are where we don't have the description we wanted.

Exercise 25.8. Show that this family of elliptic curves actually has the monodromies specified by m_1 and m_2 .

The electromagnetic lattice is described by the homology $H_1(\Sigma_u, \mathbb{Z})$ over the point u . This is a geometric representation of the monodromy: if you travel around a loop around a singular point, what you get is a different homology class.

Lecture 26.

Understanding the singular points: 12/7/17

Last time, we discussed the low-energy limit of $\mathcal{N} = 2$ super-Yang-Mills with gauge group $G = \mathrm{SU}(2)$ as proposed by Seiberg and Witten. If Λ is the scale of the theory, then the low-energy limit is locally described by a $\mathrm{U}(1)$ -super-Yang-Mills theory, but not globally. There's a quantum moduli space $\mathcal{M} = \mathbb{C}$ with coordinate u , and outside of a divisor $D = \{\pm\Lambda\}$, the effective physics around u is governed by $\mathcal{N} = 2$ super-Yang-Mills theory with group $\mathrm{U}(1)$. Moreover, the coupling τ is the modulus of an elliptic curve Σ_u , given by a family of elliptic curves over \mathcal{M} as in (25.7).

This family is smooth on $\mathcal{M} \setminus D$, and degenerates on D . The monodromy around $\pm\Lambda$ is given in (25.6), and this is something extremely concrete. The $\mathrm{SL}_2(\mathbb{Z})$ -action can be read off from the *Gauss-Manin connection* on $H_1(\Sigma_u; \mathbb{Z})$, which is a local system of lattices (concretely, the associated bundle to the principal $\mathrm{SL}_2(\mathbb{Z})$ -bundle we described last time), and has a physical meaning: the lattice of electromagnetic charges in the effective theory. Thus, the monodromy of this local system says that if you vary the parameters in the theory, an electrically charged particle may pick up a magnetic charge as well when you return to your starting point.

The map $(y, x) \mapsto x$ identifies Σ_u as a double cover of \mathbb{C} , branched at four points. The homology is generated by two explicit cycles: A around $\pm\Lambda$, and B around λ and u . Monodromy means we move u around (in \mathbb{C}), which eventually means moving A and B . Moving u towards Λ pushes part of A under the branch cut; then, redrawing the surface with a less twisted branch cut. Eventually this entails splitting B into two copies, so we get $A + 2B$ as expected, and you can make similar arguments to get m_1 . So this is a little confusing, but you can draw the picture, and everything is concrete.

We now more or less understand the description of the physics on $\mathcal{M} \setminus D$, but what about at $u = \pm\Lambda$? The existence of these singularities at all is surprising: there was no such singularity in the original, high-energy theory. So the singularity could be interpreted as a mistake that we made in our passage to the low-energy theory (which isn't a problem when $u \neq \pm\Lambda$). Specifically, we integrated out the massive fields, since they contribute to the effective action inversely proportional to the mass.

But the mass of a particle can vary as you vary the parameters of a theory. What if the mass of a particle goes to 0 as we approach a point u ? At these points the $\mathrm{U}(1)$ description does not suffice, but Seiberg and Witten also made a proposal for this. They described the effective physics as a $\mathcal{N} = 2$ super-Yang-Mills theory for $G = \mathrm{U}(1)$ again, but coupled to electrically charged matter.

In this theory, there is one extra field, and it's relatively simple. The fields are

- (P, ∇) , λ^\pm , $a \in \mathfrak{g}_{\mathbb{C}}$, and D just as in the usual $\mathcal{N} = 2$ $\mathrm{U}(1)$ -super-Yang-Mills theory,
- two fields $\psi^\pm \in \Pi\Gamma(S^\pm \otimes (V_c)_P)$, where S is the spinor bundle, $c \in \mathbb{Z}$, and V_c is the c^{th} irreducible representation of $\mathrm{U}(1)$,⁷⁶
- two fields $\chi^\pm \in \Gamma(S^\pm \otimes (V_c)_P)$, and
- a $q \in \Gamma((V_c)_P \otimes R)$.

The new fields are collectively called a *hypermultiplet*.

So you could think of this theory as “ $\mathcal{N} = 2$ super-Yang-Mills with one extra particle,” and that particle behaves like an electron. The fact that this admits a supersymmetry is nontrivial, but it does work out.

There's a single parameter m . The action has the form

$$S = S_{\mathrm{gauge}} + \int_{\mathbb{R}^4} \|\nabla q\|^2 + \|(m + c \cdot a)q\|^2 + (\text{fermionic terms}).$$

So expanded around a fixed a , q has mass $|m + ca|$: as we mentioned, the mass of q changes as we vary the parameters.

As q is a massive field, we would like to integrate it out. If we initially took τ to be constant, then after integrating (using a one-loop Feynman diagram calculation), we obtain an effective τ

$$\tau(a) = \frac{c^2}{\pi i} \log\left(\frac{m'}{\Lambda'}\right) + C$$

for some constant C . This looks a lot like the theory at infinity that was our first description of the effective theory for large u .

⁷⁶In particular, this implies V_c is one-dimensional.

Now we have theories parameterized by $a \in \mathbb{C}$. We need to be careful with $a = -m/c$, the place where q becomes massless, and if one walks in a circle around this point, the monodromy is $\tau \mapsto \tau + 2c^2$, i.e. the matrix

$$\begin{pmatrix} 1 & 2c^2 \\ 0 & 1 \end{pmatrix}.$$

If $c = 1$, this is the monodromy matrix we were trying to explain, and the other is conjugate to this, so both of the monodromies in (25.6) are explained by this fuller picture of the effective physics.

- If you want to describe the physics in a region not containing either point in D , you get pure $\mathcal{N} = 2$ super-Yang-Mills for $U(1)$ with $\tau = \tau(u)$ as before. So it's "just" supersymmetric electromagnetism. This description is $SL_2(\mathbb{Z})$ -invariant.
- If you want to describe the physics in a (small) region containing a point of D , you again have supersymmetric electromagnetism, but with an electron: the theory is $\mathcal{N} = 2$ super-Yang-Mills with a hypermultiplet of charge 1. This description is *not* $SL_2(\mathbb{Z})$ -invariant: the elliptic curve which produced the $SL_2(\mathbb{Z})$ -action degenerates over D , so one cycle (the smooth one) has been picked out.

On overlapping domains, these should be consistent, so where does the electron go? This is actually fine: there are two descriptions, and one relates to the other by integrating the electron out.

What's particularly neat about this description is that $SU(2)$ is gone: no matter where you are, you only have to talk about abelian gauge groups.

Applications to Donaldson theory. Recall that Donaldson theory has Q -closed observables $u = \text{tr } \phi^2$, a local operator, and a surface operator S . From this, one can build the generating function for the Donaldson invariants of our spacetime (Riemannian 4-manifold) X ,

$$Z_{\text{DT}}(p, n^\mp) := \langle e^{p \cdot u + n^\mp \cdot S_I} \rangle,$$

where I runs over a basis for $H^2(X; \mathbb{Z})$.

A priori, Z_{DT} is some arbitrary function. But it actually has a lot of structure, and its dependence on p is very simple.

Theorem 26.1 (Kronheimer-Mrowka [17, 18]). *For 4-manifolds of simple type and Betti numbers $b_1 = 0$ and $b_2^+ > 1$,*

$$\left(\frac{\partial^2}{\partial p^2} - 4 \right) Z_{\text{DT}}(p, n^\mp) = 0.$$

Unfortunately, this was more or less the definition of simple type! However, Kronheimer and Mrowka provided many classes of 4-manifolds that are of simple type, e.g. Kähler manifolds.

Kronheimer and Mrowka also described a splitting of the Donaldson invariants: roughly,

$$Z_{\text{DT}}(p, n^\mp) = e^{2p} Z_1(n) + e^{-2p} Z_2(n),$$

for some Z_1 and Z_2 that weren't well understood (so these correspond to $u = \pm 2$).

Witten [29] therefore conjectured that to compute Z_{DT} , one can instead use the (topologically twisted version of the) effective description of the u -plane. The answer is concentrated at the singularities $u = \pm \Lambda$, and the rest of the plane gives zero!

Thus, we can look at (topologically twisted) $\mathcal{N} = 2$ super-Yang-Mills with $G = U(1)$ and a hypermultiplet, and compute its partition function on X . As usual, this involves localization on a moduli space, but now it involves a new set of equations, the *Seiberg-Witten equations*.⁷⁷ These are not particularly complicated, but nobody had written them down before.

Let $\lambda \in H^2(X; \mathbb{Z}) + \frac{1}{2}\omega_2(X)$ and \mathcal{M}_λ denote the moduli space of $U(1)$ -connections ∇ for the line bundle L with $c_1(L)/2 = \lambda$ together with sections $M \in \Gamma(S^+ \otimes L^{1/2})$. Here, the existence of S^+ requires the choice of a Spin^c structure on X .⁷⁸ On \mathcal{M}_λ , the Seiberg-Witten equations are

$$\begin{aligned} \not{D} M &= 0 \\ F^+ &= (\overline{M} M)^+. \end{aligned}$$

⁷⁷Some people call these the *monopole equations*, but there's nothing about monopoles going on here, so that name can be confusing.

⁷⁸As with the instanton equations, the overall moduli space splits into pieces indexed by λ , and using this makes the computation easier.

Just as for the moduli space of instantons, one can calculate the dimension of the moduli space and find that

$$\dim \mathcal{M}_\lambda = \int_X \lambda^2 - \frac{2\chi(X) + 3\sigma(X)}{4}.$$

Here $\chi(X)$ is the Euler characteristic of X , and $\sigma(X)$ is its signature. Since $\lambda \in H^2(X)$, $\lambda^2 \in H^4(X)$ so we may evaluate it on the fundamental class of X to obtain an integer, and this is what is meant by $\int_X \lambda^2$.

Moreover, to compute Z_{DT} , one only needs to consider the pieces \mathcal{M}_λ with dimension zero, meaning the integral is just counting the solutions. Witten [29] conjectured a formula in this case relating Z_1 and Z_2 to these moduli spaces \mathcal{M}_λ . For 4-manifolds of simple type, this conjecture has been proven.

Remark 26.2.

- (1) This story happened in 1995. Subsequently, there was a burst of progress by forgetting about Donaldson theory and Witten's formula and just using \mathcal{M}_λ and the Seiberg-Witten equations to prove longstanding conjectures in Donaldson theory. Progress was very rapid, and had important consequences in 4-manifold topology.
- (2) There should be an analogous theory of 4-manifold invariants for any 4D $\mathcal{N} = 2$ supersymmetric quantum field theory — the way in which we applied physics to topology is not specific to $\text{SU}(2)$ -super-Yang-Mills theory. For example, one could consider other gauge groups. But that was the simplest example, so the others are more complicated, and the other examples that have been worked out produce weaker invariants than the Donaldson invariants. There are heuristic reasons why this might be, but whether you can produce anything new and interesting is an open question, and research continues, e.g. in a recent paper of Moore-Nidaiev [22]. ◀

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