

The Algebraic Theory
of
Topological Quantum Information

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Abstract

This book aims to give a comprehensive account of the algebraic theory of topological quantum information. It is intended to be accessible both to mathematicians unfamiliar with quantum mechanics and theoretical physicists unfamiliar with category theory. Additionally, this text should make a good reference for working researchers in the field. A primary focus of this text is balancing powerful algebraic generalities with concrete examples, principles, and applications.

For my mentors

0 Preface

This book is a mathematical treatment of topological quantum information, with a focus on formal algebraic aspects and a special eye towards topological quantum computation. This manuscript began as an extended set of notes from a course on topological quantum field theory given by Zhenghan Wang in the winter of 2022 at UC Santa Barbara. Through his courses, his private tutoring, and his recommendations, Zhenghan took me from a state of almost complete ignorance of mathematical physics to being a young researcher in the field. I am greatly emdebtetd to him for this, and it is certain that this book would not have existed without his guidance - he richly deserves of my apple.

This book would not have been possible without the tutelage of my esteemed mentors. Those who most directly contributed are the ones who used their time and energy to teach me topological quantum information - Dave Aasen, Mike Freedman, and Yuri Lensky. There are also those who took a chance on a young mathematician when they had every reason not to - Roald Dejean, Peter Bloomsburgh, Edward Frenkel, and Ken Ribet.

[WORK: There's some other people to thank. Andrew Sylvester for letting me try out my arguments on him. Alexei Kitaev and Daniel Ranard now for advising me.]

Great pains have been taken to make this book as pedagogical and accessable as possible. The hope is that it should be readable by both mathematicians unfamiliar with quantum mechanics as well as theoretical physicists unfamiliar with category theory. A primary focus of this text is balancing powerful algebraic generalities with concrete examples, principles, and applications. The prerequisites for this book are a undergraduate-level understanding of topology, linear algebra, and group theory, as well as a popular-science level of familarity with quantum mechanics.

There are already many great references to learn aspects of the material covered in this book. An excellently written and relatively complete book on topological quantum information from the perspective of a physicist is Steven Simon's text [Sim23]. Simon's book is algebraic, but does *not* include any category theory. The main references for the relevant category theory are Bakalov-Kirillov [BK⁺01] and Etingof-Gelaki-Nikshych-Ostrik [EGNO16]. While both excellent texts, they suffer notable shortcomings for learning topological quantum information. Bakalov-Kirillov was written in 2001, making it outdated. Etingof-Gelaki-Nikshych-Ostrik is modern, but makes no connections to physics and does not use the language of string diagrams. The manuscript most similar to this one is Kong-Zhang's preprint [KZ22]. We distinguish ourself from Kong-Zhang by our rigorous mathematical treatment, our different choice of topics, and our extended scope. Other relevant books and review articles include Wang's monograph [Wan10] and Kauffman-Lomonaco's quantum topology themed review [KL09].

[WORK: I will add a section detailing the structure of this book, and how it should be read. I have not written enough for this to be useful yet.]

Contents

0 Preface	3
1 Overview	9
1.1 Conceptual introduction	9
1.1.1 Motivation and applications	9
1.1.2 Mathematical picture	11
1.1.3 History of the subject	11
1.2 Technical introduction	14
1.2.1 Principles of topological quantum information	14
1.2.2 Defects in ordered media	17
1.2.3 The fundamental group	21
1.2.4 Topological classical computation	24
2 Quantum mechanics	31
2.1 Overview	31
2.1.1 Introduction	31
2.1.2 Experimental motivation	31
2.2 Axiomatic development	32
2.2.1 Probability theory	32
2.2.2 Basis-dependent quantum mechanics	34
2.2.3 Measurement	35
2.2.4 Incomplete measurement	37
2.2.5 Basis-independent quantum mechanics	37
2.2.6 Hamiltonians and the Schrodinger equation	41
3 Topological quantum order	45
3.1 Overview	45
3.1.1 Introduction	45
3.2 Discrete gauge theory	49
3.2.1 Ordered media on a lattice	49
3.2.2 From ordered media to gauge theory	53
3.2.3 Kitaev quantum double model	54
3.3 The toric code	57
3.3.1 Simplified Hamiltonian	58
3.3.2 Exact solution of the toric code	60
3.3.3 Toric code as a topologically ordered system	62
3.4 Anyons	62
3.4.1 Topological quantum information in excited states	62
3.4.2 Definition and principles of anyons	65
3.4.3 Anyons in the toric code	66
3.4.4 Anyons in discrete gauge theory	67
4 Category theory	70
4.1 Overview	70
4.1.1 Introduction	70
4.1.2 Definition and important obervations	71
4.2 Structures in category theory	74

4.3	Monoidal categories	75
4.3.1	Motivation, definition, and string diagrams	75
4.3.2	Braided monoidal categories	78
4.3.3	Examples, equivalences, and MacLane's coherence theorem	82
4.3.4	Pivotal monoidal categories	88
5	Modular categories	96
5.1	Overview	96
5.1.1	Introduction	96
5.1.2	Using the final product	97
5.2	First properties	98
5.2.1	Definition	98
5.2.2	Anyons in modular categories	99
5.2.3	States in modular categories and unitarity	102
5.2.4	Topological charge measurement	105
5.3	The modular category toolkit	105
5.3.1	Trace	105
5.3.2	Duality	107
5.3.3	Quantum dimension and Frobenius-Perron dimension	108
5.3.4	Twist	112
5.3.5	Functors, natural transformations, and equivalence	115
5.3.6	Deligne tensor product	116
5.4	The category of G -graded G -representations	118
5.4.1	Overview	118
5.4.2	Higher linear algebra	119
5.4.3	Spherical fusion structures	119
5.4.4	Braiding and modularity	119
5.5	The modular representation	119
5.5.1	Definition	119
5.5.2	Torus perspective	121
5.5.3	Bruguières's modularity theorem and the Verlinde formula	122
5.5.4	Verlinde formula	126
5.5.5	Proof of modularity	127
5.5.6	Vafa's theorem, unitarity of S -matrix, and the Chiral central charge .	130
5.6	Skeletonization	133
5.6.1	Principle	133
5.6.2	F -symbols	133
5.6.3	R -symbols	133
5.6.4	θ -symbols	133
5.6.5	Reconstruction theorem	133
5.7	Quantum double modular categories	133
5.7.1	The Drinfeld center	133
5.7.2	Muger's theorem	135
5.7.3	Discrete gauge theory as a quantum double and Morita equivalence .	135
5.7.4	Factorizability and time reversal symmetry	136
5.7.5	Levin-Wen model	136
5.8	Unitarity	136
5.8.1	Characterization of unitarizability	136

5.8.2	Unitary braidings	137
5.8.3	Uniqueness of unitary structure	138
5.8.4	Skeletonization of unitarity	138
5.9	Number theory in modular categories	138
5.9.1	Techniques and first results	138
5.9.2	Galois conjugation	138
5.9.3	Ocneanu rigidity	138
5.9.4	Rank-finiteness theorem	138
5.9.5	Schauenberg-Ng theorem	138
6	Further structure	142
6.1	Overview	142
6.1.1	Introduction	142
6.2	Boundaries and domain walls	143
6.2.1	Physical picture	143
6.2.2	Boundaries and domain walls in the toric code	144
6.2.3	The algebraic theory of gapped boundaries	144
6.2.4	Lagrangian algebras and algebra objects	145
6.2.5	Module categories and boundaries in the Levin-Wen model	145
6.2.6	The Drinfeld center via bimodules	153
6.2.7	Muger's theorem and Lagrange's theorem for fusion categories	155
6.2.8	Skeletonization of gapped boundaries	155
6.3	Symmetry enriched topological order	155
6.3.1	Physical picture	155
6.3.2	Haldane spin chain	156
6.3.3	Twist defects in the toric code	156
6.3.4	The algebraic theory of SET phases	156
6.3.5	Gauging symmetries	159
6.3.6	Obstruction theory, anomalies, and zesting	161
6.3.7	G -crossed modularity	161
6.3.8	Skeletonization of SET phases	161
6.4	Fermionic topological order	162
6.4.1	Physical picture	162
6.4.2	Kitaev spin chain	162
6.4.3	The algebraic theory of fermionic topological order	162
6.4.4	The 16-fold way on the trivial fermionic theory	162
6.4.5	Modular representations of fermionic topological order	162
6.4.6	Super-groups and Deligne's theorem	162
7	Topological quantum computation	167
7.1	Overview	167
7.1.1	Introduction	167
7.1.2	Universality	168
7.2	Computation with Fibonacci anyons	170
7.2.1	Methodology	170
7.2.2	The Jones invariant	170
7.2.3	Proof of universality	170
7.3	Computation with doubles of finite groups	171

7.4	Computation with the toric code	171
7.5	Computation with Ising anyons	171
7.6	Computation with Majorana zero modes	171
A	Odds and ends	174
A.1	Topological quantum field theories	174
A.2	Quasitriangular weak Hopf algebras	175
A.3	Quantum groups	176
A.4	Subfactors	176
A.5	Vertex operator algebras	176
B	Anyon data	178
B.1	Low-rank modular categories	178
B.2	Abelian modular categories	178
B.3	Group-theoretical modular categories	178
B.4	Miscellaneous examples	178

1 Overview

1.1 Conceptual introduction

1.1.1 Motivation and applications

We will take as a definition *topological quantum information* to be the study of information in topological quantum systems. A topological quantum system is some mathematical or physical system which is in a fundamental sense described by the mathematics of both quantum mechanics and topology. The term *quantum system* here is used in contrast to *classical system*. The flow of current through a conducting copper wire is described perfectly well by classical electromagnetism, whereas the flow of current through a superconducting niobium-titanium wire necessarily requires quantum mechanics for its description.

The term *topological system* is used in contrast to *geometric systems*, though the term “geometric system” is a nonstandard one. In a geometric system measurable quantities and phenomena depend on quantitative local aspects of the system - the distance between wires, the exact shape of some sample, or the curvature of some component. In a topological system measurable quantities and phenomena depend only on qualitative global aspects of the system - whether two wires cross or not, whether a sample is connected or not, whether a component curves into a ball or has a boundary.

The title of this book refers to “topological quantum information” and not “topological quantum systems” for two reasons. The first is to highlight the fact that there is more to a topological quantum system than its global topological properties. Topological quantum systems also have local geometric descriptions which are important for understanding many phenomena. However, we will mostly be ignoring these local effects in favor of focusing on global topological properties. The beauty of topological quantum systems lies exactly in the fact that this global perspective captures the essential information in the system. The second reason is to highlight this book’s eye towards topological quantum computing, the idea of making computers using topological quantum systems.

Since Peter Shor’s 1994 discovery of an efficient factoring algorithm on quantum computers [Sho94], the primary goal of quantum information theorists has been to harness quantum information sufficiently well so that it can be used to make an efficient scalable quantum computer. One of the major hurdles in achieving this goal is that quantum information is *fragile*. Small amounts of noise coming from nearby electromagnetic fields or imperfections in experimental devices are often enough to affect the information being stored, resulting in *errors* in the computation. In the early days of quantum computing it was not clear whether there was any way around this problem. Perhaps the inherent fragility of quantum information would make quantum computation impossible. This turned out to be false.

The beautiful observation is that errors are not nearly as catastrophic in *topological quantum systems*. Errors are typically local. By definition the information topological systems does not depend on local properties, and hence is not affected by local changes. Hence, under suitable conditions, topological systems are naturally error resistant! In the same way that invariants of topological spaces are supposed to be invariant under deformations in pure mathematics, information in topological systems is invariant under errors in mathematical physics. Hence, to solve the problem of noise all one has to do is make a *topological quantum computer*! This observation was made in 1997 and is due independently to Alexei Kitaev and Michael Freedman [Kit03, Fre98]. Since then topological quantum computing

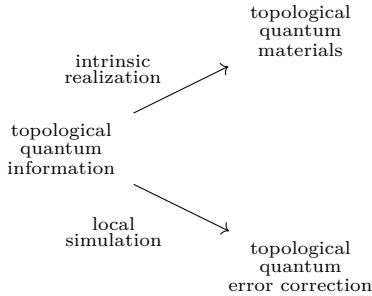


Figure 1.1: The two major branches of topological quantum information.

has grown and evolved, finding its way into almost every modern proposal for fault-tolerant quantum computing.

The first approach to topological quantum computing is to use a physical material, some literal condensed collection of atoms, which naturally behaves as a topological quantum system. These exist and have been studied for a long time. For instance, a two dimensional sheet of graphene behaves topologically when it is subjected to low temperatures (≈ 5 degrees Kelvin) and large magnetic fields (≈ 15 Teslas) [BGS⁺09]. Topological quantum materials which can be used to make scalable quantum computers require intricate experiments to operate, which has been the most prominent roadblock in this approach.

The second approach to topological quantum computing is to artificially construct a topological system within a geometric one. This sort of artificial construction typically takes place within a noisy non-topological quantum computer. The function of a quantum computer, almost by definition, is to simulate quantum systems. In particular, it can simulate *topological* quantum systems. Since topological systems are resistant to local errors, this means that the original computer which is simulating the topological system will itself become resistant to local noise! This works exactly as described as long as the simulation itself is local, that is, local effects in the original system correspond to local effects in the simulated system. This technique of simulating topological systems to inherit their error-resistant properties is known as *topological quantum error correction*. The advantage of this approach is that it works on any hardware available. The disadvantage is that to perform useful computations one must pass through the simulation involved the topological quantum error correction. This additional layer adds a hefty amount of overhead, which can eat up the majority of runtime and resources. It is for this reason that *efficient* topological quantum error correction is an important and active area of research.

Of course, the above discussion presents only one motivation for topological quantum information and only one example of an application. Topological quantum materials open a whole world of potential applications, and it seems they may play an important role in the technologies of the future [RS20]. Some proposed applications include processing classical information using topological defects in magnetic devices (with the end goal of making high-speed low-energy transmissions) [MZ21, ŠMYM18], creating highly sensitive photodetectors (with the end goal of making night-vision goggles or sensors) [CLRL17], creating technologies with high thermoelectric effect (with the end goal of making efficient fridges or electric generators) [SF18], creating highly-efficient transistors [FEC⁺21], and engineering tiny electrical components [VD14, PBD17].

This breadth of potential applications is due in part to the number of different types

of topological materials which have been discovered or theorized. This includes quantized Hall states [vKCK⁺20], topological insulators [HK10], fractional Chern insulators [RB11], Weyl/Dirac semimetals [AMV18], topological superconductors [SA17], and many more. The contents of this book certainly do not provide the entire picture for any of these materials. However, the hope is that it gives a picture of the algebraic structures within them, helping readers think both concretely and conceptually about these materials and their applications.

1.1.2 Mathematical picture

The term *topological quantum system* is broad. To get a rigorous mathematical subject, we will focus on a specific type of topological quantum system known as a *topologically ordered* quantum system. Topological order is much more precise, though there are still conflicting definitions in the literature. Specifically, we will focus on *(2+1)-dimensional* topological order - topologically ordered systems in two space dimensions and one time dimension. That is, I will be discussing a locally flat topologically ordered system. For example, a single sheet of graphene at low temperatures and large magnetic fields can exhibit a form of (2+1)D topological order, and any quantum computer running topological quantum error correction can also exhibit (2+1)D topological order.

All systems in this book are two-dimensional unless stated otherwise.

Topological quantum systems can be described in many different ways. In this book we will take an *algebraic* approach. This means we focus on the big-picture structure of the information, based on symbolic equations and relationships. The algebraic structure which houses the data of a (2+1)D topologically ordered system is known as a *modular category*. These algebraic structures are the main mathematical object of this text. Once one has a modular category it is easy to manipulate the stored information to predict the result of computations. This gives the overall schema of our discussion, illustrated visually below:

$$\begin{array}{ccc} & \text{(modular)} & \\ & \text{categories} & \\ \nearrow & & \searrow \\ \left(\begin{array}{c} \text{topologically} \\ \text{ordered systems} \end{array} \right) & \longrightarrow & \left(\begin{array}{c} \text{topological quantum} \\ \text{computing} \end{array} \right) \end{array}$$

In Chapter 3 we describe topological order. In Chapter 5 we describe topological order algebraically in terms of modular categories. In Chapter 6 we describe further algebraic structures which lie beyond plain modular categories, which allow us to describe more complex behaviors in topological order. Finally, in Chapter 7 we will use the tools we have established to detail several algorithms and procedures for topological quantum computation. Two introductory chapters are also included: Chapter 2 which establishes the theory of finite dimensional quantum systems and Chapter 4 which establishes category theory.

1.1.3 History of the subject

Like with any sufficiently rich subject, the history of topological quantum information can be traced back as far as one wants. So let us do exactly that. The first use of topology in information science was roughly 2600 BCE, with the South American *Quipu* [AA81]. Quipu are intricate knotted strings typically made out of cotton fibers. The knots in the string are used to store various types of information, typically numbers. Mathematically we say that Quipu store their information in knot invariants, and hence hold *topological* information.

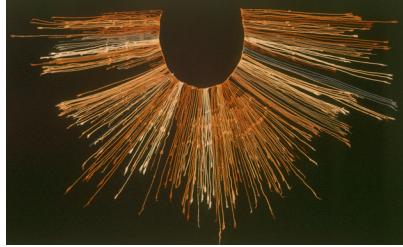


Figure 1.2: An incan quipu

Quipu were so successful that they remained the primary method of information processing in much of South America for thousands of years. They reached their peak of usage in the 15th century via the Inca empire. The Inca empire was the largest pre-Columbian empire in the western hemisphere, with over ten million subjects and spanning over two million square kilometers. Despite their intricate government, the Incas had *no written language*. This distinguished them from their contemporary empires, such as the Mali, Mongolian, or Chinese empires, which all relied on the written word. The success of the Inca empire can be seen as a testament to the versatility and power of knot invariants. The difference between the Inca and modern proposals for topological quantum computers is that instead of the strings being made out of cotton fibers they are made out of the spacetime trajectories of quasiparticles in topological systems.

Just like the history of topology in information science can be traced back to the origin of information science, the history of topology in quantum mechanics can be traced back to the origins of quantum mechanics. There is a 1931 paper of Paul Dirac [Dir31] which introduces many of the ideas which would become foundational to topological quantum mechanics. In the 1950s, explicitly topological ideas such as the Aharonov-Bohm effect [AB59] and the theory of point defects by Tony Skyrme [Sky62] were beginning to emerge. By the 1970s nontrivial abstract topological considerations were leading to novel contributions to contemporary physics, such as the theoretical description of the A-phase of superfluid Helium-3 [AT77] and the theory of phase transitions in the xy model proposed by Kosterlitz-Thouless [KT73]. These results were associated with the 1996 and 2016 Nobel prize respectively.

It was in the 1980s, however, that topology established itself as one of the leading themes in condensed matter physics. The discovery of the quantum Hall effect in 1980 [KDP80] and the subsequent discovery of the fractional quantum Hall effect in 1982 [TSG82] gave the first examples of topologically ordered systems in our modern sense of the word, and resulted in the 1985 and 1998 Nobel prizes respectively. These systems gave theorists the license to dream big about what possibilities could lie ahead. This led to major work by Frank Wilczek [Wil82b, ASWZ85], Duncan Haldane [Hal83, Hal88], and others on the theory of topological quantum systems.

The most notable of these theorists for our present story is Edward Witten, with his introduction of *topological quantum field theory* in 1988 [Wit88]. This work not only put the modern experiments within a larger context, but it also connected these developments to a parallel story which had been developing within pure mathematics. Namely, knot theory. In 1984 Vaughn Jones discovered his landmark knot invariant, which was powerful in its ability to distinguish between non-equivalent knots [Jon97]. This marked the first major progress in the field since Alexander's invariant in 1928 [Ale28]. However, Jones' construction was

steeped in opaque subfactor theory, so much so that the fact that it resulted in knot invariant felt almost like a happy accident. Hence, a widespread topic on the mind of contemporary mathematicians was how to properly interpret the Jones invariant, and how to construct other invariants like it. Witten seemed to answer both. After defining topological quantum field theory, he showed how the Jones invariant could be obtained as an observable quantity within a certain field theory [Wit89]! This shocking result gave a new interpretation of the Jones invariant in terms of mathematical physics which was appealing to experts. Seeing as the Jones invariant was constructed from a topological quantum field theory, it was natural to expect that other field theories might give new invariants which could distinguish between even more knots. This vision of invariants in low-dimensional topology constructed using topological quantum field theory became known as *quantum topology*, and evolved into its own discipline in the following years.

This brings us to 1997. Quantum topology is an active area in pure mathematics, and topological themes in condensed matter physics are at the forefront of the field. The open problem is how to construct a fault tolerant quantum computer. Peter Shor had recently discovered his factoring algorithm [Sho94], and there was debate about whether scalable quantum error correction was possible [Lan95]. This led to two independent proposals for topological quantum computation in the same year. One was by the mathematician Michael Freedman [Fre98]. His vision was clear. A recent paper had shown that computing the Jones invariant of knots was in general an NP-hard problem [JWW90]. However, by the work of Witten, the Jones invariants of knots were observables in certain topological quantum field theories. Hence, if one could construct physically a topologically ordered system which was described by Witten's topological quantum field theory then the Jones polynomial of knots could be computed efficiently by making measurements on the system. Hence, one would obtain a very powerful computer! This was Freedman's proposal.

The other proposal was made by theoretical physicist Alexei Kitaev [Kit03]. His proposal was much more precise. He gave a toy model for a certain family of topologically ordered systems. He then outlined a technique for storing and manipulating information within these systems. The deep observation was that these systems were intricate enough that they could be used to make a powerful quantum computer [Moc03].

In the subsequent years Freedman and Kitaev teamed up with collaborators Zhenghan Wang, Michael Larsen, and others to study the new field of topological quantum information and the possibility of constructing a topological quantum computer. One of the first major results was that no topological quantum computer could be more powerful than a standard quantum computer [FKW02]. This went against Freedman's original hope to solve NP-hard problems using topological quantum computers. Freedman's mistake was in asserting that topological quantum computers could compute the Jones polynomial. The measurements which give the Jones invariant in topological quantum field theory will always be *approximate*. Approximating the Jones invariant in this way is computationally easier than evaluating the Jones invariant exactly. In fact, this way of approximating the Jones invariant is *not* NP-hard - it can only be used to solve problems which could efficiently be solved using standard quantum computers.

The second major result of Freedman, Kitaev, Wang, and Larsen was the converse of their first result [FLW02]. Namely, they showed that every quantum algorithm can be efficiently run on a topological quantum computer. They do this by showing that every quantum algorithm can be efficiently reinterpreted in terms of computing the Jones invariant of some knot. In this way computing the Jones invariant is not NP-hard, but it is a *universal problem* for quantum computation. They then formalize Freedman's ideas about topological

quantum field theory, and show directly that realistic operations on a topologically ordered quantum system described by Witten's quantum field theory can be used to compute the Jones invariants of knots.

Together, these two results show in a real sense that topological quantum computing is equally powerful as standard quantum computing with circuits. This laid the groundwork for fruitful studies of fault-tolerant topological quantum computing, both using error correcting codes and physical materials. This has resulted in a great number of important results, which we will discuss at length throughout the rest of this manuscript.

1.2 Technical introduction

1.2.1 Principles of topological quantum information

In this section we will lay out the general principles of topological quantum information. As an organizational tool, these principles are introduced one by one as we construct a sample topological system. This example is meant to be representative of the systems we will encounter throughout this text, and within the broader field of topological quantum information. As a further organization tool, this example is constructed with the stated goal of obtaining a topological quantum computer.

Our system will be flat, containing only two spatial dimensions. Our system will be homogenous, essentially identical everywhere, at the exception of finitely many localized regions. These regions will differ substantially from the homogeneous bulk. These localized regions are called *quasiparticles*. The beauty of systems like these is that they behave as though the homogeneous bulk were empty, and the quasiparticles were fundamental particles. In fact, in its algebraic description, these topological systems are *identical* to ones in which the homogenous bulk is empty and the quasiparticles are fundamental particles. This is where the term quasiparticle arises. It is important however to remember that in most relevant applications the bulk is *not* empty and the quasiparticles are *not* fundamental particles. The bulk is typically some highly entangled quantum wavefunction, and the quasiparticles are emergent phenomena made up of smaller microscopic degrees of freedom.

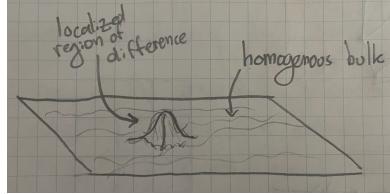


Figure 1.3: A quasiparticle in a two dimensional system

Our aim is to build a computer. In general this requires three components:

1. A method of storing information;
2. A method of manipulating information;
3. A method of reading out information.

Information is stored in the state of the system - the bulk is described by some parameters, and the details of those parameters encodes information. Our method for manipulating

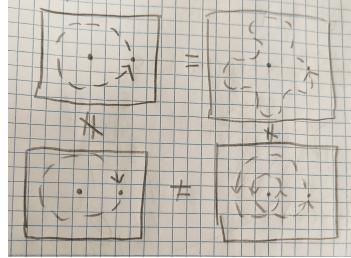


Figure 1.4: Samples of braids quasiparticles can take around one another

information is *braiding*. Braiding is the process whereby quasiparticles are moved along continuous paths around one another. There are two important points about braiding to keep in mind. The first is that braiding changes the state of the system. Even though the quasiparticles might be in identical places before and after the braid, the details of the system will change - there is more to the state of the system than just the positions of the quasiparticles. The second point is that the way that the state of the system changes *only depends on the topology of the braid*, and not the geometry. Small deformations in the path taken by the quasiparticles do not affect the result - only global changes, like whether a path is taken clockwise or counterclockwise, make a difference. This invariance is due to the fact that our system is topological. In geometric systems we expect the exact path taken by quasiparticles matters a great deal. The independence of the details of the paths is extremely specific to topological systems, and in the present setting is the *defining feature* of a topological system.

At this point we can already see we have succeeded in our goal of making our computations fault-tolerant. Noise in the system will correspond to uncontrolled perturbations in the trajectories of the quasiparticles. This uncontrolled movement won't change global properties of paths taken, and hence will not change the action of the braids on the system. That is, small errors won't affect computation! Of course, large enough errors could unintentionally make one quasiparticle wind around another. This would change the topology of the braid and hence ruin the computation. These errors are controllable, however, by moving the quasiparticles far apart and limiting the magnitude of the noise.

The final step in making our computer is to introduce a method for reading out information. This is done using *fusion*. Fusion is the process whereby two quasiparticles are brought together, resulting in a single quasiparticle. In sufficiently complicated topological systems the result of fusion depends on the details of the state of the overall system. That is, the result of fusion can be used as a way of reading out information about the state! In its most basic form, when two quasiparticles fuse they can either result in a localized region which is identical to the homogenous bulk or is different from the homogenous bulk. If they result in a localized region identical to the bulk we say that the two quasiparticles have *annihilated* each other. This can be seen as the difference between constructive and destructive interference. Two waves can either have destructive interference and annihilate each other when they meet, or they can have constructive interference and result in a new wave. Measuring whether or not two quasiparticles annihilate upon fusion gives a method for reading out information.

In some situations, the result of fusion can even be nondeterministic. In this case the fusion can be repeated multiple times, which allows one to measure the *probability* that two

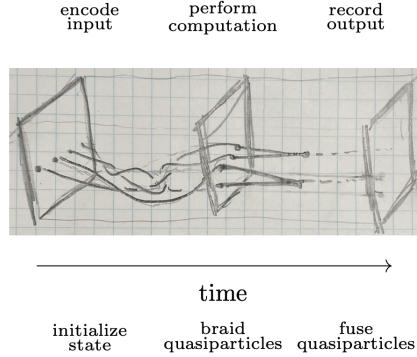
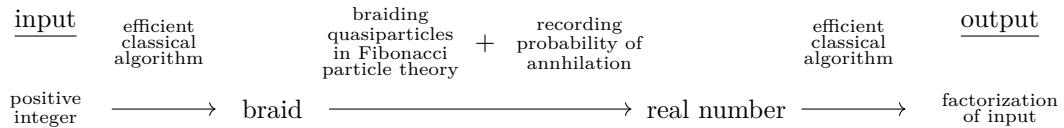


Figure 1.5: A schematic of topological quantum computing

quasiparticles will annihilate each other. These probabilities are a rich source of data, and will serve as our way of reading out information in the current setting. The fact that our system is topological implies that the result of fusion does not depend on the specifics of the path taken, and hence this method of readout preserves the invariance of our computation to noise. This gives us a full picture of topological quantum computation, as seen in figure 1.5.

To make the above discussion more concrete, we will give a worked example. In this example we use a specific topological order known as the *Fibonacci particle theory* to run Shor's efficient quantum factorization algorithm [Sho94]. The input of Shor's algorithm is a positive integer. The output of Shor's algorithm is the factorization of that integer. Shor's algorithm is *efficient* in the sense that it uses polynomially many quantum logic gates to arrive at its answer relative to the size of the input. Throughout this passage we will use *efficient* and *polynomially sized* interchangeably. The Fibonacci particle theory is a specific topological order, which describes in an algebraic fashion how the overall state changes when quasiparticles are braided and fused.

The first step in running Shor's algorithm on a Fibonacci quantum computer is to translate the positive integer input into a certain braid. This is done using an efficient classical algorithm. The second step is to run this braid on a Fibonacci quantum computer. This is done by initializing some prescribed state and then braiding its quasiparticles in the correct fashion. This initialization and braiding is performed repeatedly, and after every time two of the quasiparticles are fused. This lets us record a real number between 0 and 1, which is the probability that the two quasiparticles annihilate after the braiding. An efficient classical algorithm is then used to take this real number and obtain from it the factorization of the original input. Since all of these steps are efficient, it gives a topological quantum algorithm for factoring integers. The schematic for this process is shown below:



The magic in the above procedure is the existence of these two efficient classical algo-

rithms: a first one for encoding integers into braids and a second one for decoding real numbers into factorizations. These algorithms are nontrivial. They are due to Freedman-Larsen-Kitaev-Wang [FLW02]. In fact, Freedman-Larsen-Kitaev-Wang showed that any problem which can be efficiently solved using a quantum circuit can also be solved using the Fibonacci particle theory, via a similar method of efficient classical preprocessing and postprocessing. It is in this sense that the Fibonacci theory is *universal* for quantum computation.

The final step of this process would be to create a physical topological system which is described by the Fibonacci theory, which would serve as our quantum computer. In the realm of materials, the most promising approach seems to be specially tuned versions of the fractional quantum Hall system [ZGHS15]. While these materials are theorized to host quasiparticles described by the Fibonacci theory, the difficulty of the experiment makes them inaccessible to current technology. There has been progress made on topological quantum error correcting codes which work by simulating the Fibonacci theory [SZBV22, SBZ22, XSW⁺24]. However these codes at the current moment have structural issues and require an unbearable amount of overhead to run, making them unfeasible to use on modern computers.

Progress on topological quantum computing has thus been focused on realizing topological particle theories other than the Fibonacci theory. These other theories can be constructed in more workable materials, and can be simulated as topological quantum error correcting codes with less overhead. The drawback of these other theories is that they are typically less computationally powerful, meaning that they require more tricks and techniques to achieve universal quantum computing. There are a great number of different proposals for how to achieve universal topological quantum computing, based on different particle theories, different methods of encoding information, different methods of manipulating information, and different methods of reading out information. It is an exciting time to be a theorist in the field of topological quantum information!

1.2.2 Defects in ordered media

We will now work through a complete mathematical example of a family of topological systems. Seeing as we don't assume that the reader is familiar with quantum mechanics, our examples will be *classical* topological systems. Many of the important subtleties of topological quantum information are already present in the classical case. However, topological classical information is a smaller domain than topological quantum information - the reader should have a relatively complete grasp of the subject by the end of the chapter. Much of our discussion is taken from an excellent review article by Mermin [Mer79].

The family of systems we will describe goes by many names. In communities of experimentally focused physicists it goes by the name *ordered media*. In mathematical physics communities it goes by the name *classical field theory*. In pure mathematics it would be described as *homotopy theory*. As an input to our construction we will take a topological space M . We call M the order space of our theory. We assume throughout this chapter that M is path connected.

To describe a system in physics, the first step is to define the space of possible states of the system. In this case, states will correspond to *continuous maps* $\phi : \mathbb{R}^2 \rightarrow M$. We now give physical intuition for this choice of state space. The choice of \mathbb{R}^2 as a source represents the underlying material. We are working on an infinite flat plane. Describing a function $\phi : \mathbb{R}^2 \rightarrow M$ amounts to choosing a value $\phi(x)$ for every point $x \in \mathbb{R}^2$. In this way we

imagine our system as being made up of infinitely many objects, one placed at each point in \mathbb{R}^2 , each of which has an internal state space M . Choosing the state of the overall system amounts to choosing the state of each individual object, that is, a value in M for every point in the plane. The fact that ϕ must be continuous is a compatibility condition between the states of the objects at nearby points. It says that nearby objects must have similar states. We now list some applications of this model:

- **Classical xy model of a 2D electron gas** ([KT18]). This model describes a possible behavior electrons in a flat 2D plane. An electron can be modeled as a point particle with a magnetic dipole pointing in some direction. This magnetic dipole is known as the *spin* of the electron, and can point in any direction in the plane. The topological space of all possible directions in the plane is a circle. Hence, in this system, the order space M is the circle. The fact that nearby electrons must have similar spins is known as Hund's rule, and is the most fundamental incarnation of ferromagnetism. It is physically derived as a consequence of the Pauli exclusion principle.
- **Superfluid Helium-3** ([Lee97]). One famous example of ordered media is superfluid helium-3. Helium is an element. It has two stable isotopes: helium-3 and helium-4. The vast majority of helium on earth is helium-4, but there is also naturally occurring helium-3. At cold temperatures helium-3 undergoes a phase transition, becoming a superfluid. There are several different superfluid phases helium-3 can go into: the B-phase, dipole-locked A-phase, and dipole-free A-phase. The dipole-locked A-phase is well modeled by ordered media with order space $M = \text{SO}(3)$, where $\text{SO}(3)$ is the space of rotations in three dimensional space. The dipole-free A-phase is well modeled by ordered media with order space $M = \text{SO}(3) \times \text{SO}(3)/H$, where H is the subgroup of $\text{SO}(3) \times \text{SO}(3)$ generated by rotations across the z -axis of any degree, and simultaneous rotations of 180° on both copies of $\text{SO}(3)$ around any axis.
- **Biaxial nematics** ([Ran88]). The objects at every point in the biaxial nematic should be thought of as small rectangles with unequal side lengths. These rectangles can be oriented in any direction in three dimensional space. In practice these objects will often be molecular compounds. They will not be exactly rectangular, but they will have the same symmetry group as a rectangle which is enough for the model to be accurate. The order space M of the model is equal to the space of possible orientations of a rectangle in three dimensions. Note that in this space every orientation is equal to its 180° rotation across any axis of the rectangle, since the rectangle is symmetric under such rotations.

We now seek to extend our model of ordered media and perform a detailed analysis of it. To do this, we will need to use notions from topology such as a *continuous deformation* of a system and the *inside* of a loop in \mathbb{R}^2 . While these notions are certainly intuitive, they can be shockingly difficult to prove. The fact that every loop in \mathbb{R}^2 partitions space into an *inside* and an *outside* is known as the Jordan curve theorem. The proof of this trivial seeming statement is famously difficult [Tve80]. Thus, we adopt the following policy:

Seeing as the primary focus of this text is algebra we opt to leave our topology non-rigorous, focusing instead on principles and techniques. We only return to mathematical rigor once the subject has become sufficiently algebraic.

We now add a picture of *dynamics* into our model - how states will change through time.

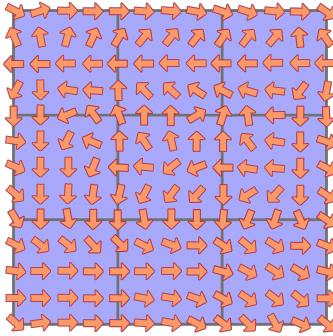


Figure 1.6: The xy-model of a 2D electron gas

In particular, we imagine that as time passes the system changes continuously. Suppose that $\phi_t : \mathbb{R}^2 \rightarrow M$ is the state of the system at time t . If t_0, t_1 are similar times we require $\phi_{t_0}(x)$ and $\phi_{t_1}(x)$ will be close. Formally, we say that a family of states ϕ_t for $0 \leq t \leq T$ form a valid trajectory in time if the map

$$\begin{aligned}\phi : & \mathbb{R}^2 \times [0, T] \rightarrow M \\ & (x, t) \mapsto \phi_t(x)\end{aligned}$$

is continuous. We also call ϕ a continuous deformation from ϕ_0 to ϕ_T , because the trajectory in time continuously changes the initial state ϕ_0 until it becomes the final state ϕ_T .

We now observe a general fact from topology: every pair of maps $\phi_0, \phi_1 : \mathbb{R}^2 \rightarrow M$ can be continuously deformed from one to the other. Intuitively, this means that as our system evolves it can transition continuously from any state to any other state. In this way, our system does not store any information which is invariant under continuous deformations. That is, it does not hold any *topologically invariant* information.

Clearly, this means that our system is not complicated enough to build a computer yet because it cannot store information. We rectify this situation by introducing quasiparticles. These quasiparticles go by many names. In the theory of ordered media they are known as defects. In field theory they are known as particles. In homotopy theory they are known as point singularities. For the sake of brevity, we use the term defect.

A defect is a point at which we drop our condition that the state $\phi : \mathbb{R}^2 \rightarrow M$ be continuous. This is done by making ϕ *undefined* at certain points. Our new system is called *ordered media with finitely many defects*. The state space consists of pairs (S, ϕ) , where $S \subset \mathbb{R}^2$ is a finite set and $\phi : \mathbb{R}^2 \setminus S \rightarrow M$ is a continuous map.

Dynamics in our new system still correspond to continuous deformations. The subtlety now is that the defects can move as the state is deformed. We call these *defect-mobile deformations*. Following our general stance on topological rigor, we omit the precise definition of defect-mobile deformation.

The vision for building our computer is that the experimenter should have control of the trajectories of the defects, but no control over the rest of the system. This means that the system will transform under defect-mobile deformations with definite defect paths

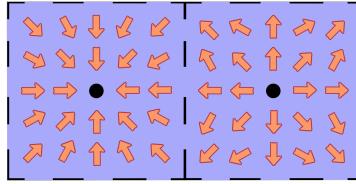


Figure 1.7: Two defects in the xy-model

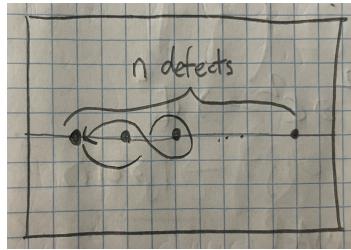


Figure 1.8: Defects on a line, with a sample trajectory of one particle around the others.

chosen by the experimenter, but the details of the rest of the deformation is arbitrarily and uncontrollable. The fact that the details of the deformation are uncontrollable comes from the fact that we expect the individual objects making up the ordered media to be noisy - they are prone to uncontrolled transformations and errors. The defects, however, are made up of many objects and are stabilized by the assumption that there is a physical mechanism which forces most nearby objects to have similar states.

We can now outline the big idea of how the computer will work. We will arrange n defects on a line in the plane. We keep these defects still, so that the system is changing only by deformations which keep the defects in place. We call these *defect-fixed deformations*. We store our information in the possible topologically-distinct states of this system:

$$(\text{information storage space}) = \binom{\text{states with } n \text{ defects}}{\text{arranged in a line}} / (\text{defect-fixed deformation})$$

The way we act on this information is by moving the defects around each other. This movement of defects induces some defect-mobile deformation. The space we are storing our information in is invariant under defect-fixed deformation, but not defect-mobile deformations. Hence, moving the defects around non-trivial paths will have non-trivial action on the stored information. This action on stored information is exactly how we perform our computations.

Finally, we must introduce a method for reading out information. This is done via fusion. Two defects can be brought together and fused. The result of this fusion is a topologically invariant quantity, and we will assume that it can be measured by an experimenter. In its most simple form, this amounts to detecting whether two defects annihilated or not. This gives us some information about the state, which is the output of our computation. This process is described visually in Figure 1.5

In the rest of this chapter we will describe exactly what the space we are storing our information in looks like, how braids act on that information, and how this can be used to make

a functioning computer. This will give a detailed picture of how topological computation works.

1.2.3 The fundamental group

To understand topological computation in ordered media we will need to put in some real work in analysing the system, and make some non-trivial observations. The structure of this analysis will be largely the same as the analysis which will be taking place throughout the rest of this book. We recall the overall outline of this text, which goes as follows:

[WORK: add outline.]

In this section we will do a very similar thing. We will take our physical model, ordered media, and take its algebraic description. That algebraic description can then be used for making a computer. Luckily for us, the algebraic theory underlying ordered media is much simpler than the algebraic theory underlying topological order: it's group theory. In particular, we will even assume that all relevant groups are *finite*. In this way, modular categories can be seen as vast quantum generalizations of finite groups. The schematic in our case is shown below:

[WORK: add outline w/ group theory instead of MTC.]

The way we go from defects in ordered media to group theory is by using a construction known as the *fundamental group* from homotopy theory.

The fundamental group is derived from a careful analysis of loops in topological spaces. We first clarify what we mean by *loop*. Loops, for our purposes, are always *oriented* and are allowed arbitrary self intersections. Examples of loops around a point are shown below:

[WORK: add pictures of loops.]

Formally, we define a loop in a topological space M to be a continuous map $\alpha : [0, 1] \rightarrow M$ such that $\phi(0) = \phi(1)$. Our main goal is to understand topological information. Topological information is stored in properties which are invariant under deformations. Hence, we are naturally interested in the space of loops up to deformation.

In the plane \mathbb{R}^2 with a point removed, points up to deformation are classified by their *winding number*. This winding number is an integer which says how many times the loop went around the point. This winding number is an integer in \mathbb{Z} , with positive numbers corresponding to counterclockwise rotations and negative numbers corresponding to clockwise rotations. The loops in figure [ref] have their winding numbers given as an illustration of the concept.

The key ingredient we are missing is the *group structure*. We want to get groups out of topological spaces, but so far all we have is a set. The group structure comes from composition. Given two loops we can compose them by first following one loop and then following the other. In this language, we see the topologists' winding-number version of $1 + 1 = 2$ below:

[WORK: add 1+1=2 with winding numbers.]

This definition has a big problem though. To compose, we need to choose a point to start and stop two loops being composed at. This special starting/stopping point is known as a *basepoint*. This deal of choosing basepoints is very important for the theory. Formally, a loop in M based at $m \in M$ is a map $\alpha : [0, 1] \rightarrow M$ such that $\phi(0) = \phi(1) = m$. We define the composition of two loops α_0, α_1 in M based at $m \in M$ to be

$$(\alpha_1 \circ \alpha_0)(t) = \begin{cases} \alpha_0(2t) & 0 \leq t \leq 1/2 \\ \alpha_1(2(t - 1/2)) & 1/2 < t \leq 1. \end{cases}$$

The reason we need to add the factors of 2 is to ensure that the domain of the loop is still the unit interval $[0, 1]$. Intuitively, to fit two loops in the same amount of time we had to speed-up both loops by a factor of 2.

We are now almost ready to define the fundamental group: we have defined based loops, and we have defined a rule for their composition. The last subtlety is in discussing what it means to deform based loops. In particular, should deformations be allowed to move the basepoint? The issue that we want to be able to compose our loops. To compose loops they need to have the same basepoint. If the basepoint moves then we will lose out composition structure. Hence, for the time being, we should only work with deformations which preserve the basepoint. With this subtlety out of the way, we can finally define the fundamental group. Given any connected topological space M and any point $m \in M$, we define the *fundamental group of M based at m* to be the group

$$\pi_1(M, m) := (\text{loops in } M \text{ based at } m) / (\text{basepoint preserving deformations})$$

whose group structure is given by the composition of based loops. As an example, our earlier comments about loops around points can be summarized as the statement that $\pi_1(\mathbb{R}^2 \setminus \{p\}, b) \cong \mathbb{Z}$ for any distinct points $b, p \in \mathbb{R}^2$. The identity element in the fundamental group is the trivial loop which stays at its basepoint and doesn't move (formally, the constant map $\alpha : [0, 1] \rightarrow M$), and inverses are given by reversing orientation (formally, the inverse of $\alpha : [0, 1] \rightarrow M$ is $\alpha^{-1}(t) = \alpha(1 - t)$.)

We can now start to use the fundamental group to analyse defects in ordered media. The first major insight is that loops in physical space yield loops in order space. Let S be a finite set of defects and let $\phi : \mathbb{R}^2 \setminus S \rightarrow M$ be a state. Given any loop α in $\mathbb{R}^2 \setminus S$ based at $b \notin S$, postcomposing with ϕ gives a loop in M :

$$(\phi \circ \alpha) : [0, 1] \xrightarrow{\alpha} \mathbb{R}^2 \setminus S \xrightarrow{\phi} M.$$

This loop has basepoint $(\phi \circ \alpha)(0) = (\phi \circ \alpha)(1) = \phi(b)$. This gives an element of $\pi_1(M, \phi(b))$. Given any state ϕ and given any loop α based at b , we call the corresponding element of $\pi_1(M, \phi(b))$ the *winding number of ϕ along α* . This sort of winding number generalizes the standard notion of a winding number of a loop around a point discussed before.

Now, consider the system with n defects arranged in a line. We can choose a basepoint b above all of the defects. We add loops α_i based at b for each $1 \leq i \leq n$, each of which go directly around defect i counterclockwise exactly once. This is depicted visually below:

[WORK: add picture.]

Given any ordered state ϕ on this system, we can take the winding numbers of each loop $\{\alpha_i\}_{i=1}^n$. These winding numbers all live in $\pi_1(M, \phi(b))$. Hence, to each state we can assign an element in the n -fold Cartesian product $\pi_1(M, \phi(b))^n$.

Fantastically, the values in $\pi_1(M, \phi(b))^n$ change in a well-behaved way under braids. We use a 2-defect system to illustrate the principle.

[WORK: add good pictures to show why making g_1 go under g_2 maps (g_1, g_2) to (g_2, g_1) . Talk a bit about it in words too]

This gives us a picture for how our computer works: information is stored in the winding numbers of the defects, and braiding acts by conjugating the winding numbers by each other.

There are still a few lingering points that need to be sorted out before we can start building our computer. The first is the issue of whether our information being stored is

actually invariant under defect-fixed deformations. Unfortunately, it is *not*. The problem is that deformations in general have no reason to preserve the value of $\phi(b)$. The deformations will hence change the basepoints. However, elements of the fundamental group are only defined up to basepoint-preserving deformations! Hence, deformations of the state will change its winding numbers. In general we have the following key relation:

$$(\text{loops in } M \text{ based at } m) / (\text{basepoint preserving deformations}) = \pi_1(M, m)$$

$$(\text{loops in } M \text{ based at } m) / (\text{arbitrary deformations}) = (\text{conjugacy classes in } \pi_1(M, m)).$$

The intuition for this above fact is as follows. Let α be a loop based at b . Let α' be the same loop but with a different choice of basepoint b' . Let ϵ be the portion of the loop between b and b' . Going along α is the same as first going along ϵ to get to b' , then going along α' , and then going along ϵ^{-1} to get back to b . Hence, we have $\alpha = \epsilon^{-1} \circ \alpha' \circ \epsilon$. Hence, choosing different basepoints amounts to conjugation. This is illustrated below:

[WORK: add diagram]

All this is to say that the winding numbers in an n -defect ordered media state ϕ are not preserved up to defect-fixed deformations. Properly dealing deformations requires properly keeping track of conjugacy classes versus group elements. This sort of effort, however, is unnecessary because it does not change any of the key takeways or any of the important concepts. Hence, *we will assume that all of our deformations do not change the value of $\phi(b)$* . We will choose some element $m \in M$ and assume $\phi(b) = m$ is fixed. For physical motivation, one can imagine moving the basepoint far away towards infinity. Since all of our physics is local we can imagine that the magnitude of the deformations go to zero away from the origin and hence the point at infinity is preserved by all local nosie.

Our last subtlety to discuss is reading out information. As we have just discussed, the values of the winding numbers in $\pi_1(M, m)$ are very dependent on the choice of basepoint. This means that the information encoded in this winding number is spread out over the whole region between the defect and the basepoint. This nonlocal nature makes it hard to measure, especially when the basepoint is taken away towards infinity. The readily measurable local information is the non basepoint-preserving winding number of the loop around defects. That is, the conjugacy classes in $\pi_1(M, m)$ associated to the defects. These conjugacy classes should be measurable in a reasonable experimental setup.

However, as we braid, these conjugacy classes do not change, and hence the outcomes of our measurements won't be affected. In a way, this is the point - braiding can change the spread-out global topological information, but will not change local quantities, like the conjugacy class in $\pi_1(M, m)$. In this way we can think of the conjugacy class as a well-defined *type* of the defect, whereas the exact value in $\pi_1(M, m)$ is a global quantity which depends on the choice of basepoint

The way to get around this issue is to fuse the defects before braiding. Fusing defects together amounts to bringing them close together until they act like a single defect. If the defects have winding number g_1, g_2 , then their fused defect will have winding number $g_1 g_2$ as illustrated by the below diagram:

[WORK: add diagram.]

Hence, given a state (g_1, g_2, \dots, g_n) , the measurable quantities are the conjugacy classes of products of adjacent defects. For instance, fusing all of the defects one-by-one to the left

would allow one to measure the conjugacy classes of g_1 , g_1g_2 , $g_1g_2g_3$, all the way up to $g_1g_2g_3\dots g_n$.

This completes our analysis of defects in ordered media based on the fundamental group.

1.2.4 Topological classical computation

We are now ready to describe the theory of topological classical computation. In the previous section, we showed how all of the topological information of defects in ordered media is controlled by the fundamental group $G = \pi_1(M, m)$ of the order space M relative to some basepoint $m \in M$. This is the heart of the algebraic theory of topological computing. Even though our physical model is complicated, the algebraic data can be succinctly summarized as a single group $G = \pi_1(M, m)$. In this way, we will formulate our discussion of topological classical computation in a way which does not make reference to the order space at all. We will choose an abstract group G and make a computer using it.

On a purely mathematical level we don't need to make any restrictions on the group G . It is a theorem from homotopy theory that every group is the fundamental group of some topological space. However, we will make restrictions on our choice of group coming from physical concerns. We will ask that our group be *finite*. A first reason for this is error protection. Suppose that G were some continuous group, like $G = \mathbb{R}$. If the stored information was $\pi = 3.1415\dots$, it would be difficult for this information to not drift to $3.1416\dots$. Even though there aren't any continuous deformations which would change a winding number, small non-continuous deformations could still make this sort of jump. For this reason it is preferable to work with *discrete* groups - groups whose natural topology is discrete. Discreteness implies a degree of separation, which means that winding numbers will not spontaneously jump from one element to the next. This gives the system error resistance.

The choice to make the group finite is more subtle. For instance, $G = \mathbb{Z}$ appears in many physically reasonable contexts. Our choice mainly stems from the practical consideration that finite groups are simpler to work with than infinite ones, and the general physical principle that symmetry groups should be compact. Any compact discrete group must be finite. This compactness argument will become especially relevant once we pass to quantum mechanics.

In summary, the algebraic theory of topological classical information we present is really just a special case of finite group theory, which is a well understood subject.

Our systems will consist of n defects on a straight line. These defects will be labeled with group elements $g_i \in G$, for $1 \leq i \leq n$. We recall that these labels correspond to winding numbers around loops:

[WORK: add diagram.]

Braiding g_i under g_{i+1} amounts to replacing g_i with g_{i+1} , and replacing g_{i+1} with $g_{i+1}^{-1}g_ig_{i+1}$. Fusing the defects g_i and g_{i+1} amounts to replacing them with a single defect labeled g_ig_{i+1} . The *type* of a defect is the conjugacy class of its label in G . These types are local observables which can be measured by the experimenter.

[WORK: add little pictures summarizing these rules]

This leads to the very natural question: *for which groups G can I make a full classical computer?* This is a design problem in finite group theory.

Our first step towards answering it is making a closer analysis of how braiding works. We define the braid group as our main tool:

[WORK: B_n = ways of braiding n points around each other/endpoint-preserving deformations. Define P_n too. Also give examples. Maybe this is where to introduce the graphical

language, things going thru time?]

Every braid can be made up piece-by-piece using individual swaps. Let σ_i denote the swap which sends the point at position i under the point at position $i + 1$. The group B_n is generated by the σ_i , for $1 \leq i \leq n - 1$. We observe that the following two braids can be deformed from one to the other:

[WORK: add Yang-Baxter braid]

Algebraically, this is the identity $\sigma_i\sigma_{i+1}\sigma_i = \sigma_{i+1}\sigma_i\sigma_{i+1}$. Additionally, if i and j are not adjacent, that is $|i - j| \geq 2$, then $\sigma_i\sigma_j = \sigma_j\sigma_i$ as seen in the below picture:

[WORK: far-commutativity picture.]

This yields the algebraic fact that

$$B_n = \langle \sigma_1, \sigma_2, \dots, \sigma_{n-1} \mid \begin{matrix} \sigma_i\sigma_{i+1}\sigma_i = \sigma_{i+1}\sigma_i\sigma_{i+1}, \\ \sigma_i\sigma_j = \sigma_j\sigma_i \end{matrix} \forall 1 \leq i, j \leq n-1, |i - j| \geq 2 \rangle.$$

It is at this point that we can move past our non-rigorous topological considerations for braiding. The lack of rigor in our treatment of braid groups and our lack of rigor in our treatment of defect trajectories cancel, giving us our first well-formed mathematical proposition. It encodes the fact that moving defects by braids acts on the stored information:

Proposition 1.1. *Let G be a finite group, and let $n \geq 1$ be an integer. The map*

$$\begin{aligned} \rho_n : B_n &\rightarrow \text{Sym}(G^n) \\ \sigma_i &\mapsto ((g_1 \dots g_i, g_{i+1} \dots g_n) \mapsto (g_1 \dots g_{i+1}, g_{i+1}^{-1} g_i g_{i+1} \dots g_n)) \end{aligned}$$

defines a homomorphism of groups between the braid group B_n and the group of set-wise permutations of the Cartesian product G^n , $\text{Sym}(G^n)$.

Proof. .[WORK: do proof] □

The dream would be that braids alone would be enough for universal classical computation. For instance, suppose that G is a finite group with a conjugacy class C that has order 2. Identifying $C \cong \{0, 1\}$ as a set, we can identify $C^n \cong \{0, 1\}^n$ with n -bit strings. Suppose the homomorphism $\rho_n : B_n \rightarrow \text{Sym}(C^n) \cong \text{Sym}(\{0, 1\}^n)$ were surjective. This would mean that every element of $\text{Sym}(\{0, 1\}^n)$ could be written in terms of braids on G . In the language of computer science, maps $\{0, 1\}^n \rightarrow \{0, 1\}^n$ are called *boolean functions*. The surjectivity of the braid group homomorphism would say that every bijective boolean function could be implemented by braiding defects, and hence every bijective computation could be done using braiding. Bijective boolean functions in the world of computation are known as *reversible*. Most problems we care about are not reversible. For instance, adding integers is not reversible because multiple summands could give the same sum as demonstrated in $1 + 3 = 2 + 2 = 4$.

This leads to two problems to deal with:

1. Braids can only ever give reversible computations. This is not a major issue because it's known that there is an efficient clever way of turning every non-reversible computation problem into a reversible one [Ben73]. However, reversibility is still cumbersome and annoying.
2. There aren't any groups G which have conjugacy classes C for which the map $\rho_n : B_n \rightarrow \text{Sym}(C^n)$ is surjective. This is because, for instance, choosing some $x \in C^n$ we

could have the whole input be a string of all x : $(x)_{i=1}^n \in C^n$. Since x commutes by itself and the braid group acts by conjugation, every braid will act trivially on this input. Hence, $\rho_n(\beta)((x)_{i=1}^n) = (x)_{i=1}^n$ for every braid $\beta \in B_n$. There is no way to make a nontrivial computation from braiding if all inputs commute.

The way around the problem of commuting inputs is to use pair-creation. Just like how at any time two defects (g_1, g_2) to make the defect g_1g_2 , at any time this process could reverse and out of nothing the pair of defects (g, g^{-1}) could be created. The important part of this process is that g can be chosen so that it does not commute with the inputs.

[WORK: add a bit of explanation about this, introducing all the ingredients we need]

We now have all of our ingredients, we can finally state our main results about making computers. Clearly, if G is abelian then it won't be useful for making a computer. Our computer is based on acting by conjugation, and conjugation is trivial in abelian groups. However, just being nonabelian is not enough. The group needs to be *sufficiently non-abelian* so that conjugation is powerful enough to implement any computation possible.

We define the subgroup

$$[G, G] = \{\text{subgroup of } G \text{ generated by elements of the form } g_0g_1g_0^{-1}g_1^{-1} \text{ for } g_0, g_1 \in G\}.$$

The subgroup $[G, G]$ is called the *commutator subgroup* of G , and the elements $g_0g_1g_0^{-1}g_1^{-1}$ are called commutators. Clearly, if G is abelian then all of its commutators are trivial and hence $[G, G] = 0$. To make a good computer, we need sufficiently many commutators. In fact, we need the following condition: $[G, G] = G$. For technical reasons it also easiest to assume that G is *simple*. That is, it has no proper nontrivial normal subgroups.

We observe that every non-abelian simple group is automatically perfect. This follows from the fact that the commutator subgroup $[G, G]$ is a normal subgroup of G . Since G is nonabelian the commutator subgroup $[G, G]$ is nontrivial, and since the only nontrivial subgroup of G is itself we get that $[G, G] = G$.

In this case we have the following important result:

Theorem 1.1 (Mochon). . [WORK: I want a good statement of Mochon's theorem, need to think more about it and read Mochon's paper again. State it as "for every non-abelian simple group G ".]

Proof. . [WORK: do proof.] □

We can now give a full example of how this works. [WORK: give a full example, using $G = A_5$. Of course, full is a tricky word here. Enough to make what's going on clear to people.]

Now that we know how to make a universal computer, we can analyse how the power of computation changes as G is chosen to be more or less abelian. The general theme is that if a group is more non-abelian then it will have more computational power. Of course, being "more" or "less" abelian is not a well-defined term. We introduce here some formal notions from group theory which measure abelianness.

If a group is very nonabelian it should have a big commutator subgroup. That is, $[G, G]$ should be large in a certain sense. One way for this to be true is for the group to be perfect, that is, $[G, G] = G$. A weaker condition is that G should have a perfect subgroup - a subgroup $H \leq G$ such that $[H, H] = H$. Intuitively, it makes sense that any group with a perfect subgroup should be useable for universal topological computation. You can just

focus on the perfect subgroup and use Mochon's theorem, and forget about the rest of the group. A group with a perfect subgroup is called *non-solvable*.

There is another useful step to consider between non-solvable and abelian. Some groups G have subgroups H such that $[H, H]$ might be small, but $[H, G]$ is bigger. When H is allowed to take commutators with elements of G you get potentially more elements, so $[H, H] \leq [H, G]$. When H is a normal subgroup, we find that $[H, G] \leq H$ since for all $h \in H, g \in G$, the commutator

$$ghg^{-1}h^{-1} = (ghg^{-1})h \in H$$

since $(ghg^{-1}) \in H$ by the normality of H and $h \in H$ by assumption. If G has a normal subgroup H such that $[H, G] = H$, we call G *non-nilpotent*. Clearly we have the following inclusions:

$$\left(\begin{smallmatrix} \text{non-solvable} \\ \text{groups} \end{smallmatrix} \right) \subset \left(\begin{smallmatrix} \text{non-nilpotent} \\ \text{groups} \end{smallmatrix} \right) \subset \left(\begin{smallmatrix} \text{non-abelian} \\ \text{groups} \end{smallmatrix} \right).$$

This induces a hierarchy of adjectives, from most abelian to least abelian:

$$\begin{array}{cccc} (\text{abelian}) & \left(\begin{smallmatrix} \text{nilpotent} \\ \text{non-abelian} \end{smallmatrix} \right) & \left(\begin{smallmatrix} \text{solvable} \\ \text{non-nilpotent} \end{smallmatrix} \right) & (\text{non-solvable}) \\ \text{more} \leftarrow \text{abelian} & & & \rightarrow \text{less abelian} \end{array}$$

The following phenomenon presents itself. We find that less abelian a group is the more computational power it has. Additionally, the more computational power it has the harder it is to create topological systems in the lab which are algebraically described that group. It is harder to make systems which make good computers.

This phenomenon is especially well developed in the quantum case. Given a finite group G , we can describe a classical system of ordered media based on G . This system can be *quantized*. This turns it into a topological quantum system whose behavior is still governed by the group G . This quantized system is known as the *quantum double* of G , and is denoted $\mathfrak{D}(G)$. These systems behave very similarly to the ones discussed in this chapter, just made quantum. These quantum doubles and the algebraic theory describing them and their generalizations will be the topic of much of this book.

In the case of quantum doubles we can make a table detailing exactly the relationship between level of abelianness, computational power, and experimental status:

[WORK: need to introduce universality as a concept]

Abelianness	Smallest Example	Computational power of $\mathfrak{D}(G)$	Experimental Status
non-solvable	$G = A_5$, alternating group with $ A_5 = 60$	Straightforwardly universal. (Chapter [ref], [Moc03])	Fundamental limitations coming from intensive circuit-depth requirements, and the size of the smallest example. ([BKK22]).
solvable non-nilpotent	$G = S_3$, symmetric group with $ S_3 = 6$	Universal with tricks. (Chapter [ref], [Moc04])	Has yet to be done. There are some inherent difficulties involved. ([TVV23])
nilpotent non-abelian	$G = D_4$, dihedral group with $ D_4 = 8$?	Preliminary experiments have been successful. ([ITV ⁺ 24]) ¹
abelian	$G = \mathbb{Z}_2$, cyclic group with $ \mathbb{Z}_2 = 2$	Universal schemes seem to be impossible. Non-topological methods seem to be required. (Chapter [ref], [BK13, EK09])	Widely used in most applications. ([BCG ⁺ 24, HDSHL24] [BKK ⁺ 24, goo23])

We make a few comments about this table.

1. We notice that difficulty inherent to experimentally realizing topological phases is *not* completely controlled by the size of that group. The quantum double $\mathfrak{D}(D_4)$ is simpler to realize than $\mathfrak{D}(S_3)$ because it is nilpotent and S_3 is not, despite the fact that $|D_4| = 8$ is larger than $|S_3| = 6$.
2. All of the experimental results cited come from the side of topological quantum error correction and not topological quantum materials. This is because most topological quantum materials are described by algebraic theories which are not doubles of finite groups. Doubles of finite groups are primarily used in topological quantum error correction theory.
3. This table details a general programme. Given an algebraic theory of topological information, there is the question of how to make a universal quantum computer. The culmination of this book is Chapter [ref], where we show describe six different families of algebraic theories and show how to make a universal quantum computer out of all of them.

This concludes our overview of the subject!

Exercises:

- 1.1. .[WORK: show that simulability \implies *efficient* simulability. Do the work in the case $G = A_5$]
- 1.2. .[WORK: show that nilpotent \implies polynomial growth? Can reference the more general picture of size of braid group images.]

2 Quantum mechanics

2.1 Overview

2.1.1 Introduction

In this chapter we will give an introduction to quantum mechanics. The goal of this book is to give an exposition of topological quantum information. So far we have described topological *classical* information - all that's missing now is quantum mechanics!

One of the difficulties of quantum mechanics is that it is physically unintuitive to most uninitiated learners. Conversely, one of the advantages of quantum mechanics is that it is mathematically basic. Quantum mechanics is mathematically linear algebra. The mathematical intricacies of quantum mechanics often arise from complications from working in infinite dimensional spaces. In topological quantum information, however, all of the spaces of interest are finite dimensional and hence the mathematics involved is quite straightforward: finite dimensional linear algebra is largely a solved subject. In this chapter we will give a dictionary between the physical language of quantum mechanics and the mathematical language of linear algebra.

The first physical principle about quantum mechanics to know is that it is typically used to describe small objects. A natural question is *why*. If quantum mechanics is correct, then it should equally well apply to small and large objects. The answer to this question is subtle, and brings us back to the thesis of this book.

Large scale macroscopic phenomena are emergent from coherent small scale microscopic phenomena. The word *coherent* is used intentionally. It is used to mean “held together”, “integrated”, or “organized”. Sometimes collections of microscopic degrees of freedom fail to form observable macroscopic degrees of freedom. This failure is known as *decoherence*. It is an empirically observed fact that microscopic quantum degrees of freedom typically decohere. It is the ubiquity of decoherence which makes the macroscopic world seem classical.

It is exactly for this reason that topological quantum systems are so special. They are essentially unique in the fact that they can coherently hold quantum information at macroscopic length and time scales. This is because decoherence is caused by repeated noise from the environment, which corrupts fragile quantum information. Topological quantum systems are defined by the property that their stored information is not affected by local changes. Hence, if noise is sufficiently local and sufficiently controlled, the information in topological quantum systems will remain coherent.

This makes topological quantum matter a fantastic place to first learn quantum theory. The mathematics is simple because all spaces involved are finite dimensional, and the quantum effects are more dominant than in almost any other macroscopic phenomena! It is an exciting and rich subject.

2.1.2 Experimental motivation

Before diving into a formal treatment of quantum mechanics, let us first motivate why quantum mechanics has to be like it is. The most famous aspect of quantum mechanics is its probabilistic nature. As Einstein famously said, “*God does not play dice*”. If quantum mechanics was just probabilistic, however, it wouldn't bother physicists nearly as much as it does. Quantum mechanics is a sort of twisted probability theory:

“What happens if you try to come up with a theory that’s *like* probability theory, but based on the 2-norm instead of the 1-norm?... Quantum mechanics is what inevitably results.” - Scott Aaronson²

Throughout this introduction to quantum mechanics we will take the lens of comparing quantum mechanics with classical probability theory. Some properties of quantum mechanics, like *superposition* and *entanglement*, are already clearly present in the world of probability. Other properties, like *interference*, are not. To make this clear, we will present a few experiments which demonstrate the probabilistic nature of quantum mechanics, and the ways in which quantum mechanics goes beyond probability theory.

[WORK: which experiments should I chose? Double slit? Polarized light? Pairs of entangled photons? It would be cool to get experiments which are relevant to topological matter if possible. It would also be cool to get experiments which almost immediately motivate the exact form of quantum mechanics. I’m not a physicist though - need to get someone else more knowledgeable to give me a lecture.]

2.2 Axiomatic development

2.2.1 Probability theory

Seeing as quantum mechanics is a modified probability theory, before axiomatizing quantum mechanics we will first axiomatize probability theory in terms of linear algebra. The goal is to highlight what an axiomatization of a physical theory should look like, so that the jump to quantum mechanics is as predictable as possible.

Intuitively, we all know what probability theory is. We start with some set S which represents the possible outcomes of our probability theory. States in the probabilistic system are probability distributions on S . That is, assignments of probabilities (positive real numbers) to each elements of S such that the total probability is 1. We will focus entirely on *finite* probability spaces. This greatly simplifies our analysis. Finite probability spaces require only basic linear algebra to describe, whereas infinite probability spaces requires measure theory.

For example, suppose we are flipping a coin. The space of possible outcomes is $S = \{\text{head}, \text{tails}\}$. A fair coin flip would have $50\% = 1/2$ probability of giving heads, and $50\% = 1/2$ probability of giving tails.

A convenient notation for probability distribution is the language of weighted sums. The state $\sum_{x \in S} p_x |x\rangle$ denotes the state with probability $p_x \geq 0$ of having outcome $|x\rangle$, where $\sum_{x \in S} p_x = 1$. In the case of heads and tails, we would write

$$|\text{fair flip}\rangle := \frac{1}{2} |\text{heads}\rangle + \frac{1}{2} |\text{tails}\rangle.$$

The notation $|\cdot\rangle$ for states is known as a *ket*. This is part of so-called *Dirac notation*, which is widespread in quantum theory. We use it here to help ease our transition from probability theory to quantum mechanics.

Mathematically a formal sum is an element of a vector space. That is, the weighted sums corresponding to probability distributions are elements of the vector space

$$\mathbb{R}[S] := \text{span} \{ |x\rangle | x \in S \}.$$

²Page 112 of Aaronson’s “Quantum Computing since Democritus” [Aar13]

For convenience we will refer to elements of $\mathbb{R}[S]$ of the form $\sum_{x \in S} p_x |x\rangle$ with $p_x \geq 0$, $\sum_{x \in S} p_x = 1$ as *normalized vectors*. Our discussion can be summarized as saying that probability distributions on S correspond to normalized vectors in $\mathbb{R}[S]$.

We now move on to discussing the way that probability spaces can evolve, or be related to one another. Certainly, a relation between a probability space with outcomes S and a probability space with outcomes S' will be some function

$$(\text{normalized vectors in } \mathbb{R}[S]) \rightarrow (\text{normalized vectors in } \mathbb{R}[S'])$$

which gives a rule for going from probability distributions on S to probability distributions on S' . However, not every function will give a valid assignment. For example, suppose we are studying the outcomes of lottery tickets. Ticket 1 has an 80% chance of being a winner, and Ticket 2 has a 40% of being a winner. You haven't scratched your ticket yet, so you know you have a 50% chance of having Ticket 1 and a 50% chance of having Ticket 2. What is the probability that you win the lottery? The standard way of computing it would be as follows:

$$\begin{aligned} \text{result}(|\text{your ticket}\rangle) &= \text{result}\left(\frac{1}{2}|\text{Ticket 1}\rangle + \frac{1}{2}|\text{Ticket 2}\rangle\right) \\ &= \frac{1}{2}\text{result}(|\text{Ticket 1}\rangle) + \frac{1}{2}\text{result}(|\text{Ticket 2}\rangle) \\ &= \frac{1}{2}\left(\frac{4}{5}|\text{win}\rangle + \frac{1}{5}|\text{lose}\rangle\right) + \frac{1}{2}\left(\frac{2}{5}|\text{win}\rangle + \frac{3}{5}|\text{lose}\rangle\right) \\ &= \frac{3}{5}|\text{win}\rangle + \frac{2}{5}|\text{lose}\rangle. \end{aligned}$$

Hence, you have a $3/5 = 60\%$ chance of winning. The key insight in this computation was that probabilistic processes are *linear*. That is, “result” induces a linear map from $\mathbb{R}[\{\text{Ticket 1}, \text{Ticket 2}\}]$ to $\mathbb{R}[\text{win, lose}]$. More generally, given finite sets S, S' any linear map $\mathbb{R}[S] \rightarrow \mathbb{R}[S']$ which sends normalized vectors to normalized vectors could represent some valid probabilistic process.

The final topic to tackle before giving the full axiomatization is the question of *joining* probabilistic systems. In this book we will mostly be constructing systems out of a lot of smaller constituent parts, so the question of fitting together smaller systems to make one larger system is of utmost importance. Suppose we have two smaller systems with possible outcomes S, S' . To describe a state in the joined system, it is necessary and sufficient to describe how that state restricts to each subsystem. In this way, possible outcomes of the joined system will correspond to pairs (x, x') where $x \in S$ is the portion of the overall state in S and $x' \in S'$ is the portion of the overall state in S' . This means the space of outcomes in the joined system is the Cartesian product $S \times S'$.

We are now ready to state the full axioms of probability theory:

Definition (Axioms of probability theory).

1. (Systems) A probabilistic system is a real vector space of the form $\mathbb{R}[S]$, where S is a finite set. Valid states are normalized vectors in $\mathbb{R}[S]$, which we call probability distributions on S .
2. (Processes) A probabilistic process going from a system S to a system S' is a linear map $\mathbb{R}[S] \rightarrow \mathbb{R}[S']$ which sends normalized vectors to normalized vectors.

3. (Joining systems) If S and S' are two finite sets, the system obtained by joining $\mathbb{R}[S]$ and $\mathbb{R}[S']$ is $\mathbb{R}[S \times S']$.
4. (Measuring systems) Given a normalized vector $\sum_{x \in S} p_x |x\rangle \in \mathbb{R}[S]$, measurement corresponds to collapsing onto an outcome, where we collapse into each $x \in S$ with probability p_x .

□

2.2.2 Basis-dependent quantum mechanics

The basis-dependent version of quantum mechanics can be established by copying the axioms of probability theory almost verbatim, replacing the 1-norm with the 2-norm.

Given a finite set S , a normalized vector in $\mathbb{R}[S]$ is one of the form $\sum_{x \in S} p_x |x\rangle$, where $p_x \geq 0$ and $\sum_{x \in S} p_x = 1$. This quantity $\sum_{x \in S} p_x$ is known as the *1-norm* of the vector $p = (p_x)_{x \in S}$.

In quantum mechanics we re-define the notation of normalized vector. A normalized vector in quantum mechanics is a state $\sum_{x \in S} c_x |x\rangle$, where $c_x \in \mathbb{C}$ are arbitrary complex numbers and $\sum_{x \in S} |c_x|^2 = 1$. The root of the sum of norm-squares $\sqrt{\sum_{x \in S} |c_x|^2}$ is known as the *2-norm* of the vector $c = (c_x)_{x \in S}$. In this way, the norm-squares $|c_x|^2$ form a probability distribution on S .

Thus, given some finite set S , states in the quantum system based on S correspond to normalized vectors in $\mathbb{C}[S]$. As a matter of convention, normalized vectors in $\mathbb{R}[S]$ will always refer to the 1-norm definition and normalized vectors in $\mathbb{C}[S]$ will always refer to the 2-norm definition. We are now ready to state the basic axioms of quantum theory, with the caveat that it does not give the full picture of measurement:

Definition (Axioms of quantum mechanics, basis dependent version).

1. (Systems) A quantum system is a complex vector space of the form $\mathbb{C}[S]$, where S is a finite set. The normalized vectors in $\mathbb{C}[S]$ correspond to quantum states on S . Here, a *normalized* vector $v = \sum_{x \in S} c_x |x\rangle$ is one for which $\sum_{x \in S} |c_x|^2 = 1$, where $|c_x|^2$ denotes the norm square.
2. (Processes) A quantum process going from a system S to a system S' is a linear map $\mathbb{C}[S] \rightarrow \mathbb{C}[S']$ which sends normalized vectors to normalized vectors.
3. (Joining systems) If S and S' are two finite sets, the system obtained by joining $\mathbb{C}[S]$ and $\mathbb{C}[S']$ is $\mathbb{C}[S \times S']$.
4. (Measuring systems) Given a normalized vector $\sum_{x \in S} c_x |x\rangle \in \mathbb{C}[S]$, measurement corresponds to collapsing into a pure state, where we collapse into each $x \in S$ with probability $|c_x|^2$.

□

We now relate these axioms to the previous discussion and introduce terminology. The formal sums $\sum_{x \in S} c_x |x\rangle$ are not probability distributions. They are called *wavefunctions*. Every state in quantum mechanics is encoded in a wavefunction. Treating the possible outcomes in S as positions, we get the analogy

- Wave = multiple positions, spread-out $= \sum_{x \in S} c_x |x\rangle \in \mathbb{C}[S]$;

- Particle = single positions, definite = $|x\rangle$, $x \in S$.

By axiom (4), measuring of wavefunction collapses it into a single particle. This is the essence of wave-particle duality in quantum mechanics. The numbers c_x are not probabilities. They are called *amplitudes*. If a state $|\psi\rangle = \sum_{x \in S} c_x |x\rangle$ has non-zero amplitude at $x, y \in S$, then we say that $|\psi\rangle$ is in a *superposition* of being in state $|x\rangle$ and $|y\rangle$.

Within this framework it is easy to demonstrate the phenomenon of interference. Define the transformation $M : \mathbb{C}[S] \rightarrow \mathbb{C}[S]$ by

$$M(|0\rangle) = \frac{1}{\sqrt{2}} |0\rangle + \frac{1}{\sqrt{2}} |1\rangle,$$

$$M(|1\rangle) = \frac{1}{\sqrt{2}} |0\rangle - \frac{1}{\sqrt{2}} |1\rangle.$$

Applying M to $|0\rangle$ and measuring gives 0 and 1 with equal probability, and same with applying M to $|1\rangle$. When we apply M to the equal superposition of 0 and 1, however, this results in the state

$$H\left(\frac{1}{\sqrt{2}}|0\rangle + \frac{1}{\sqrt{2}}|1\rangle\right) = \frac{1}{\sqrt{2}}\left(\frac{1}{\sqrt{2}}|0\rangle + \frac{1}{\sqrt{2}}|1\rangle\right) + \frac{1}{\sqrt{2}}\left(\frac{1}{\sqrt{2}}|0\rangle - \frac{1}{\sqrt{2}}|1\rangle\right) = |0\rangle.$$

We can summarize this as saying that there was *constructive interference* in the $|0\rangle$, and *destructive interference* in the $|1\rangle$. The amplitudes had the same signs in the $|0\rangle$ causing the probability of measuring 0 to add, and the amplitudes had opposite signs in the $|1\rangle$, causing the probabilities of measuring 1 to cancel.

2.2.3 Measurement

The axioms in the previous section are all accurate, but they do not give a complete picture of measurement in quantum theory. In particular, the type of measurement which takes a state $\sum_{x \in S} c_x |x\rangle$ and collapses it to $|x\rangle$ with probability $|c_x|^2$ is only a special type of measurement. There are key subtleties that are ignored in our naive treatment:

1. It is possible to measure with respect to bases other than the standard basis;
2. Measurements can be incomplete, meaning that they do not collapse a wavefunction all the way down to a particle;
3. Measurements always have *observables* associated with them.

The easiest point to discuss is observables. Every time you measure something in a laboratory, there is always a real number output associated with the measurement:

- If you measure the velocity of a particle, the output is a speed in meters/second;
- If you measure the relative position of two objects, the output is a distance in meters;
- If you measure the intensity of a light source, the output is a luminescence in candela/square meter;
- etc, etc...

Seeing as these real numbers are the only quantities which we actually get to record as experiments, we have to incorporate them into our theory. For example, consider some finite set S with associated quantum system $\mathbb{C}[S]$. Suppose we measure the energy of the system in joules (J). Since S is finite there are finitely many possibilities for the energy, say 1J, 5J, 10J. In a quantum system, measuring with respect to energy will produce some output (1J, 5J, or 10J) and collapse the system onto a state with a well-defined energy.

A crucial point is that these states with well-defined energy have *absolutely no reason* to be the same as the elements of S . Different observables can have different collections of states with well-defined values of those observables. A state with a well-defined value of some observable is called an *eigenstate* of that observable. This will connect back to our usual notation of eigenvector from linear algebra.

As an example, suppose $S = \{0, 1\}$. We define an observable called energy. We say that the state $\frac{1}{\sqrt{2}}|0\rangle + \frac{1}{\sqrt{2}}|1\rangle$ has energy 2J and the state $\frac{1}{\sqrt{2}}|0\rangle - \frac{1}{\sqrt{2}}|1\rangle$ has energy 3J. The state $|0\rangle$ can be decomposed as

$$|0\rangle = \frac{1}{\sqrt{2}} \left(\frac{1}{\sqrt{2}}|0\rangle + \frac{1}{\sqrt{2}}|1\rangle \right) + \frac{1}{\sqrt{2}} \left(\frac{1}{\sqrt{2}}|0\rangle - \frac{1}{\sqrt{2}}|1\rangle \right).$$

We see here that $|0\rangle$ is in an equal superposition of the state with energy 2J and the state with energy 3J. When we measure this state, it will collapse onto some energy eigenstate. It will collapse onto $\frac{1}{\sqrt{2}}|0\rangle + \frac{1}{\sqrt{2}}|1\rangle$ with probability 1/2 and it will collapse onto $\frac{1}{\sqrt{2}}|0\rangle - \frac{1}{\sqrt{2}}|1\rangle$ with probability 1/2, depending on the value of energy that was measured.

It is important that one needs to take care when defining observables to make sure that no contradictions appear. For instance, once the values of the observable are specified on a basis then the rest of the values of the observable follow by linearity. A more subtle restriction is seen in the following example. Suppose that $|0\rangle$ is given energy 2J and $\frac{1}{\sqrt{2}}|0\rangle + \frac{1}{\sqrt{2}}|1\rangle$ is given energy 3J. Then, we can write

$$|1\rangle = -\sqrt{2}(|0\rangle) + \sqrt{2} \left(\frac{1}{\sqrt{2}}|0\rangle + \frac{1}{\sqrt{2}}|1\rangle \right).$$

In this way, $|1\rangle$ has energy 2J with amplitude $-\sqrt{2}$ and energy 3J with amplitude $+\sqrt{2}$. Clearly, the norm squares of these amplitudes does not give a valid probability distribution. The key algebraic requirement is *orthogonality*. Namely, we have an *inner product* on $\mathbb{C}[S]$ defined by

$$\left\langle \sum_{x \in S} c_x |x\rangle \middle| \sum_{x \in S} c'_x |x\rangle \right\rangle = \sum_{x \in S} c_x \overline{c'_x}.$$

Two states in $\mathbb{C}[S]$ are called *orthogonal* if their inner product is 0. If the values of an observable are specified with respect to a basis in which every basis vector is normalized and every pair of basis vectors is orthogonal, then this observable can be extended to all normalized vectors in $\mathbb{C}[S]$ without issues. Before stating this axiom formally, we introduce some notation. If a basis of $\mathbb{C}[S]$ consists of normalized pairwise orthogonal vectors, we call it *orthonormal*. An *observable* on $\mathbb{C}[S]$ is a pair (B, v) where $B \subset \mathbb{C}[S]$ is an orthonormal basis and $v : B \rightarrow \mathbb{R}$ is a set function.

This gives us our next version of the axioms of quantum mechanics. There are issues that arise when v is not injective, so we state our axioms with a restriction on v for now:

- 3'. (Measuring systems) Let (B, v) be an observable for which v is injective. The system $\mathbb{C}[S]$ can be measured with respect to (B, v) . When $|\psi\rangle = \sum_{b \in B} c_b |b\rangle \in \mathbb{C}[S]$ is measured with respect to (B, v) , the state collapses to each $|b\rangle$, $b \in B$, with probability $|c_b|^2$. In the case that $|\psi\rangle$ collapses onto $|b\rangle$, we say that the outcome of the measurement is $v(b) \in \mathbb{R}$.

We will verify that the values $|c_b|^2$ indeed form a probability distribution later in the section.

2.2.4 Incomplete measurement

The above discussion is still missing some generality. Namely, it ignores the fact that measurements can be *incomplete*. Incomplete measurements arise when two linearly independent vectors have the same value of an observable. When the observable is measured, it doesn't know which of those two linearly independent vectors to collapse to! In this situation, we say that the observable is *degenerate*. The term degeneracy here comes from its general mathematical usage, whereby it used to describe edge cases where not-necessarily-equal values happen to be equal.

Instead of collapsing all the way down to an eigenstate, the measurement of degenerate observables will project a state onto the subspace spanned by the eigenstates with the measured value of the observable. For example, let $S = \{0, 1, 2\}$. Suppose that the state $|0\rangle$ has energy 5J, and that the states $|1\rangle$ and $|2\rangle$ have energy 10J. Suppose further that we measure the state

$$\frac{1}{\sqrt{3}}|0\rangle + \frac{1}{\sqrt{3}}|1\rangle - \frac{i}{\sqrt{3}}|2\rangle$$

with respect to energy, and the observed value is 10J. This will collapse the state onto $\frac{1}{\sqrt{2}}|1\rangle - \frac{i}{\sqrt{2}}|2\rangle$. The projection respects phases, but scales the absolute value of the state so that it becomes normalized. Formally, this is an orthogonal projection. To state this axiom it is good to introduce some notation. Let S be a finite set and let $|\psi\rangle$, $|\varphi\rangle$ be states in $\mathbb{C}[S]$. We use the notation

$$\langle |\psi\rangle | |\varphi\rangle \rangle := \langle \psi | \varphi \rangle .$$

This gives us a complete description of measurement in quantum mechanics:

- 3''. (Measuring systems) Let (B, v) be an observable. The system $\mathbb{C}[S]$ can be measured with respect to (B, v) . Let $|\psi\rangle = \sum_{b \in B} c_b |b\rangle \in \mathbb{C}[S]$ be a state, and let $\lambda \in \mathbb{R}$ be a real number. The probability that the outcome of the measurement is equal to λ is $\sum_{v(b)=\lambda} |c_b|^2$. In this case, the state $|\psi\rangle$ will collapse onto

$$\left(\sum_{v(b)=\lambda} c_b |b\rangle \right) / \left(\sum_{v(b)=\lambda} |c_b|^2 \right).$$

2.2.5 Basis-independent quantum mechanics

From our discussion of measurement it is clear that, unlike probabilistic systems, quantum systems do not have a favored choice of basis. However, our definition of quantum system is still woefully basis-dependent. Namely, it starts by choosing a distinguished basis S of

$\mathbb{C}[S]$. What would be better if we could remove this choice, and make a quantum system simply a vector space.

This poses some immediate problems however. The first is that vector spaces have no notion of norm. Hence, we cannot speak of normalized vectors, and hence we cannot speak of states. What's more, measurements are required to use an orthonormal basis. To define orthogonality we used the canonical inner product on $\mathbb{C}[S]$. Without a basis there is no distinguished choice of inner product. However, in a real sense that is the *only* piece of information we need about our basis - its inner product. This means that we can state the axioms of quantum mechanics for any vector space with a distinguished choice of inner product. We define what it means for a space to have an inner product below:

[WORK: define Hilbert space]

In any Hilbert space V , we can define the 2-norm of a vector $|\psi\rangle \in V$ to be

$$||\psi\rangle| = \sqrt{\langle\psi|\psi\rangle}$$

A normalized vector in a Hilbert space is any state for which $||\psi\rangle| = 1$. Observe that this agrees with our previous definition of normalized vector. If B is any orthonormal basis of V and $|\psi\rangle = \sum_{b \in B} c_b |b\rangle$, then

$$\begin{aligned} \langle\psi|\psi\rangle &= \left\langle \sum_{b \in B} c_b |b\rangle \middle| \sum_{b \in B} c_b |b\rangle \right\rangle \\ &= \sum_{b_0, b_1 \in B} c_{b_0} \bar{c}_{b_1} \langle b_0 | b_1 \rangle \\ &= \sum_{b \in B} |c_b|^2. \end{aligned}$$

Thus, $||\psi\rangle| = 1$ if and only if $\sum_{b \in B} |c_b|^2 = 1$ relative to any (equivalently, all) orthonormal bases.

The quantum process and quantum measurement axioms are obvious to state in any Hilbert space. The difficulty is in the joining axiom. It's here that we observe that for any finite sets S, S' , there is a canonical isomorphism

$$\begin{aligned} \mathbb{C}[S \times S'] &\cong \mathbb{C}[S] \otimes \mathbb{C}[S'] \\ |(x, x')\rangle &\mapsto |s\rangle \otimes |s'\rangle \end{aligned}$$

where \otimes is the tensor product. For those unfamiliar with the tensor product, this could be taken as the *definition* of it. We note that the tensor product of two Hilbert spaces $(V, \langle \cdot | \cdot \rangle_V), (V', \langle \cdot | \cdot \rangle_{V'})$ is a Hilbert space. The inner product on $V \otimes V'$ is given by

$$\langle(v \otimes v')|(w \otimes w')\rangle_{V \otimes V'} = \langle v | w \rangle_V \cdot \langle v' | w' \rangle_{V'}.$$

This leads us to the following basis independent formulation of the axioms of quantum mechanics:

Definition (Axioms of quantum mechanics, basis independent version).

1. (Systems) A quantum system is a complex Hilbert space V

2. (Processes) A quantum process going from a system V to a system W is a unitary transformation from V to W
3. (Joining systems) If V and W are two quantum systems, the system obtained by joining V and W is $V \otimes W$.
4. (Measuring systems) Let (B, v) be an observable. The system V can be measured with respect to (B, v) . Let $|\psi\rangle = \sum_{b \in B} c_b |b\rangle \in V$ be a state, and let $\lambda \in \mathbb{R}$ be a real number. The probability that the outcome of the measurement is equal to λ is $\sum_{v(b)=\lambda} |c_b|^2$. In this case, the state $|\psi\rangle$ will collapse onto

$$\left(\sum_{v(b)=\lambda} c_b |b\rangle \right) \Bigg/ \left(\sum_{v(b)=\lambda} |c_b|^2 \right).$$

□

Now that we have stated our final version of the axioms of quantum mechanics, we make some technical comments which aid in our future endeavors. The first is that operators which send normalized states to normalized states have a very concise characterization in terms of the *conjugate transpose*. Of course, without a basis we have no way of identifying linear operators with matrices, and hence no way of defining the transpose. Given a Hilbert space V and a linear map $M : V \rightarrow V$ there may be no way to define the transpose but there *is* a way of defining the component-wise conjugate transpose of V . This conjugate transpose is denoted M^\dagger , and is defined by the inner-product formula

$$\langle U\psi | \varphi \rangle = \langle \psi | U^\dagger \varphi \rangle.$$

It is verified in Exercise [ref] that this formula always specifies a unique well-defined operator, and that this operator is equal to the conjugate transpose of V relative to any orthonormal basis. Here is our characterization of maps which send normalized vectors to normalized vectors:

Proposition 2.1. *Let V be a Hilbert space, and let $U : V \rightarrow V$ be a linear transformation. The following are equivalent:*

1. U sends normalized vectors to normalized vectors;
2. $U^\dagger = U^{-1}$.

If either of these two equivalent conditions are met, we call U a unitary transformation.

Proof. We observe that if $U^\dagger = U^{-1}$, then for any normalized vector $|\psi\rangle$

$$|U|\psi\rangle| = \langle U\psi | U\psi \rangle = \langle \psi | U^\dagger U\psi \rangle = \langle \psi | \psi \rangle = 1.$$

Hence, (2) \implies (1). To show the other direction, suppose that U sends normalized vectors to normalized vectors. By scaling, we observe that $|U|\psi\rangle| = ||\psi\rangle|$ for all $|\psi\rangle \in V$. We now show that U sends orthogonal vectors to orthogonal vectors. Let $|\psi\rangle, |\varphi\rangle$ be orthogonal vectors. We wish to show that $U|\psi\rangle$ and $U|\varphi\rangle$ are orthogonal as well. We compute:

$$\begin{aligned}
||\psi\rangle|^2 + ||\varphi\rangle|^2 &= \langle\psi + \varphi|\psi + \varphi\rangle \\
&= \langle U(\psi + \varphi)|U(\psi + \varphi)\rangle \\
&= \langle U\psi|U\psi\rangle + \langle U\varphi|U\varphi\rangle + \langle U\psi|U\varphi\rangle + \langle U\varphi|U\psi\rangle \\
&= ||\psi\rangle|^2 + ||\varphi\rangle|^2 + 2\Re(\langle U\psi|U\varphi\rangle)
\end{aligned}$$

where $\Re(\cdot)$ denotes the real part of a complex number. Thus, we conclude that $\Re(\langle U\psi|U\varphi\rangle) = 0$. However, chaning $|\varphi\rangle$ by a phase, we can assume without loss of generality that $\langle U\psi|U\varphi\rangle$ is real, and hence we conclude that $\langle U\psi|U\varphi\rangle = 0$. Thus, we conclude that $\langle U\psi|U\varphi\rangle = \langle\psi|\varphi\rangle$ whenever ψ and φ are equal or orthogonal. Letting ψ, φ run over an orthonormal basis, we thus conclude that the equation $\langle U\psi|U\varphi\rangle = \langle\psi|\varphi\rangle$ holds on a basis. Extending via linearity we conclude it holds everywhere, which is exactly the statement that $U^\dagger = U^{-1}$, as desired. \square

Our second comment is in its heart a way of compact packaging the data of an observable. Given a Hilbert space V , instead of working with a choice of orthonormal basis B and a function $v : B \rightarrow \mathbb{R}$ we can work instead with a single operator $H : V \rightarrow V$. This is done by defining

$$H(b) = v(b) \cdot b$$

for all $b \in B$. The set B can now be recovered as the eigenvectors of H , and the measured results of the observable correspond to the eigenvalues. It is from this repackaging that the states in B get the name eigenstate. This packaging is useful because the space of linear operators $H : V \rightarrow V$ has more structure than the space of orthonormal bases of B paired with functions $v : B \rightarrow \mathbb{R}$. For example, we can now add two observables together, or tensor two observables on smaller systems to obtain an observable on a larger system. These sorts of operations will be very important going forward. In fact, the operator H will often have a simple form, and even computing what the elements of B are can be highly complex.

In a similar vein to our characterization of unitary operators, we give a characterization of those linear operators which arrise from observables:

Proposition 2.2 (Spectral theorem). *Let $H : V \rightarrow V$ be a linear transformation. The following are equivalent:*

1. *There exists an observable (B, v) such that $H(b) = v(b) \cdot b$ for all $b \in B$;*
2. $H = H^\dagger$.

If any of the three equivalent conditions are met, we call H a Hermitian matrix.

Proof. We do (1) \implies (2) first. From Exercise [ref], we know that H^\dagger can be computed as the conjugate transpose relative to any orthonormal basis. Choosing the orthonormal basis B , H is a real diagonal matrix. Hence, it is clearly equal to its own conjugate transpose.

We now prove the converse. We consider the map $\langle \cdot | \cdot \rangle$ as defined in the proof of Proposition 2.1. Since \mathbb{C} is algebraically closed the characteristic polynomial of H must have a root, hence we know that H has some eigenvector e , with eigenvalue λ . Scaling e if neccecary, we can assume without loss of generality that $\langle e|e\rangle = 1$. Let V be the subspace of vectors

$x \in \mathbb{C}[S]$ such that $\langle e|x \rangle = 0$. This space has dimension one less than V . We know from the definition of conjugate transpose that

$$\langle x|Hy\rangle = \langle Hx|y\rangle \quad \forall x, y \in \mathbb{C}[S].$$

In particular, if $\langle e|x \rangle = 0$ then

$$\langle e|Hx\rangle = \langle He|x\rangle = \lambda \langle e|x\rangle = 0.$$

Thus, H restricts to a map on V . Continuing this process of picking eigenvectors and restricting H to the subspace of vectors orthogonal to it, we find that V has an orthonormal basis of eigenvectors. Moreover, all of these eigenvectors satisfy

$$\lambda \langle e|e \rangle = \langle H(e)|e \rangle = \langle e|H(e) \rangle = \bar{\lambda} \langle e|e \rangle,$$

so their eigenvalues $\lambda = \bar{\lambda}$ are real. Thus, (2) \implies (1) as desired. \square

This concludes our treatment of the basic axioms of quantum mechanics.

2.2.6 Hamiltonians and the Schrodinger equation

We now know the basic rules of quantum mechanics. Suppose, however, that we are given some quantum mechanical system in a lab. How will it evolve in time? Certainly it will evolve by a unitary transformation, as per the axioms. But *which* unitary? The answer to this question is the Schrodinger equation. It gives us time dynamics in quantum mechanics. Once the initial state of the universe was set, the rest of time was just an evolution by the Schrodinger equation.

At the heart of the Schrodinger equation is the *Hamiltonian* of a quantum system. The Hamiltonian is an observable. The physical quantity it corresponds to is *total energy*. States with definite total energy are known as energy eigenstates, and their energy is some real number. In line with general principles established in the previous subsection, we will think of the Hamiltonian as being a linear operator $H : V \rightarrow V$. The Schrodinger equation is defined as follows:

Definition. (Schrodinger equation) Let V be a Hilbert space, corresponding to a quantum system. Let H be a Hermitian operator, corresponding to the Hamiltonian of V . Let $|\psi(t)\rangle$ denote the state of the system at time t . We have the formula

$$|\psi(t)\rangle = e^{-iHt} |\psi(0)\rangle$$

where $e^M = \sum_{n=0}^{\infty} \frac{M^n}{n!}$ is the matrix exponential. \square

This equation deserves several comments. First, we comment on terminology. Initially words *state of the system at time t* currently have no meaning. In fact time itself is at the current moment undefined. In this way, the Schrodinger equation is defining what time is in quantum mechanics (a one dimensional real parameter) and what it means for a system to be in a state at a time. We still do need to verify that the Schrodinger equation is consistent with our intuitive notion of time. For instance, if we first evolve the system in forward by t time units and then by s time units is that the same as evolving the system forward by $t + s$ time units? Under the Schrodinger equation, this is the same as verifying the equation

$$e^{-iH(t+s)} |\psi(0)\rangle \stackrel{?}{=} e^{-iHt} e^{-iHs} |\psi(0)\rangle.$$

This formula follows from the well known fact about matrix exponentials, which we will not prove:

Proposition 2.3. *If A and B are commuting operators, then*

$$e^A e^B = e^{A+B}.$$

Second, we make sure that the equation as stated is consistent with the axioms of quantum mechanics as we have previously defined them. In other words, is the map $e^{-iHt} : V \rightarrow V$ really a unitary operator for every $t \in \mathbb{R}$? This follows from the following important computation:

$$\begin{aligned} (e^{-iHt})^\dagger &= \left(\sum_{n=0}^{\infty} \frac{(-iHt)^n}{n!} \right)^\dagger \\ &= \sum_{n=0}^{\infty} \frac{((-iHt)^\dagger)^n}{n!} \\ &= \sum_{n=0}^{\infty} \frac{(iHt)^n}{n!} \\ &= e^{iHt}. \end{aligned}$$

The operators e^{-iHt} and e^{iHt} are inverses by Proposition [ref].

A third comment to make about the Schrodinger equation is about units. Both time and energy, austensibly, should have units. However, we have treated them as dimensionless mathematical quantities. How can this be? The answer is that implicitly we *did* choose units. When different choices of units are made, different constants need to be put into the Schrodinger equation. The version of the Schrodinger equataion which includes units is

$$|\psi(t)\rangle = e^{-iHt/\hbar} |\psi(0)\rangle$$

where \hbar is the normalized plank constant. In our original statement of the Schrodinger equation we have simply decided to use units in which the normalized plank constant is equal to 1.

[WORK: talk to a physicist who can say why the Schrodinger equation is true. I only have vague waffle.]

The Schrodinger equation tells us that all we need to do to understand the dynamics of a quantum system is solve the Schrodinger equation. Suppose now that $|\psi(0)\rangle$ is some initial state in a quantum system with extended state space V and Hamiltonian H . Suppose that we have a decomposition $|\psi(0)\rangle = \sum_{x \in B} c_x |x\rangle$ where B is the set of energy eigenstates of H . Then, the Schrodinger equation would tell us that

$$|\psi(t)\rangle = \sum_{x \in B} e^{-iv(b)t} c_x |x\rangle$$

where $v(b)$ is the eigenvalue corresponding to b . In this way, we see that by writing $|\psi(0)\rangle$ in terms of an energy eigenbasis we can exactly solve the Schrodinger equation.

In this way, solving quantum dynamics correponds exactly to finding the eigenvectors of the Hamiltonian. Or, in other words, diagonalizing the Hamiltonian. This task, while conceptually easy, can be very difficult in specific cases. Diagonalizing matrices has never been so exciting!

History and further reading:

[WORK: there are people who can do a history of quantum mechanics way better than me]

A fantastic place to first learn about quantum mechanics and its principles is the popular science book “Quantum computing since Democritus” [Aar13]. A more formal, but still excellent, introduction to finite-dimensional quantum theory is Nielsen-Chuang’s book “Quantum computation and quantum information” [NC10]. Past this there are many great textbooks which go into full depth on infinite-dimensional quantum theory and advanced properties of quantum systems. A good physics-oriented text is Shankar’s “Principles of quantum mechanics” [Sha12], and a good math-oriented text is Hall’s “Quantum theory for mathematicians” [Hal13].

Exercises:

2.1. .[WORK: show that the adjoint really is the conjugate transpose] [WORK: change verbiage above from “conjugate transpose” to “adjoint”]

[WORK: need to add somewhere that global phases don’t matter, clear up this ambiguity]

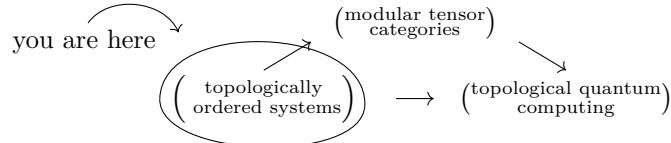
3 Topological quantum order

3.1 Overview

3.1.1 Introduction

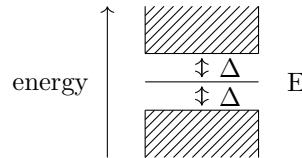
[WORK: balance generality and concreteness. I'll want to work with the definition of TO given by Bravyi and Hastings. This has the problem that it doesn't include the most physically realistic TO Hamiltonian, the honeycomb model! The honeycomb model only gives TC in its 4-th order perturbation theory. Something to deal with in the treatment - the big ideas are more important than the details. However, without the details things can easily get way off track.]

In this chapter we will be giving a detailed analysis of topological quantum order, a particular type of topological quantum system. We recall below how this fits into the general framework of this book:



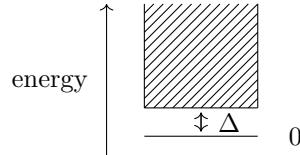
Topological quantum systems are distinguished by the fact that their states don't depend on local properties - they depend only on global topological properties of the system. One way of getting this sort of topological invariance is through *discreteness*. If a system is discrete, all of its parts are in a sense *far away* from each other. Things which are far away cannot be continuously deformed from one to another - local changes can't change discrete objects. A real-number valued invariant could move all over the place and depend heavily on local properties of a system, but an integer-valued invariant is *necessarily* topological invariant.

We demonstrate this below in its most basic form. Suppose that V is a Hilbert space and $H : V \rightarrow V$ is a Hermitian operator. This represents a quantum system and its Hamiltonian. Let $|\psi\rangle$ be an energy eigenstate with energy E . Suppose further that the E -eigenspace of H is one dimensional, and that every other eigenvalue E' on H satisfies $|E' - E| \geq \Delta$ for some real number $\Delta > 0$. This situation is demonstrated in the below graph:



This energy gap around E adds a sort of discreteness to the spectrum of H . Suppose that the system is in state $|\psi\rangle$ and we distort it a small amount. Typically, *this will not affect the state*. The state $|\psi\rangle$ would need to jump all the way to some other state, but all other states have significantly different energies. In particular, if the perturbation applied to $|\psi\rangle$ has magnitude significantly less than Δ , then $|\psi\rangle$ cannot change. This connection between gaps in energy spectra and topological states is so essential that many physicists use the terms *topological system* and *gapped system* interchangeably.

So, in practice, how do we make sure that the perturbations being applied to $|\psi\rangle$ are always much smaller than Δ ? We make the system *cold*. Roughly we say that a system has *temperature T* if the states of the Hamiltonian being occupied all have energy $< T$, and perturbations from the environment have magnitude $\approx T$. We renormalize our Hamiltonian so that the lowest energy eigenstate has energy 0. We call the lowest energy eigenstates the *ground states* of the system. We now assume that the ground state space is one dimensional, so there is a unique ground state. We assume that the next lowest energy eigenvalue is $\Delta > 0$. This gives us a new picture:

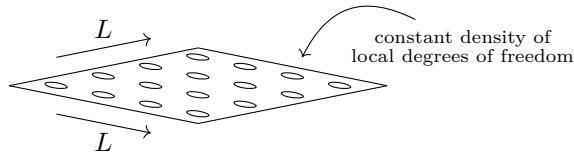


So long as the temperature is much smaller than the energy gap ($T \ll \Delta$), then our system will remain in the ground state. We say that the system (V, H) *becomes topologically ordered at low temperature*.

[WORK: here is where I should introduce TO properly]

[WORK: The issue of topological order at nonzero temperature is actually quite subtle. Seeing as in two dimensions there are no self-correcting codes, we can conclude that all topological order is *unstable* at nonzero temperature - it needs the external probes to drive it into the ground state: [Has11]. Another good discussion of topological order at nonzero temperature is [NO09].]

Of course, there's a big problem in our above discussion. *Every* finite dimensional quantum system is gapped. The Hamiltonian has finitely many eigenvalues, so its spectrum is necessarily discrete. What we should really be imagining is an infinite family of systems, parameterized by some real number $L > 0$ called the *linear system size*. Working in a two dimensional system, this will look like the below picture:



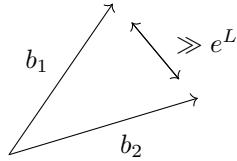
Letting $\dim(V)$ denote the dimension of V , this gives us an asymptotic formula $\dim(V) \sim e^{(\text{const}) \cdot L^2}$ where the constant in the exponent depends on the density of quantum degrees of freedom in the system. Let Δ_L be the lowest nonzero energy of the Hamiltonian in the size- L system. Of course, we will always have a gap $\Delta_L > 0$. What's important is that we require that $\Delta_L > \Delta$ for some uniform $\Delta > 0$. In most quantum systems this will *not* be the case - as the system size gets larger there will be states with smaller and smaller nonzero energies.

The issue with our discussion up to now is that it is *no use* for making a topological quantum computing. There is only a single ground state, so there is no non-trivial topologically protected information. There's just a point. To make a quantum computer we will need to introduce *degeneracy* into the ground states - make the lowest energy eigenspace higher dimensional. This degenerate ground space is where we will store our information.

If we do this naively, there's an immediate issue which appears. What if a perturbation of the system keeps the vectors in the ground space, but perturbs exactly which vector in

the ground space is being stored. Wouldn't this corrupt the data? The trick is choose the ground space correctly so that this does not happen. The way this works is by choosing a ground space which has a basis consisting of vectors which are in a certain sense "far apart". Because they are far apart, they cannot easily be distorted from one to another.

More explicitly, let us choose standard basis S for V , inducing an isomorphism $V = \mathbb{C}[S]$. This basis should correspond to the physical degrees of freedom underlying the system. If V is made up of an L by L grid of some repeating quantum sub-system, then choosing some arbitrary basis D for that subsystem a good choice for S is D^{L^2} , coming from the isomorphism $\mathbb{C}[S] \cong \mathbb{C}[D]^{\otimes L^2}$ where \otimes in the exponent denotes repeated tensor product. The canonical metric on \mathbb{C} induces a product metric on $V = \mathbb{C}[S]$. It is with respect to this topology that we want our basis for the ground space to be far apart. That is, we require a basis B for the ground space such that for every $b_1, b_2 \in B$, $|b_1 - b_2|$ is large. The exact scale of large depends on the topological system. At the very least it should tend to infinity with system size. In this case we will require an exponential scaling, $|b_1 - b_2| > e^{(\text{const})L}$:



[WORK: this stuff about distance is totally bogus. The real point is that if you differ at a large number of sites then it necessarily takes a large number of local errors to make a difference! Probability gets exponentially suppressed. Global feature \implies touches $> (\text{const}) \cdot L$ sites.]

Explicitly, this condition says that there exists a basis $|\psi_i\rangle$ such that

$$\langle \psi_i | \mathcal{O} | \psi_j \rangle = 0$$

for all $i \neq j$ and local operators \mathcal{O} . Paired with the error correcting code property, these conditions can be succinctly summarized as

$$\langle \psi_i | \mathcal{O} | \psi_j \rangle = \lambda \delta_{i,j}$$

Here's a good reference which talks about when all this stuff is physically possible and is fault tolerant - [KZL⁺16].]

This allows us to state a full picture of how to store topological information in a gapped system. Suppose we have some gapped system as before with distinguished geometric basis S , Hilbert space $V = \mathbb{C}[S]$, Hamiltonian H , temperature T , topological energy gap Δ , and linear system size L . We suppose $T \ll \Delta$, $L \gg 0$. Suppose further that the information we wish to store is the ground state

$$|\psi\rangle = \sum_{x \in S} c_x |x\rangle .$$

As time goes, we imagine the coefficients c_x continuously varying due to noise. This noise should have magnitude $\cong T$. We control our information by repeatedly measuring with respect to H . This measurement continually projects the our information back into an eigenstate. This is a mathematical mechanism for *cooling* - keeping the energy low. A few things could happen when H is measured.

1. Typically, after measuring the state will be projected back into the ground state space. The stored information will change a small continuous amount. The magnitude of this change is on the order of $T/e^{(\text{const})L}$. This is because basis vectors in the ground space are on the scale of $e^{(\text{const})L}$ times further apart than the basis vectors of $\mathbb{C}[S]$. Hence, the metric on the ground space is dialated by a factor of $e^{(\text{const})L}$, which has the effect of dampening the magnitude of the drift. Even though our stored information is always being corrupted by noise, the magnitude of this noise is tiny. Making the system size large, we can efficiently make the drift arbitrarily small. For any polynomial-length algorithm, the total amount of drift is still suppressed to large enough degree that the errors are tolerable. This means that our information is *topologically protected* in this case.
2. After measurement, the state could get projected onto an energy eigenstate which is *not* a ground state. This corresponds to a spontaneous jump in energy. The probability of such a jump is suppressed by the magnitude of the gap, giving a probability on the order of T/Δ . Choosing $T \ll \Delta$, we can make this probability small. However, we cannot make it arbitrarily small, and errors of this type need to be dealt with as they will surely appear in any sufficiently long algorithm. The upside is that when these errors occur it is entirely detectable - the outcome of the measurement of H is some observable energy, and it can be detected when that energy becomes nonzero. When it is detected that the energy is nonzero, then the experimenters can project the system back into the ground space by applying some external probe. The experiments can choose this projection carefully so that it sends the state to the nearest ground state, keeping the information drift on the order $T/e^{(\text{const})L}$. The details of how experimenters project non-ground states into ground states depends from topological system to topological system, and is often the heart of a proposal for topological quantum computing.

All in all, we find that following the procedures outlined above we can store topological information with essentially no errors. This is topological quantum memory.

The question now is how to make a *computer* of this. How do you act on the information stored this way in a gapped system? How do we go from one state to another in a topologically protected way? There are lots of different ways to do this, each of which have many equivalent descriptions. Here I will present a framework similar to the one introduced by Aasen-Wang-Hastings [AWH22]. In this framework, we perform computations by slowly transforming which Hamiltonian H we use to cool the system.

Suppose we have some state $|\psi\rangle$ we want to perform our computation on. We will choose some a family of Hamiltonians H_t , one for each time $t \in [0, 1]$. We will require that $H_0 = H_1 = H$ is our original Hamiltonian. We will continuously transform which Hamiltonian we use to cool the system. That is, at every time step t , we measure the system with respect to the Hamiltonian H_t . Assuming that the Hamiltonians vary slowly enough, our comments above apply. Namely, at time t either the state will stay a ground state of H_t with minimal drift or it will spontaneously jump to an excited state. In the case that it jumps to an excited state, we can apply an external probe to project it back into a ground state. Letting $|\psi(t)\rangle$ denote the state at time t , we find that $|\psi(1)\rangle$ will be some new ground state of H , which is well-defined up to errors on the scale $T/e^{(\text{const})L}$.

The beautiful observation is that $|\psi(1)\rangle$ does not need to be equal to $|\psi(0)\rangle = |\psi\rangle$. If the path taken by the Hamiltonians is non-trivial it can have a non-trivial action on the ground states, and serve as a source of computation. This is topological quantum computation. This sort of continuous evolution of a Hamiltonian while keeping a state in the ground state is

known as an *adiabatic* evolution of the Hamiltonian. An important point to emphasize is that for the above procedure to work, the Hamiltonians H_t must all have energy gaps, and these gaps must all be bounded below. Namely, $> \Delta$ for a fixed Δ . This model of computation can be summarized as saying that computations are performed by adiabatically transforming the Hamiltonian along non-trivial paths in the configuration space of all possible gapped Hamiltonians.

This already allows us to make interesting comments about the nature of topological quantum computing. To make a powerful quantum computer, there needs to be a lot of different loops that the Hamiltonian can go around, corresponding to a lot of possible different gates that can be applied. This means that the path-connected component of the original Hamiltonian in the configuration space of all possible gapped Hamiltonians has to have lots of non-trivial loops - its fundamental group needs to be large. Choosing gapped Hamiltonians whose path connected component in the space of gapped Hamiltonians has interesting topology is the art of topological quantum computing. It is here that we can get the definition of what a topological order is. It is a path connected component in the configuration space of gapped Hamiltonians. Or, equivalently, an equivalence class of gapped Hamiltonian up to continuous deformation.

Note that the exact definition of gapped Hamiltonian is subtle, because really we are talking about infinite families of Hamiltonians parameterized by system size, and so our above definitions of topological order are only approximate. The point is that topological order captures the inherent algebraic structure and nontrivial topology with a gapped Hamiltonian, while forgetting the details of how that Hamiltonian is defined.

[WORK: How should I define topological order, as opposed to simply “gapped Hamiltonian”? What am I missing? Is this something I even want to define it? Add a subsection?]

[WORK: The papers [CDH⁺20, BHM10, BH11] all agree on two axioms of TQO, TQO-1 and TQO-2. The exact implementation of these axioms are different, but their philosophy is here. Bring them in.

TQO-1 = ground states are error correcting code (topological protection)

TQO-2 = local ground state coincides with global one (allows for quasiparticle picture)

]

3.2 Discrete gauge theory

3.2.1 Ordered media on a lattice

Above we defined topological order. The best way to demonstrate the general principles of topological order is to give a good family of examples. The examples we will give in this section come from *discrete gauge theory*. At its heart, discrete gauge theory is a quantum version of the notion of ordered media we defined in Chapter [ref] section [ref]. While mathematically unnecessary, the next two subsections give physical motivation for why the formulas for discrete gauge theory have to be like they are, and why their analysis behaves like it does. Those who feel comfortable working with unmotivated formulas should skip to subsection [ref].

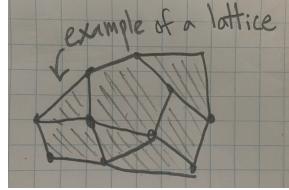
We will go from ordered media to discrete gauge theory in two steps:

Step 1: Put ordered media on a lattice;

Step 2: Make it quantum.

This first subsection is focused on Step 1. We will do Step 2 in the next subsection.

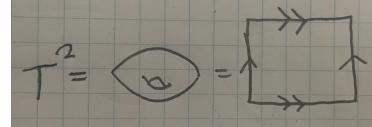
The first natural question to ask is *what is a lattice*. For our purpose a lattice is something like the picture below:



A lattice is a collection of vertices, edges, and faces connected in some way. To keep in line with the terminology common in topological quantum information, we refer to the faces of our lattices using the French term *plaquette*. Formally, by lattice we mean “simplicial 2-complex” though there is no need to go into details because we will never be dealing with the subtleties in the definition. Often times we will need to deal with *directed* lattices. These are lattices in which every edge has a direction, which we represent as an arrow on that edge.

Before putting ordered media on a lattice, a good question is *why* we would want to do this. There are two primary reasons. The first is that this will make this Hilbert spaces involved all finite dimensional. This is very important because we have only established quantum mechanics in the finite dimensional case, and working with the continuum limit can be highly complex. The second reason is that in practice, many of the systems physicists deal with are on lattices. For example, the chip of a quantum computer will store its information at finitely many sites, which can correspond to the vertices of some lattice. Many topological systems also arrise from materials which have crystal structures, which are modeled well by a lattice with atoms at the vertices and edges representing the geometry of the crystal.

The best setting for putting our ordered media on a lattice is by first putting on a torus. This helps for several reasons. Firstly, a torus is compact and hence it will add even more finiteness to the problem. Secondly, a torus has nontrivial topology which is useful for seeing the characteristic phenomenona of topological order. Thirdly, a torus has no boundary, which helps because boundaries in topological order are subtle and require more work to describe. We denote the torus by T^2 , and identify it with a square having its opposite sides glued:



Ordered media on the torus corresponds to continuous maps $\phi : T^2 \rightarrow M$ where M is some fixed order space. The steps to transforming a state ϕ into a lattice version of itself go as follows:

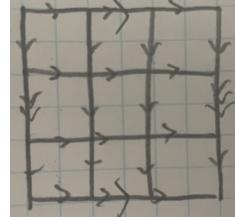
Step 1(a): Choose a directed lattice on the torus;

Step 1(b): Choose a basepoint $m \in M$. Make *local twists* around each vertex so that $\phi(v) = m$ for all vertices v in the lattice.

Step 1(c): On every edge, write down the winding number of ϕ along that edge, as an element of $\pi_1(M, m)$;

Step 1(d): Forget ϕ , and remember only the assignment of group elements in $\pi_1(M, m)$ to edges in the lattice.

These steps deserve explanation. Step 1(a) is clear: we choose an arbitrary lattice on the torus. Typically we will choose the square lattice on the torus:

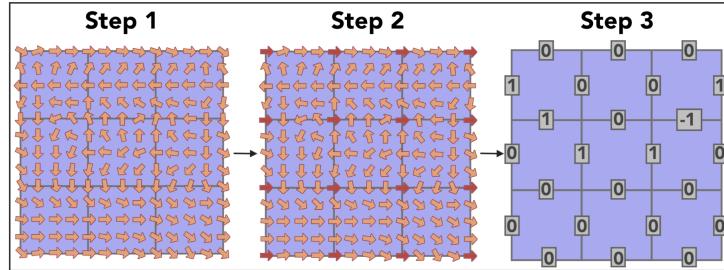


Step 1(b) requires more explanation. The picture to imagine is that we take the state ϕ and twist its values in small neighborhoods around each vertex to enforce the condition $\phi(v) = m$. Formally, this means choosing another state $\tilde{\phi}$ such that $\tilde{\phi}(v) = m$ for every vertex v of the lattice, and $\tilde{\phi} = \phi$ outside of some chosen small neighborhoods around each vertex. The fact that we can always choose such as state $\tilde{\phi}$ is a consequence of general mathematical principles in homotopy theory. Of course, different choices of $\tilde{\phi}$ will change the final result of our lattice encoding. However because any two choices of $\tilde{\phi}$ can only differ by local changes they can't be *too* different, in a way we will quantify later in the subsection.

Step 1(c) is straightforward. Every edge can be thought of as a path. Pushing forward with ϕ , this gives us a path in M . Since the edge starts and ends at vertices and ϕ sends all vertices to m , this means that the push forward of our edge gives a loop in M based at m . Hence, it gives an element of $\pi_1(M, m)$. We can record this element and attach it as a piece of data associated to the edge.

Step 1(d) is entirely book keeping. It records the fact that we have successfully transformed our continuous data ($\phi : T^2 \rightarrow M$) into discrete data (an assignment of group elements to edges in a lattice).

A worked example is shown below in the case that $M = S^1$ is the circle:



We now analyse our encoding of states in ordered media into assignments of group elements in $\pi_1(M, m)$ to edges in the lattice. The first fact from homotopy theory we will use is that these group elements determine the state ϕ exactly up to deformations localized

within each face. Taking a limit of denser and denser lattices, this means that the group elements will specify ϕ up to increasingly local deformations. The intuition is that by taking an infinite lattice limit we should recover ϕ up to “infinitely local deformations”, i.e., we recover it exactly. In this way we did a good job with our lattice encoding.

We observe that not every assignment of group elements to edges appears in our construction. There are implicit conditions. In particular, imagine taking the product of the group elements on edges along some contractible loop, taking inverses appropriately so that all the arrows are pointing in the same direction. This product will be equal to the group element associated with the loop around this whole path. The winding number along any contractible path under a continuous map should be trivial. Hence, the product of these group elements should be trivial. In particular, given any plaquette, the ordered product of group elements along its edges should be zero:

$$g_1 \xrightarrow{g_2} g_3 \text{ fillable} \Leftrightarrow [g_1 g_4 g_3^{-1} g_2^{-1}] = 1$$

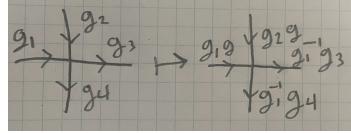
Moreover, *any* coloring of the edges of the lattice by elements of $\pi_1(M, m)$ will come from some map ϕ so long as it satisfies the condition above. This is one of the key formulas of the theory. It is in a real sense a lattice version of the continuity condition, since it is *equivalent* to the condition of continuity in the infinite lattice limit. This lattice version of continuity is called *flatness*. Flatness conditions are the most common sort of compatibility conditions which appear when you have local degrees of freedom valued in some group, making this lattice situation very general.

The last thing to deal with in analysing our system is deformation. When analysing states in ordered media, a huge amount of our time was spent on performing continuous deformations. Topological information is defined to be information which is invariant under continuous deformation. What does this correspond to in the lattice model?

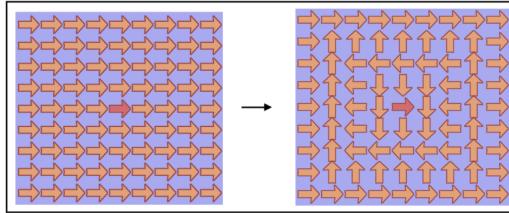
Suppose we are given an ordered media state ϕ and its corresponding lattice coloring. If we deform ϕ in some small neighborhood within a face, this will not change the values along the edges and hence will not change the coloring. If we deform ϕ in some small neighborhood around the interior of some edge this also won't change the coloring, because this will correspond to deforming the loop in M induced by going along that edge, and elements of the fundamental group are invariant under deformations of this sort. Another way of seeing that the coloring can't change is that flatness must be preserved - if the group element on the deformed edge changed, it would ruin flatness on the faces it bounds.

Finally, we can consider deforming ϕ around some vertex. This certainly *can* impact the coloring. An easy way to compute how it must impact the coloring is by using the fact that the flatness condition must be preserved. Suppose that an incoming edge labeled by g_1 changes to $g_1 g$ after the deformation. Enforcing flatness along all of the faces touching the vertex allows one to conclude that all incoming edges g_k will get changed to $g_k g$, and all outgoing edges g_k will get changed to $g^{-1} g_k$, as shown below:

Another way of seeing this result is by analysing what a deformation of ϕ does. The value $\phi(v)$ can move along some loop, starting and ending at m . This loop induces some element of the fundamental group, $g \in \pi_1(M, m)$. Performing this deformation exactly acts by



precomposing/postcomposing the adjoint edges with g/ g^{-1} accordingly. We can see below a concrete example for $G = S^1$:



Hence, we have a picture for ordered media on the lattice: states correspond to flat colorings of elements of $\pi_1(M, m)$ on a fixed lattice, and continuous deformations correspond to certain vertex actions by elements of $\pi_1(M, m)$.

3.2.2 From ordered media to gauge theory

In the previous subsection we showed how to put ordered media on a lattice. In this section we show how to make it quantum, turning it from a classical field theory to a quantum gauge theory. The idea of this jump is as follows. In Section [ref] we obtained an equivalence

$$(\text{topological information})_{\text{in ordered media}} = (\text{states}) / (\text{continuous deformation}).$$

It is necessary to mod out by continuous deformation because there is topological information in states, but also local degrees of freedom. For instance, the group element assigned to any individual edge in ordered media on a lattice can be changed by a gauge transformation and hence is not topologically invariant. The idea of going from ordered media to gauge theory is as follows: gauge theory is what results from ordered media when quantum fluctuations become so strong that local degrees of freedom are completely washed out and only the topology remains.

The fluctuations are quantum because we will imagine that our states will evolve in such a way that they are in a superposition of gauge transformations having been applied and not having been applied. Our states in gauge theory will be *equal superpositions over all possible deformations* of a given state. In this way, we are using quantum mechanics as a physical mechanism for quotients. Equivalence classes under deformation will be physically realized as equal superpositions over all possible representatives.

This can all be made completely rigorous. Choose a lattice on the torus, an order space M , and a basepoint m . We define a Hilbert space

$$\mathcal{N} = \bigotimes_{\text{edges}} \mathbb{C}[G].$$

We canonically identify the standard basis of \mathcal{N} with G -colorings of the lattice. Let C be an equivalence class of flat G -colorings of \mathcal{N} up to gauge transformations. There is a corresponding state

$$|C\rangle = \sum_{\gamma \in C} \frac{1}{\sqrt{|C|}} |\gamma\rangle.$$

This state is a normalized equal superposition of representatives of C . This defines a sub-Hilbert space

$$\mathcal{C} = \text{span} \{ |C\rangle \mid C \in (\text{flat } G\text{-colorings}) / (\text{gauge transformations}) \}.$$

This Hilbert space \mathcal{C} stores the information in our gauge theory.

So far our system is relatively trivial - it is just a Hilbert space, with no Hamiltonian. We connect it back to our original picture of topological order. The space \mathcal{C} is the collection of ground states in a topologically ordered system. Above it there is a whole spectrum of other states. This fuller picture with a Hamiltonian adds all of the subtlety and intrigue to the system.

In particular, we observed in Chapter [ref] section [ref] we observed the imports of quasiparticles in ordered media. These formed the heart of our information processing. Similarly, in gauge theory there will be quasiparticles as well which appear higher up in the spectrum of the Hamiltonian. Some of these quasiparticles will correspond to the classical quasiparticles in ordered media, but others are entirely new features of the system which did not exist before. We will analyse all this in more in the subsection that follows.

3.2.3 Kitaev quantum double model

[WORK: not sure if this is readable to someone who skipped the first two sections, but it should be. Something to keep an eye on.]

[WORK: Use $\mathfrak{D}(G)$ as notation for the doubled quantum order associated to G .]

In this section we will give the Hamiltonian formulation of discrete gauge theory. Seeing as we have moved passed ordered media, we will no longer be working with order spaces and base points. Instead, we will choose an abstract finite group G which replaces $\pi_1(M, m)$. The general picture for creating our Hamiltonian is simple, and follows a very general pattern in quantum theory: instead of enforcing properties rigidly as conditions, we will enforce them enforce properties energetically as terms in a Hamiltonian. The formulation we give below is known as the *Kitaev quantum double model of discrete gauge theory*. It was introduced in Kitaev's semiminal paper on topological quantum information [ref]. It has been studied extensively in the literature by many authors [add more refs].

Choose a directed lattice on the torus. Let

$$\mathcal{N} = \bigotimes_{\text{edges}} \mathbb{C}[G]$$

be the Hilbert space of our quantum system. The space \mathcal{N} has a canonical basis given by $\prod_{\text{edges}} G$, which we identify with G -colorings of the lattice. Given a G -coloring γ , we will denote the corresponding state in \mathcal{N} by $|\gamma\rangle$. For every plaquette p in the lattice, we define an operator on \mathcal{N} by

$$B_p |\gamma\rangle = \begin{cases} |\gamma\rangle & \gamma \text{ flat at } p \\ 0 & \text{otherwise.} \end{cases}$$

We observe immediately that

$$\sum_{\text{plaquettes } p} (1 - B_p) |\gamma\rangle = 0 \iff |\gamma\rangle \text{ is flat.}$$

It is in this way that we can enforce properties energetically by adding them as terms to a Hamiltonian. If we chose the Hamiltonian to be $\sum_{\text{plaquettes } p} (1 - B_p)$, then the lowest energy eigenspace would exactly correspond to the space spanned by flat G -colorings. For every vertex v and group element $g \in G$, we define an operator on \mathcal{N} by

$$A_{v,g} |\gamma\rangle = |\gamma \text{ acted on by the } g \text{ gauge action at } v\rangle.$$

For any $|\psi\rangle \in \mathcal{N}$, we call $|\psi\rangle$ *gauge invariant at v* if $A_{v,g} |\psi\rangle = |\psi\rangle$ for all $g \in G$. We call $|\psi\rangle$ *gauge invariant* if it is gauge invariant at v for all vertices v . We define

$$A_v = \frac{1}{|G|} \sum_{g \in G} A_{v,g}.$$

We define the Hamiltonian of our system to be

$$H = \sum_{\text{vertices } v} (I - A_v) + \sum_{\text{plaquettes } p} (I - B_p)$$

where I is the identity operator. We summarize the basic properties of this Hamiltonian below:

Proposition 3.1. *The following properties of the Kitaev quantum double Hamiltonian hold:*

- (a) *The operators A_v , B_p , and H are Hermitian for all vertices v and plaquettes p ;*
- (b) *The formula $A_{v,g}^\dagger = A_{v,-g}$ holds for all vertices v and $g \in G$;*
- (c) *All of the operators in the set $\{A_v, B_p\}_{v \in \text{vertices}, p \in \text{plaquettes}}$ commute with every other operator in the set;*
- (d) *The eigenstates of H are simultaneous eigenstates of the operators A_v , B_p ;*
- (e) *The eigenvalues of the A_v , B_p are all 0 or 1;*
- (f) *The lowest eigenvalue of H is 0, and the 0-eigenspace of H is*

$$\mathcal{C} = \text{span}\{ |C\rangle \mid C \in (\text{flat } G\text{-colorings}) / (\text{gauge transformations}) \}.$$

where for we define the ket

$$|C\rangle = \sum_{\gamma \in C} \frac{1}{\sqrt{|C|}} |\gamma\rangle$$

for any equivalence class C of G -colorings of the lattice up to gauge transformations.

Proof. .[WORK: do proof]

□

In particular, the above proposition tells us exactly that we have achieved our goal of realizing a Hamiltonian whose ground states capture the topological information in a lattice-version of ordered media. The term “double” in the Kitaev quantum double model refers to the fact that there are two families of terms in H - one family of type A_v and one family of type B_p . We can readily compute the dimension of the ground space as follows:

Proposition 3.2. *Choose a vertex v in the lattice. Every G -coloring of the lattice induces an assignment of lattice loops on the torus based at v to elements of G , based on taking the oriented winding number along that loop relative to the coloring. This restricts to a map*

$$(\text{flat } G\text{-colorings}) \rightarrow \text{Hom}(\pi_1(T^2, v), G)$$

where $\text{Hom}(\cdot, \cdot)$ denotes the space of group homomorphisms between two groups. Any two flat G -colorings which differ by gauge transformations will induce the same map in $\text{Hom}(\pi_1(T^2, v), G)$, up to global conjugation by an element of G . This induces a bijection

$$(\text{flat } G\text{-colorings}) / (\text{gauge transformations}) \rightarrow \text{Hom}(\pi_1(T^2, v), G) / (\text{conjugation})^{(\text{simultaneous})}.$$

The set of vectors $|C\rangle_{C \in (\text{flat } G\text{-colorings}) / (\text{gauge transformations})}$ is linearly independent. Hence, there is a canonical isomorphism

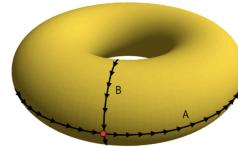
$$\mathcal{C} \rightarrow \mathbb{C}[\text{Hom}(\pi_1(T^2, v), G) / (\text{conjugation})]$$

given by taking winding numbers.

Proof. .[WORK: give proof. I'm scared this could be too hard. It's already] \square

The final step in using the above formula is to compute the fundamental group of the torus:

Proposition 3.3. $\pi_1(T^2, v) \cong \mathbb{Z}^2$ for any vertex v . The two loops shown below are generators for $\pi_1(T^2, v)$:



Proof. .[WORK: proof] \square

One last observation to make about this ground space is that its ground states really are globally different:

Proposition 3.4. *Let L be length of the shortest non-contractible loop on the lattice. Let γ_0, γ_1 be non-gauge equivalent flat G -colorings of the lattice. There are at least L edges at which γ_0 and γ_1 assign different values.*

Proof. .[WORK: do proof] \square

In particular, if we choose the square lattice on the torus, then the length L of the shortest non-contractible loop is obviously a good measure of linear system size. Proposition [ref] tells us that the number of local changes requires to go from one ground state to another is on the order of L . This is exactly the sort of condition we needed in Section [ref] to conclude topological protection in the ground states. Of course, the smallest non-zero eigenvalue of H is at least 1, which is bounded away from zero and hence there is a system-size independent gap between the ground states and the other states. Hence, we see that H is a good topologically ordered Hamiltonian.

The excited states of H will be described localized excitations with quasiparticle behavior. Given a state $|\psi\rangle \in \mathcal{N}$, we will say that a state has an *excitation at vertex v* if $A_v |\psi\rangle = 0$ and we will say that it is *unoccupied at p* if $A_v |\psi\rangle = 1$. We say that $|\psi\rangle$ has an *excitation at plaquette p* if $B_p |\psi\rangle = 0$ and that it is *unoccupied at p* if $B_p |\psi\rangle = 1$. By Proposition [ref], every energy eigensate is either occupied or unoccupied at every vertex/plaquette. The regions in which $|\psi\rangle$ is unoccupied are all essentially identical, leading to a homogenous bulk. The sites at which $|\psi\rangle$ is occupied are different, and behave as quasiparticles. We will define operators which move these excitations around.

[WORK: Add something about local indistinguishability of ground states - reinforce this “homogenous bulk” idea.]

More than this, it is important to note that earlier we are giving a rigorous definition of topological order. It is not immediately obvious that the KQDM satisfies this definition. This is the main content of the paper [CDH⁺20]. Should I include a proof? At the very least there should be some mention of how this fits into the definition. In fact, this should be a big point. The KQDM is being introduced with the main goal of giving an example of TO. Needs to talk about how it is topological.]

[WORK: Maybe also reinforce that this could be done on *any manifold*, and the ground states would be the same?]

3.3 The toric code

[WORK: Maybe this section can be re-done. We know that the total space \mathcal{N} can be decomposed as a direct sum

$$\mathcal{N} = \bigoplus_{\lambda} \mathcal{N}_{\lambda}$$

as a direct sum over syndromes, by general principles of diagonalizable matrices. To prove that all of the \mathcal{N}_{λ} have an even # of excited terms in H of both A_v and B_p type is easy. Proving that they all have the same dimension involves the simple observation that applying σ_X and σ_Z between two excited terms will make them both ground states. Simple counting recovers the fact that the ground state is 4-dimensional.

The beauty of this approach is that it is immediately grounded. We have a Hamiltonian, we want to solve it - i.e. we want to compute the dimensions of the \mathcal{N}_{λ} , and explicitly have a way of creating those basis states. The proof in a real sense is using the quasiparticle nature of the A_v and B_p excitations. Namely they are being moved along paths to annihilate with one another.

Every operator can be decomposed as a sum of Pauli operators. Hence, understanding how Paulis act on \mathcal{N} lets understand how every operator acts on \mathcal{N} . Paulis act on \mathcal{N} by creating/moving/fusing vertex/plaquette excitations. Hence, understanding vertex/plaquette excitations tells you everything you need to know about the toric code. Saying it this way

makes everything feel very grounded, and it doesn't bring in anyons unneccarily early into the picture. We can talk about anyons after. Highlight the fact that they behave like quasi-particles and that they will become objects of independent interest, but that isn't the point yet.

]

3.3.1 Simplified Hamiltonian

In this section we move on to analyzing the Kitaev quantum double model for $G = \mathbb{Z}_2$, which is known as the *toric code*. The name toric code comes from the fact that the toric code was first introduced as an error correcting code, and was only later recast as a topologically ordered system [refs]. The toric code is still the basis for many of the most popular error correcting codes [refs]. In a real sense the toric code is the simplest nontrivial topological order. It is a fantastic example which demonstrates almost all of the phenomena of topological order with relatively little work involved. The toric code, and more generally \mathbb{Z}_2 discrete gauge theories, can be found in all sorts of systems such as [WORK: give examples. The ones that jump to mind are dimer models - [MS01, MSP02]].

We describe the model now. Because $G = \mathbb{Z}_2$ is abelian, we will switch to additive notation for our group operation. We choose a *non-oriented* lattice structure on the torus. This lattice does not need to be oriented because changing the direction of edges in the lattice corresponds to taking inverses, and $g = g^{-1}$ for every element $g \in \mathbb{Z}_2$. We define

$$\mathcal{N} = \bigotimes_{\text{edges}} \mathbb{C}[\mathbb{Z}_2] = \bigotimes_{\text{edges}} \mathbb{C}^2.$$

Here, we identify $\mathbb{C}[\mathbb{Z}_2]$ with \mathbb{C}^2 for convenience, endowing \mathbb{C}^2 with a canonical basis $\{|0\rangle, |1\rangle\}$. We call \mathbb{C}^2 a *qubit*, in analogy to “bits” for classical computing. It is a standard two-level quantum system. Most quantum computers are based on qubits, which makes the toric code especially accessible to practical implementation as an error correcting code. The definition of the Hilbert space \mathcal{N} can be summarized as putting a qubit on every edge of the lattice. The Hamiltonian is

$$H = \sum_{\text{vertices } v} (1 - A_v) + \sum_{\text{plaquettes } p} (1 - B_p).$$

We unpack the general definitions of A_v and B_p for the toric code. The operator $A_{v,0}$ is the identity. The operator $A_{v,1}$ acts by a gauge transformation,

Defining

$$\begin{aligned} \sigma_X : \mathbb{C}^2 &\rightarrow \mathbb{C}^2 \\ |0\rangle &\mapsto |1\rangle \\ |1\rangle &\mapsto |0\rangle \end{aligned}$$

we thus find that

$$A_{v,1} = \bigotimes_{\substack{\text{edges} \\ \text{touching } v}} \sigma_X, \quad A_v = \frac{1}{2} (I + A_{v,1}).$$

Moving on to B_p , we recall that

$$B_p \left| \begin{array}{c} \text{square} \\ \text{edges} \\ g_1, g_2, g_3, g_4 \end{array} \right\rangle = \begin{cases} 1 & \text{if } g_1 + g_2 + g_3 + g_4 = 0 \\ 0 & \text{otherwise} \end{cases}$$

In the present case, B_p has a more workable expressing that is symmetric to our description of B_p . Define

$$\begin{aligned} \sigma_Z : \mathbb{C}^2 &\rightarrow \mathbb{C}^2. \\ |0\rangle &\mapsto |0\rangle \\ |1\rangle &\mapsto -|1\rangle \end{aligned}$$

Philosophically, it is useful to interpret σ_Z as acting as $|g\rangle \mapsto \chi(g)|g\rangle$ where $\chi : \mathbb{Z}_2 \rightarrow \mathbb{C}^\times$ is the unique nontrivial character of \mathbb{Z}_2 , $\chi(0) = 1$, $\chi(1) = -1$. Since χ is a group isomorphism, for any $g_1, g_2, g_3, g_4 \in G$ we have an equivalence

$$g_1 + g_2 + g_3 + g_4 = 0 \iff \chi(g_1)\chi(g_2)\chi(g_3)\chi(g_4) = 1.$$

Defining an auxillary $B_{p,1}$, we thus find the following expression for B_p :

$$B_{p,1} = \bigotimes_{\substack{\text{edges} \\ \text{bounding } p}} \sigma_Z, \quad B_p = \frac{1}{2} (I + B_{p,1}).$$

For simplicity, we will often rewrite the Hamiltonian as

$$H = \frac{1}{2} \sum_{\text{vertices } v} (1 - A_{v,1}) + \frac{1}{2} \sum_{\text{plquettes } p} (1 - B_{p,1}).$$

The matrices σ_X and σ_Z we defined are known as *Pauli matrices*. They are extremely common across formulae in quantum mechanics - this is another reason that the toric code is so ammenable to error correction applications. The basic properties of these matrices are summarized below:

Proposition 3.5.

- (a) The operators σ_X and σ_Z are simultaneously unitary and Hermitian;
- (b) $\sigma_X^2 = \sigma_Z^2 = I$;
- (c) $\sigma_X \sigma_Z = -\sigma_Z \sigma_X$;

Proof. .[WORK: do proof] □

An important thing to note is that $A_{v,1}$ and $B_{p,1}$ commute, despite the fact that σ_X and σ_Z anticommute. The fact that they commute follows from Proposition [ref], though it is fruitful to reevaluate that proposition in this present context. The important fact is that given any vertex v on the exterior of any face touching p , there are an *even number* of edges which both touch v and bound p . Hence, the number of tensor factors in which $A_{v,1}$ and $B_{p,1}$ anticommute is even, and hence overall they commute.

The last step in reinterpreting our general theory of Kitaev quantum double models to the toric code is computing the ground space. We observe that since \mathbb{Z}_2 is abelian acting by conjugation does nothing, and hence

$$\text{Hom}(\pi_1(T^2, v), \mathbb{Z}_2) / (\text{conjugation}) = \text{Hom}(\pi_1(T^2, v), \mathbb{Z}_2).$$

Seeing as we are no longer modding out by conjugation, the group operation on \mathbb{Z}_2 extends to a group operation on $\text{Hom}(\pi_1(T^2, v), \mathbb{Z}_2)$. Hence this space forms an abelian group, which we denote

$$H^1(T^2, \mathbb{Z}_2) = \text{Hom}(\pi_1(T^2, v), \mathbb{Z}_2) = (\text{flat } \mathbb{Z}_2\text{-colorings}) / (\text{gauge transformations}).$$

[WORK: maybe set notation and write out four elements explicitly? Might be too much.]

This is the *cohomology group of T^2 with coefficients in \mathbb{Z}_2* . Since $\pi_1(T^2, v) \cong \mathbb{Z}^2$, we conclude that

$$H^1(T^2, \mathbb{Z}_2) \cong \mathbb{Z}_2^2.$$

Hence, we obtain the following:

Proposition 3.6. *The 0-eigenspace of H is four dimensional. It is spanned by the vectors*

$$|C\rangle = \frac{1}{\sqrt{|C|}} \sum_{\gamma \in C} |\gamma\rangle$$

for $C \in H^1(T^2, \mathbb{Z}_2)$.

Proof. .[WORK: do proof] □

3.3.2 Exact solution of the toric code

When given a quantum system, the first thing to do with it is to *solve it*. This means diagonalizing the Hamiltonian. In this case of the toric code, the diagonalization of \mathcal{N} is the direct sum decomposition

$$\mathcal{N} = \bigoplus_{E \in \mathbb{R}} \mathcal{N}_E$$

where \mathcal{N}_E is the energy E eigenspace,

$$\mathcal{N}_E = \{|\psi\rangle |, H|\psi\rangle = E|\psi\rangle\}.$$

Solving the toric code amounts to explicitly describing \mathcal{N}_E for each E . In particular, this means computing the dimension of each space. The Hamiltonian for the toric code is

$$H = \sum_v (1 - A_v) + \sum_p (1 - B_p).$$

Since the A_v and B_p all commute with each other, they are *simultaneously diagonalizable*. This is a huge help in our analysis. We introduce some notation to take advantage of this insight. We define a *syndrome* on the toric code to be a map

$$\lambda : (\text{faces}) \sqcup (\text{vertices}) \rightarrow \{\pm 1\}.$$

We define the syndrome λ subspace of \mathcal{N} to be

$$\mathcal{N}_\lambda = \{|\psi\rangle \in \mathcal{N} \mid A_{v,1} |\psi\rangle = \lambda(v) |\psi\rangle, B_{p,1} |\psi\rangle = \lambda(p) |\psi\rangle \forall v, p\}$$

We define the energy E_λ of a syndrome λ by the formula

$$E_\lambda = \sum_v \frac{1}{2}(1 - \lambda(v)) + \sum_p \frac{1}{2}(1 - \lambda(p)).$$

The fact that the operators A_v , B_p , and H are simultaneously diagonalizable is codified in the following observations:

Proposition 3.7. *We have that*

$$\mathcal{N} = \bigoplus_{\text{syndromes } \lambda} \mathcal{N}_\lambda, \quad \mathcal{N}_E = \bigoplus_{\substack{\text{syndromes } \lambda \\ E_\lambda = E}} \mathcal{N}_\lambda.$$

Proof. This follows immediately from the above discussion. \square

We can now solve the toric code:

Proposition 3.8. *We have that*

$$\dim(\mathcal{N}_\lambda) = \begin{cases} 4 & \text{if } \prod_v \lambda(v) = \prod_p \lambda(p) = 1 \\ 0 & \text{otherwise.} \end{cases}$$

Proof. .[WORK: do proof] \square

Corollary 3.1. *We have that*

$$\dim(\mathcal{N}_E) = [\text{WORK : write formula}]$$

Proof. . [WORK: do proof] \square

[WORK: this section will have some commentary about how Pauli operators are being used to fuse X -type and Z -type excitations. The way this commentary sounds should depend on the way the proof looks. I might want to include this lemma before the proof:

Given an edge e in the lattice and an operator $U : \mathbb{C}^2 \rightarrow \mathbb{C}^2$, denote by $(U)_e : \mathcal{N} \rightarrow \mathcal{N}$ the operator which applies U on the tensor factor of \mathbb{C}^2 at edge e . We compute the following:

Lemma 3.1. *For any vertex v , edge e , plaquette p , we have*

$$A_v(\sigma_X)_e = (\sigma_X)_e A_v, \quad B_p(\sigma_Z)_e = (\sigma_Z)_e B_p,$$

$$A_v(\sigma_Z)_e = \begin{cases} -(\sigma_Z)_e A_v & \text{if } e \text{ touches } v \\ (\sigma_Z)_e A_V & \text{otherwise} \end{cases}$$

and

$$B_p(\sigma_X)_e = \begin{cases} -(\sigma_X)_e B_p & \text{if } e \text{ bounds } p \\ (\sigma_X)_e B_p & \text{otherwise} \end{cases}$$

]

3.3.3 Toric code as a topologically ordered system

.[WORK: prove that toric code satisfies axioms of TO.]

[WORK: This section will really dig into how the Pauli algebra acts on states]

[WORK: this section can also show that *global* errors (e.g. σ_X all the way around a torus) acts non-trivially. Maybe give a proof? Make a little quantum computer? If I do, I should reformulate it in the “adiabatically changing Hamiltonian” way to tie it into the way I talked about it in the introduction.]

3.4 Anyons

3.4.1 Topological quantum information in excited states

[WORK:

I’ve realized that I was wrong about anyons. I thought that they were localized excitations in a topologically ordered system. I was wrong. Anyons are *stable* localized excitations in a topologically ordered system. This stability is what ensures that they will have a well-defined type. This correct definition should make the presentation a lot easier!]

Let us recap our current picture of topological quantum information. We introduced topological order, and argued as a general principle that its ground states are invariant under local deformations. That is, if we start with a ground state in a topologically ordered system, apply some local operator, and then project back onto ground states, then we will get back to the state we started with up to some phase. This lets us store information in ground states. This information can be used as a place to store stable information, and can be acted on by global operators to perform computations such as in proposition [ref].

We now take this perspective to its logical extreme. The dimension of the ground state space depends only on the topological order of the system and the global topology of the physical space. For instance, toric code topological order on a torus will always have a four dimensional ground space, independent of system size or choice of microscopic lattice.

To make a computer, however, we need to be able to store arbitrarily large amounts of information. This means that either we will need to work with increasingly complex topological orders or increasingly complex surfaces. There are only a handful of topological orders known to be physically realizable, so working with increasingly complex topological orders is out of the question. Hence, one needs to work with increasingly complex surfaces.

By working on high-genus surfaces, the ground state space can be made as large as possible. This gives the following picture for a topological quantum computer:

[WORK: add picture, a-la Freedman et al.]

This approach is potentially possible, and we will explore it more in section [ref]. However, no matter the implementation it adds a great deal of complexity. Adding genus onto a computer chip is hard work. It would be much better to work with a state with the least complex topology possible, preferably a plane, sphere, or torus.

It is for this reason that we turn to a key idea in the theory of topological order: *being careful, it is possible to store topologically protected information in the excited states of a topologically ordered system.*

[WORK: picture of spectrum diagram, with excited states boxed. “Use these to store information!”]

Using the excited states of a topologically ordered system to store information comes with several obstacles:

1. It requires careful study. This isn’t a fundamental problem, but it is relevant to us since we are about to embark on that study;
2. It requires a more powerful control over the topologically ordered system than working only with ground states. This can make the approach impossible in some physical systems;
3. It introduces non-topological local degrees of freedom. This non-topological information needs to be worked around so that proper fault-tolerant computations can be performed.

We now recall our procedure for storing information in ground states. The key point is that we have a projector $\mathcal{P} : \mathcal{N} \rightarrow \mathcal{N}_{g.s.}$. This projector satisfies the relation that $\mathcal{P}\mathcal{O}\mathcal{P} = \lambda \cdot \mathcal{P}$ for any local operator \mathcal{O} , for some $\lambda \in \mathbb{C}$. Our physical picture for topological quantum computing is that we are continually measuring with respect to \mathcal{P} , and hence constantly projecting into the ground space. The formula $\mathcal{P}\mathcal{O}\mathcal{P} = \lambda \cdot \mathcal{P}$ says that even if noise is applied between rounds of projection it is okay because our state will only change by a phase and hence will still store the same information.

As we leave the ground states, this protocol breaks down. We cannot continually project into ground states because that would destroy any information we are storing in excited states! If we don’t have any sort of projector, however, errors will accumulate and we will lose all of the fault-tolerance of topological quantum computing. The answer is to find a new subspace $\mathcal{N}' \subset \mathcal{N}$ to store our information in, so that the orthogonal projector $\mathcal{P}_{\mathcal{N}'} : \mathcal{N} \rightarrow \mathcal{N}'$ satisfies a formula similar to $\mathcal{P}_{\mathcal{N}'}\mathcal{O}\mathcal{P}_{\mathcal{N}'} = \lambda \cdot \mathcal{P}_{\mathcal{N}'}$.

A first guess might to constantly measure with respect to the Hamiltonian, as before, and then project into a different energy eigenspace. We have a decomposition of \mathcal{N} by diagonalizing the Hamiltonian, $\mathcal{N} = \bigoplus_{E \in \mathbb{R}} \mathcal{N}_E$. We can try storing our information in \mathcal{N}_E , for some fixed $E > 0$ independent of system size.

We demonstrate how this fails using the toric code, working with $E = 4$. A generic state $|\psi\rangle$ in $\mathcal{N}_{E=4}$ might look like this one:

[WORK: add picture, two sites with X -excitations, two sites with Z -excitations.]

The fact that there are exactly four terms in the Hamiltonian that are violated corresponds exactly to the fact that the energy of the state is $E = 4$. We now consider applying σ_X to $|\psi\rangle$ at an edge adjacent to one of the X -excitations:

[WORK: add picture. Applying $(\sigma_X)_e$ moves the X -excitation.]

This new state with the X -excitation moved still has the same number of anyons, and hence $(\sigma_X)_e |\psi\rangle$ still leaves in \mathcal{N}_E . Hence, $\mathcal{P}_E(\sigma_X)_e \mathcal{P}_E |\psi\rangle = (\sigma_X)_e |\psi\rangle$ does *not* differ by a scalar. Even worse, despite the constant application of the projector \mathcal{P}_E , local errors can accumulate to become global errors. If we keep applying σ_X on edges around some loop then this will never be detected by \mathcal{P}_E , and this will be the same as acting by some non-trivial global operator.

[WORK: add picture of local noise accumulating to be a global error.]

In summary, storing information in an excited energy level of a topologically ordered system does not result in information which is resistant to local noise. The problem is that the terms which violate the ground-state condition in the Hamiltonian can *drift*, moving around the physical space in an unrolled fashion without changing the energy of the system.

The solution to this problem is to *constrain the drift*. This works as follows. We work in the Kitaev quantum double model based on some finite group G . Define a *region* in a space to be a compact connected subset of it. Choose n disjoint regions $R_1 \dots R_n$ on the torus T^2 . We define the space

$$\mathcal{N}_{R_1, R_2 \dots R_N} = \left\{ |\psi\rangle \in \mathcal{N} \mid A_v |\psi\rangle = B_p |\psi\rangle = |\psi\rangle, \forall v, p \notin \bigcup_{i=1}^n R_i \right\}.$$

This space consists of states which satisfy the condition to be in the ground state of $H = \sum_v (1 - A_v) + \sum_p (1 - B_p)$ outside of the regions $R_1 \dots R_n$, but is allowed to do whatever it wants within the regions. This is shown visually below:

[WORK: add picture]

The space $\mathcal{N}_{R_1, R_2 \dots R_N}$ contains all ground states, but also a large amount of excited states. The regions $R_1 \dots R_n$ constrain the drift of excitations. To illustrate the power of this space, consider the example of the toric code. Suppose we have our same state $|\psi\rangle$ as below with energy $E = 4$, but we consider it instead as a subspace of $\mathcal{N}_{R_1, R_2, R_3, R_4}$ for some regions R_1, \dots, R_4 containing the four excited operators. Applying σ_X to an edge can still change the state, and so we don't have the formula $\mathcal{P}_{N'} \mathcal{O} \mathcal{P}_{N'} = \lambda \cdot \mathcal{P}_{N'}$, but local errors *cannot* accumulate to become global errors! If the excited term starts to drift, it will eventually leave the region in started in and hence will no longer be in the subspace $\mathcal{N}_{R_1, R_2, R_3, R_4}$, so projecting back into it will fix the error. This is shown pictorially below:

[WORK: add picture]

In this way, local operators can change the information stored in the system but it can only change it to a controlled degree. There is still global information within $\mathcal{N}_{R_1, R_2 \dots R_n}$ which is roughly invariant under local operators. Getting this global information out in such a way that the answer does not depend on whatever local operators are applied to the system is non-insurmountable challenge. This is the heart of point (3) in the obstacles of working with excited states: it introduces non-topological degrees of freedom which need to be worked around.

[WORK: notation is a bit junky because I'm working with $R_1, R_2 \dots R_n$ all the time. It would be nicer in some ways if I worked with just one region R , and I dropped the condition that it be connected. This makes the definition of "anyon" a bit more janky though. Not sure what to do.]

[WORK: So far I haven't defined what an anyon is for topological orders other than the KQDM. One easy way to do this is by putting some set L of sites, choosing $d > 0$, defining $\mathcal{N} = \bigotimes_{\ell \in L} \mathbb{C}^d$. Then we have $H = \sum_i H_i$, $[H_i, H_j] = 0$, $H_i^2 = 1$, H_i localized to some

region U_i . This is a commuting local projector Hamiltonian. It is easy to define anyons in a model like this one. I'm not sure if this would be informative or distracting.

Actually, on second thought, something like this might actually be necessary for the definition of TO. Hence, we might have it ready-to-use for this section.]

3.4.2 Definition and principles of anyons

With the discussion in the previous section, we can start to see how our picture for topological quantum computing is similar to the picture for topological classical computing described in chapter [ref]. The regions $R_1, R_2 \dots R_n$ are localized regions of difference within a homogenous bulk. The bulk is homogenous because the wavefunction is groundstate of the Hamiltonian in those regions, and TQO-2 implies that all of these ground states are locally indistinguishable. The regions $R_1, R_2 \dots R_n$ are different because they are allowed to be excited. Hence, the regions $R_1 \dots R_n$ behave as *quasiparticles*. We will show that these quasiparticles can be pair-created, braided, and fused, just like in topological classical computing. The major difference between our scheme for topological quantum computing and topological classical computing is that instead of our quasiparticles being defects in ordered media, they are localized excitations in a topological order. This leads us to the following definition:

Definition. An *anyon* is a localized excitation in a topologically ordered system.

In our present context, this localization is best viewed not as an unavoidable physical reality but as a design decision for building a topological quantum computer. States $|\psi\rangle$ in topologically ordered systems have terms in the Hamiltonian that the violate and ones they don't. By circling regions around the terms they don't and constraining the excitation to those regions by repeatedly applying the projector $\mathcal{P}_{\mathcal{N}_{R_1, R_2 \dots R_n}}$, we localize the excitation.

To store coherent topological quantum information in anyons, there are a few key principles one must follow. The first principle is straightforward:

Anyons should be kept far apart

This is motivated as follows. Suppose that the anyons were kept in close proximity to one another. Then local noise could affect two anyons at once:

[WORK: add picture of R_1, R_2 with some local noise operator \mathcal{O} touching both]

Our picture of our system is that there is constantly local noise being applied, and that we are constantly projecting into the space $\mathcal{N}_{R_1, R_2 \dots R_n}$. The fault-tolerant information we want to get at is exactly the information which is invariant under this noisy picture. That is, information which does not change when local operators from $\mathcal{N}_{R_1, R_2 \dots R_n}$ to itself is applied. We call this topological information.

Some of this topological information can be measured using local observables. Because physics is local, any realistic observable should be local. Suppose we have a local Hermitian operator $\pi : \mathcal{N}_{R_1, R_2 \dots R_n} \rightarrow \mathcal{N}_{R_1, R_2 \dots R_n}$ which commutes with every noise operator $\mathcal{N}_{R_1, R_2 \dots R_n} \rightarrow \mathcal{N}_{R_1, R_2 \dots R_n}$. Then π can be physically measured *and* the result of that measurement is invariant under noise. Hence, it gives topological information. We call such observables π which are locally supported and commute with every noise operator *local topological observables*. We imagine that the outcomes of local topological measurements are readily available to experimenters - they can be measured with physical devices in a way that does not depend on noise.

This allows us to break-down the information in our system:

[WORK:

Four-square diagram for information in $\mathcal{N}_{R_1, R_2 \dots R_n}$.

Columns: measurable by local topological observables. (yes/no)

Rows: invariant under local noise. (yes/no)

C-yes R-no: Impossible.

C-no R-no: Non-topological information.

C-yes R-yes: Classical information / local topological information

C-no R-no: Topological quantum information / global information

Add little arrows going to each box explaining it.]

[WORK:

Measuring under all topological observables should be called a “topological charge measurement”. This leads to harmony with the language used in the MTC section.

] We now break down this general picture into exact mathematical statements.

[WORK: at this point I need to read Kitaev’s argument in more detail.

The first step is to go $\mathcal{N}_{R_1, \dots, R_n} = \bigoplus_{i=1}^N \mathcal{N}_i$ where \mathcal{N}_i indexes over classical information.

This is easy. The non-trivial step is to observe that there exists some finite set \mathcal{L} such that the terms \mathcal{N}_i can be rearranged as

$$\mathcal{N}_{R_1, \dots, R_n} = \bigoplus_{(A_i)_{i=1}^n \in \mathcal{L}^n} \mathcal{N}_{A_1, A_2 \dots A_n}.$$

Each A_i is the result of a measurement localized around R_i . The set \mathcal{L} is the set of anyon types, and given some state $|\psi\rangle \in \mathcal{N}_{A_1, A_2 \dots A_n}$, we call A_i the type of the anyon at R_i . Anyon types = topological classical information.

Now, the next step is to observe that there is a non-canonical tensor decomposition

$$\mathcal{N}_{A_1, A_2 \dots A_n} = \mathcal{N}_{A_1, A_2 \dots A_n}^{loc} \otimes \mathcal{N}_{A_1, A_2 \dots A_n}^{top}$$

of $\mathcal{N}_{A_1, A_2 \dots A_n}$ into a local part and a topological part. It satisfies the condition that every local noise operator \mathcal{O} on $\mathcal{N}_{A_1, A_2 \dots A_n}$ can be decomposed as $\mathcal{O} = \mathcal{O}' \otimes \text{id}$, and hence the information in the topological part remains unchanged. Moreover, it is maximal subject to this condition. Moreover, we have a splitting

$$\mathcal{N}_{A_1, A_2 \dots A_n}^{loc} = \mathcal{N}_{A_1}^{loc} \otimes \mathcal{N}_{A_2}^{loc} \dots \otimes \mathcal{N}_{A_n}^{loc}.$$

Information of which direct summand I’m in = topological classical information

Information left over = tensor product of local and topological. Cannot be distinguished by the non-canonical nature of this tensor decomposition. Needs subtle techniques (i.e. fusion of anyons) to be measured. It would be fantastic if all of this could be written up correctly and codified into propositions.]

[WORK: I want to get to the fact that anyons are moved by unique operators, and hence we can ignore the specific choice of operator and just move the anyons.]

[WORK: The notion of anyon types is explained well by Kawagoe and Levin [KL20]:

quote: We begin with the idea of “anyon types.” ... anyon excitation of type a, type b, type c, etc.]

3.4.3 Anyons in the toric code

[WORK: do it first with toric code. Everything is painfully simple and obvious here.

The problem is that adding more anyons does not encode more information, so it’s hard to get the points of TQC across.

Hopefully this is a very short subsection.]

3.4.4 Anyons in discrete gauge theory

[WORK:

Here we go from KQDM to G -crossed G -representations. By the end we should have the category $\mathfrak{D}(G)$ as a set, with morphisms motivated.

A lot of the intricate setup has been done in the previous sections, so I think this can be relatively contained. It would be nice to have proofs. Even if this section is more intricate that should be included in most lectures series, I'd say its nice to have nonetheless as a reference.]

History and further reading:

The term topological order was first used in 1972 by Kosterlitz and Thouless to describe topological classical systems of the sort discussed in Chapter [ref] [KT18]. The term has since evolved, and was re-coined in 1989 by Xiao-Gang Wen to describe the sort of topological classical systems defined in this chapter [Wen89].

The history anyons is distinct from the history of topological order. It was first noted in 1976 in a paper of Leinass and Myrheim that the classification of particles in terms of fermions and bosons broke down in two dimensions [LM77]. The subject of anyons was then taken over by Wilczek who published a series of seminal papers on the topic [Wil82a, Wil82b, ASW84]. It was in these papers that Wilczek observed that anyons were present in the quantum Hall effect, and hence connected the theory of anyons and topological order together.

[WORK: what is the history of gauge theory, and when was it introduced to the picture? A great reference is the de Wild Propitius and Bais survey. Also should mention Kitaev's paper again.]

Exercises:

3.1. For vertices v and plaquettes p , define

$$A'_{v,1} = \bigotimes_{\substack{\text{edges} \\ \text{touching } v}} \sigma_Z, \quad A'_v = \frac{1}{2} (I + A'_{v,1}),$$

$$B'_{p,1} = \bigotimes_{\substack{\text{edges} \\ \text{bounding } p}} \sigma_X, \quad B'_p = \frac{1}{2} (I + B'_{p,1}),$$

and

$$H' = \sum_{\text{vertices } v} (1 - A'_v) + \sum_{\text{plaquettes } p} (1 - B'_p).$$

Define $M : \mathbb{C}^2 \rightarrow \mathbb{C}^2$ by $M(|0\rangle) = \frac{1}{\sqrt{2}}(|0\rangle + |1\rangle)$ and $M(|1\rangle) = \frac{1}{\sqrt{2}}(|0\rangle - |1\rangle)$. Show that

$$\sigma_X = M\sigma_Z M^{-1}, \quad \sigma_Z = M\sigma_X M^{-1},$$

and show that H and H' are similar in the sense that $H' = MHM^{-1}$. Use this to conclude that all basis independent properties of the toric code are formally symmetric by replacing σ_X with σ_Z . For example, conclude that the codespace of H' is 4 dimensional.

4 Category theory

4.1 Overview

4.1.1 Introduction

There is a lot of math in the world. The development of the subject has spanned thousands of years, and has enjoyed a large uptick in the last two hundred or so. This has given ample time for the most important ideas to rise to the top. Among these important concepts there is one which is the focus of chapter: **composition**.

Let A, B, C be sets. Let $f : A \rightarrow B$ and $g : B \rightarrow C$ be functions. The *composition* of f and g is the function $g \circ f : A \rightarrow C$ defined by the formula $(g \circ f)(x) = g(f(x))$ for all $x \in A$. More generally, composition is the act of doing one process followed by a second process. Composition is distinguished in its importance for two reasons:

1. Composition is ubiquitous;
2. A very large number of more complicated structures can be described in terms of composition.

These two primary reasons of importance lead to several emergent applications of composition:

1. It's a good organization principle - thinking in terms of composition gives a unified approach to disparate subjects, which highlights the universality latent within mathematics;
2. It's a good compression technique - in a composition-forward approach there's no need to remember details about objects or functions between them, only the way that those functions compose is used;
3. Sometimes composition rules are the only data we have, making a composition-forward technique the only approach possible.

This third point is the situation we find ourselves in with the algebraic theory of topological quantum information. We're trying to give a usable mathematical description of topologically ordered systems. The way that we find ourselves doing this is by focusing on anyons (the local quasiparticle excitations in topological order). Doing this we run into three important points:

1. Describing anyons exactly is hard. They are emergent phenomena, found within highly-entangled energy eigenstates of arbitrarily complicated gapped Hamiltonians.
2. Describing the possible ways that anyons can transform is hard. This involved specifying intricate unitary operators on high-dimensional Hilbert spaces.
3. Describing the ways that these transformations compose with one another is always relatively simple. It can be done using explicit-to-describe rules, which are independent of system size or choice of gapped Hamiltonian.

What to do in this situation is clear: we will take a composition-first approach to anyons.

We give some examples to demonstrate our point. Suppose we want to discuss braiding anyons in the toric code. We can abstractly talk about a syndrome of the toric code in which there is one X -type particle and one Z -type particle:

[WORK: write out state.]

On these states we can talk about braiding. We use the same sorts of spactime diagrams as before to represent these transformations:

[WORK: write out state.]

Without talking about the fact that transformations of this type are realized explicitly using Pauli operators, we can still abstractly discuss the way they compose with each other:

[WORK: write out composition line]

[WORK: add more complicated example coming from whatever case of Kitaev quantum double I describe explicitly in the TO chapter]

The mathematical objects which allows one to speak intelligently about composition-first approaches is known as a *category*. The composition-first approach to mathematics is known as *category theory*. Of course, to describe anyons we will need more than just the structure of composition. We will also need a way to encode what happens when we put anyons together, braid them, and fuse them. These structures are all completely compatible with the composition-first approach, and correspond to adding extra structures onto the category. The categories describing anyons will have all their extra structures known as a *modular tensor category*, and will be the subject of much of this book. This chapter deals with introducing category theory, as well as some of the structures which will be important for discussing anyons and modular tensor categories.

4.1.2 Definition and important observations

As discussed before, a category is the structure which allows for a composition-first approach to map. Before going forward let's give a formal definition of category:

Definition (Category). A category is the following data:

1. (Objects) A set \mathcal{C} .
2. (Morphisms) A set $\text{Hom}(A, B)$ for all $A, B \in \mathcal{C}$
3. (Composition) Functions

$$\circ : \text{Hom}(B, C) \times \text{Hom}(A, B) \rightarrow \text{Hom}(A, C)$$

for all $A, B, C \in \mathcal{C}$.

Such that:

1. $(h \circ g) \circ f = h \circ (g \circ f)$, for all morphisms $f \in \text{Hom}(A, B)$, $g \in \text{Hom}(B, C)$, $h \in \text{Hom}(C, D)$, and objects $A, B, C, D \in \mathcal{C}$.
2. (Identity) For all objects $A \in \mathcal{C}$ there exists a morphism $\text{id}_A : A \rightarrow A$ such that for all $B \in \mathcal{C}$, $f \in \text{Hom}(A, B)$, and $g \in \text{Hom}(B, A)$,

$$f \circ \text{id}_A = f, \quad \text{id}_A \circ g = g.$$

□

The rest of this section will contain a loosely-related series of five observations about this definition:

Observation 1: *The structure of this definition is very typical of algebra.*

Roughly, algebra is defined to be the study of algebraic structures. An algebraic structure, roughly, is defined to be some collection structures on some space, with rules outlining how the structures interact with each other. The general way of defining algebraic structures is to first list the structures, and then list the axioms of how these structures interact with each other. We will see many definitions of this sort throughout the rest of the book, so it is good to get used to it now.

Observation 2: *In this text we have already seen many examples of categories.*

We list some of them below:

- **Set**, the category of sets. The objects are sets and the morphisms are functions.
- **Top**, the category of topological spaces. The objects are topological spaces and the morphisms are continuous functions.
- **Vec_k**, the category of finite dimensional vector spaces over a field k . The objects are finite dimensional vector spaces over k and the morphisms are linear operators.
- **Grp**, the category of finite groups. The objects are finite groups and the morphisms are group homomorphisms.
- **Hilb**, the category of quantum systems. The objects are finite dimensional Hilbert spaces and the morphisms are unitary operators.
- **Prob**, the category of probability spaces. The objects are finite dimensional real vector spaces with distinguished bases and the morphisms are operators which send normalized vectors to normalized vectors.
- **Ord_M**, the category associated with ordered media with order space M . The objects are continuous maps $\phi : \mathbb{R}^2 \rightarrow M$ and the morphisms are continuous deformations.
- **D(G)**, the category associated with discrete gauge theory based on the finite group G . The objects are G -graded G -representations and the morphisms are linear maps which respect both the G -grading and the G -action.

Observation 3: *The objects and morphisms of a category do not have much complexity implicit to them. All of the interesting structure is encoded within the composition structure.*

This is despite the fact that when we listed our examples in Observation 2 we only described the objects and morphisms, and not the composition structure. The reason for this is that the composition structure between morphisms in all of our examples is clear. In all our examples the objects are sets with extra structure, and the morphisms are maps of sets. The composition structure is inherited from the composition structure on functions between sets.

Going further, however, we observe that objects in abstract categories are *not* required to be sets and the morphisms are *not* required to be functions of sets. Most of our examples

of categories will have objects which are sets and morphisms which are functions of sets, but there will be notable counterexamples. It is important to remember that there are some categories for which there is no interpretation of morphisms as functions between sets [Fre70].

Observation 4: *A category isn't just a space with a good notion of composition - it also has identity maps.*

These identity maps are important, and we include them in the definition purposefully. There are two primary reasons: firstly that all of the relevant examples of categories will have identity maps, and secondly that most interesting properties of categories only make sense because of the identity maps. Hence if we didn't require identity maps then we would find ourselves constantly requiring them as a condition, which is a waste of space.

It is important to take a closer look at what the identity map means, though. The identity map is trying to capture a very general phenomenon about transformations: *there is always the trivial transformation which results from doing nothing*. This do-nothing map is the identity. In the category of sets, the identity maps on the set A is given by the formula $\text{id}_A(x) = x$ for all $x \in A$. The fact that these maps are the identities in the category of sets is the reason that the identity axiom for categories is defined like it is. Really, there is an implicit lemma hidden in the definition of category:

Lemma 4.1. *Let A be a set. For all sets B and for all $f : A \rightarrow B$, $g : B \rightarrow A$ we have*

$$f \circ \text{id}_A = f, \quad \text{id}_A \circ g = g.$$

*In particular, id_A satisfies the axiom of an identity in the category of sets, and hence **Set** forms a category.*

Proof. .[WORK: do proof] □

These sorts of implicit lemmas are everywhere in category theory. Whenever a composition-forward definition is given in category theory, there is the assumption that it agrees with the standard definition at least in the category of sets. For instance, we make the following definition:

Definition (Isomorphism). Let \mathcal{C} be a category, let $A, B \in \mathcal{C}$ be objects, and let $f : A \rightarrow B$ be a morphism. We say that f is an *isomorphism* if there exists a morphism $f^{-1} : B \rightarrow A$ such that $f^{-1} \circ f = \text{id}_A$ and $f \circ f^{-1} = \text{id}_B$. We call f^{-1} the *inverse* of f . In this case, we say that A and B are *isomorphic objects*.

□

The implicit lemma in this definition is as follows:

Lemma 4.2. *Let A, B be sets, and let $f : A \rightarrow B$ be a function. The map f is a bijection if and only if there exists a function $f^{-1} : B \rightarrow A$ such that $f^{-1} \circ f = \text{id}_A$ and $f \circ f^{-1} = \text{id}_B$. In particular, a function f in the category **Set** is an isomorphism if and only if it is a bijection.*

Proof. .[WORK: do proof] □

Observation 5: *Statements in category theory can be very broadly applied.*

This is in some sense obvious by the fact that there are so many different examples of categories, but it's good to state the observation explicitly. Here's a good example:

Proposition 4.1. *Let \mathcal{C} be a category. Identities in \mathcal{C} are unique. Explicitely, let $A \in \mathcal{C}$ be an object and let $\text{id}_A, \tilde{\text{id}}_A : A \rightarrow A$ be morphisms satisfying the identity axiom. We have that $\text{id}_A = \tilde{\text{id}}_A$.*

Proof. . Using the fact that $\text{id}_A \circ f = f$ and $f \circ \tilde{\text{id}}_A = f$ for any $f : A \rightarrow A$, we compute that

$$\text{id}_A = \text{id}_A \circ \tilde{\text{id}}_A = \tilde{\text{id}}_A$$

as desired. \square

This has broad application. For instance: why are identity elements in groups unique? Let G be a group and let $1, 1' \in G$ be identity elements. We find that $1 = 1 \cdot 1' = 1'$ as desired. Going further, here is another proposition in category theory:

Proposition 4.2. *Let \mathcal{C} be a category. Let A, B be objects and let $f : A \rightarrow B$ be an isomorphism. The inverse of f is unique. That is, let f^{-1}, \tilde{f}^{-1} be morphisms satisfying the definition of the inverse of f . We have that $f^{-1} = \tilde{f}^{-1}$.*

Proof. Using the associativity axiom, we compute

$$f^{-1} = f^{-1} \circ \text{id}_B = f^{-1} \circ (f \circ \tilde{f}^{-1}) = (f^{-1} \circ f) \circ \tilde{f}^{-1} = \text{id}_A \circ \tilde{f}^{-1} = \tilde{f}^{-1}$$

as desired. \square

This is very general. Why are inverses unique in groups? Why are inverses of matrices unique? Abstractly, why should the inverse of any reversible process be unique? Proposition 4.2 gives the answer.

4.2 Structures in category theory

[WORK: this section should include all of the structures which are neccecary for the rest of the book, and are too cumbersome to define on-site. It should also read as an introducing to how to think in the language of categories. Here is the running list of neccecary topics

- Products/coproducts/biproducts;
- zero objects;
- \mathbb{C} -linear structure;
- Functors, natural equivalence, equivalence of categories, NOT Yoneda lemma;
- Opposite category. This is relevant for the discussions of time-reversal which will be at play during the book. Namely, given an MTC I want to define $\overline{\mathcal{C}}$ so that $Z(\mathcal{C}) \cong \mathcal{C} \boxtimes \overline{\mathcal{C}}$. Also, when talking about module categories, the category \mathcal{M}^{op} is going to show up when we reverse the direction of the boundary.

]

[WORK: maybe use homotopy theory as a reccuring motivating example?]

Definition (\mathbb{C} -linear category). A \mathbb{C} -linear category is the following data:

1. A category \mathcal{C} ;

2. The structure of a \mathbb{C} -vector space on $\text{Hom}(A, B)$ for all $A, B \in \mathcal{C}$.

Such that:

1. The composition maps $\circ : \text{Hom}(B, C) \times \text{Hom}(A, B) \rightarrow \text{Hom}(A, C)$ are bilinear maps of vector spaces for all $A, B, C \in \mathcal{C}$.

□

Definition (\mathbb{C} -linear functor). A \mathbb{C} -linear functor between \mathbb{C} -linear categories \mathcal{C}, \mathcal{D} is a functor $F : \mathcal{C} \rightarrow \mathcal{D}$ such that $F : \text{Hom}_{\mathcal{C}}(A, B) \rightarrow \text{Hom}_{\mathcal{D}}(F(A), F(B))$ is a linear map of vector spaces for all $A, B \in \mathcal{C}$.

□

[WORK: do I need to define \mathbb{C} -linear natural transformation?]

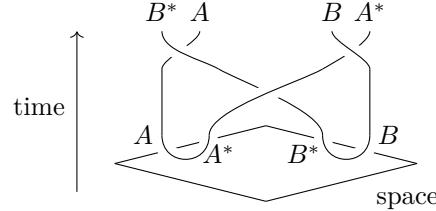
4.3 Monoidal categories

4.3.1 Motivation, definition, and string diagrams

[WORK: an early reference for string diagrams is [Mou84]. Maybe I should cite it?]

The goal of this section is to introduce the language necessary for a proper detailed discussion of modular tensor categories. Despite the fact that the language of composition is very useful for MTCs, there are still many concepts in MTC theory which require more structure than just composition. In the sections that follow we will introduce these structures one-by-one, giving motivation and proving basic properties along the way.

By the end of our discussion, we will be able to discuss situations like these, where we create and braid quasiparticles:



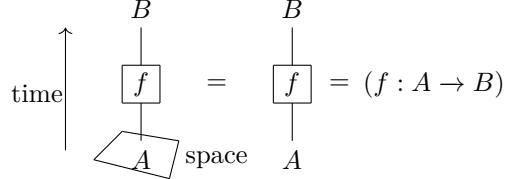
This sort of diagram will take place in some category \mathcal{C} . The labels A, B, A^*, B^* represent objects in \mathcal{C} . The objects (A, A^*) form a particle/antiparticle pair, and the objects (B, B^*) form a particle/antiparticle pair. Naively, we could interpret this diagram as follows. To begin, there are no particles. Then, we have creation maps create_{A,A^*} and $\text{create}_{B^*,B}$ which pairs of particles and their antiparticles. Then we have three different braids, braid_{A^*,B^*} , braid_{A,B^*} , and $\text{braid}_{A^*,B}$. The overall process is the composition of these:

$$(\text{braid}_{A^*,B}) \circ (\text{braid}_{A,B^*}) \circ (\text{braid}_{A^*,B^*}) \circ (\text{create}_{B^*,B}) \circ (\text{create}_{A,A^*}).$$

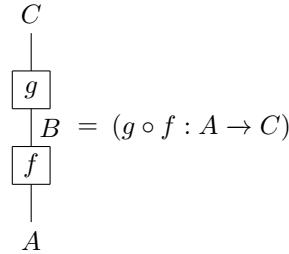
These creation maps and braiding maps are exactly the sort of maps which we will be introducing as extra structures on our category during this chapter. One lingering problem, however, is that the naive approach to formalising these diagrams in category theory results in long chains of composition. These long chains of composition hide the real structure of the problem, and make processes like the one in diagram [ref] much harder to parse. It is

for this reason that we introduce a *graphical language for category theory*. This graphical language makes the diagrams like [ref] rigorous mathematical notion which describe well-defined morphisms. These diagrams are known as *string diagrams*.

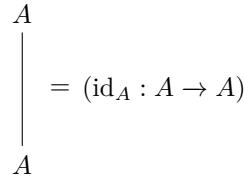
The main principle of string diagrams is that morphisms are represented as follows:



The direction of time going from bottom to up and the space being two-dimensional slices is the same in every diagram, and hence is left implicit from here on out. Composition can be expressed cleanly in this language as stacking, for all $f : A \rightarrow B, g : B \rightarrow C$:



Accordingly, the identity map has a simple implementation:



We now give our first major example of adding structure, and how that structure can be interpreted in terms of string diagrams. This structure is that of a *monoidal category*. For technical reasons we only define *strict* monoidal categories for now - we will come back to the general definition later. Monoidal categories give a way to put objects together. For instance, in diagram [ref] we had four particles all together. We need a way to discuss composite-particle systems. In quantum mechanics, forming a composite system is done by taking the tensor product. Hence, we will use the notation \otimes for joining particles in our current setting. We will even use the term “tensor product” to discuss it. In general, joining two systems is one way of going from pairs of systems to individual systems:

$$\begin{aligned} & (\text{systems}) \times (\text{systems}) \rightarrow (\text{systems}). \\ & (\text{system 1}, \text{system 2}) \mapsto (\text{system 1}) \otimes (\text{system 2}) \end{aligned}$$

In the world of category theory, we only require some basic properties of this joining. Namely, it should be functorial and satisfy some simple conditions:

Definition (Strict monoidal category). A strict monoidal category is the following data:

1. A category \mathcal{C} ;
2. (Tensor product) A functor $\otimes : \mathcal{C} \times \mathcal{C} \rightarrow \mathcal{C}$;
3. (Unit) A distinguished element $\mathbf{1} \in \mathcal{C}$;

Such that:

1. (Unit axiom) Let $A, A' \in \mathcal{C}$ be objects and let $f : A \rightarrow A'$ be a morphism. We have

$$A \otimes \mathbf{1} = \mathbf{1} \otimes A = A, \quad f \otimes \text{id}_{\mathbf{1}} = \text{id}_{\mathbf{1}} \otimes f = f.$$

2. (Associativity) Let $A, B, C, A', B', C' \in \mathcal{C}$ be objects, and let $f : A \rightarrow A'$, $g : B \rightarrow B'$, $h : C \rightarrow C'$ be morphisms. We have

$$(A \otimes B) \otimes C = A \otimes (B \otimes C), \quad (f \otimes g) \otimes h = f \otimes (g \otimes h).$$

□

The object $\mathbf{1} \in \mathcal{C}$ is important. Just like how groups of symmetries always include the “do-nothing” symmetry, strict monoidal categories should always include the unit. In this case, $\mathbf{1} \in \mathcal{C}$ represents the empty particle - no particle at all. In every particle theory there should be the possibility of not having any particles. Joining the empty particle with any other particle should obviously do nothing, hence the axiom $\mathbf{1} \otimes A = A \otimes \mathbf{1} = A$.

We can now work strict monoidal categories into our graphical language. The tensor product of two objects is represented by putting two lines adjacent to one another. For instance, let \mathcal{C} be a strict monoidal category, let $A, B, C, D \in \mathcal{C}$ be four objects, and let $f : A \rightarrow C$, $g : B \rightarrow D$ be morphisms. We have

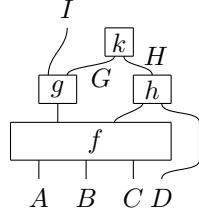
$$\begin{array}{ccc} C & & D \\ \downarrow & & \downarrow \\ \boxed{f} & \quad \boxed{g} & = (f \otimes g : A \otimes B \rightarrow C \otimes D) \\ \downarrow & & \downarrow \\ A & & B \end{array}$$

The monoidal unit $\mathbf{1}$ is distinguished in monoidal categories, and hence is represented with a special line. We will either use a dotted line, or no line at all:

$$\begin{array}{ccc} A & & A \\ \downarrow & & \downarrow \\ \boxed{f} & = & \boxed{f} = (f : \mathbf{1} \rightarrow A) \\ \vdots & & \\ \mathbf{1} & & \end{array}$$

We *do not* require that the lines drawn in string diagrams be straight. They can curve any amount so long as it is clear that they are directly connecting an output to an input. The lines cannot cross each other or double back. Additionally, when it is clear from context, we *do not* require ourselves to include every label. For example, the following is a valid

diagram in all strict monoidal categories \mathcal{C} , where $A, B, C, D, E, F, G, H \in \mathcal{C}$ are objects, and $f : A \otimes B \otimes C \rightarrow E \times F$, $g : E \rightarrow I \otimes G$, $h : F \otimes D \rightarrow H$, and $k : G \otimes H \rightarrow \mathbf{1}$ are morphisms:



4.3.2 Braided monoidal categories

We continue our definitions of structures on monoidal categories, and their expression in the language of string diagrams. Our next definition is that of a strict *braided* monoidal category:

Definition (Strict braided monoidal category). A strict braided monoidal category is the following data:

1. A strict monoidal category \mathcal{C} ;
2. (Braiding) Isomorphisms $\beta_{A,B} : A \otimes B \rightarrow B \otimes A$ for all $A, B \in \mathcal{C}$ which form a natural isomorphism between the functors $\mathcal{C} \times \mathcal{C} \rightarrow \mathcal{C}$ given by $(A, B) \mapsto A \otimes B$ and $(A, B) \mapsto B \otimes A$.

Such that for all $A, B, C \in \mathcal{C}$, the diagrams

$$\begin{array}{ccc}
 A \otimes B \otimes C & & \\
 \downarrow \beta_{A,B \otimes C} & \searrow \beta_{A,B} \otimes \text{id}_C & \\
 B \otimes A \otimes C & & \\
 & \swarrow \text{id}_B \otimes \beta_{A,C} & \\
 B \otimes C \otimes A & &
 \end{array}$$

and

$$\begin{array}{ccc}
 A \otimes B \otimes C & & \\
 \downarrow \beta_{B \otimes C, A}^{-1} & \searrow \beta_{B,A}^{-1} \otimes \text{id}_C & \\
 B \otimes A \otimes C & & \\
 & \swarrow \text{id}_B \otimes \beta_{C,A}^{-1} & \\
 B \otimes C \otimes A & &
 \end{array}$$

commute.

□

The idea for how to implement braided monoidal categories in the language of string diagrams is to introduce a special symbol for the braiding map $\beta_{A,B}$. Namely, we define graphically overcrossing and undercrossing as follows:

$$\begin{array}{c} B & A \\ \diagup & \diagdown \\ A & B \end{array} = \beta_{A,B}, \quad
 \begin{array}{c} B & A \\ \diagdown & \diagup \\ A & B \end{array} = \beta_{B,A}^{-1}.$$

The fact that overcrossing and undercrossing are related by an inverse encodes the following key fact that

$$\begin{array}{c} A & B \\ \diagup & \diagdown \\ \text{overcrossing} \end{array} = \beta_{A,B}^{-1} \circ \beta_{A,B} = \text{id}_{A \otimes B} = \begin{array}{c} A & B \\ | & | \\ A & B \end{array}$$

We can now describe the conditions on a strict braided monoidal category in a graphical way. The fact that β is a natural transformation can be reinterpreted as follows:

Lemma 4.3. *Let \mathcal{C} be a strict braided monoidal category. For all $A, B, C, D \in \mathcal{C}$ and $f : A \rightarrow C, g : B \rightarrow D$, we have the following equality of string diagrams:*

$$\begin{array}{c} D & C \\ \diagup & \diagdown \\ f & g \\ | & | \\ A & B \end{array} = \begin{array}{c} D & C \\ \square g & \square f \\ \diagup & \diagdown \\ A & B \end{array}$$

The same formula holds replacing overcrossing with undercrossing on both sides.

Proof. Consider the morphism $(f, g) : (A, B) \rightarrow (C, D)$ in $\mathcal{C} \times \mathcal{C}$. The naturality of β implies the following commutative square:

$$\begin{array}{ccc} A \otimes B & \xrightarrow{f \otimes g} & C \otimes D \\ \downarrow \beta_{A,B} & & \downarrow \beta_{C,D} \\ B \otimes A & \xrightarrow{g \otimes f} & D \otimes C \end{array}$$

Expanding this square in diagrammatic language gives the first part of the proposition. Reversing the direction of the arrows by taking inverses gives the second part. \square

The coherence axiom can be stated diagrammatically as follows,

$$\begin{array}{c} B & C & A \\ \diagup & \diagdown & \diagdown \\ \beta_{A,B \otimes C} & A & B & C \end{array} = \begin{array}{c} B & C & A \\ \diagup & \diagdown & \diagdown \\ \beta_{A,B} & A & B & C \\ \beta_{A,C} & \text{---} & \text{---} \end{array}$$

and similarly with replacing overcrossing with undercrossing. The importance of this axiom is that it means that our graphical language can express braid diagrams without other ambiguity. We can safely deform strings behind braids and not need to worry about whether we are applying $\beta_{A,B \otimes C}$ or $(\text{id } B \otimes \beta_{A,C}) \circ (\beta_{A,B} \otimes \text{id}_C)$.

A fundamental result about braided monoidal categories is that the coherence condition given is enough to rearrange braids at will. In particular, we have the following key proposition:

Proposition 4.3 (Yang-Baxter equation). *Let \mathcal{C} be a strict braided monoidal category. Let $A, B, C \in \mathcal{C}$ be objects. We have*

$$\begin{array}{ccc}
 \begin{array}{c}
 \text{Diagram showing three strands labeled } A, B, \text{ and } C. \\
 \text{Strand } C \text{ has a dashed arrow } \beta_{B,C} \text{ pointing right.} \\
 \text{Strand } A \text{ has a dashed arrow } \beta_{A,C} \text{ pointing left.} \\
 \text{Strand } B \text{ has a dashed arrow } \beta_{A,B} \text{ pointing right.} \\
 \text{The strands cross in a standard braid pattern.}
 \end{array}
 & = &
 \begin{array}{c}
 \text{Diagram showing the same strands } A, B, \text{ and } C. \\
 \text{Strand } C \text{ has a dashed arrow } \beta_{A,C} \text{ pointing left.} \\
 \text{Strand } A \text{ has a dashed arrow } \beta_{B,C} \text{ pointing right.} \\
 \text{Strand } B \text{ has a dashed arrow } \beta_{A,B} \text{ pointing right.} \\
 \text{The strands cross in a different braid pattern, illustrating the Yang-Baxter equation.}
 \end{array}
 \end{array}$$

Proof. We offer a graphical proof, using first the coherence condition and then naturality:

$$\begin{array}{c}
\text{Diagram 1: } \begin{array}{ccc} C & B & A \\ \swarrow & \curvearrowright & \downarrow \\ A & B & C \end{array} = \begin{array}{ccc} C & B \otimes A & \\ \swarrow & \curvearrowright & \downarrow \\ \boxed{\beta_{A,B}} & A & B & C \end{array} \\
\\
\text{Diagram 2: } \begin{array}{ccc} C & B \otimes A & \\ \swarrow & \curvearrowright & \downarrow \\ \boxed{\beta_{A,B}} & A & B & C \end{array} = \begin{array}{ccc} C & B & A \\ \swarrow & \curvearrowright & \downarrow \\ A & B & C \end{array}.
\end{array}$$

□

We get the following corrolary:

Corollary 4.1. *Let \mathcal{C} be a strict braided monoidal category. Let $A \in \mathcal{C}$ be an object. The map*

$$\begin{aligned}
B_n &\rightarrow \text{Aut}(A^{\otimes n}) \\
\sigma_i &\mapsto \text{id}_{A^{\otimes i-1}} \otimes \beta_{A,A} \otimes \text{id}_{A^{n-i-1}}
\end{aligned}$$

is a homomorphism of groups.

Proof. We saw in Proposition [ref] that the relations on B_n are generated by the conditons $\sigma_{i+1}\sigma_i\sigma_{i+1} = \sigma_i\sigma_{i+1}\sigma_i$. These conditions are satisfied by the braiding by Proposition [ref]. Hence, the map is a homomorphism of groups. □

4.3.3 Examples, equivalences, and MacLane's coherence theorem

In this section we will give concrete examples of monoidal categories and braided monoidal categories. What we will find, however, is that these examples will all demonstrate the same subtle problem. For example, here is a category which we would want to give as an example of a monoidal category:

$$\mathcal{C} = \mathbf{Set}, \otimes = \text{Cartesian product}.$$

The Cartesian product is certainly functorial. Namely, given morphisms $f : A \rightarrow C$ and $g : B \rightarrow D$ we get a morphism

$$(f \times g) : A \times B \rightarrow C \times D,$$

$$(a, b) \mapsto (f(a), g(b))$$

However we get a key issue $(A \times B) \times C \neq A \times (B \times C)$. We have an isomorphism

$$\alpha : (A \times B) \times C \rightarrow A \times (B \times C),$$

$$((a, b), c) \mapsto (a, (b, c))$$

but this isomorphism is *not* an equality. This means that **Set** does not satisfy the definition of a strict monoidal category! In general, all the examples we would want to give of monoidal categories fail to be strict monoidal categories. In this section we discuss a method for loosening the definition of monoidal category so that **Set** and other examples can be included in the definition.

For this reason, we give the following warning: **This section is not necessary for a conceptual understanding of the subject matter. It is material of technical importance, and thus of interest to those who want a correct formal understanding of the mathematics at play.**

This is because, despite the fact that we will loosen the notation of strict monoidal category to a more general sort of possibly non-strict category, we will do the following:

We assume monoidal categories are strict whenever it is convenient.

The fact that this does not cause issues is a corollary of MacLane's coherence theorem. We will discuss these issues in detail in this section.

The most naive way of loosening the definition of monoidal category is to only enforce the condition $(A \otimes B) \otimes C \cong A \otimes (B \otimes C)$ instead of equality. However, this leads to a problem. The associativity axiom on morphisms $(f \otimes g) \otimes h = f \otimes (g \otimes h)$ no longer makes sense because there is no way of comparing morphisms on $(A \otimes B) \otimes C$ and $A \otimes (B \otimes C)$. In general category theory fashion, we should choose specific isomorphisms $\alpha_{A,B,C} : (A \otimes B) \otimes C \xrightarrow{\sim} A \otimes (B \otimes C)$ and require that those isomorphisms satisfy certain coherence conditions. This leads us to our definition of a non-strict monoidal category:

Definition (Monoidal category). A monoidal category is the following data:

1. A category \mathcal{C} .
2. (Tensor product) A functor $\otimes : \mathcal{C} \times \mathcal{C} \rightarrow \mathcal{C}$.

3. (Unit) A distinguished element $\mathbf{1} \in \mathcal{C}$.

4. (Associator) A natural isomorphism

$$\alpha : - \otimes (- \otimes -) \xrightarrow{\sim} (- \otimes -) \otimes -,$$

where $- \otimes (- \otimes -)$ denotes the functor $\mathcal{C} \times \mathcal{C} \times \mathcal{C} \rightarrow \mathcal{C}$ sending (A, B, C) to $A \otimes (B \otimes C)$, and similarly for $(- \otimes -) \otimes -$.

- 5. (Left unit) A natural isomorphism $\lambda : \mathbf{1} \otimes - \xrightarrow{\sim} -$, where $\mathbf{1} \otimes -$ denotes the functor $\mathcal{C} \rightarrow \mathcal{C}$ sending A to $\mathbf{1} \otimes A$, and $-$ denotes the identity.
- 6. (Right unit) A natural isomorphism $\rho : - \otimes \mathbf{1} \xrightarrow{\sim} -$, where $- \otimes \mathbf{1}$ is the functor $\mathcal{C} \rightarrow \mathcal{C}$ sending A to $A \otimes \mathbf{1}$.

Additionally, a monoidal category is required to satisfy the following properties:

1. (Triangle identity) The diagram

$$\begin{array}{ccc} (A \otimes \mathbf{1}) \otimes B & \xrightarrow{\alpha_{A, \mathbf{1}, B}} & A \otimes (\mathbf{1} \otimes B) \\ & \searrow \rho_A \otimes \text{id}_B & \swarrow \text{id}_A \otimes \lambda_B \\ & A \otimes B & \end{array}$$

commutes for all $A, B \in \mathcal{C}$.

2. (Pentagon identity) The diagram

$$\begin{array}{ccccc} & & (A \otimes B) \otimes (C \otimes D) & & \\ & \nearrow \alpha_{A \otimes B, C, D} & & \swarrow \alpha_{A, B, C \otimes D} & \\ ((A \otimes B) \otimes C) \otimes D & & & & A \otimes (B \otimes (C \otimes D)) \\ \downarrow \alpha_{A, B, C} \otimes \text{id}_D & & & & \uparrow \text{id}_A \otimes_{B, C, D} \\ (A \otimes (B \otimes C)) \otimes D & \xrightarrow{\alpha_{A, B \otimes C, D}} & & & A \otimes ((B \otimes C) \otimes D) \end{array}$$

commutes for all $A, B, C, D \in \mathcal{C}$.

□

With this more general definition, we now have many examples of monoidal categories:

Proposition 4.4. *The following collections of data form monoidal categories*

- (i) The category $\mathcal{C} = \mathbf{Set}$, with tensor product $\otimes =$ Cartesian product, monoidal unit $\mathbf{1} = \{\ast\}$, associator

$$\begin{aligned}\alpha_{A,B,C} : A \times (B \times C) &\xrightarrow{\sim} (A \times B) \times C, \\ (a, (b, c)) &\mapsto ((a, b), c)\end{aligned}$$

and unitors

$$\begin{aligned}\lambda : \mathbf{1} \otimes A &\rightarrow A \\ (*, a) &\mapsto a\end{aligned}\qquad\qquad\qquad\begin{aligned}\rho : A \otimes \mathbf{1} &\rightarrow A. \\ (a, *) &\mapsto a\end{aligned}$$

- (ii) The plain category $\mathcal{C} = \mathbf{Vec}_{\mathbb{C}}$, with its standard tensor product, monoidal unit $\mathbf{1} = \mathbb{C}$, associator

$$\begin{aligned}\alpha_{A,B,C} : A \times (B \times C) &\xrightarrow{\sim} (A \times B) \times C, \\ a \otimes (b \otimes c) &\mapsto (a \otimes b) \otimes c\end{aligned}$$

and unitors

$$\begin{aligned}\lambda : \mathbf{1} \otimes A &\rightarrow A \\ 1 \otimes a &\mapsto a\end{aligned}\qquad\qquad\qquad\begin{aligned}\rho : A \otimes \mathbf{1} &\rightarrow A. \\ a \otimes 1 &\mapsto a\end{aligned}$$

- (iii) The category $\mathcal{C} = \mathbf{Set}$ with tensor product $\otimes = \text{Disjoint union}$ and $\mathbf{1} = \{\}$, with a standard choice of associators and unitors;

- (iv) The category $\mathcal{C} = \mathbf{Vec}_{\mathbb{C}}$ with tensor product $\otimes = \text{Direct sum}$, and $\mathbf{1} = 0$, with a standard choice of associators and unitors.

Proof. These facts are straightforward to verify, and are left as an exercise to the reader. \square

In expanding our definition from strict monoidal category to monoidal category, however, we have introduced a subtle problem. The diagram

$$\begin{array}{ccc} A & B & C \\ | & | & | \\ A & B & C \end{array} = \text{id}_{A \otimes B \otimes C}$$

no longer makes sense! The map $\text{id}_{A \otimes B \otimes C}$ no longer exists, because $A \otimes B \otimes C$ no longer exists. One must make a choice of $(A \otimes B) \otimes C$ or $A \otimes (B \otimes C)$. These maps may be isomorphic, but they have no need to be equal! The correct diagram would be

$$\begin{array}{ccc} A & B & C \\ | & | & | \\ \boxed{\alpha_{A,B,C}} & & \\ | & | & | \\ A & B & C \end{array}$$

All string diagrams would now need α maps thrown in at key points to make a well-defined language. This is exceedingly complicated, and has deep issues that need to be addressed. Hence, we maintain that our graphical language only applies to strict monoidal categories.

This means that we haven't really gotten anywhere. We defined the notion of a non-strict monoidal category so that we could include our favorite examples, but then we observed that string diagrams still fail to describe those examples! This seemingly bad situation is rectified by the following theorem, which we first state informally.

MacLane's coherence theorem: *every monoidal category is equivalent to a strict monoidal category.*

This gives us a workflow for the book. We will frame our discussion so that it applies to arbitrary monoidal categories. That way, all our usual examples are included. Then, when we want to use string diagrams, we use MacLane's coherence theorem to pass to an equivalent strict category, in which our diagrams make sense. Then, when we are done using the diagram, we pass the conclusion of the argument through the equivalence! We will be using this subtle technique repeatedly throughout the book. To save time and energy, we won't explicitly mention it. We will implicitly pass to an equivalent strict category without making any special note.

Sometimes we will want to pass to a strict monoidal category even before string diagrams come into play. For instance, in Proposition [ref] we proved that every strict braided monoidal category \mathcal{C} comes paired with a group homomorphism

$$B_n \rightarrow \text{Aut}(A^{\otimes n})$$

for all $A \in \mathcal{C}$, $n \geq 1$. Once we generalize strict braided monoidal categories to possibly non-strict braided monoidal categories, this proposition will become false. The object $A^{\otimes n}$ does not exist - a choice of parenthesization needs to be made. Every time that an element of the braid group acts on $A^{\otimes n}$, the parentheses need to be re-arranged using associators, then the braiding map β can be applied, and then the parentheses need to be re-arranged back into their original position using associators again.

Not only does this non-strict version of Proposition [ref] take more time and space to set-up, but it also leads to potential thorny issues. There are multiple ways to rearrange parentheses from one starting point to the other. How do we know that they will all give the same map, and hence into a well-defined homomorphism from B_n ? It follows from general combinatorial principles and a repeated application of the pentagon identity.

This is indicative of the general feeling of working with non-strict monoidal categories. Statements and proofs which were obvious for strict monoidal categories become needlessly unintuitive for non-strict monoidal categories. Hence, it is much better to pass to a strict monoidal category using MacLane's coherence theorem at our first convenience.

Of course, all of this discussion rests on the notion of *equivalent* in MacLane's coherence theorem being well defined, so that we can pass information back and forth through the equivalence. Our notation of equivalence is modeled after the more general notation of equivalence of categories - a pair of functors whose compositions are both naturally isomorphic to the identity. To translate to the present setting, we need a good notion of monoidal functor and monoidal natural transformation so that the equivalence can pass through information about the monoidal structure.

Definition (Monoidal functor). A monoidal functor between monoidal categories $(\mathcal{C}, \otimes_{\mathcal{C}}, \alpha_{\mathcal{C}}, \lambda_{\mathcal{C}}, \rho_{\mathcal{C}}, \mathbf{1}_{\mathcal{C}})$ and $(\mathcal{D}, \otimes_{\mathcal{D}}, \alpha_{\mathcal{D}}, \lambda_{\mathcal{D}}, \rho_{\mathcal{D}}, \mathbf{1}_{\mathcal{D}})$ is the following data:

1. A functor $F : \mathcal{C} \rightarrow \mathcal{D}$.
2. A morphism $\epsilon : 1_{\mathcal{D}} \rightarrow F(1_{\mathcal{C}})$.
3. A natural isomorphism $\mu : F(-) \otimes_{\mathcal{D}} F(-) \xrightarrow{\sim} F(- \otimes_{\mathcal{C}} -)$.

Additionally, a monoidal functor is required to satisfy the following properties:

1. (Associativity) The diagram

$$\begin{array}{ccc}
 (F(A) \otimes_{\mathcal{D}} F(B)) \otimes_{\mathcal{D}} F(C) & \xrightarrow{\alpha_{\mathcal{D}; F(A), F(B), F(C)}} & F(A) \otimes_{\mathcal{D}} (F(B) \otimes_{\mathcal{D}} F(C)) \\
 \downarrow \mu_{A,B} \otimes \text{id}_{F(C)} & & \downarrow \text{id}_{F(A)} \otimes \mu_{B,C} \\
 F(A \otimes_{\mathcal{C}} B) \otimes_{\mathcal{D}} F(C) & & F(A) \otimes_{\mathcal{D}} F(B \otimes_{\mathcal{C}} C) \\
 \downarrow \mu_{A \otimes_{\mathcal{C}} B, C} & & \downarrow \mu_{A,B} \otimes_{\mathcal{C}} C \\
 F((A \otimes_{\mathcal{C}} B) \otimes_{\mathcal{C}} C) & \xrightarrow{F(\alpha_{\mathcal{C}; A, B, C})} & F(A \otimes_{\mathcal{C}} (B \otimes_{\mathcal{C}} C))
 \end{array}$$

commutes for all $A, B, C \in \mathcal{C}$.

2. (Unitality) The diagrams

$$\begin{array}{ccc}
 1_{\mathcal{D}} \otimes_{\mathcal{D}} F(A) & \xrightarrow{\epsilon \otimes \text{id}_{F(A)}} & F(1_{\mathcal{C}}) \otimes F(A) \\
 \downarrow \lambda_{\mathcal{C}; F(A)} & & \downarrow \mu_{1_{\mathcal{C}}, A} \\
 F(A) & \xleftarrow{F(\lambda_{\mathcal{C}; A})} & F(1_{\mathcal{C}} \otimes A)
 \end{array}$$

and

$$\begin{array}{ccc}
 F(A) \otimes_{\mathcal{D}} 1_{\mathcal{D}} & \xrightarrow{\text{id}_{F(A)} \otimes \epsilon} & F(A) \otimes_{\mathcal{D}} F(1_{\mathcal{C}}) \\
 \downarrow \rho_{\mathcal{C}; F(A)} & & \downarrow \mu_{A, 1_{\mathcal{C}}} \\
 F(A) & \xleftarrow{F(\rho_{\mathcal{C}; A})} & F(1_{\mathcal{C}} \otimes A)
 \end{array}$$

commute for all $A \in \mathcal{C}$.

□

Definition (Monoidal natural transformation). A monoidal natural transformation between two monoidal functors (F_0, μ_0, ϵ_0) and (F_1, μ_1, ϵ_1) between monoidal categories $(\mathcal{C}, \otimes_{\mathcal{C}}, 1_{\mathcal{C}})$ and $(\mathcal{D}, \otimes_{\mathcal{D}}, 1_{\mathcal{D}})$ is a natural transformation η between the underlying functors F_0, F_1 . Additionally, a monoidal natural transformation is required to satisfy the following properties:

1. (Compatibility with tensor product) For all objects $A, B \in \mathcal{C}$, the diagram

$$\begin{array}{ccc}
 F_0(A) \otimes_{\mathcal{D}} F_1(B) & \xrightarrow{\eta_A \otimes \eta_B} & F_1(A) \otimes_{\mathcal{D}} F_1(B) \\
 \downarrow \mu_{0; A, B} & & \downarrow \mu_{1; A, B} \\
 F_0(A \otimes_{\mathcal{C}} B) & \xrightarrow{\eta_{A \otimes B}} & F_1(A \otimes_{\mathcal{C}} B)
 \end{array}$$

commutes.

2. (Compatibility with unit) The diagram

$$\begin{array}{ccccc}
 & & 1_{\mathcal{D}} & & \\
 & \swarrow \epsilon_0 & & \searrow \epsilon_1 & \\
 F_0(1_{\mathcal{C}}) & \xrightarrow{\eta_{1_{\mathcal{C}}}} & F_1(1_{\mathcal{C}}) & &
 \end{array}$$

commutes.

□

We can now define monoidal equivalence. A *monoidal equivalence* between two monoidal categories \mathcal{C}, \mathcal{D} is a pair of monoidal functors $F : \mathcal{C} \rightarrow \mathcal{D}$, $G : \mathcal{D} \rightarrow \mathcal{C}$ such that $G \circ F$ is monoidally naturally isomorphic to $\text{id}_{\mathcal{C}}$ and $F \circ G$ is monoidally naturally isomorphic to $\text{id}_{\mathcal{D}}$. We say two categories are monoidally equivalent if there is a monoidal equivalence between them. We can now state MacLane's coherence theorem:

Theorem 4.1 (MacLane's coherence theorem,). *Every monoidal category is monoidally equivalent to a strict monoidal category.*

As we add more structure, it will be a non-trivial task to verify that we can still apply MacLane's coherence theorem. In particular, we will need to strengthen our notion of equivalence to make sure it is strong enough to pass through information about our additional structures. We can see this in the case of braidings already.

Definition (Braided monoidal category). A braided monoidal category is the following data:

1. A monoidal category $(\mathcal{C}, \otimes, \alpha, \mathbf{1})$.
2. (Braiding) A natural isomorphism β between the functor $\mathcal{C} \times \mathcal{C} \rightarrow \mathcal{C}$ sending (A, B) to $A \otimes B$, and the functor sending (A, B) to $B \otimes A$.

Additionally, a braided monoidal category is required to satisfy the following properties:

1. (Hexagon identities) The diagrams

$$\begin{array}{ccccc}
 A \otimes (B \otimes C) & \xrightarrow{\alpha_{A,B,C}} & (A \otimes B) \otimes C & \xrightarrow{\beta_{A \otimes B,C}} & C \otimes (A \otimes B) \\
 \downarrow \text{id}_A \otimes \beta_{B,C} & & & & \downarrow \alpha_{B,C,A} \\
 A \otimes (C \otimes B) & \xrightarrow{\alpha_{A,C,B}} & (A \otimes C) \otimes B & \xrightarrow{\beta_{A,C} \otimes \text{id}_B} & (C \otimes A) \otimes B
 \end{array}$$

and

$$\begin{array}{ccccc}
 (A \otimes B) \otimes C & \xrightarrow{\alpha_{A,B,C}^{-1}} & A \otimes (B \otimes C) & \xrightarrow{\beta_{A,B \otimes C}} & (B \otimes C) \otimes A \\
 \downarrow \beta_{A,B} \otimes \text{id}_C & & & & \downarrow \alpha_{B,C,A}^{-1} \\
 (B \otimes A) \otimes C & \xrightarrow{\alpha_{B,A,C}^{-1}} & B \otimes (A \otimes C) & \xrightarrow{\text{id}_B \otimes \beta_{A,C}} & B \otimes (C \otimes A)
 \end{array}$$

commute for all $A, B, C \in \mathcal{C}$.

□

Definition (Braided monoidal functor). A braided monoidal functor between braided monoidal categories $(\mathcal{C}, \otimes_{\mathcal{C}}, \beta_{\mathcal{C}})$, $(\mathcal{D}, \otimes_{\mathcal{D}}, \beta_{\mathcal{D}})$ is a monoidal functor $(F, \mu) : \mathcal{C} \rightarrow \mathcal{D}$ such that the diagram

$$\begin{array}{ccc} F(A) \otimes_{\mathcal{D}} F(B) & \xrightarrow{\beta_{\mathcal{D}; F(A), F(B)}} & F(B) \otimes_{\mathcal{D}} F(A) \\ \downarrow \mu_{A,B} & & \downarrow \mu_{B,A} \\ F(A \otimes_{\mathcal{C}} B) & \xrightarrow{F(\beta_{\mathcal{C}; F(A), F(B)})} & F(B \otimes_{\mathcal{C}} A) \end{array}$$

commutes for all $A, B \in \mathcal{C}$.

□

Thankfully, there is no notion of braided monoidal natural transformation - any monoidal natural transformation will automatically respect the braiding. Hence, we can define two braided monoidal categories to be equivalent if there are braided monoidal functors between them which have compositions which are naturally isomorphic to the identity. Hence we can state a braided MacLane coherence theorem:

Theorem 4.2 (Braided MacLane coherence theorem). *Every braided monoidal equivalent is equivalent as a braided monoidal category to a strict braided monoidal category.*

Proof. .[WORK: I bet the proof is not that bad. Include it if so.] □

As we go through this text, we will define increasingly more structure on monoidal categories. We will be implicitly assuming theorems which assert that every structured monoidal categories is equivalent to a structure monoidal category whose underlying monoidal category is strict. Importantly, we will assume that this equivalence respects the relevant structure. We will not state these theorems as we go along the way, but they are true and necessary for our discussion. [WORK: Is this accurate? What is a good reference for this sort of coherence theorem? Will I talk about it more in the “Yoneda perspective” chapter?]

4.3.4 Pivotal monoidal categories

So far we have defined a language for putting particles together and braiding them. The next frontier is to introduce a language for creating and fusing particles/antiparticles. Categories with a mechanism for creating and fusing particles/antiparticles is known as a *pivotal monoidal category*.

In any realistic system, every particle will have a dual *antiparticle*. Particle/antiparticle pairs can always spontaneously be created from the vacuum. Often, particles/antiparticles can annihilate each other to go back to the vacuum. This process of annihilation is much more subtle however, because a particle/antiparticle pair could also fuse to make a particle which is *not* the vacuum. We delay the subtleties of fusion to our chapter on modular tensor categories. In a category with duals, every object $A \in \mathcal{C}$ will have a dual object which we denote $A^* \in \mathcal{C}$. For now, we introduce categories with maps for annihilation/creation which we call *evaluation* and *coevaluation* respectively:

Definition (Right-rigid monoidal category). A right-rigid monoidal category is the following data:

1. A monoidal category \mathcal{C} .
2. Objects A^* for all $A \in \mathcal{C}$.
3. Morphisms $\text{ev}_A : A \otimes A^* \rightarrow 1$, and $\text{coev}_A : 1 \rightarrow A^* \otimes A$ for all $A \in \mathcal{C}$.

Such that $(\text{ev}_A \otimes \text{id}_A) \circ (\text{id}_A \otimes \text{coev}_A) = \text{id}_A$ and $(\text{id}_{A^*} \otimes \text{ev}_A) \circ (\text{coev}_A \otimes \text{id}_{A^*}) = \text{id}_{A^*}$ for all $A \in \mathcal{C}$. \square

We implement right-rigid monoidal categories in string diagrams as follows. We denote evaluation and coevaluation as follows:

$$\begin{array}{ccc} \text{---} & & \text{---} \\ \text{---} & \curvearrowright & \text{---} \\ A & & A^* \end{array} = \text{ev}_A, \quad \begin{array}{ccc} A^* & & A \\ & \curvearrowleft & \\ A^* & & A \end{array} = \text{coev}_A.$$

The compatibility conditions are stated graphically as follows:

$$\begin{array}{ccc} \text{---} & & \text{---} \\ \text{---} & \curvearrowright & \text{---} \\ A & & A^* \\ \text{---} & \curvearrowright & \text{---} \\ A^* & & A \end{array} = \text{id}_A, \quad \begin{array}{ccc} A^* & & A \\ & \curvearrowleft & \\ A^* & & A \end{array} = \text{id}_{A^*}$$

We now note that particle/antiparticle pairs could also be created on the other side. This gives a left-rigid monoidal category, defined similarly:

Definition (Left-rigid monoidal category). A left-rigid monoidal category is the following data:

1. A monoidal category \mathcal{C} .
2. Objects A^* for all $A \in \mathcal{C}$.
3. Morphisms $\text{ev}_A : A^* \otimes A \rightarrow 1$, and $\text{coev}_A : 1 \rightarrow A \otimes A^*$ for all $A \in \mathcal{C}$.

Additionally, a rigid category is required to satisfy the property that $(\text{id}_A \otimes \text{ev}_A) \circ (\text{coev}_A \otimes \text{id}_A) = \text{id}_A$ and $(\text{ev}_A \otimes \text{id}_{A^*}) \circ (\text{id}_{A^*} \otimes \text{coev}_A) = \text{id}_{A^*}$ for all $A \in \mathcal{C}$. \square

In a left-rigid monoidal category, graphical cups and caps can be defined just like in right-rigid monoidal categories.

This leads us to our main definition of the section. We want to discuss categories which have a full theory of particles/antiparticles. This means that they should be able to create particle/antiparticle pairs on both sides, leading to a left-rigid and right-rigid structure on \mathcal{C} . As per usual, there should be some compatibility conditions between these two rigid structures. We give this full definition now:

Definition (Pivotal monoidal category). A pivotal monoidal category is the following data:

1. A monoidal category \mathcal{C} ;
2. A right-rigid structure $(\text{ev}^R, \text{coev}^R)$ on \mathcal{C} ;
3. A left-rigid structure $(\text{ev}^L, \text{coev}^L)$ on \mathcal{C} .

Such that:

1. The right-duals and left-duals of all objects are equal;
2. For all $A, B \in \mathcal{C}$, we have an equality of morphisms $B^* \otimes A^* \rightarrow (A \otimes B)^*$,

$$\begin{array}{ccc} (A \otimes B)^* & & \\ \text{coev}_{A \otimes B}^R & \xrightarrow{\quad \text{ev}_B^R \quad} & (A \otimes B)^* \\ B^* \otimes A^* & & \end{array} = \begin{array}{ccc} \text{ev}_B^L & & (A \otimes B)^* \\ \text{ev}_A^L & \xrightarrow{\quad \text{coev}_{A \otimes B}^L \quad} & B^* \otimes A^* \\ B^* & & \end{array}$$

3. For all $A, B \in \mathcal{C}$ and $f : A \rightarrow B$,

$$\begin{array}{ccc} A^* & & \\ \text{ev}_B^R & \xrightarrow{\quad f \quad} & \text{ev}_B^L \\ B^* & \text{coev}_A^R & B^* \\ & & \text{coev}_A^L \end{array} = \begin{array}{ccc} A^* & & \\ B & \xrightarrow{\quad f \quad} & B \\ \text{coev}_A^L & & B^* \end{array}$$

□

Now that we have given our main definitions, we prove some basic properties of rigid and pivotal categories.

The first thing to observe is that even though there is a lot of structure involved in the definition of a rigid monoidal category, most of it is in a real sense inessential. That is, we could have chosen different duals and the result would have been essentially the same:

Proposition 4.5. *Let \mathcal{C} be right (resp. left) rigid monoidal category. Let $A \in \mathcal{C}$ be an object, and let $(\tilde{A}^*, \tilde{\text{ev}}_A, \tilde{\text{coev}}_A)$ be another triple satisfying the axioms of rigidity. There is a unique morphism $i : A^* \rightarrow \tilde{A}^*$ making the diagram*

$$\begin{array}{ccc} & A^* \otimes A & \\ \text{coev}_A \nearrow & \downarrow \sim & \searrow \text{coev}_A \\ 1 & & \\ \text{co}\tilde{\text{ev}}_A \searrow & & \end{array}$$

$$A \otimes \tilde{A}^*$$

commute (resp. reverse order of tensor factors). This unique morphism is an isomorphism, and it is given by

$$i = \begin{array}{c} \tilde{A}^* \\ \text{co}\tilde{\text{ev}}_A \\ \text{ev}_A \\ A^* \end{array}$$

Remark. This proposition can be summarized by saying that *duals are unique up to unique isomorphism*.

Proof. By the computation

$$\begin{array}{ccc}
 \begin{array}{c} \tilde{A}^* \\ \boxed{i} \\ A^* \end{array} & = & \begin{array}{c} \tilde{A}^* \\ \boxed{i} \\ A^* \end{array} \\
 & & \text{---} \\
 & = & \begin{array}{c} \tilde{A}^* \\ \boxed{i} \\ A^* \end{array} = \begin{array}{c} \tilde{A}^* \\ \text{coev}_A \\ \tilde{A}^* \end{array} = \begin{array}{c} \tilde{A}^* \\ \text{coev}_A \text{ } \tilde{\text{ev}}_A \\ \tilde{A}^* \end{array} = \begin{array}{c} \tilde{A}^* \\ \tilde{\text{ev}}_A \\ A^* \end{array}
 \end{array}$$

we find that i is unique, and it has the desired formula. To prove that i is an isomorphism we observe that the map

$$\begin{array}{c} A^* \\ \text{coev}_A \\ \tilde{A}^* \end{array} = \begin{array}{c} \tilde{A}^* \\ \tilde{\text{ev}}_A \\ \tilde{A}^* \end{array}$$

serves as an inverse. This gives a proof of the result. \square

We now discuss the correct notion of functors between rigid and pivotal categories. Let $F : \mathcal{C} \rightarrow \mathcal{D}$ be a functor between pivotal categories. Given an object $A \in \mathcal{C}$, the evaluation and coevaluation maps naturally extend through the functor to endow $F(A^*)$ with the structure of a dual for A . Thus, by Proposition [ref], we have a canonical isomorphism $F(A^*) \cong F(A)^*$. This isomorphism exists without needing to add any extra conditions on F . In this way, the correct notion of functor between right/left rigid categories is just functor! There is, however, extra an compatibility condition needed for pivotal category. Both the left-rigid and right-rigid structures induce isomorphisms $F(A^*) \cong F(A)^*$. These induced isomorphisms should be the same. This is known as a pivotal functor.

Another important thing to know about rigid monoidal categories is that duality is *functorial*. That is, the duals of objects induce functors:

Proposition 4.6. *Let \mathcal{C} be right (resp. left) rigid monoidal category. Define a monoidal category $\overline{\mathcal{C}}$ as follows. The underlying category on $\overline{\mathcal{C}}$ is the opposite category for \mathcal{C} . The tensor product is given by $A \overline{\otimes} B = B \otimes A$, and the monoidal unit is $\mathbf{1} \in \mathcal{C}$. This gives a well-defined monoidal category.*

- (i) *The right (resp.left) rigid structure on \mathcal{C} induces a left (resp. right) rigid structure on $\overline{\mathcal{C}}$;*
- (ii) *Given any morphism $f : A \rightarrow B$ in \mathcal{C} , define*

$$f^* = \begin{array}{c} A^* \\ \curvearrowleft B \\ f \\ \curvearrowright A \\ B^* \end{array}$$

to be the dual for f (resp. same diagram using left rigidity). The assignment $A \mapsto A^*$, $f \mapsto f^*$ induces a functor from \mathcal{C} to $\overline{\mathcal{C}}$ which we denote $(-)^*$.

(iii) Given any objects $A, B \in \mathcal{C}$, define the map

$$\begin{array}{c} (A \otimes B)^* \\ \downarrow \\ \text{coev}_{A \otimes B} \\ \curvearrowleft \\ B^* A^* \end{array}$$

from $B^* \otimes A^*$ to $(A \otimes B)^*$ (resp. same diagram using left rigidity). These maps endows $(-)^*$ with the structure of a monoidal functor.

(iv) The functor $(-)^*$ is fully faithful. If \mathcal{C} is a pivotal category, then the functor above is an equivalence of monoidal categories between \mathcal{C} and $\overline{\mathcal{C}}$.

Remark. This proposition can be used to motivate the axioms of a pivotal category. Both the right and left rigid structures in a pivotal category induce functors $\mathcal{C} \rightarrow \overline{\mathcal{C}}$. The coherence condition is that these two functors should be equal.

Proof. We do only the proofs for right-rigid categories. The left-rigid proof is identical.

- (i) This follows immediately from the definitions;
- (ii) Functoriality is the condition that $(f \circ g)^* = g^* \circ f^*$. This follows from the following argument:

$$g^* \circ f^* = \begin{array}{c} A^* \\ \curvearrowleft B \\ f \\ \curvearrowright A \\ B^* \\ \downarrow \\ C^* \end{array} = \begin{array}{c} A^* \\ \curvearrowleft C \\ g \\ \curvearrowright f \\ A \\ \downarrow \\ C^* \end{array} = (f \circ g)^*$$

- (iii) This is an unlightening and straightforward computation;
- (iv) We first prove that $(-)^*$ is fully faithful. Given any objects $A, B \in \mathcal{C}$ and any morphism $g : B^* \rightarrow A^*$, the morphism
[WORK: add formula.]

has the property that $f^* = g$. Hence, $(-)^*$ is bijective on hom-spaces as desired.

We now show that $(-)^*$ is an equivalence of categories with \mathcal{C} is pivotal. By part (i), $\overline{\mathcal{C}}$ is a pivotal monoidal category. Hence duality once again induces a monoidal functor, this time $\overline{\mathcal{C}} \rightarrow \overline{\overline{\mathcal{C}}}$. Clearly, by our definition of $\overline{\mathcal{C}}$, $\overline{\overline{\mathcal{C}}} = \mathcal{C}$. Hence we have a pair of functors $\mathcal{C} \rightarrow \overline{\mathcal{C}}$ and $\overline{\mathcal{C}} \rightarrow \mathcal{C}$, each given by duality. Proving this proposition hence amounts to showing that the double dual map is monoidally naturally isomorphic to the identity.

To do this, we define a natural isomorphism explicitly by the isomorphisms $i : A \xrightarrow{\sim} A^{**}$

$$\begin{array}{ccc} A^{**} & & A^{**} \\ \boxed{i} & = & \text{ev}_A^R \curvearrowleft A^* \curvearrowright \text{coev}_{A^*}^L \\ A & & A \end{array}$$

for all $A \in \mathcal{C}$. To show that these morphisms induce a natural transformation, we observe that for all $f : A \rightarrow B$

$$\begin{array}{c} B^{**} \\ \boxed{f^{**}} \\ \boxed{i} \\ A \end{array} = \begin{array}{c} B^{**} \\ \boxed{B} \\ \boxed{f} \\ \boxed{A} \\ A^{**} \end{array} = \begin{array}{c} B^{**} \\ \boxed{B} \\ \boxed{f} \\ \boxed{A} \\ A^{**} \end{array} = \begin{array}{c} B^{**} \\ \boxed{B} \\ f \\ A \end{array} = \begin{array}{c} B^{**} \\ \boxed{i} \\ f \\ A \end{array}$$

The fact that \mathcal{C} is compatible with the tensor product is a straightforward computation, using the fact that computing the tensor product using right-rigidity and left-rigidity gives the same answer, and compatibility of \mathcal{C} with the unit is immediate. \square

As a key part of Proposition [ref], we showed that every pivotal structure on a right-rigid monoidal category induces a natural isomorphism between the identity functor and the double dual functor. This gives an alternate description of pivotal categories which is useful in some applications:

Corollary 4.2. *Let \mathcal{C} be a right-rigid monoidal category. Let $i : \text{id}_{\mathcal{C}} \xrightarrow{\sim} (-)^{**}$ be a monoidal natural isomorphism between the identity functor and the double dual functor. The maps*

$$\text{coev}_A^L \quad = \quad \begin{array}{c} A \ A^* \\ \text{coev}_{A^*}^R \end{array}, \quad \text{coev}_A^L \quad = \quad \begin{array}{c} \text{coev}_{A^*}^R \\ A \ A^* \end{array}$$

induce a pivotal structure on \mathcal{C} . Moreover, this assignment induces a bijection between pivotal structures on \mathcal{C} and monoidal natural isomorphisms $\text{id}_{\mathcal{C}} \xrightarrow{\sim} (\text{id}_{\mathcal{C}})^{**}$.

Proof. Proving that the maps provided satisfy the axioms of a left-rigid structure is immediate. Proving that they satisfy the axioms of a pivotal structure comes from running the arguments in the proof of proposition [ref] in reverse. The operations of inducing a monoidal natural isomorphism from a pivotal structure and inducing a pivotal structure from a monoidal natural isomorphism are inverses of one another. Hence, they induce a bijection between the two types of structures as desired. \square

History and further reading:

Category theory was first introduced and formalized by Saunders Mac Lane and Samuel Eilenberg in 1945 [EM45]. Of course, the ideas underlying category theory were present earlier and can be traced back arbitrarily far. In the subsequent decades the formalism of category theory spread far and wide, bringing with it the discovery of many deep theorems. The first major explicit appearance of category theory in physics was Vladimir Drinfeld's work on so-called *quantum groups* in the early 1980s [Dri86]. Quantum groups are certain kinds of mathematical objects rightly related to content in this book. They were introduced as tools to help generate exactly-solvable models in condensed matter physics. Very quickly quantum groups were absorbed into the theory of the ideas of string theory of topological quantum field theory, which were both new at the time [BPZ84, Wit88]. The physics in this area has since become and remained extremely categorical in nature [Lur08, BDSPV15].

There are many excellent introductory texts to category theory. Some authors find it fruitful to reformulate all of quantum mechanics, and especially quantum information, in terms of category theory. A good source outlining this approach and introducing category theory through it is Coecke-Kissinger's textbook [CK18]. The Kong-Zhang textbook [KZ22] gives an introduction to category theory in the context of topological order. A good general-purpose textbook on category theory is Fong-Spivak [FS19], and a classical but slightly dated reference is [ML13].

Exercises:

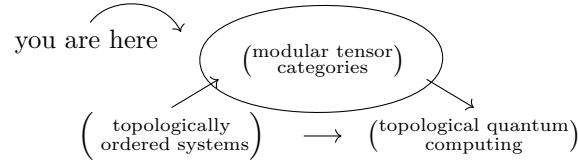
- 4.1. .[WORK: If \mathcal{C} is a category with products, then the product forms a monoidal structure (with a good **1** given of course), and same for coproducts.]
- 4.2. .[WORK: Show that endofunctors form a *strict* monoidal category.]
- 4.3. .[WORK: Add an exercise giving some compatibility conditions between monoidal/rigid structures and direct sums. Namely, they distribute nicely.]

5 Modular categories

5.1 Overview

5.1.1 Introduction

In this chapter we will be giving a detailed analysis of modular categories, the abstract algebraic structures used to describe anyons in topological order. We recall below how this fits into the general framework of this book:



Describing exactly what an anyon and how it can transform in terms of states and unitary operators on a Hilbert space can be difficult. However, describing abstractly how these transformations compose with one another can be done relatively simply. Hence we take a composition-first category-theoretic approach to anyons. We will make heavy use of the diagrammatic language of braided monoidal categories established in Chapter [ref]. Concretely, we will think of a modular category as being the category with the following data:

$$\begin{pmatrix} \text{objects: finite collections of anyons} \\ \text{morphisms: motions/behaviors of anyons} \end{pmatrix}$$

[WORK: is the above picture correct? It's subtle and I should think through it.]

[WORK: I am using nonstandard language. I am using the term “modular category” to refer to what most authors call “modular tensor categories”. I do this because modular category is shorter than modular tensor category and I want to be succinct. For many people (notably Zhenghan) modular category means potentially non-semisimple modular category. For instance, Turaev’s original definition did not include the semisimplicity axiom and he called them modular categories. The term tensor is used to highlight semisimplicity by many authors, so I can highlight semisimplicity with a footnote.]

Up to topological equivalence, there are not that many things that a collection of anyons can do. The most basic thing is to move anyons around each other - this is known as braiding. If the anyons touch each other then they can congeal into a single composite anyon - this is known as fusion. Even if there are no anyons in a system, however, there is always something possible. Anyons can be spontaneously created, so long as every anyon which is created comes along with its corresponding antiparticle. This is known as *pair-creation*. These three operations are the fundamental structures which we will build into modular categories:

1. braiding;
2. fusion;
3. pair-creation.

One potentially useful way of thinking about modular categories comes from analogy with classical physics. We saw in Chapter [ref] that topological classical systems have an algebraic description in terms of finite groups. Namely, quasiparticles in the system of ordered media with order space M is algebraically characterized by the fundamental group $\pi_1(M, m)$ of M relative to some basepoint $m \in M$. Seeing as topological order is a vast quantum generalization of classical ordered media, we can think of modular categories as being a vast quantum generalization of finite groups. Every finite group induces a modular category, by first constructing the Kitaev quantum double model based on that finite group and then describing its anyons. Most modular categories, however, lie beyond this description.

[WORK: There's an issue in this treatment. There is one piece of data beyond the scope of modular categories - the chiral central charge. This is a remnant of the fact that the bulk-to-boundary correspondence is not exact because the boundary can have stacked E_8 phases, see [Bon21]. Of course this is not something to dive into now. However, I want to be maximally honest - point out that there is a unique piece of topologically invariant information beyond the scope of modular categories. Maybe include somewhere (as an exercise?) the treatment of chiral central charge mod 8? I think this would make for a good footnote. I'm realizing now that the fact that the chiral central charge is a root of unity is part of Vafa's theorem. So, it makes the most sense for this to be in the number theory section.]

5.1.2 Using the final product

Before developing the theory of modular categories, it is good to get a feel for what using the final product is like. A modular category itself will be a big infinite thing, with infinitely many objects and infinitely many morphisms between those objects. However, all modular categories are in a real sense *finitely generated*. What we mean by this is that plugging in a finite number of objects and morphisms, the rest of the objects and morphisms can be recovered by the abstract rules encoded in the formalism. For example, consider the 3-strand braid group B_3 . This group has infinitely many elements and the group operation $\cdot : B_3 \times B_3 \rightarrow B_3$ takes a-priori an infinite amount of data to describe. However, the presentation

$$B_3 = \langle \sigma_1, \sigma_2 \mid \sigma_1\sigma_2\sigma_1 = \sigma_2\sigma_1\sigma_2 \rangle$$

gives completely finite description of B_3 . It is important to note, however, that this presentation would *not* have been enough to recover B_3 if we had just been told that B_3 is a monoid. The fact that B_3 is a group implied the existence of elements σ_1^{-1} , σ_2^{-1} , and defined how they interacted with σ_1, σ_2 . We see in this way that the axioms of a group not only serve as a restriction on what mathematical objects are allowed to be groups, but they also serve as a compression technique. They give the rules by which a minimal collection of data can be used to generate the rest.

In a similar way, the axioms of modular categories are not only necessary by the fact that they restrict which categories can be modular categories, but they are also vital in the fact that they allow us to generate a full description of anyons from a minimal collection of data. For practically-minded readers, this can be viewed as the main motivation for defining modular categories at all, instead of just working with important examples.

The final challenge in going from modular categories to their description in terms of a finite set of data is in coming up with an efficient standard way of describing morphisms in a modular category. This is done using skeletonization, as discussed in section [ref].

In the end, the data of a modular category will look like what we have below for the toric code:

[WORK: add toric code modular category data]

Or, for a more complicated example, we can consider the data for $G = S_3$:

[WORK: add $G = S_3$ modular category data]

A large table of these descriptions are found in Chapter [ref]. We now give a worked example of how this data is used to compute observable quantities.

[WORK: add good example, computing some probability of annihilation]

5.2 First properties

5.2.1 Definition

In this section we finally define modular categories, which are the main mathematical content of this book. Seeing as lots of data is involved, we spread out the definition over a series of steps as to not overload the senses. These intermediate definitions are also important in their own right, because they will be used in other places in the algebraic theory of topological phases.

Definition (Fusion category). A fusion category is the following data:

1. A category \mathcal{C} ;
2. The structure of a right-rigid monoidal category on \mathcal{C} ;
3. The structure of a \mathbb{C} -linear category on \mathcal{C} .

Such that:

1. The tensor product functor $\otimes : \mathcal{C} \times \mathcal{C} \rightarrow \mathcal{C}$ induces bilinear maps hom-spaces;
2. There is an equivalence $\mathcal{C} \cong \mathbf{Vec}_{\mathbb{C}}^n$ as \mathbb{C} -linear categories;
3. $\text{End}_{\mathcal{C}}(\mathbf{1}) \cong \mathbb{C}$ as \mathbb{C} -vector spaces

□

A fusion category is part of the way towards having all of the requisite structures of a modular category: it has a method for fusion inherited from the tensor product, and it has half of a method for pair-creation coming from right-rigidity. The \mathbb{C} -linearity allows us to think of hom-spaces as vector spaces, which allows us to treat hom-spaces as quantum systems. The condition (1) is a compatibility between the \mathbb{C} -linear structure and the monoidal structure. The conditions (2)-(3) are strong niceness and finiteness conditions - we will explain them in detail later. We now move one step closer to our definition of modular category:

Definition (Spherical fusion category). A spherical fusion category is the following data:

1. A fusion category \mathcal{C} ;
2. A left-rigid structure on \mathcal{C} .

Such that:

1. The left-rigid and right-rigid structures on \mathcal{C} satisfy the axioms of a pivotal structure on \mathcal{C} ;
2. For every object $A \in \mathcal{C}$ and for every morphism $f : A \rightarrow A$, we have

$$\begin{array}{c} B \\ \text{---} \\ f \\ \text{---} \\ A \end{array} = \begin{array}{c} B \\ \text{---} \\ f \\ \text{---} \\ A \end{array}$$

□

A spherical fusion category now has a structure for fusion, and a full structure for pair-creation. The 2nd condition is known as the *spherical axiom*. We will explain this axiom in more detail later.

Adding on a braiding, we can get all of the structures of a modular category. However, adding this structure still misses one key structure of being a modular category. Hence, we call it *pre-modular*:

Definition (Pre-modular category). A pre-modular category is the following data:

1. A spherical fusion category \mathcal{C} ;
2. A braided structure on \mathcal{C} .

No extra compatibility conditions are required.

□

This definition now has all of the structure we wanted it to have: fusion, pair-creation, and braiding. The final axiom is a non-degeneracy condition. It is subtle in its interpretation, and we will explain it several different ways throughout this chapter:

Definition (Modular category). A modular category is a pre-modular category satisfying the following condition. Let $A \in \mathcal{C}$ be an object. If

$$\begin{array}{c} A & B \\ \text{---} \\ \text{---} \\ A & B \end{array} = \begin{array}{c} A & B \\ | & | \\ A & B \end{array}$$

for all $B \in \mathcal{C}$, then $A \cong \mathbf{1}$.

□

5.2.2 Anyons in modular categories

Modular categories are supposed to be theories of anyons in topological order. So, now that we have the definition of modular category, it is natural to ask: what do anyons mathematically correspond to, in modular categories? The answer lies within the condition in a fusion category \mathcal{C} that there is an equivalence $\mathcal{C} \cong \mathbf{Vec}_{\mathbb{C}}^n$ as \mathbb{C} -linear categories. We explore the importance of this condition.

Suppose we are given an object $V = (V_1, V_2, \dots, V_n) \in \mathbf{Vec}_{\mathbb{C}}^n$. For all $1 \leq i \leq n$, let $C_i \in \mathbf{Vec}_{\mathbb{C}}^n$ denote the object which has dimension zero in every index $j \neq i$ and is equal to \mathbb{C} in index i . We observe the isomorphism

$$\begin{aligned}
V &\cong \bigoplus_{i=1}^n (0 \dots V_i \dots 0) \\
&\cong \bigoplus_{i=1}^n (0 \dots \mathbb{C}^{\dim(V_i)} \dots 0) \\
&\cong \bigoplus_{i=1}^n \dim(V_i) \cdot \mathbb{C}_i
\end{aligned}$$

where $\dim(V_i) \cdot (\mathbb{C}_i) = \mathbb{C}_i \oplus \mathbb{C}_i \dots \oplus \mathbb{C}_i$, $\dim(V_i)$ many times. This computation shows that any object in $\mathbf{Vec}_{\mathbb{C}}^n$ can be decomposed into irreducible components \mathbb{C}_i . These objects \mathbb{C}_i are in a real sense the building blocks of $\mathbf{Vec}_{\mathbb{C}}^n$. They will correspond physically to anyons. More concretely, we make the following definition:

Definition. A *simple object* A in a fusion category \mathcal{C} is an object which has no direct sum decomposition into smaller objects. That is, $A \not\cong B \oplus C$ for any non-zero objects $B, C \in \mathcal{C}$ where \oplus denotes the biproduct in \mathcal{C} .

Our physics-math dictionary is that anyon types correspond to isomorphism classes of simple objects.

We now state the basic proposition which ensures that the necessary properties from $\mathbf{Vec}_{\mathbb{C}}^n$ follow through the equivalence of categories.

Proposition 5.1. *Let \mathcal{C} be a fusion category. The biproduct of any two elements in \mathcal{C} exists, and \mathcal{C} has a zero object. Let \mathcal{L} denote the set of isomorphism classes of simple objects in \mathcal{C} . The set \mathcal{L} is finite. Choose an object $X \in \mathcal{C}$. There exist unique nonnegative integers $c_{[A]}$, $[A] \in \mathcal{L}$ such that*

$$X \cong \bigoplus_{[A] \in \mathcal{L}} N_{[A]} \cdot A.$$

Proof. We first show that \mathcal{C} has biproducts. Let $F : \mathcal{C} \rightarrow \mathbf{Vec}_{\mathbb{C}}^n$, $G : \mathbf{Vec}_{\mathbb{C}}^n \rightarrow \mathcal{C}$ be a pair of functors which induces an equivalence of categories, for some $n \geq 1$. Let $A, B \in \mathcal{C}$ be objects. Since G and F are fully faithful, the universal property of the direct sum $F(A) \oplus F(B)$ guarantees that $G(F(A) \oplus F(B))$ will be a direct sum of $G(F(A))$ and $G(F(B))$. Since $G(F(A)) \cong A$ and $G(F(B)) \cong B$, we conclude that \mathcal{C} has biproducts. The object $G(0)$ is a zero object for \mathcal{C} .

We now prove that \mathcal{C} has finitely many isomorphism classes of simple objects. It is clear that an object $A \in \mathcal{C}$ is simple if and only if $F(A) \in \mathbf{Vec}_{\mathbb{C}}^n$ is simple. Thus, since G serves as an inverse, F establishes a bijection between isomorphism classes of simple objects in \mathcal{C} and isomorphism classes of simple objects in $\mathbf{Vec}_{\mathbb{C}}^n$. Every simple object in $\mathbf{Vec}_{\mathbb{C}}^n$ will be isomorphic to \mathbb{C}_i for some $1 \leq i \leq n$. Hence, there are n simple objects in $\mathbf{Vec}_{\mathbb{C}}^n$. Hence, there are n simple objects in \mathcal{C} , which is finite.

The unique direct sum decomposition is clearly true in $\mathbf{Vec}_{\mathbb{C}}^n$. It is immediate that it passes to a unique direct sum decomposition in \mathcal{C} . \square

The set of simple objects has an alternative description, known as Schur's lemma:

Proposition 5.2 (Schur's Lemma). *Let \mathcal{C} be a fusion category. An object $A \in \mathcal{C}$ is simple if and only if its endomorphism ring $\text{End}(A)$ is one-dimensional. Additionally, if $A, B \in \mathcal{C}$ are nonisomorphic simple objects then $\text{Hom}(A, B) = 0$.*

Proof. Let $F : \mathcal{C} \rightarrow \mathbf{Vec}_{\mathbb{C}}^n$, $G : \mathbf{Vec}_{\mathbb{C}}^n \rightarrow \mathcal{C}$ be a pair of \mathbb{C} -linear functors which establishes an equivalence between \mathcal{C} and $\mathbf{Vec}_{\mathbb{C}}^n$ as \mathbb{C} -linear categories. The simple objects in $\mathbf{Vec}_{\mathbb{C}}^n$ are all isomorphic to \mathbb{C}_i for some $1 \leq i \leq n$. We compute that

$$\dim \left(\text{Hom}_{\mathbf{Vec}_{\mathbb{C}}^n} \left(\bigoplus_{i=1}^n n_i \mathbb{C}_i, \bigoplus_{i=1}^n m_i \mathbb{C}_i \right) \right) = \sum_{i=1}^n n_i m_i.$$

As a corollary of this formula, we find that if $A = \bigoplus_{i=1}^n n_i \mathbb{C}_i$ then $\dim(\text{End}_{\mathbf{Vec}_{\mathbb{C}}^n}(A)) = \sum_{i=1}^n n_i^2$. Clearly, this dimension is equal to one if $A = \mathbb{C}_i$ for some $1 \leq i \leq n$, and is greater than one otherwise. As a second corollary, we compute that $\text{Hom}(\mathbb{C}_i, \mathbb{C}_j) = 0$ whenever $i \neq j$.

The functor G induces a bijection between isomorphism classes of simple objects in $\mathbf{Vec}_{\mathbb{C}}^n$ and isomorphism classes of simple objects in \mathcal{C} , and it induces vector space isomorphisms on hom spaces. This means that the results for $\mathbf{Vec}_{\mathbb{C}}^n$ translate to the desired result on \mathcal{C} . \square

As an immediate application of Schur's lemma, we observe that the monoidal unit $\mathbf{1}$ is a simple object in every fusion category. By our physics-math dictionary, this means that $\mathbf{1}$ corresponds to an anyon type. This type is the *vacuum* type - empty space. The anyon $\mathbf{1}$ is the trivial no-anyon type.

Another application of Schur's lemma is to make a first verification that simple objects are a good choice of mathematical characterization of anyons. If A, B are distinct anyon types, then there should not be any physical process which goes from one to another. There is no physical mechanism for locally turning one anyon type into another. This is captured by the formula $\text{Hom}(A, B) = 0$. Similarly, given an anyon A , there is no nontrivial action that can be locally performed on A . This comes from the fact that information is topologically protected, and thus cannot be changed by acting on a single particle - topological information processing requires global braiding between multiple particles. This is encoded in the fact that $\text{Hom}(A, A) \cong \mathbb{C}$ is one dimensional and hence consists only of trivial phase gates.

Expanding our physics-math dictionary, we say that for every anyon A its *antiparticle* is the dual A^* which comes from right-rigidity. This gives a valid anyon type by the following computation:

Proposition 5.3. *Let \mathcal{C} be a fusion category. If $A \in \mathcal{C}$ is a simple object, then so is A^* .*

Proof. By Proposition [ref] duality induces a bijection on hom-spaces. Since composition is bilinear, this bijection is an isomorphism of vector spaces. Hence, for all $A \in \mathcal{C}$ there is an isomorphism $\text{Hom}(A, A) \cong \text{Hom}(A^*, A^*)$. Hence, $\dim \text{Hom}(A, A) = 1$ if and only if $\dim \text{Hom}(A^*, A^*) = 1$ so the result follows by Schur's lemma. \square

An important part of understanding simple objects in modular categories is making sense of the direct sum decompositions coming from Proposition [ref]. Let \mathcal{C} be a fusion category with simple objects $A, B \in \mathcal{C}$. Consider the decomposition

$$A \otimes B \cong \bigoplus_{[C] \in \mathcal{L}} N_C^{A,B} \cdot C$$

where $N_C^{A,B} \geq 0$ are nonnegative integers, and \mathcal{L} is the set of isomorphism classes of simple objects. The integers $\{N_C^{A,B}\}_{[C] \in \mathcal{L}}$ are known as fusion coefficients, because they specify the behavior of A and B when they fuse.

The tensor product \otimes physically corresponds to joining anyons, forming a composite anyon configuration. The object $A \otimes B$ corresponds to the configuration with one A -type anyon and one B -type anyon. The direct sum decomposition is physically interpreted as saying that when $A \otimes B$ are fused, the possible results of that fusion are all of the anyon types $[C] \in \mathcal{L}$ for which $N_C^{A,B} \neq 0$.

[WORK: add nontrivial example from Kitaev quantum double model]

A more detailed understanding of the physical meaning of the direct sum will have to wait for later.

This concludes our basic picture of anyons in fusion categories.

5.2.3 States in modular categories and unitarity

[WORK: This whole section only discusses *pure* states. For a good discussion of mixed states and entropy, see [BKP17]]

It is now worth reflecting on what exactly states correspond to in modular categories. In particular, objects in modular categories are *not* quantum systems. They don't have vector space structure. The spaces with vector space structure are the hom-spaces, by \mathbb{C} -linearity. Objects will correspond to anyon configurations. States will correspond to normalized vectors in certain hom-spaces. In particular:

$$\left(\begin{array}{l} \text{states of topological order } \mathcal{C} \\ \text{on the sphere} \\ \text{with anyon configuration } A_1, A_2 \dots A_n \end{array} \right) = \left(\begin{array}{l} \text{normalized vectors in the Hilbert space} \\ \text{Hom}_{\mathcal{C}}(\mathbf{1}, A_1 \otimes A_2 \dots \otimes A_n) \end{array} \right)$$

where by "anyon configuration $A_1, A_2 \dots A_n$ " we mean that the state has anyons present in n sites, arranged left to right on a one dimensional subspace of \mathbb{R}^2 , with corresponding anyon type $A_1, A_2 \dots A_n$. For the sake of concreteness, one can imagine that at the point $(i, 0) \in \mathbb{R}^2$ the state has an anyon of type A_i .

The remainder of this subsection is a series of loosely-related observations about this choice of state space:

Observation 1: *The physical space is a sphere.*

It is not immediately clear where in the formula $\text{Hom}(\mathbf{1}, A_1 \otimes A_2 \dots \otimes A_n)$ we chose the sphere as the physical space. To make it make sense, we observe as a special case of the general formula that

$$\dim \left(\begin{array}{l} \text{Hilbert space of topological order } \mathcal{C} \\ \text{on the sphere } S^2 \\ \text{with exactly one anyon of type } A \end{array} \right) = \dim \text{Hom}(\mathbf{1}, A) = \begin{cases} 1 & A = \mathbf{1} \\ 0 & \text{otherwise.} \end{cases}$$

Hence, the state formula tells us that if the sphere has exactly one anyon on it then that anyon type must be trivial. Moreover, there is unique state on the sphere with no anyons.

This is consistent with our general principles about topological order on the sphere.

[WORK: sketch nice argument for why there is a unique ground state on the sphere. What I'm struggling with here is why an anyon type in a region must necessarily be detectable by its surrounding region.]

This gives us a framework for extending our state-space/hom-space correspondances to physical spaces other than the sphere. In particular, we find the following:

$$\left(\begin{array}{l} \text{states of topological order } \mathcal{C} \\ \text{on the infinite flat plane } \mathbb{R}^2 \\ \text{with anyon configuration } A_1, A_2 \dots A_n \end{array} \right) = \left(\begin{array}{l} \text{normalized vectors in the Hilbert space} \\ \text{Hom}_{\mathcal{C}}\left(\bigoplus_{[B] \in \mathcal{L}} B, A_1 \otimes A_2 \dots \otimes A_n\right) \end{array} \right)$$

where \mathcal{L} is the set of isomorphism classes of simple objects in \mathcal{C} . Replacing $\mathbf{1}$ with $\bigoplus_{[B] \in \mathcal{L}} B$ reflects the differences between the sphere and the infinite flat plane.

[WORK: sketch nice argument for why states on infinite flat plane are determined by their overall charge. The subtlety here is exactly the same as the one for the sphere. Think about it then put it down.]

[WORK: What happens for higher genus surfaces? I should add a few words about them. Zhenghan says all of this is contained in Turaev's book about the Reshetikhin-Turaev construction.]

Observation 2: *The anyons are always assumed to be arranged in a line.*

The anyon configurations are always assumed to be linear. The main reason to do this is because it makes the mathematics much simpler. If we kept track of the positions of each of the anyons in two dimensional space it would add more pieces of data and structures to keep track of. Seeing as every anyon configuration can be pushed onto a one-dimensional space, only working with a one-dimensional configuration does not affect the generality of the answers and hence it is very much preferred.

Observation 3: *The formula $\text{Hom}_{\mathcal{C}}(\mathbf{1}, A_1 \otimes A_2 \dots \otimes A_n)$ encodes the fact that states can be specified by their history.*

A good first question to ask when seeing the Hilbert space $\text{Hom}_{\mathcal{C}}(\mathbf{1}, A_1 \otimes A_2 \dots \otimes A_n)$ is *why* this should describe a state with anyon configuration $A_1 \dots A_n$. The answer is that states can be described their history. [WORK: give good example of making a state by specifying its history; argue why it has to be this way in general].

Observation 4: *In the definition of an modular category hom-spaces are vector spaces and not Hilbert spaces, so this choice of physics-math correspondance is incorrect as literally written.*

To make this definition work, all of the hom-spaces of the modular category \mathcal{C} should be equipped with Hilbert space structures. Furthermore, the natural operators we wish to perform like braiding should all be unitary with respect to these inner products. This amounts to adding a large number of compatibility conditions on the Hilbert space structures. A modular category with this choice of structure is known as a *unitary* modular category. We give the formal definition below:

Definition (Unitary fusion category). A unitary fusion category is the following data:

1. An fusion category \mathcal{C} .
2. (Conjugation) A linear map $\dagger : \text{Hom}(A, B) \rightarrow \text{Hom}(B, A)$ for all $A, B \in \mathcal{C}$.

Such that:

1. (Unitarity) Given $f : A \rightarrow A$ an endomorphism of $A \in \mathcal{C}$, define

$$\text{tr}(f) = \text{ev}_A \circ (\text{id}_{A^*} \otimes f) \circ (\text{ev}_A)^\dagger.$$

The map $\langle \cdot | \cdot \rangle : \text{Hom}(A, B) \times \text{Hom}(A, B) \rightarrow \mathbb{C}$ defined by $\langle f | g \rangle = \text{tr}(f^\dagger \circ g)$ is an inner product, endowing $\text{Hom}(A, B)$ with the structure of a Hilbert space.

2. $(f^\dagger)^\dagger = f$ for all $f \in \text{Hom}(A, B)$, $A, B \in \mathcal{C}$.

3. $(f \circ g)^\dagger = g^\dagger \circ f^\dagger$ for all $f \in \text{Hom}(B, C), g \in \text{Hom}(A, B), A, B, C \in \mathcal{C}$.
4. $(f \otimes g)^\dagger = f^\dagger \otimes g^\dagger$ for all $f \in \text{Hom}(A, B), g \in \text{Hom}(C, D), A, B, C, D \in \mathcal{C}$.
5. $(\text{coev}_A)^\dagger \circ (f \otimes \text{id}_{A^*}) \circ \text{coev}_A = \text{tr}(f)$ for all $A \in \mathcal{C}$

□

Unitary fusion categories make for a pleasant object of study because the distinguished maps $(\text{ev}_A)^\dagger : 1 \rightarrow A^* \otimes A$ and $(\text{coev}_A)^\dagger : A \otimes A^* \rightarrow 1$ induce a pivotal structure:

Proposition 5.4. *Let \mathcal{C} be a unitary fusion category. The maps $\text{ev}_A^L = (\text{coev}_A)^\dagger$ and $\text{coev}_A^L = (\text{ev}_A)^\dagger$ give a left-rigid structure on \mathcal{C} . This left-rigidity endows \mathcal{C} with the structure of a spherical fusion category.*

Proof. [WORK: This is actually subtler than I expected. Either give a proof here, or postpone it to the unitarity section.] □

We can now define unitary modular categories. The compatibility conditions for the twist are chosen so that the definition of trace as a modular category and the definition of trace as a unitary fusion category coincide.

Definition (Unitary pre-modular category). A unitary modular category is the following data:

1. A modular category \mathcal{C} ;
2. (Conjugation) A linear map $\dagger : \text{Hom}(A, B) \rightarrow \text{Hom}(B, A)$ for all $A, B \in \mathcal{C}$.

Such that:

1. Forgetting the left-rigid structure and braiding, (\mathcal{C}, \dagger) forms a unitary fusion category.
2. $(\text{ev}_A^R)^\dagger = \text{coev}_A^L$;
3. $(\text{coev}_A^R)^\dagger = \text{ev}_A^L$;
4. $(\beta_{A,B})^\dagger = \beta_{B,A}^{-1}$.

□

Definition (Unitary modular category). A unitary modular category is a unitary pre-modular category which satisfies the non-degeneracy axiom of a modular category.

For this reason, the correct algebraic structure to underlie the theory of topological order is not a modular category, but a unitary modular category. We have chosen to not emphasize this before because the difference between unitary modular categories and non-unitary modular categories is very small. [WORK: talk about uniqueness + positive q.d. criterion this will make more sense once we write the actual section about unitarity. A good thing to emphasize is that unitary modular categories don't let you use less data in your definition, and you can still do essentially everything you want to do. It's just way more cumbersome. They're all equivalent but you still have to choose, c.f. the fact that the category of vector spaces and Hilbert spaces with linear maps as morphisms are equivalent].

5.2.4 Topological charge measurement

When two anyons are fused together, they will form a superposition of other anyon types. Measuring the result of the fusion will collapse the answer into a specific anyon type. The outcome of this measurement is an observable quantity, which allows for the measurement of topological quantum information. In many cases this is the *only* local observable quantity. We give the formalism behind computing these probabilities now.

[WORK: do this right - I don't know it well but it shouldn't be hard to learn. Don't introduce anything too general, like trace or whatnot. Just quantum dimension, which should already have been introduced in previous chapter.]

[WORK: The correct reference for this subsection is [Bon21]. The paper [CCW17] claims to introduce the term topological charge measurement and gives a nice formal treatment. Clarifying the situation seems important.]

5.3 The modular category toolkit

In this section, we will introduce and prove the basic facts about the most important structures in the theory of modular categories. These facts and structures are the tools used for solving problems about the algebraic theory of anyons.

[WORK: I don't have a section on fusion coefficients yet. I guess this isn't a problem, because there isn't that much to say. I would like to have the associativity of fusion coefficients and the fact that braiding \Rightarrow commutative said somewhere explicitly, though. Find a place?]

5.3.1 Trace

The first structure to define in the theory of modular categories is the *trace*. Let \mathcal{C} be a spherical fusion category. Given any object $A \in \mathcal{C}$ and any endomorphism $f : A \rightarrow A$, we define the *trace of f* by the following formula:

$$\text{tr}(f) = A^* \begin{array}{c} A \\ \square \\ A \end{array}$$

Initially, the trace is a morphism, $\text{tr}(f) : \mathbf{1} \rightarrow \mathbf{1}$. However, we will choose to think of the trace of a morphism as a *complex number*, $\text{tr}(f) \in \mathbb{C}$. This can be done because the definition of a fusion category $\text{End}(\mathbf{1}) \cong \mathbb{C}$. This isomorphism can be made canonical by identifying an endomorphism $g \in \text{End}(\mathbf{1})$ with the unique $\lambda \in \mathbb{C}$ such that $g = \lambda \cdot \text{id}_{\mathbf{1}}$.

The trace is used mainly as a tool for linearization. Morphisms and objects are hard to describe, but the trace is a complex number.

Proposition 5.5. *Let \mathcal{C} be a spherical fusion category. For all $A, B \in \mathcal{C}$, $f \in \text{End}(A)$ the following claims are all true:*

1. $\text{tr} : \text{End}(A) \rightarrow \mathbb{C}$ is a linear map of vector spaces,
2. $\text{tr}(f^*) = \overline{\text{tr}(f)}$,
3. $\text{tr}(f \oplus g) = \text{tr}(f) + \text{tr}(g)$ for all $g \in \text{End}(B)$,

4. $\text{tr}(f \otimes g) = \text{tr}(f) \cdot \text{tr}(g)$ for all $g \in \text{End}(B)$,
5. $\text{tr}(h \circ g) = \text{tr}(g \circ h)$ for all $g : A \rightarrow B$, $h : B \rightarrow A$.
6. Trace is preserved by functors. That is, let \mathcal{C}, \mathcal{D} be spherical categories with traces $\text{tr}_{\mathcal{C}}, \text{tr}_{\mathcal{D}}$ respectively. Let $F : \mathcal{C} \rightarrow \mathcal{D}$ be a pivotal functor. We have that $\text{tr}_{\mathcal{C}}(f) = \text{tr}_{\mathcal{D}}(F(f))$;

Proof. We prove the claims one by one.

1. This follows immediately from the bilinearity of composition.
2. This is a straightforward computation.
3. [WORK: do proof. This uses facts about the direct sum we haven't established yet.]
4. Using Proposition [ref] we compute

$$\begin{aligned} \text{tr}(f \otimes g) &= \begin{array}{c} A \otimes B \\ \text{---} \\ \boxed{f \otimes g} \\ \text{---} \\ A \otimes B \end{array} = \begin{array}{c} A \quad B \\ \text{---} \\ \boxed{f} \quad \boxed{g} \\ \text{---} \\ A \quad B \end{array} \\ &= \begin{array}{c} A \quad B \\ \text{---} \\ \boxed{f} \quad \boxed{g} \\ \text{---} \\ A \quad B \end{array} = \text{tr}(f) \cdot \text{tr}(g). \end{aligned}$$

5. Using Proposition [ref] we find that

$$\text{tr}(f \circ g) = \begin{array}{c} A \\ \text{---} \\ \boxed{g} \\ \text{---} \\ \boxed{f} \\ \text{---} \\ A \end{array} = \begin{array}{c} B \\ \text{---} \\ f^* \\ \text{---} \\ g \\ \text{---} \\ B \end{array} = \begin{array}{c} A \\ \text{---} \\ \boxed{f} \\ \text{---} \\ \boxed{g} \\ \text{---} \\ A \end{array} = \text{tr}(g \circ f).$$

6. [WORK: do the proof. It's not very hard, but involves a diagram and uses pivotality of the functor.]

This completes the proof. \square

With these properties in hand, we can explicitly compute the trace using a straightforward procedure:

Corollary 5.1. Let $f : A \rightarrow A$ be an endomorphism in a fusion category \mathcal{C} . Fix a decomposition $A \cong \bigoplus_{i \in I} A_i$ of A into simple objects A_i . Moreover, we take the decomposition such that if $A_i \cong A_j$ then $A_i = A_j$. We can decompose

$$\text{Hom}(A, A) \cong \text{Hom}\left(\bigoplus_{i \in I} A_i, \bigoplus_{i \in I} A_i\right) = \bigoplus_{i \in I, j \in I} \text{Hom}(A_i, A_j).$$

Let M be the matrix whose columns and rows are labeled by I , and whose (i, j) entry is 0 if $A_i \not\cong A_j$ and $\lambda \cdot d_{A_i}$ if $A_i = A_j$, where $\lambda \in \mathbb{C}$ is the unique value such that the $\text{Hom}(A_i, A_j)$ component of f is $\lambda \cdot \text{id}_{A_i}$. We have that

$$\text{tr}_{\mathcal{C}}(f) = \text{tr}_{\mathbf{Vec}}(M).$$

Proof. Suppose that $A = A_0 \oplus A_1$ is the direct sum of two objects, not necessarily simple. By proposition [ref] we have a canonical decomposition

$$(A_0 \oplus A_1) \otimes (A_0 \oplus A_1)^* \cong (A_0 \otimes A_0^*) \oplus (A_0 \otimes A_1^*) \oplus (A_1 \otimes A_0^*) \oplus (A_1 \otimes A_1^*).$$

Suppose that $h : A \rightarrow A$ is an endomorphism. We can decompose $h = h_{A_0, A_0} + h_{A_0, A_1} + h_{A_1, A_0} + h_{A_1, A_1}$ as a sum of morphisms which restrict to maps $A_i \rightarrow A_j$. We find that $\text{coev}_{A \oplus B}$ restricts to a map whose codomain is $(A_0 \otimes A_0^*) \oplus (A_1 \otimes A_0)^*$ and similarly ev restricts to a map whose domain is $(A_0 \otimes A_0^*) \oplus (A_1 \otimes A_0)^*$, since $\text{coev}_{A \oplus B} = \text{coev}_A \oplus \text{coev}_B$ and $\text{ev}_{A \oplus B} = \text{ev}_A \oplus \text{ev}_B$.

Hence, in the definition of trace, we find that the cross terms $h_{A_0, A_1} + h_{A_1, A_0}$ act by zero since they send the codomain of $\text{coev}_{A \oplus B}$ to elements with no effect on the map $\text{ev}_{A \oplus B}$. Moreover, we compute in this way that $\text{tr}(h) = \text{tr}(h_{A_0, A_0}) + \text{tr}(h_{A_1, A_1})$.

In this way, the trace splits over direct sums and only picks out diagonal elements. Applying this result inductively reduces the proof to the case that A is a simple object. This follows directly from the definition of quantum dimension. \square

5.3.2 Duality

Duality is baked into our definition of modular categories as a fundamental part of the structure. It acts in a very controlled way on fusion coefficients:

Proposition 5.6. *Let \mathcal{C} be a fusion category and let $A, B, C \in \mathcal{C}$ be simple objects. We have the following:*

- (i) (Anti-involution) $N_C^{A,B} = N_{C^*}^{B^*,A^*}$;
- (ii) (Frobenius reciprocity) $N_C^{A,B} = N_B^{A^*,C} = N_A^{C,B^*}$.

Proof. Part (i) follows from the fact that the duality functor is fully faithful and monoidal from proposition [ref], so

$$N_C^{A,B} = \dim \text{Hom}(C, A \otimes B) = \dim \text{Hom}(A^* \otimes B^*, C^*) = N_{C^*}^{B^*,A^*}.$$

Part (ii) follows from the following computation. Consider the map

$$i : \text{Hom}(A, B \otimes C) \longrightarrow \text{Hom}(A^* \otimes C, B)$$

$$\begin{array}{ccc} AB & & B \\ \boxed{\alpha} & \longmapsto & \boxed{\alpha} \\ C & & A^* C \end{array}$$

Since composition is bilinear, i is a linear map. The map

$$\begin{array}{ccc} \mathrm{Hom}(A^* \otimes C, B) & \longrightarrow & \mathrm{Hom}(A, B \otimes C) \\ \begin{array}{c} B \\ \square \\ \alpha \\ \square \\ A^*C \end{array} & \mapsto & \begin{array}{c} A \quad B \\ \cup \\ \square \\ \alpha \\ \square \\ C \end{array} \end{array}$$

serves as an inverse for i by rigidity. Hence, we conclude that

$$N_C^{A,B} = \dim \mathrm{Hom}(C, A \otimes B) = \dim \mathrm{Hom}(A^* \otimes C, B) = N_B^{A^*,C}.$$

The third equality in Frobenius reciprocity follows from an identical argument, and hence we conclude the proof. \square

In particular, we can describe the fusion rules of a simple object with its dual:

Corollary 5.2. *Let \mathcal{C} be a fusion category. Let $A, B \in \mathcal{C}$ be simple objects. We find that*

$$N_{\mathbf{1}}^{A,B} = N_{\mathbf{1}}^{B,A} = \begin{cases} 1 & B \cong A^* \\ 0 & \text{otherwise.} \end{cases}$$

Proof. This follows from Frobenius reciprocity and Schur's lemma:

$$N_{\mathbf{1}}^{A,B} = N_A^{A^*,\mathbf{1}} = \dim(\mathrm{Hom}(A, A^*)) = \begin{cases} 1 & B \cong A^* \\ 0 & \text{otherwise.} \end{cases}$$

\square

Specializing even more, we get the following corollary:

Corollary 5.3. *If \mathcal{C} is a fusion category, then $A \cong A^{**}$ for all $A \in \mathcal{C}$.*

Proof. Since $N_{\mathbf{1}}^{A^*,A^{**}} > 0$, we conclude that $A^{**} \cong A$ by the (iii) \implies (i) implication in proposition [ref]. \square

However, despite this corollary, we *cannot* conclude that every fusion category admits a pivotal structure. The isomorphism $A \cong A^{**}$ may fail to form a monoidal natural transformation. It is an open problem whether or not every fusion category admits a pivotal structure, and it is furthermore an open problem whether every pivotal fusion category admits a spherical structure [ENO05].

5.3.3 Quantum dimension and Frobenius-Perron dimension

.[WORK: include something about global quantum dimension \mathcal{D} .]

Our next tool to discuss is the *quantum dimension*. Given any spherical fusion category \mathcal{C} and any object $A \in \mathcal{C}$, we define its quantum dimension using the following formula:

$$d_A = A \bigcirc A^*$$

As usual, we identify d_A with a complex number via the canonical isomorphism $\mathrm{End}(\mathbf{1}) \cong \mathbb{C}$. The quantum dimension is clearly equal to the trace of the identity map on A , $d_A = \mathrm{tr}(\mathrm{id}_A)$. The first properties of quantum dimension follow from our general analysis of trace:

Proposition 5.7. *For every spherical fusion category \mathcal{C} and any objects $A, B \in \mathcal{C}$, we have the following formulas:*

- (i) *If $A \cong B$, then $d_A = d_B$;*
- (ii) *$d_{A \oplus B} = d_A + d_B$;*
- (iii) *$d_{A \otimes B} = d_A \cdot d_B$;*
- (iv) *$d_{A^*} = d_A$.*
- (v) *$d_A \neq 0$.*

Proof.

- (i) Let $f : A \cong B$ be an isomorphism. Using Proposition [ref] we find

$$d_A = \text{tr}(\text{id}_A) = \text{tr}(f^{-1} \circ f) = \text{tr}(f \circ f^{-1}) = \text{tr}(\text{id}_B) = d_B.$$

- (ii) This follows from proposition [ref].
- (iii) This follows from proposition [ref].
- (iv) This follows from proposition [ref].
- (v) From proposition [ref], we know that $A \otimes A^* \cong \mathbf{1} \oplus X$ for some $X \in \mathcal{C}$ which does not have any factors of $\mathbf{1}$ in its direct sum decomposition. The map $\text{coev}_A^R : \mathbf{1} \rightarrow A \otimes A^*$ is thus a non-zero scalar times the inclusion $\mathbf{1} \hookrightarrow \mathbf{1} \oplus X$, and the map $\text{ev}_A^L : A \otimes A^* \rightarrow \mathbf{1}$ is a non-zero scalar times the projection $\mathbf{1} \oplus X \rightarrow \mathbf{1}$. Since inclusion composed with projection is the identity, we find that $\text{ev}_A^L \circ \text{coev}_A^R$ is a non-zero scalar times the identity, as desired.

□

The above propositions tell us that the values d_A as A ranges over isomorphism classes of simple objects determines all the other values of d_A . Moreover, proposition [ref] tells us that the quantum dimensions of simple objects determines the trace of *every* endomorphism! Hence, computing d_A for each isomorphism class $[A] \in \mathcal{L}$ is an important step in analysing a modular category. The following formula and its linear-algebraic reformulation are the primary insight in performing the computation:

Proposition 5.8. *Let \mathcal{C} be a spherical fusion category.*

- (i) *Let $A, B \in \mathcal{C}$ be simple objects. We have that*

$$d_A d_B = \sum_{[C] \in \mathcal{L}} N_C^{A,B} d_C.$$

- (ii) *Let $A \in \mathcal{C}$ be a simple object. Define an operator*

$$\begin{aligned} N^A : \mathbb{C}[\mathcal{L}] &\rightarrow \mathbb{C}[\mathcal{L}], \\ |[B]\rangle &\mapsto \sum_{[C] \in \mathcal{L}} N_C^{A,B} |[C]\rangle \end{aligned}$$

Define $\mathbf{d} = \sum_{[B] \in \mathcal{L}} d_B |[B]\rangle \in \mathbb{C}[\mathcal{L}]$. We have that

$$N^A \mathbf{d} = d_A \mathbf{d}.$$

Proof. From proposition [ref], we have an isomorphism

$$A \otimes B \cong \bigoplus_{[C] \in \mathcal{L}} N_C^{A,B} \cdot C$$

and thus

$$\text{tr}(\text{id}_{A \otimes B}) = \text{tr}(\text{id}_{\bigoplus_{[C] \in \mathcal{L}} N_C^{A,B} \cdot C}).$$

Expanding using the rules in proposition [ref] gives part (i). Part (ii) follows from expanding the definition of the linear operator and applying part (i). \square

We now make commentary about the above proposition. It tells us that d_A is an eigenvalue of N^A . Since N^A is an operator with integer coefficients, this immediately tells us that d_A is the root of polynomial with integer coefficients. Namely, the characteristic polynomial of N^A . We can even be more precise about the nature of d_A :

[WORK: On the MathOverflow question “Modular Tensor Categories: Reasoning behind the axioms”, a commentator said “You also have to assume that the categorical dimensions arising from the pivotal structure are all real. This is called the spherical axiom”. Why is this the same thing as the spherical axiom? What is the motivation? I should include this as a remark on the below proposition.]

Corollary 5.4. *Let \mathcal{C} be a spherical fusion category. The quantum dimensions of all simple objects in \mathcal{C} are real numbers.*

Proof. .[WORK: do proof.] \square

The question is whether or not the quantum dimensions are *positive* real numbers. We recall that we defined a unitarizable spherical fusion category to be one in which the quantum dimensions are all positive. It is at this point that this becomes relevant. In particular, if \mathcal{C} is unitarizable then its quantum dimensions are eigenvalues of N^A , and their corresponding eigenvector \mathbf{d} has positive real entries. There is a theorem about eigenvalues of non-negative matrices with positive eigenvectors:

Theorem 5.1 (Frobenius-Perron theorem, [EGNO16]). *Let B be a square matrix with nonnegative real entries.*

- (i) *B has a non-negative real eigenvalue. The largest non-negative real eigenvalue $\lambda(B)$ of B dominates the absolute values of all other eigenvalues μ of B : $|\mu| \leq \lambda(B)$. Moreover, there is an eigenvector of B with non-negative entries and eigenvalue $\lambda(B)$.*
- (ii) *If B has strictly positive entries then $\lambda(B)$ is a simple positive eigenvalue, and the corresponding eigenvector can be normalized to have strictly positive entries. Moreover, $|\mu| < \lambda(B)$ for any other eigenvalue μ of B .*
- (iii) *If a matrix B with non-negative entries has an eigenvector v with strictly positive entries, then the corresponding eigenvalue is $\lambda(B)$.*

We call the largest positive real eigenvalue of a matrix its *Frobenius-Perron eigenvalue*. The Frobenius-Perron theorem tells us the following:

Corollary 5.5. *Let \mathcal{C} be a unitarizable spherical fusion category. Let $A \in \mathcal{C}$ be a simple object. The quantum dimension d_A is equal to the Frobenius-Perron eigenvalue of N^A .*

Proof. Since \mathcal{C} is unitarizable, the vector $\mathbf{d} = \sum_{[B] \in \mathcal{L}} d_B |[B]\rangle \in \mathbb{C}[\mathcal{L}]$ has positive entries and has eigenvalue d_A . Hence, d_A is the Frobenius-Perron eigenvalue of N^A as desired. \square

In this chapter we will mostly work with generic spherical fusion categories with no conditions on unitarizability. Hence, it is useful to make the following definition. Let $A \in \mathcal{C}$ be a simple object in a spherical fusion category. We define

$$\text{FPdim}(A) = (\text{Frobenius-Perron eigenvalue of } N^A).$$

When \mathcal{C} is unitarizable, $\text{FPdim}(A) = d_A$. Many formulas about quantum dimension in the unitary world apply to the Frobenius-Perron dimension in the non-unitary world. An interesting observation is that the definition of quantum dimension strongly uses the spherical structure on \mathcal{C} . However, the Frobenius-Perron dimension only uses the fusion coefficients, and those are well-defined in any fusion category. Hence, the Frobenius-Perron dimension also derives utility from being applicable in a broader set of situations than the quantum dimension.

We now give an alternate interpretation of the Frobenius-Perron dimension in terms of growth in tensor powers. This sort of alternate perspective of dimension applies to several types of objects outside the scope of tensor category theory [COT24].

Proposition 5.9. *Let \mathcal{C} be a fusion category, and let $A \in \mathcal{C}$ be a simple object.*

- (i) $\text{FPdim}(A) = \lim_{n \rightarrow \infty} \dim(\text{Hom}(A^{\otimes n}, A^{\otimes n}))^{1/(2n)}$
- (ii) $\text{FPdim}(A) = \lim_{n \rightarrow \infty} \dim(\text{Hom}(\mathbf{1}, A^{\otimes n}))^{1/n}$
- (iii)

$$\text{FPdim}(A) = \lim_{n \rightarrow \infty} (\# \text{ of simple objects in the direct sum decomposition of } A^{\otimes n})^{1/n}.$$

Proof. [WORK: I can do a good part of this when \mathcal{C} is unitarizable, so that its largest eigenvalue is strictly larger than all of the others. When there are multiple large eigenvalues all of the same size then the proofs go wrong. Is there something about the structure of N^A I can exploit? Are these theorems true for fusion categories, or do I need to pass to unitarizable fusion categories?] \square

This proposition can be interpreted as saying that the simple object A has $\text{FPdim}(A)$ internal degrees of freedom “on average”. Elements of the vector space $\text{Hom}(\mathbf{1}, A^{\otimes n})$ correspond to states in the system with n anyons of type A arranged in a line. If the internal configuration space of each anyon was $\text{FPdim}(A)$ -dimensional, then the overall dimension would be $\text{FPdim}(A)^n$. By Proposition [ref], $\text{FPdim}(A)^n$ is approximately $\dim(\text{Hom}(\mathbf{1}, A^{\otimes n}))$ for large n . Hence, each anyon has approximately $\text{FPdim}(A)$ internal degrees of freedom. Of course, $\text{FPdim}(A)$ has no reason to be an integer! In the Fibonacci theory $\text{FPdim}(\tau) = \phi = 1.61\dots$ Frobenius-Perron dimension just gives an average amount for large values.

[WORK: re-do this explanation way better + add diagram for it.]

5.3.4 Twist

In this section we will discuss *twists*. The twist is a subtle concept, which we have not explicitly mentioned up to now. The idea is that anyons can *rotate in place*. Since the space of endomorphisms of an anyon is one dimensional, this rotation must act by a phase. This phase is physically relevant, and can be measured in experiment.

For example, consider the Y -type on the toric code. It consists of the fusion of an X -type anyon and a Z -type anyon, as shown below:

[WORK: add figure of Y as a thick X and Z together; could be hard to draw these nice]

Twisting Y in place will correspond to twisting X and Z around each other. This twisting thus results in a phase of -1 . In general, we can imagine anyons as having some thickness to them. Anyons are not localized at points - they are localized at small regions. Twisting this region all the way around can be viewed visually as

[WORK: twisted anyon.]

This is the twist. One way of working with the twist is to work with thickened diagrams, where strings are replaced with ribbons. While popular in some parts of the literature, we will continue to work with string diagrams for simplicity. The key observation is that the twist can be constructed using string diagrammatic structures we already have as follows:

[WORK: twist as a swirl diagram, compared with ribbon.]

Hence, letting \mathcal{C} be a pre-modular fusion category, we *define* the twist θ_A of an object $A \in \mathcal{C}$ to be

$$\begin{array}{c} A \\ | \\ \boxed{\theta_A} \\ | \\ A \end{array} = \begin{array}{c} A \\ | \\ \circlearrowleft \\ | \\ A \end{array}$$

For every simple object $A \in \mathcal{C}$, the map $\theta_A \in \text{End}(A)$ can be identified with the unique complex number λ such that $\theta_A = \lambda \cdot \text{id}_A$. Equivalently, we can identify θ_A with the complex number $\lambda = \text{tr}(\theta_A)/d_A$ which gives the graphical formula

$$\theta_A = \frac{1}{d_A} \bigcirc A$$

We can reinterpret all other twist-like maps in terms of θ :

Lemma 5.1. *Let \mathcal{C} be a pre-modular fusion category. We have that*

$$\begin{array}{ccc} \begin{array}{c} A \\ | \\ \circlearrowleft \\ | \\ A \end{array} & = & \begin{array}{c} A \\ | \\ \boxed{\theta_A} \\ | \\ A \end{array}, \quad \begin{array}{c} A \\ | \\ \circlearrowright \\ | \\ A \end{array} & = & \begin{array}{c} A \\ | \\ \boxed{\theta_A^{-1}} \\ | \\ A \end{array} \\ \begin{array}{c} A \\ | \\ \circlearrowright \\ | \\ A \end{array} & = & \begin{array}{c} A \\ | \\ \boxed{\theta_A} \\ | \\ A \end{array}, \quad \begin{array}{c} A \\ | \\ \circlearrowleft \\ | \\ A \end{array} & = & \begin{array}{c} A \\ | \\ \boxed{\theta_A^{-1}} \\ | \\ A \end{array} \end{array}$$

Proof. To begin we show that

$$\begin{array}{c} A \\ \text{---} \\ \text{---} \\ A \end{array} = \begin{array}{c} A \\ | \\ \text{---} \\ A \end{array}.$$

When A is simple, this follows from the spherical axiom. Taking the trace of both sides gives the same formula for θ_A as a figure-eight. Additionally, pushing through duals it is clear that both sides in the above proposed equality are natural isomorphisms. Natural isomorphisms are determined by their action on simple objects because they commute with direct sums. Hence, we conclude that the sides are equal for all objects.

To get that the two reversed formulas are equal to θ_A^{-1} , it suffices to compose with θ_A and use string-diagram manipulations to show that it results in the identity. This is a simple exercise and is left as an exercise to the reader. \square

We now characterize the key properties of the twist:

Proposition 5.10. *Let \mathcal{C} be a pre-modular fusion category. The twists θ induce a monoidal natural isomorphism $\text{id}_{\mathcal{C}} \xrightarrow{\sim} \text{id}_{\mathcal{C}}$. Additionally, θ satisfies the identity*

$$\theta_{A \otimes B} = \beta_{B,A} \circ \beta_{A,B} \circ (\theta_A \otimes \theta_B)$$

for all $A, B \in \mathcal{C}$, and $\theta_{A^*} = (\theta_A)^*$.

Proof. Naturality of θ follows from pushing through duals. The formula $\theta_{A \otimes B} = \beta_{B,A} \circ \beta_{A,B} \circ (\theta_A \otimes \theta_B)$ comes from manipulating string diagrams to get the equation

$$\begin{array}{c} AB \\ \text{---} \\ \text{---} \\ AB \end{array} = \begin{array}{c} A \ B \\ \text{---} \\ \text{---} \\ A \ B \end{array}$$

Finally, $\theta_{A^*} = (\theta_A)^*$ comes from the string-diagram manipulation and proposition [ref]:

$$\begin{array}{c} A^* \\ | \\ \boxed{(\theta_A)^*} \\ | \\ A^* \end{array} = \begin{array}{c} A^* \\ \text{---} \\ \text{---} \\ A^* \end{array} = \begin{array}{c} A^* \\ | \\ \text{---} \\ A^* \end{array} = \begin{array}{c} A \\ | \\ \boxed{\theta_A} \\ | \\ A \end{array}$$

as desired. \square

The naive reason to care about twists is that they describe a physically relevant quantity and hence should be studied. The more subtle reason to care about twists is that they are the most efficient way of encoding the spherical structure on \mathcal{C} . A spherical structure is first and foremost a pivotal structure, meaning that it has a right and left rigid structure which

are compatible. Given a spherical structure one can always obtain twists. Conversely, given a right-rigid structure and twists one can recover the left-rigid structure via the formulas

$$A^* \underset{\text{---}}{\cup} A = \begin{array}{c} A^* \quad A \\ \text{---} \\ \theta_A \end{array} \quad \text{and} \quad A \underset{\text{---}}{\cap} A^* = \begin{array}{c} \theta_A^{-1} \\ \text{---} \\ A \quad A^* \end{array}$$

In this way, giving a spherical structure on a right-rigid monoidal category is the *same* as giving a twist structure. This is codified in the following lemma:

Proposition 5.11 (Deligne's twisting lemma, [Yet92]). *Let \mathcal{C} be a right-rigid braided monoidal category. Every pivotal structure on \mathcal{C} naturally gives a twist natural transformation $\theta : \text{id}_{\mathcal{C}} \rightarrow \text{id}_{\mathcal{C}}$. This assignment induces a canonical bijection between the set of pivotal structures on \mathcal{C} and the set of natural isomorphisms $\theta : \text{id}_{\mathcal{C}} \rightarrow \text{id}_{\mathcal{C}}$ satisfying $\theta_{A \otimes B} = \beta_{B,A} \circ \beta_{A,B} \circ (\theta_A \otimes \theta_B)$ for all $A, B \in \mathcal{C}$.*

Moreover, restricting the assignment to the space of spherical structures on \mathcal{C} induces a canonical bijection between the set of spherical structures on \mathcal{C} and the set of isomorphisms $\theta : \text{id}_{\mathcal{C}} \rightarrow \text{id}_{\mathcal{C}}$ satisfying $\theta_{A \otimes B} = \beta_{B,A} \circ \beta_{A,B} \circ (\theta_A \otimes \theta_B)$ for all $A, B \in \mathcal{C}$ and $\theta_{A^*} = (\theta_A)^*$.

Proof. We already showed in proposition [ref] that every spherical category gives a twist natural transformation satisfying the desired axioms. Restricting the proof to only a possibly non-spherical pivotal category still gives a twist natural transformation satisfying $\theta_{A \otimes B} = \beta_{B,A} \circ \beta_{A,B} \circ (\theta_A \otimes \theta_B)$ for all $A, B \in \mathcal{C}$. The heart of the proof is showing that the formulas [ref] induce pivotal and spherical structures with the twist satisfies the right axioms. The process of inducing a pivotal structure and inducing a twist are inverses to one another because

$$\begin{array}{ccc} A & & A \\ \downarrow & \text{---} & \downarrow \\ \theta_A & = & \text{---} \\ \downarrow & \text{---} & \downarrow \\ A & & A \end{array}.$$

To begin, we assume that $\theta_{A \otimes B} = \beta_{B,A} \circ \beta_{A,B} \circ (\theta_A \otimes \theta_B)$ and we seek to prove that the corresponding ev^L , coev^L maps induce a pivotal structure. We first axiom of pivotality follows from use of the axiom $\theta_{A \otimes B} = \beta_{B,A} \circ \beta_{A,B} \circ (\theta_A \otimes \theta_B)$:

$$\begin{array}{c}
B^* \otimes A^* \quad A \otimes B \\
\text{---} = \quad \text{---} = \quad \text{---} \\
\text{---} = \quad \text{---} = \quad \text{---} = \quad \text{---}
\end{array}$$

The second axiom of pivotality follows from the use of the naturality of θ :

$$\begin{array}{c}
A^* \quad \quad \quad A^* \\
\text{---} = \quad \text{---} = \quad \text{---} \\
\text{---} = \quad \text{---} = \quad \text{---}
\end{array}$$

Finally, we assume that $(\theta_A)^* = \theta_{A^*}$ and we seek to prove the spherical axiom. Taking the dual of theta we can get all of the equalities in Lemma [ref]. Applying them we get that

$$\begin{array}{c}
f \\
\text{---} = \quad \text{---} = \quad \text{---} = \quad \text{---} \\
\text{---} = \quad \text{---} = \quad \text{---} = \quad \text{---}
\end{array}$$

as desired. \square

5.3.5 Functors, natural transformations, and equivalence

In this section, we will talk about functors, natural transformations, and equivalences between fusion, spherical, pre-modular, and modular categories. Given a topological order, there is *not* a unique modular category describing it. There is a unique modular category *up to equivalence*. Hence, the notion of equivalence of categories is baked into our physics-math correspondance so it is important that we state it explicitly.

Functors which do not induce equivalences of categories are also physically relevant. In certain contexts, a functor $F : \mathcal{C} \rightarrow \mathcal{D}$ is used to model a *phase transition* from \mathcal{C} to \mathcal{D} . We will see a lot more functors and natural transformations between modular categories throughout the book, especially in chapter [ref].

Even though structures in categories require a lot of compatibility conditions, the conditions on the functors do not. This means that we have the following:

- The correct notion of functor between fusion categories is \mathbb{C} -linear monoidal functor. There is no compatibility condition required between the \mathbb{C} -linear structure and the monoidal structure. The correct notion of natural transformation between \mathbb{C} -linear monoidal functors is a monoidal natural transformation.
- The correct notion of functor between spherical fusion categories is \mathbb{C} -linear pivotal monoidal functor. There is no compatibility condition required between the \mathbb{C} -linear structure and the pivotal monoidal structure. The correct notion of natural transformation is monoidal natural transformation.
- The correct notion of functor between pre-modular categories is \mathbb{C} -linear pivotal braided monoidal functor. There is no compatibility condition required between the \mathbb{C} -linear structure, pivotal monoidal structure, or braided monoidal structure. The correct notion of natural transformation is monoidal natural transformation.
- The correct notions of functors/natural transformations for modular categories are the same as for pre-modular categories.

[WORK: this section is very short. I don't have much to say, actually. Should this be moved? Maybe I keep a very short section? I don't know.]

5.3.6 Deligne tensor product

In the theory of any class of mathematical object, an important consideration is the ways in which examples can be put together to give new examples. In the case of fusion categories, this basic operation is known as the *Deligne tensor product*. Given any fusion categories \mathcal{C}, \mathcal{D} , their Deligne tensor product $\mathcal{C} \boxtimes \mathcal{D}$ is a new fusion category. The Deligne tensor product of spherical fusion categories will be equipped with the structure of a spherical fusion category, and the Deligne tensor product of (pre-)modular categories will be equipped with the structure of a (pre-)modular category.

Physically, the Deligne tensor product corresponds to *stacking*. Consider two sheets of material. We choose two modular categories \mathcal{C}, \mathcal{D} . We endow the top sheet with the structure of a topologically ordered quantum system described by \mathcal{C} and we endow the bottom with the structure of a topologically ordered quantum system described by \mathcal{D} . The algebraic description of this bilayer system is $\mathcal{C} \boxtimes \mathcal{D}$. This can be viewed as the physical definition of $\mathcal{C} \boxtimes \mathcal{D}$.

[WORK: add bilayer system diagram]

We now mathematically define the Deligne tensor product.

Definition. Let \mathcal{C}, \mathcal{D} be a \mathbb{C} -linear categories, isomorphic as a \mathbb{C} -linear categories to $\mathbf{Vec}_{\mathbb{C}}^n, \mathbf{Vec}_{\mathbb{C}}^m$ respectively. We define a Deligne tensor product of \mathcal{C} and \mathcal{D} to be the following data:

1. A \mathbb{C} -linear category $\mathcal{C} \boxtimes \mathcal{D}$;
2. A \mathbb{C} -linear functor $\mathcal{C} \times \mathcal{D} \rightarrow \mathcal{C} \boxtimes \mathcal{D}$.

Such that:

1. Every object $X \in \mathcal{C} \boxtimes \mathcal{D}$ has a direct sum decomposition

$$X \cong \bigoplus_{i=1}^n A_i \boxtimes B_i$$

for some $n \geq 1$, $A_i \in \mathcal{C}$, $B_i \in \mathcal{D}$.

2. There is an equality of vector spaces

$$\text{Hom}_{\mathcal{C} \boxtimes \mathcal{D}}(A \boxtimes B, A' \boxtimes B') = \text{Hom}_{\mathcal{C}}(A, A') \otimes \text{Hom}_{\mathcal{D}}(B, B').$$

3. Given any $A, A', A'' \in \mathcal{C}$, $B, B', B'' \in \mathcal{D}$, $f : A \rightarrow A'$, $f' : A' \rightarrow A''$, $g : B \rightarrow B'$, $g' : B' \rightarrow B''$, the diagram

$$\begin{array}{ccccc} A \boxtimes B & \xrightarrow{f \boxtimes g} & A' \boxtimes B' & \xrightarrow{f' \boxtimes g'} & A'' \boxtimes B'' \\ & \searrow & & \nearrow & \\ & & (f' \circ f) \boxtimes (g' \circ f) & & \end{array}$$

commutes.

We now state the main existence/uniqueness result about the Deligne tensor product:

Proposition 5.12. *Let \mathcal{C}, \mathcal{D} be \mathbb{C} -linear categories isomorphic as \mathbb{C} -linear categories to $\mathbf{Vec}_{\mathbb{C}}^n$ and $\mathbf{Vec}_{\mathbb{C}}^m$ respectively. There exists a Deligne product $\mathcal{C} \boxtimes \mathcal{D}$ for \mathcal{C} and \mathcal{D} . Moreover, given any other deligne tensor product $\mathcal{C} \boxtimes' \mathcal{D}$ of \mathcal{C} and \mathcal{D} there exists a unique functor $F : \mathcal{C} \boxtimes \mathcal{D} \rightarrow \mathcal{C} \boxtimes' \mathcal{D}$ making the diagram*

$$\begin{array}{ccc} \mathcal{C} \times \mathcal{D} & \longrightarrow & \mathcal{C} \boxtimes \mathcal{D} \\ & \searrow & \downarrow F \\ & & \mathcal{C} \boxtimes' \mathcal{D} \end{array}$$

commute. This functor is an equivalence of categories.

Proof. It is clear that $\mathbf{Vec}_{\mathbb{C}}^n \boxtimes \mathbf{Vec}_{\mathbb{C}}^m = \mathbf{Vec}_{\mathbb{C}}^{nm}$. Every equivalence of categories $\mathcal{C} \rightarrow \mathcal{C}'$ induces an equivalence of categories $\mathcal{C} \boxtimes \mathcal{D} \rightarrow \mathcal{C}' \boxtimes \mathcal{D}$. Hence, since \mathcal{D} and \mathcal{D} are equivalent to $\mathbf{Vec}_{\mathbb{C}}^n$ and $\mathbf{Vec}_{\mathbb{C}}^m$ respectively, their Deligne tensor product exists and is equivalent to $\mathbf{Vec}_{\mathbb{C}}^{nm}$.

Any functor making the diagram commute must send $A \boxtimes B$ to $A \boxtimes' B$. The definition of Deligne tensor produce tells us this is enough to conclude that the map is an equivalence of categories, since axiom 3 this map is always a functor, axiom 2 implies it is fully faithful, and axiom 1 implies it is essentially surjective, and hence we can apply proposition [ref]. \square

Now that we have defined the Deligne tensor product of \mathbb{C} -linear categories equivalent to $\mathbf{Vec}_{\mathbb{C}}^n$, we move on to defining the Deligne tensor product of fusion categories, spherical fusion categories, pre-modular categories, and modular categories.

Proposition 5.13. *The following claims are all true.*

- (i) Let \mathcal{C}, \mathcal{D} be fusion categories. On the level of objects, define a monoidal structure $\mathcal{C} \boxtimes \mathcal{D}$ by the formula

$$(A \boxtimes B) \otimes (A' \boxtimes B') = (A \otimes A') \boxtimes (B \otimes B').$$

Along with a natural choice of action of the tensor product on morphisms, unit $\mathbf{1}_{\mathcal{C} \boxtimes \mathcal{D}} = \mathbf{1}_{\mathcal{C}} \boxtimes \mathbf{1}_{\mathcal{D}}$, and a natural choice of associator and unitors, this induces the structure of a monoidal category on \mathcal{C} .

Define a right-rigid structure on $\mathcal{C} \boxtimes \mathcal{D}$ as follows. The dual of an object $A \boxtimes B$ is $A^* \boxtimes B^*$. Define $\text{ev}_{A \boxtimes B} = \text{ev}_A \boxtimes \text{ev}_B$, $\text{coev}_{A \boxtimes B} = \text{coev}_A \boxtimes \text{coev}_B$. This induces a well-defined right-rigid structure on $\mathcal{C} \boxtimes \mathcal{D}$.

The above definitions induce the structure of a fusion category on $\mathcal{C} \boxtimes \mathcal{D}$.

- (ii) Let \mathcal{C}, \mathcal{D} be spherical fusion categories. The evaluation and coevaluation maps $\text{ev}_{A \boxtimes B}^L = \text{ev}_A^L \boxtimes \text{ev}_B^L$ and $\text{coev}_{A \boxtimes B}^L = \text{coev}_A^L \boxtimes \text{coev}_B^L$ induce a left-rigid structure on $\mathcal{C} \boxtimes \mathcal{D}$. Along with the canonical structure of a fusion category on $\mathcal{C} \boxtimes \mathcal{D}$, this induces the structure of a spherical fusion category on $\mathcal{C} \boxtimes \mathcal{D}$.
- (iii) Let \mathcal{C}, \mathcal{D} be pre-modular categories. The braiding map $\beta_{\mathcal{C} \boxtimes \mathcal{D}} = \beta_{\mathcal{C}} \boxtimes \beta_{\mathcal{D}}$ induces the structure of a pre-modular category on $\mathcal{C} \boxtimes \mathcal{D}$. The product $\mathcal{C} \boxtimes \mathcal{D}$ is modular if and only if \mathcal{C}, \mathcal{D} are both modular.

Proof. Given any of the above structures, all of the axioms on $\mathcal{C} \boxtimes \mathcal{D}$ immediately follow from their respective axioms on \mathcal{C} and \mathcal{D} . Hence, the proof is an exercise in recalling definitions which we omit. \square

5.4 The category of G -graded G -representations

5.4.1 Overview

We've talked about a lot of general theory of modular categories. It's time for us to focus on our main family of *examples*. Namely, the categories $\mathfrak{D}(G)$ of G -graded G -representations. These categories describe discrete gauge theory based on the finite group G .

Before we can prove that $\mathfrak{D}(G)$ is a modular category, we need to endow $\mathfrak{D}(G)$ with the necessary structures. In particular, we will endow $\mathfrak{D}(G)$ with \mathbb{C} -linear, monoidal, braided, right-rigid, and left-rigid structures. We will need to show that all of these structures are compatible with each other in the necessary ways, and that $\mathfrak{D}(G)$ satisfies the non-degeneracy axiom. We will use this as an opportunity to introduce tools of general use for proving that categories satisfy the axioms of a modular category.

Additionally, we will also study two categories similar to $\mathfrak{D}(G)$ which will serve as extra examples to get our grip on definitions. These categories will also appear later as relevant in and of themselves. The first is \mathbf{Vec}_G , the category of G -graded vector spaces. It is defined as follows:

[WORK: define \mathbf{Vec}_G in terms of objects and composition.]

Our second structure of interest is $\text{Rep}(G)$, the category of G -representations. It is defined as follows:

[WORK: define $\text{Rep}(G)$ in terms of objects and composition.]

We will show that both \mathbf{Vec}_G and $\text{Rep}(G)$ can be naturally equipped with the structures of spherical fusion categories. We then show that $\text{Rep}(G)$ admits a braiding which turns it

into a pre-modular category. This braiding is symmetric in the sense that $\beta_{B,A} \circ \beta_{A,B} = \text{id}_A \otimes \text{id}_B$ for all $A, B \in \mathcal{C}$, and hence $\text{Rep}(G)$ is not a modular category. The category \mathbf{Vec}_G is shown to not admit a braiding whenever G is non-abelian.

5.4.2 Higher linear algebra

.[WORK: In this section we define the \mathbb{C} -linear structures on \mathbf{Vec}_G , $\text{Rep}(G)$, and $\mathfrak{D}(G)$. Our goal is to show that they are all equivalent to $\mathbf{Vec}_{\mathbb{C}}^n$ for some $n \geq 1$.

It seems like the best approach is through higher linear algebra. Namely, we show that if \mathcal{C} is abelian, \mathbb{C} -linear, semisimple, and has finitely many isomorphism classes of simple objects then it must be isomorphic to $\mathbf{Vec}_{\mathbb{C}}^n$. Its a good time to wax philosophical about higher linear algebra and 2-vector spaces. However, its not clear that this approach actually helps at all. It might be easier to immediately note that everybody is the direct sum of irreducibles, prove a Schur's lemma, and call it a day. Of course these approaches are all equivalent but its not clear what's best.

Here's a lemma I would like to use. Since \times and \boxtimes look too similar, I should use \boxplus to denote the Cartesian product and call it the direct sum. I should then prove this:

Lemma 5.2. *Consider the category $\mathbf{2Vec}$ whose objects are \mathbb{C} -linear categories equivalent to $\mathbf{Vec}_{\mathbb{C}}^n$ for some $n \geq 0$, and whose morphisms are \mathbb{C} -linear functors. The direct sum \boxplus is a biproduct in $\mathbf{2Vec}$*

Proof. .[WORK: do proof]

□

]

5.4.3 Spherical fusion structures

.[WORK: show that the categories have duals and monoidal structure. This should be pretty easy and painless. Pentagon identity should follow from the pentagon identity on $\mathbf{Vec}_{\mathbb{C}}$.]

5.4.4 Braiding and modularity

.[WORK: Introduce braidings. Show that $\text{Rep}(G)$ is symmetric. Show that \mathbf{Vec}_G does not admit a braiding if G is not abelian and does admit a symmetric braiding if G is abelian. Show that $\mathfrak{D}(G)$ admits a non-degenerate braiding.]

5.5 The modular representation

5.5.1 Definition

In this chapter we are going to talk about the *modular representations* of modular categories. Here's the point. Let \mathcal{C} be a modular category. Let \mathcal{L} be the set of isomorphism classes of simple objects of \mathcal{C} . We will define a group homomorphism

$$\rho_{\mathcal{C}} : \text{SL}_2(\mathbb{Z}) \rightarrow \text{Aut}(\mathbb{C}[\mathcal{L}])$$

associated to \mathcal{C} , where $\text{SL}_2(\mathbb{Z})$ is the group of 2-by-2 matrices with integer coefficients and unit determinant. The group $\text{SL}_2(\mathbb{Z})$ is sometimes known as the *modular group*, due to its connection with moduli spaces of elliptic curves. Hence, $\rho_{\mathcal{C}}$ is known as the *modular representation* of \mathcal{C} .

The goal of this chapter is to introduce $\rho_{\mathcal{C}}$, show it is well defined, and then prove a series of theorems related to $\rho_{\mathcal{C}}$.

Before defining $\rho_{\mathcal{C}}$, we recall the basic group theory of $\mathrm{SL}_2(\mathbb{Z})$. It is generated by the matrices

$$s = \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix}, \quad t = \begin{pmatrix} 1 & 1 \\ 0 & 1 \end{pmatrix}.$$

These two matrices satisfy the relations $s^2 = -1$ and $(st)^3 = -1$, where 1 is used to represent the identity matrix. These relations generate $\mathrm{SL}_2(\mathbb{Z})$, in the sense that we have the following presentation:

Proposition 5.14. *The following presentation is valid:*

$$\mathrm{SL}_2(\mathbb{Z}) = \langle s, t \mid s^4 = 1, (st)^3 = s^2 \rangle.$$

Proof. This is a standard fact about $\mathrm{SL}_2(\mathbb{Z})$. See for instance [WORK: ref]. \square

Hence, to define a homomorphism $\rho_{\mathcal{C}} : \mathrm{SL}_2(\mathbb{Z}) \rightarrow \mathrm{Aut}(\mathbb{C}[\mathcal{L}])$ it suffices to choose automorphisms $\rho_{\mathcal{C}}(s)$, $\rho_{\mathcal{C}}(t)$ of $\mathbb{C}[\mathcal{L}]$, and show that they satisfy the relations $\rho_{\mathcal{C}}(s)^4 = 1$ and $(\rho_{\mathcal{C}}(s)\rho_{\mathcal{C}}(t))^3 = \rho_{\mathcal{C}}(s)^2$. Since $\mathbb{C}[\mathcal{L}]$ has a canonical basis, we can think of its automorphisms as being matrices with rows and columns labeled by \mathcal{L} . We define an operator $S : \mathbb{C}[\mathcal{L}] \rightarrow \mathbb{C}[\mathcal{L}]$ via the matrix coefficients

$$S_{A,B} = A^* \circlearrowleft B^*$$

We next define the matrix $T : \mathbb{C}[\mathcal{L}] \rightarrow \mathbb{C}[\mathcal{L}]$ to be the diagonal matrix with $([A], [A])$ -entry θ_A , for all $[A] \in \mathcal{L}$.

As currently stated, the S and T matrices defined do not satisfy $S^4 = 1$ and $(ST)^3 = S^2$. They only satisfy these formula up to phases in \mathbb{C} . They still need to be normalized before we can define $\rho_{\mathcal{C}}$. The normalization factors come in terms of the *Gauss sums*,

$$p_{\mathcal{C}}^{\pm} = \sum_{[A] \in \mathcal{L}} \theta_A^{\pm 1} d_A^2.$$

We can now state the main theorem of this chapter:

Theorem 5.2. *Let \mathcal{C} be a modular category. The values $p_{\mathcal{C}}^+$ and $p_{\mathcal{C}}^-$ are nonzero, and the map*

$$\begin{aligned} \rho_{\mathcal{C}} : \mathrm{SL}_2(\mathbb{Z}) &\rightarrow \mathrm{Aut}(\mathbb{C}[\mathcal{L}]) \\ s &\mapsto \frac{1}{D} \cdot S \\ t &\mapsto (p_{\mathcal{C}}^- / p_{\mathcal{C}}^+)^{1/6} \cdot T \end{aligned}$$

is a group homomorphism.

We will prove this theorem and motivate why it should be true over the course of this chapter. We will also prove key facts about the image and kernel of this representation, as well as other formulas of interest relating to twists, S-matrix entries, and Gauss sums.

5.5.2 Torus perspective

It's good to reflect on why MCs have $\mathrm{SL}_2(\mathbb{Z})$ representations associated with them in the first place. Not only does the representation exist, but it is so fundamental to the modular category that it is chosen as the namesake. This begs the question. What's going on?

The answer has to do with topological phases on the torus.

[WORK: add torus]

Every modular category \mathcal{C} is supposed to describe a topological order. Up to now we have only considered what happens when this topological order is applied to an infinitely large flat sheet. We have not examined what happens when this topological order is put on a space with nontrivial topology. For instance, the torus. Suppose we analyse the system of \mathcal{C} applied to the torus. This amounts to breaking up the torus into some microscopic lattice and applying some Hamiltonian. This Hamiltonian will have group states $V_{\text{g.s.}}^{T^2}$, which are independent of the choice of microscopic realization of \mathcal{C} .

Suppose we start with a torus, cut it across, twist one of its legs, then glue it back together, as shown below:

[WORK: add Dehn twist picture.]

If the initial torus has some state $|\psi\rangle \in V_{\text{g.s.}}^{T^2}$ on it, then applying this procedure would give back another group state, though possibly a different one. The key phenomenon is that continues transformations on physical space correspond to linear transformations on state space:

[WORK: add schematic.]

We can make this more formal as follows. We define the *mapping class group* of a topological space X as follows:

$$\mathrm{MCG}(X) = (\text{homeomorphisms } X \rightarrow X) / (\text{continuous deformations}).$$

If two homeomorphisms can be continuously deformed from one another then they will act the same on the ground states $V_{\text{g.s.}}^{T^2}$. This is because ground states are topologically protected and hence slowly changing the diffeomorphism cannot affect them. Hence, we get a well-defined group homomorphism

$$\rho_{\mathcal{C}}^{T^2} : \mathrm{MCG}(T^2) \rightarrow \mathrm{Aut}\left(V_{\text{g.s.}}^{T^2}\right).$$

This homomorphism connects back to our modular representation as follows:

- **Claim 1:** $\mathrm{MCG}(T^2) \cong \mathrm{SL}_2(\mathbb{Z})$;
- **Claim 2:** $V_{\text{g.s.}}^{T^2} \cong \mathbb{C}[\mathcal{L}]$;
- **Claim 3:** $\rho_{\mathcal{C}}^{T^2} \cong \rho_{\mathcal{C}}$, passing through the identifications in claims 1 and 2.

In general, we see that associated to every modular category \mathcal{C} there should not only be a modular representation, but also a representation of $\mathrm{MCG}(\Sigma)$ for many other choices of topological space Σ . For instance, if $\Sigma = \Sigma_g$ is the g -holed torus then putting \mathcal{C} on Σ_g we get a map

$$\rho_{\mathcal{C}}^{\Sigma_g} : \mathrm{MCG}(\Sigma_g) \rightarrow \mathrm{Aut}(V_{\text{g.s.}}^{\Sigma_g}).$$

[WORK: maybe say a few words about these representations. I'm sure they must have an explicit description in terms of generators and relations. A good reference (though a bit early) is this one: [Lyu95]]

We now examine and motivate claims 1-3.

Claim 1: $\text{MCG}(T^2) \cong \text{SL}_2(\mathbb{Z})$. This claim is best seen by thinking of the torus as a gluing diagram,

[WORK: add gluing diagram]

[WORK: add diagram with s acting by rotating by 90 degrees. Clearly, $s^4 = 1$.]

[WORK: add diagram with t as a shift. $(st)^3 = s^2$ can be left as an exercise.]

[WORK: writing presentation for $\text{MCG}(T^2)$, note that it is the same as $\text{SL}_2(\mathbb{Z})$.]

Claim 2: $V_{\text{g.s.}}^{T^2} \cong \mathbb{C}[\mathcal{L}]$.

[WORK: explain this. Cut into cylinder, label by charge on boundary]

Claim 3: $\rho_{\mathcal{C}}^{T^2} \cong \rho_{\mathcal{C}}$.

[WORK: Showing that the Dehn twist acts diagonally by θ_A is obvious. θ_A and Dehn twist are both defined as a 2π twist. For S we need another argument, more subtle but not too hard. I think Simon has it.]

[WORK: Finish by saying this is something like TQFTs. TQFT = bundled collection of mapping class group representations. Link this to TQFT appendix.]

5.5.3 Bruguières's modularity theorem and the Verlinde formula

In this section we prove Bruguières's modularity theorem. This theorem asserts that, given a pre-modular category \mathcal{C} , the S -matrix S is invertible if and only if \mathcal{C} is modular. Historically, this theorem is backwards. The original definition of modular category included that the S -matrix should be invertible. This was the only definition of modular category, until Bruguières proved in [ref] that the invertability of the S -matrix is equivalent to \mathcal{C} having the non-degenerate braiding property that if

$$\begin{array}{c} A & B \\ \swarrow & \searrow \\ A & B \end{array} = \begin{array}{c} A & B \\ | & | \\ A & B \end{array}$$

for all $B \in \mathcal{C}$ then $A \cong \mathbf{1}$. We are thus stating a historically incorrect definition of modular category, and Bruguières's modularity theorem tells us that this is equivalent to the original definition. The proof of the modularity theorem relies on the *Verlinde algebra* of \mathcal{C} . This algebra will be of use for us in proving other theorems in the future, in particular the Verlinde formula in section [ref].

We define an *algebra* over \mathbb{C} to a vector space V paired with a bilinear map $\cdot : V \times V \rightarrow V$ called multiplication, such that multiplication is associative and has a unit. An algebra is called *commutative* if its multiplication is commutative.

We define the Verlinde algebra $K_{\mathbb{C}}(\mathcal{C})$ of \mathcal{C} as follows:

$$K_{\mathbb{C}}(\mathcal{C}) = \left\{ \mathbb{C}[\mathcal{L}] \text{ with algebra structure } |[A]\rangle \cdot |[B]\rangle = \sum_{[C] \in \mathcal{L}} N_C^{A,B} |[C]\rangle \right\}.$$

We additionally define the function algebra

$$\mathbb{C}[\mathcal{L}]^{\text{func}} = \left\{ \mathbb{C}[\mathcal{L}] \text{ with algebra structure } \left(\sum_{[A] \in \mathcal{L}} c_A |[A]\rangle \right) \cdot \left(\sum_{[A] \in \mathcal{L}} c'_A |[A]\rangle \right) = \sum_{[A] \in \mathcal{L}} c_A c'_A |[A]\rangle \right\}.$$

Lemma 5.3. Both $K_{\mathbb{C}}(\mathcal{C})$ and $\mathbb{C}[\mathcal{L}]^{\text{func}}$ are commutative algebras.

Proof. The fact that $K_{\mathbb{C}}(\mathcal{C})$ is associative follows from the associativity of the tensor product. Its unit is $|[1]\rangle$. It is commutative because \mathcal{C} is braided. The fact that $\mathbb{C}[\mathcal{L}]^{\text{func}}$ is a commutative algebra is a standard exercise in algebra. Its unit is $\sum_{[A] \in \mathcal{L}} |[A]\rangle$. \square

We now state and prove the core theorem which underlies the core properties of the S matrix:

Proposition 5.15. *The map*

$$\begin{aligned}\mathcal{S} : K_{\mathbb{C}}(\mathcal{C}) &\rightarrow \mathbb{C}[\mathcal{L}]^{\text{func}} \\ |[A]\rangle &\mapsto \sum_{[B] \in \mathcal{L}} \frac{1}{d_B} S_{B,A} |[B]\rangle\end{aligned}$$

is a morphism of algebras.

Proof. Since it was defined on a basis, \mathcal{S} is clearly a linear map. We now verify that \mathcal{S} preserves multiplication. In the below computation, we identify endomorphisms of simple objects with the unique scalar they are times the identity. We let A, B, D be simple objects.

$$\begin{aligned}\left(\frac{1}{d_D} S_{D,A}\right) \left(\frac{1}{d_D} S_{D,B}\right) &= \begin{pmatrix} D \\ | \\ A \circlearrowleft \\ D \end{pmatrix} \cdot \begin{pmatrix} D \\ | \\ B \circlearrowleft \\ D \end{pmatrix} = \begin{pmatrix} D \\ | \\ B \circlearrowleft \\ A \circlearrowleft \\ D \end{pmatrix} \\ &= \begin{pmatrix} D \\ | \\ A \otimes B \circlearrowleft \\ D \end{pmatrix} = \sum_{[C] \in \mathcal{L}} N_C^{A,B} \cdot \begin{pmatrix} D \\ | \\ C \circlearrowleft \\ D \end{pmatrix} \\ &= \sum_{[C] \in \mathcal{L}} N_C^{A,B} \left(\frac{1}{d_D} S_{D,C} \right).\end{aligned}$$

Note our key use of the fact that

$$B \oplus C \circlearrowleft = B \circlearrowleft + C \circlearrowleft$$

which follows from the facts that $\text{id}_{B \oplus C}$ can be decomposed as projection onto B plus projection onto C , and composition is bilinear. We now conclude that

$$\begin{aligned}
\mathcal{S}(|[A]\rangle) \cdot \mathcal{S}(|[B]\rangle) &= \sum_{[D] \in \mathcal{L}} \left(\frac{1}{d_D} S_{D,A} \right) \left(\frac{1}{d_D} S_{D,B} \right) |[D]\rangle \\
&= \sum_{[D] \in \mathcal{L}} \left(\sum_{[C] \in \mathcal{L}} N_C^{A,B} \left(\frac{1}{d_D} S_{D,C} \right) \right) |[D]\rangle \\
&= \mathcal{S}(|[A]\rangle \cdot |[B]\rangle)
\end{aligned}$$

as desired. \square

By Proposition [ref], we have constructed a map of algebras $\mathcal{S} : K_{\mathbb{C}}(\mathcal{C}) \rightarrow \mathbb{C}[\mathcal{L}]^{\text{func}}$. As a map of vector spaces, \mathcal{S} is equal to the S -matrix up to a rescaling of rows by nonzero factors. Hence, it is clear that the S -matrix is invertible if and only if the algebra map \mathcal{S} is invertible. We now use special properties of the map \mathcal{S} to prove the main theorem of the section:

Theorem 5.3 (Bruguières's modularity theorem). *Let \mathcal{C} be a pre-modular category. The braiding on \mathcal{C} satisfies the non-degenerate braiding axiom if and only if the S -matrix is invertible.*

Proof. We observe that \mathcal{C} has a degenerate braiding if and only if there exists some $A \not\cong \mathbf{1}$ such that

$$\begin{array}{ccc}
D & & D \\
| & & | \\
A \bigcirc \text{---} & = d_A \cdot & | \\
| & & | \\
D & & D
\end{array}$$

for all $D \in \mathcal{D}$. If such an element A exists, then clearly $\mathcal{S}(|[A]\rangle) = d_A \mathcal{S}(|[1]\rangle)$. Hence, two linearly independent vectors map to linearly dependent vectors and thus \mathcal{S} is not invertible. Thus, the invertibility of the S -matrix implies that the braiding is non-degenerate.

We now prove the converse, and hence we suppose that \mathcal{C} has nondegenerate braiding. The proof is in two main steps. First, we prove that $|[1]\rangle$ is in the image of \mathcal{S} . Then, we use the fact that $|[1]\rangle$ is in the image of \mathcal{S} to construct the rest of the image, which proves that \mathcal{S} is surjective hence invertible.

Part 1: $|[1]\rangle$ is in the image of \mathcal{S} . Since \mathcal{C} has nondegenerate braiding, for all simple objects $A \not\cong \mathbf{1}$ there exists some simple object \tilde{A} such that

$$\begin{array}{ccc}
A & & A \\
| & & | \\
\tilde{A} \bigcirc \text{---} & \neq d_{\tilde{A}} \cdot & | \\
| & & | \\
A & & A
\end{array}$$

Thus, the vector $\mathcal{S}(|[\tilde{A}]\rangle) - \frac{S_{A,\tilde{A}}}{d_A} \mathcal{S}(|[1]\rangle)$ has a coefficient zero of for $|[A]\rangle$ but a non-zero coefficient for $|[1]\rangle$. Thus, using the product structure on $\mathbb{C}[\mathcal{L}]^{\text{func}}$, we find that the vector

$$\prod_{\substack{[A] \in \mathcal{L} \\ A \not\cong 1}} \left(\mathcal{S}(|[\tilde{A}]\rangle) - \frac{S_{A,\tilde{A}}}{d_A} \mathcal{S}(|[1]\rangle) \right)$$

has a coefficient of zero for all $|[A]\rangle$, $A \not\cong \mathbf{1}$, but a non-zero coefficient of $|[\mathbf{1}]\rangle$. Hence, it is a scalar multiple of $|[\mathbf{1}]\rangle$. Since \mathcal{S} is a morphism of algebras, it is in the image of \mathcal{S} . Hence, $|[\mathbf{1}]\rangle$ is in the image of \mathcal{S} .

This completes the first part of the proof. We now use the fact that $|[\mathbf{1}]\rangle$ is in the image of \mathcal{S} to construct the rest of the vectors.

Part 2: \mathcal{S} is surjective. Let $\omega = \sum_{[A] \in \mathcal{L}} \omega_A |[A]\rangle \in K_{\mathbb{C}}(\mathcal{C})$ be a vector such that $\mathcal{S}(\omega) = |[\mathbf{1}]\rangle$, which exists by part 1 of the proof. We now compute the quantity

$$h_{X,Y} = \sum_{[A] \in \mathcal{L}} \omega_A \cdot \text{tr} \left(A \begin{array}{c} X & Y \\ \vdash & \dashv \\ - & \end{array} \right)$$

two ways, for all simple objects $X, Y \in \mathcal{C}$. The first way follows by expanding $X \otimes Y$ as a direct sum and using the fact that $\mathcal{S}(\omega) = |[\mathbf{1}]\rangle$:

$$\begin{aligned} h_{X,Y} &= \sum_{[A] \in \mathcal{L}} \sum_{[B] \in \mathcal{L}} \omega_A N_B^{X,Y} d_B \cdot \begin{array}{c} B \\ \vdash \\ A \begin{array}{c} \text{---} \\ \vdash \end{array} \\ B \end{array} \\ &= \sum_{[B] \in \mathcal{L}} N_B^{X,Y} d_B \sum_{[A] \in \mathcal{L}} \omega_A \cdot \begin{array}{c} B \\ \vdash \\ A \begin{array}{c} \text{---} \\ \vdash \end{array} \\ B \end{array} \\ &= N_{\mathbf{1}}^{X,Y} d_1 = \begin{cases} 1, & X \cong Y^* \\ 0, & \text{otherwise.} \end{cases} \end{aligned}$$

In our second way of computing $h_{X,Y}$, we relate the string diagram trace to S -matrix values:

$$\begin{aligned} \text{tr} \left(A \begin{array}{c} X & Y \\ \vdash & \dashv \\ - & \end{array} \right) &= X^* \begin{array}{c} A \\ \text{---} \\ \vdash \end{array} Y^* \\ &= X^* \begin{array}{c} A \\ \text{---} \\ \vdash \end{array} \begin{array}{c} A^* \\ \text{---} \\ \vdash \end{array} Y^* \\ &= \frac{1}{d_A} S_{X,A} S_{Y,A} \end{aligned}$$

and thus, combining our two computations, we find

$$h_{X,Y} = \sum_{[A] \in \mathcal{L}} \frac{\omega_A}{d_A} S_{X,A} S_{Y,A} = \begin{cases} 1 & X \cong Y^* \\ 0 & \text{otherwise.} \end{cases}$$

We now define the vector

$$\omega^{(X)} = \sum_{[A] \in \mathcal{L}} \frac{\omega_A}{d_A} S_{A,X^*} |[A]\rangle$$

for all simple objects $X \in \mathcal{C}$. We compute

$$\begin{aligned} \mathcal{S}(\omega^{(X)}) &= \sum_{[A] \in \mathcal{L}} \frac{\omega_A}{d_A} S_{A,X^*} \left(\sum_{[Y] \in \mathcal{L}} \frac{1}{d_Y} S_{A,Y} |[Y]\rangle \right) \\ &= \sum_{[Y] \in \mathcal{L}} \frac{1}{d_Y} \left(\sum_{[A] \in \mathcal{L}} \frac{\omega_A}{d_A} S_{A,X^*} S_{A,Y} \right) |[Y]\rangle \\ &= \sum_{[Y] \in \mathcal{L}} \frac{h_{X^*,Y}}{d_Y} |[Y]\rangle = \frac{1}{d_X} |[X]\rangle. \end{aligned}$$

Hence $|[X]\rangle$ is in the image of \mathcal{S} for $[X] \in \mathcal{L}$, as desired. \square

5.5.4 Verlinde formula

In this section, we prove the *Verlinde formula*. This formula was first conjectured by Verlinde [Ver88], and proven the following year by Moore-Seiberg [MS89]. There are now many Verlinde-type formulas. Most importantly, there is one for vertex operator algebras [Hua08] and one in algebraic geometry [Fal94]. With proposition [ref] in hand, the proof is very quick:

Theorem 5.4 (Verlinde formula). *Let \mathcal{C} be a modular category.*

(i) *For all simple objects $A, B, C \in \mathcal{C}$,*

$$N_C^{A,B} = \sum_{[E] \in \mathcal{L}} \frac{S_{A,E} S_{B,E} (S^{-1})_{C,E}}{d_E}$$

where $(S^{-1})_{C,E}$ denotes the (C, E) -coefficient of the inverse of the S matrix.

(ii) *For all simple objects $A \in \mathcal{C}$, the matrix*

$$D^A = S N^A S^{-1}$$

is diagonal with $([B], [B])$ -entry $S_{A,B}/d_B$, where $N^A = (N_C^{A,B})_{([B], [C]) \in \mathcal{L}^2}$ is the fusion matrix of A .

Proof. We begin by proving part (ii). The main observation of the proof is that the operator $N^A : \mathbb{C}[\mathcal{L}] \rightarrow \mathbb{C}[\mathcal{L}]$ is exactly left multiplication by A in $K_{\mathbb{C}}(\mathcal{C})$. Proposition [ref] says that \mathcal{S} is a morphism of algebras, and hence we can commute N^A past \mathcal{S} , and turn it into multiplication by A in $\mathbb{C}[\mathcal{L}]^{\text{func}}$. Hence, using the appropriate multiplication in the appropriate algebra, we find

$$\begin{aligned}
(\mathcal{S}N^A\mathcal{S}^{-1})|[B]\rangle &= \mathcal{S}(|[A]\rangle \cdot_{K_{\mathbb{C}}(\mathcal{C})} \mathcal{S}^{-1}(|[B]\rangle)) \\
&= \mathcal{S}(|A\rangle) \cdot_{\mathbb{C}[\mathcal{L}]^{\text{func}}} |[B]\rangle \\
&= \frac{S_{A,B}}{d_B} |[B]\rangle.
\end{aligned}$$

Thus, $\mathcal{S}N^A\mathcal{S}^{-1}$ is diagonal with $([B], [B])$ entry $S_{A,B}/d_B$. Scaling rows of \mathcal{S} does not change the effect of diagonalization. Hence, we conclude that $SN^A\mathcal{S}^{-1}$ is diagonal as well, with the same entries, and thus our proof of (ii) is complete.

We now move on to proving part (i). Expanding the formula $N^A = S^{-1}D^AS$, we find

$$\begin{aligned}
N^A |[B]\rangle &= S^{-1}D^AS |[B]\rangle \\
&= S^{-1} \left(\sum_{[E] \in \mathcal{L}} \frac{S_{A,E}S_{B,E}}{d_E} |[E]\rangle \right) \\
&= \sum_{[E] \in \mathcal{L}} \frac{S_{A,E}S_{B,E}}{d_E} \left(\sum_{[C] \in \mathcal{L}} (S^{-1})_{C,E} |[C]\rangle \right).
\end{aligned}$$

Comparing coefficients with the definition of N^A , we conclude the result. \square

5.5.5 Proof of modularity

In this section we prove that the S -matrix and T -matrix indeed give a representation of the modular group. That is, we will prove Theorem [ref]. At its heart, the fact that the modular representation of modular category is a homomorphisms comes down to proving a series of relations between the coefficients of the S -matrix and the coefficients of the T -matrix. That is, we are proving a series of relations between braiding and twisting. The general method is to take traces of certain diagrams, and then compute those traces in two ways. One way will involve more twists and the other will involve more braiding. This will give some algebraic relation, and choosing the right diagrams we will get enough algebraic relations to deduce Theorem [ref].

We begin with the most fundamental relationship between S -matrix and T -matrix entries:

Lemma 5.4. *Let \mathcal{C} be a pre-modular category. We have that*

$$S_{A,B} = \theta_A^{-1}\theta_B^{-1} \sum_{[C] \in \mathcal{L}} N_C^{A,B} \theta_C d_C$$

and

$$S_{A^*,B} = \theta_A \theta_B \sum_{[C] \in \mathcal{L}} N_C^{A,B} \theta_C^{-1} d_C.$$

Proof. By Proposition [ref], we have $\beta_{B,A} \circ \beta_{A,B} = (\theta_A^{-1} \otimes \theta_B^{-1}) \circ \theta_{A \otimes B}$. Taking the trace of this formula we get

$$S_{A,B} = \theta_A^{-1} \theta_B^{-1} \operatorname{tr}(\theta_{A \otimes B}).$$

Now, since θ is a natural transformation it splits over direct sums. Traces split over direct sums as well by proposition [ref] and hence $\operatorname{tr}(\theta_{A \otimes B}) = \sum_{[C] \in \mathcal{L}} N_C^{A,B} \theta_C d_C$. This concludes the proof of the first formula.

For the second formula, we observe that replacing overcrossings with undercrossings in the definition of $S_{A,B}$ has the effect of taking the dual of one of elements, replacing it with $S_{A^*,B}$. Hence $\operatorname{tr}(\beta_{A,B}^{-1} \circ \beta_{B,A}^{-1}) = S_{A^*,B}$. Thus, taking the trace of the formula $\beta_{A,B}^{-1} \circ \beta_{B,A}^{-1} = \theta_{A \otimes B}^{-1} \circ (\theta_A \otimes \theta_B)$ yields the desired result. \square

Before continuing to our proof of theorem [ref] we observe a key lemma:

Lemma 5.5. *Let \mathcal{C} be a pre-modular category. Let $A \in \mathcal{C}$ be a (possibly non-simple) object. We have that*

$$\sum_{[B] \in \mathcal{L}} d_B \theta_B \cdot \left(\begin{array}{c} A \\ | \\ B \text{ } \bigcirc \\ | \\ A \end{array} \right) = p_{\mathcal{C}}^+ \cdot \boxed{\theta_A^{-1}} \quad \begin{array}{c} A \\ | \\ \boxed{\theta_A^{-1}} \\ | \\ A \end{array}$$

and

$$\sum_{[B] \in \mathcal{L}} d_B \theta_B^{-1} \cdot \left(\begin{array}{c} A \\ | \\ B \text{ } \bigcirc \\ | \\ A \end{array} \right) = p_{\mathcal{C}}^- \cdot \boxed{\theta_A} \quad \begin{array}{c} A \\ | \\ \boxed{\theta_A} \\ | \\ A \end{array}$$

Proof. We only prove the first formula - the second follows by a formally dual argument. We restrict to the case that A is simple. Seeing as both sides are linear with respect to direct sums, the case that A is simple will immediately imply the general case. Since A is simple, it suffices to prove that the traces of both sides are equal. The trace on the left hand side has the effect of replacing the diagram with $S_{A,B}$. Hence, we compute as follows using Lemma [ref] and the fact that $\sum_{[C] \in \mathcal{L}} N_C^{A,B} d_B = d_A d_C$ from proposition [ref]:

$$\begin{aligned} \sum_{[B] \in \mathcal{L}} d_B \theta_B S_{A,B} &= \sum_{[B] \in \mathcal{L}} d_B \theta_B \left(\theta_A^{-1} \theta_B^{-1} \sum_{[C] \in \mathcal{L}} N_C^{A,B} \theta_C d_C \right) \\ &= \theta_A^{-1} \sum_{[C] \in \mathcal{L}} \theta_C d_C \left(\sum_{[B] \in \mathcal{L}} N_C^{A,B} d_B \right) \\ &= \theta_A^{-1} d_A \sum_{[C] \in \mathcal{L}} \theta_C d_C^2 = p_{\mathcal{C}}^+ \theta_A^{-1} d_A. \end{aligned}$$

This result is exact the trace of the right hand sides of the lemma. Hence, the proof is complete. \square

We now give the heart of the proof of theorem [ref]:

Theorem 5.5. *Let \mathcal{C} be a pre-modular category. Define the charge conjugation operator $\check{C} : \mathbb{C}[\mathcal{L}] \rightarrow \mathbb{C}[\mathcal{L}]$ to the matrix with $([A], [A^*])$ coefficient 1 for all $[A] \in \mathcal{L}$, and all other coefficients zero.*

$$(i) \check{C}S = S\check{C}, \check{C}T = T\check{C}, \text{ and } \check{C}^2 = 1;$$

$$(ii) (ST)^3 = p_{\mathcal{C}}^+ S^2;$$

$$(iii) (ST^{-1})^3 = p_{\mathcal{C}}^- S^2 \check{C}.$$

If \mathcal{C} is modular, then

$$(iv) S^2 = p_{\mathcal{C}}^+ p_{\mathcal{C}}^- \check{C}.$$

Proof. Part (i) follows from the fact that $S_{A^*, B^*} = S_{A, B}$, $\theta_{A^*} = \theta_A$, and $A^{**} \cong A$. Parts (ii) and (iii) have formally dual proofs, which arise from replacing θ with θ^{-1} at every opportunity. Part (iv) follows algebraically from combining formulas (ii) and (iii) whenever S is invertible, which is always the case when \mathcal{C} is modular by theorem [ref].

Hence, it suffices to prove part (ii). The proof comes from computing the quantity

$$h_{X,Y} = \sum_{[A] \in \mathcal{L}} d_A \theta_A \cdot \text{tr} \left(A \begin{array}{c|c} X & Y \\ \hline | & | \\ - & - \\ | & | \\ X & Y \end{array} \right)$$

two ways.

In the first way of computing $h_{X,Y}$, we use lemma [ref]. We find the following:

$$h_{X,Y} = p_{\mathcal{C}}^+ \cdot \text{tr} \left(\begin{array}{c|c} X & Y \\ \hline \boxed{\theta_{X \otimes Y}^{-1}} & \\ X & Y \end{array} \right) = p_{\mathcal{C}}^+ \theta_X \theta_Y S_{X^*, Y}$$

In our second way of computing $h_{X,Y}$, we use computation of the trace of two lines through a loop in the proof of Theorem [ref]. We find this way that

$$h_{X,Y} = \sum_{[A] \in \mathcal{L}} \theta_A S_{X,A} S_{Y,A}.$$

Thus, we find that $\sum_{[A] \in \mathcal{L}} \theta_A S_{X,A} S_{Y,A} = p_{\mathcal{C}}^+ \theta_X \theta_Y S_{X^*, Y}$. Thinking of these quantities as the $([X], [Y])$ entries in operators $\mathbb{C}[\mathcal{L}] \rightarrow \mathbb{C}[\mathcal{L}]$, we get the equation

$$STS = p_{\mathcal{C}}^+ T S T \check{C}.$$

.[WORK: This formula is WRONG. It should be

$$STS = p_{\mathcal{C}}^+ T^{-1} S T^{-1}.$$

From this we get

$$(ST)^3 = p_{\mathcal{C}}^+ S^2$$

as desired. I'm not sure where I went wrong, but something is off in here.] \square

We now know that the S and T matrices give a modular representation *up to phase!* We still need to work out the details of the phases. A key part is the following computation:

Corollary 5.6. *Let \mathcal{C} be a modular category. The quantities $p_{\mathcal{C}}^+$ and $p_{\mathcal{C}}^-$ are nonzero, and*

$$p_{\mathcal{C}}^+ p_{\mathcal{C}}^- = \mathcal{D}^2.$$

Proof. The values $p_{\mathcal{C}}^+$ and $p_{\mathcal{C}}^-$ must be nonzero because S is invertible and $S^2 = p_{\mathcal{C}}^+ p_{\mathcal{C}}^- \check{C}$. The formula $S^2 = p_{\mathcal{C}}^+ p_{\mathcal{C}}^- \check{C}$, when expanded, says that

$$\sum_{[C] \in \mathcal{L}} S_{C,A} S_{C,B} = \begin{cases} p_{\mathcal{C}}^+ p_{\mathcal{C}}^- & A \cong B^* \\ 0 & \text{otherwise.} \end{cases}$$

Applying this formula to $A = B = \mathbf{1}$, we find

$$\sum_{[C] \in \mathcal{L}} S_{C,\mathbf{1}} S_{C,\mathbf{1}} = \sum_{[C] \in \mathcal{L}} d_C^2 = p_{\mathcal{C}}^+ p_{\mathcal{C}}^-$$

as desired. \square

Now, it is clear that we can normalize the representation appropriately and conclude theorem [ref], as desired.

5.5.6 Vafa's theorem, unitarity of S -matrix, and the Chiral central charge

In this section we discuss some finer points of the structure of the modular representation. In particular, we will prove that modular representation of every modular category is *unitary*. That is, the S and T matrices are both unitary operators on $\mathbb{C}[\mathcal{L}]$ when it is endowed with its canonical inner product coming from its basis

We begin with the matrix T . For a diagonal matrix to be unitary, it is necessary and sufficient for its diagonal entries to have absolute value 1. We will prove something even stronger: that all of the entries are roots of unity! We recall that a number $z \in \mathcal{C}$ is called a root of unity if $z^n = 1$ for some integer $n \geq 1$. We begin with a key topological lemma which will underscore our proof:

Lemma 5.6 (Lantern identity). *Let \mathcal{C} be a pre-modular category. Let $A, B, C \in \mathcal{C}$ be objects. As maps $A \otimes B \otimes C \rightarrow A \otimes B \otimes C$, we have the identity*

$$\theta_{A \otimes B} \circ \theta_{A \otimes C} \circ \theta_{B \otimes C} = \theta_{A \otimes B \otimes C} \circ (\theta_A \otimes \theta_B \otimes \theta_C).$$

Proof. In the language of string diagrams, this formula becomes

[WORK: add diagram.]

It is a matter of elementary manipulations to convince one's self that these two diagrams are equal. An alternate algebraic approach is to expand the relation both sides using the formula $\theta_{X \otimes Y} = (\beta_{Y,X} \circ \beta_{X,Y}) \circ (\theta_X \otimes \theta_Y)$, cancel $\theta_A^2 \otimes \theta_B^2 \otimes \theta_C^2$, and compare the resulting braids using the hexagon and naturality. \square

We state one more linear-algebraic lemma necessary for the proof:

Lemma 5.7. Let $n \geq 1$ be an integer and let $M : \mathbb{R}^n \rightarrow \mathbb{R}^n$ be a linear operator whose (a, b) entry is $M_{a,b}$. Suppose that all off diagonal entries of M are positive, and that $\sum_{b=1}^n M_{a,b} < 0$ for all $1 \leq a \leq n$. Then, M is invertible.

Proof. .[WORK: This follows from the Gershgorin circle theorem, but that's way to general for our purposes. Find a nice proof which is appropriate for the situation.] \square

Theorem 5.6 (Vafa). Let \mathcal{C} be a pre-modular category. The values θ_A are roots of unity for all simple object $A \in \mathcal{C}$.

Remark. This theorem is interesting in part due to its history. [WORK: add history.]

Proof. To begin, we choose a simple object A . We observe that every endomorphism of $A \otimes A^* \otimes A$ induces a linear map

$$\text{Hom}(A, A \otimes A^* \otimes A) \rightarrow \text{Hom}(A, A \otimes A^* \otimes A)$$

by postcomposition. We now consider the lantern identity (proposition [ref]) on the stands A, A^*, A ,

$$\theta_{A \otimes A^*} \circ \theta_{A \otimes A} \circ \theta_{A^* \otimes A} = \theta_{A \otimes A^* \otimes A} \circ (\theta_A \otimes \theta_{A^*} \otimes \theta_A).$$

viewed as an equality of linear operators on $\text{Hom}(A, A \otimes A^* \otimes A)$. We compute the determinant of both sides. We first compute the determinant of $\theta_{A \otimes A^*}$. The eigenvalues of the operator $\theta_{A \otimes A^*}$ are the twists θ_B . The dimension of the θ_B eigenspace is exactly the dimension of the subspace of $\text{Hom}(A, A \otimes A^* \otimes A)$ in which $A \otimes A^*$ fuse to B . The dimension of this space is $N_B^{A,A^*} N_A^{B,A} = (N_B^{A,A^*})^2$. The determinant is equal to the product of the eigenvalues counted with multiplicity, and hence

$$\det \theta_{A \otimes A^*} = \prod_{[B] \in \mathcal{L}} \theta_B^{(N_B^{A,A^*})^2}.$$

Continuing this way for all the other twists and plugging them into the lantern identity we get

$$\prod_B \theta_B^{2(N_B^{A,A^*})^2 + (N_B^{A,A})^2} = \theta_A^{4 \cdot \dim \text{Hom}(A, A \otimes A^* \otimes A)}.$$

We now define the coefficients

$$M_{A,B} = \begin{cases} 2(N_B^{A,A^*})^2 + (N_B^{A,A})^2 & A \not\cong B \\ 2(N_B^{A,A^*})^2 + (N_B^{A,A})^2 - 4 \cdot \dim \text{Hom}(A, A \otimes A^* \otimes A) & A \cong B. \end{cases}$$

We observe that all of the off-diagonal entries of $M_{A,B}$ are positive and

$$\begin{aligned} \sum_{[B] \in \mathcal{L}} M_{A,B} &= \sum_{[B] \in \mathcal{L}} 2(N_B^{A,A^*})^2 + (N_B^{A,A})^2 - 4 \cdot \dim \text{Hom}(A, A \otimes A^* \otimes A) \\ &= -\dim \text{Hom}(A, A \otimes A^* \otimes A) < 0. \end{aligned}$$

Thus, by lemma [ref], we conclude that the matrix $M = (M_{A,B})_{([A],[B]) \in \mathcal{L}^2} : \mathbb{C}[\mathcal{L}] \rightarrow \mathbb{C}[\mathcal{L}]$ is invertible. Let \tilde{M} be the adjugate of M . That is, an integer valued matrix such that $M \cdot \tilde{M} = \tilde{M} \cdot M = n$ where $n = \det M$. We find for all simple objects A that

$$\theta_A^n = \prod_{[C] \in \mathcal{L}} \left(\prod_{[B] \in \mathcal{L}} \theta_B^{M_{C,B}} \right)^{\tilde{M}_{A,C}} = 1,$$

and hence θ_A is a root of unity as desired. \square

We immediately get several corollaries from this theorem. The first is that braiding enough times gets you back where you started:

Corollary 5.7. *Let \mathcal{C} be a pre-modular category. There exists an integer $n \geq 1$ such that*

$$(\beta_{B,A} \circ \beta_{A,B})^n = \text{id}_{A \otimes B}$$

for all $A, B \in \mathcal{C}$

Proof. Choose $n \geq 1$ so that $\theta_C^n = 1$ for all simple objects $C \in \mathcal{C}$, which exists by Vafa's theorem (theorem [ref]). Seeing that $\beta_{B,A} \circ \beta_{A,B} = \theta_{A \otimes B} \circ (\theta_A^{-1} \otimes \theta_B^{-1})$, in the direct sum decomposition $A \otimes B \cong \bigoplus_{[C] \in \mathcal{L}} N_C^{A,B} |[C]\rangle$ the transformation $\beta_{B,A} \circ \beta_{A,B}$ acts by the scalar $\theta_C / (\theta_A \theta_B)$ on every $|[C]\rangle$ summand. Thus, $(\beta_{B,A} \circ \beta_{A,B})^n$ acts by $(\theta_C / (\theta_A \theta_B))^n = 1$ and thus $(\beta_{B,A} \circ \beta_{A,B})^n$ is the identity as desired. \square

Corollary 5.8. *Let \mathcal{C} be a modular category. The quantity $p_{\mathcal{C}}^-/p_{\mathcal{C}}^+$ is a root of unity.*

Proof. We take determinants. From $S^2 = p_{\mathcal{C}}^+ p_{\mathcal{C}}^- \check{C}$ we find that $\det(S)^2 = \pm p_{\mathcal{C}}^+ p_{\mathcal{C}}^-$, since $\det \check{C} = \pm 1$. From the formula $(ST)^3 = p_{\mathcal{C}}^+ S^2$ we find that

$$(p_C^+ / \det(S))^2 = \det(T)^6$$

so $p_C^+ / p_C^- = \pm \det(T)^6$. By Vafa's theorem $\det(T)$ is a root of unity. Hence, we conclude that p_C^+ / p_C^- is a root of unity as desired. \square

We take a moment to observe that the quantity $p_{\mathcal{C}}^-/p_{\mathcal{C}}^+$ is of great interest to physicists. It is related to the *chiral central charge* c_{top} of the theory by the formula

$$p_{\mathcal{C}}^-/p_{\mathcal{C}}^+ = e^{-\frac{2\pi i c_{\text{top}}}{4}}.$$

[WORK: probably say a bit more? Maybe there will be a better spot somewhere else.]

We now move on to proving that the S matrix is unitary. The main technical ingredient is as follows:

Proposition 5.16. *Let \mathcal{C} be a modular category. For all simple objects $A, B \in \mathcal{C}$ we have that $S_{A^*,B} = \overline{S_{A,B}}$.*

Proof. By the Verlinde formula [ref], we know that for all simple objects A there exists a vector $\mathbf{v}_B \in \mathbb{C}[\mathcal{L}]$ such that

$$N^A \mathbf{v}_B = \frac{S_{A,B}}{d_B} \mathbf{v}_B$$

for all simple objects A . Namely, \mathbf{v}_B is the $[B]$ -column of the S matrix. Let \mathbf{v}_B^* be the row vector which is the Hermitian adjoint to \mathbf{v}_B . We have that

$$\mathbf{v}_B N^A \mathbf{v}_B = \frac{S_{A,B}}{d_B} |\mathbf{v}_B|^2.$$

Now, we observe that by Frobenius reciprocity (proposition [ref]) $(N^A)^\dagger = N^{A^*}$. Hence,

$$\begin{aligned} \mathbf{v}_B N^A \mathbf{v}_B &= \left(N^{A^*} \mathbf{v}_B \right)^* \mathbf{v}_B \\ &= \left(\frac{S_{A^*,B}}{d_B} \mathbf{v}_B \right)^* \mathbf{v}_B \\ &= \overline{\frac{S_{A^*,B}}{d_B}} |\mathbf{v}_B|^2. \end{aligned}$$

Comparing, we get the desired result. \square

We now get the following theorem:

Theorem 5.7 (Etingof-Nikshych-Ostrik). *Every matrix in the image of the modular representation $\mathrm{SL}_2(\mathbb{Z}) \rightarrow \mathrm{Aut}(\mathbb{C}[\mathcal{L}])$ is a unitary operator on $\mathbb{C}[\mathcal{L}]$.*

Proof. We already know by theorem [ref] and corollary [ref] that the normalized T -matrix is unitary. It thus suffices to show that the modular S matrix is unitary. From theorem [ref] we have that $(\frac{1}{D} S)^{-1} = (\frac{1}{D} S) \cdot \check{C}$. By proposition [ref] $(\frac{1}{D} S) \cdot \check{C} = (\frac{1}{D} S)^\dagger$ and thus $\frac{1}{D} S$ is unitary as desired. \square

5.6 Skeletonization

5.6.1 Principle

.[WORK: lots of big choices need to be made here. Do I call this the skeletonization, or do I call it something else? Do I work with multiplicity-free categories, or do I allow multiplicity? I don't know what the correct statements are or what the proofs look like so this section might be a tough one.]

A good mathematical reference is [DHW13].]

5.6.2 F -symbols

5.6.3 R -symbols

5.6.4 θ -symbols

5.6.5 Reconstruction theorem

5.7 Quantum double modular categories

5.7.1 The Drinfeld center

A quantum double modular category is a special type of modular category. They are particularly important because many of the constructions of topological order only deal

with quantum double modular categories. For instance, there are constructions of modular categories/topological order coming from the theory of tensor networks [ref], subfactors [ref], vertex operator algebras [ref], [WORK: add more sources]. All of these constructions only give quantum double modular categories. Hence, understanding quantum doubles is key to understanding how topological order work in practice.

At the heart of quantum doubles is a construction known as the *Drinfeld center*. In its most basic form the Drinfeld center induces an assignment

$$\mathcal{Z} : (\text{monoidal categories}) \rightarrow (\text{braided monoidal categories}).$$

In our context, we care about a more structured version of the Drinfeld center. It is a theorem of Muger that the Drinfeld center induces an assignment as follows:

$$\mathcal{Z} : (\text{spherical fusion categories}) \rightarrow (\text{modular categories}).$$

This theorem is fantastic because it allows one to construct modular categories using much less data than would otherwise be necessary. Without needing a braiding, non-degenerate or otherwise, the Drinfeld center allows one to construct a modular category. This makes the Drinfeld center an abundant source of modular categories. We call an modular category \mathcal{C} a quantum double if it is of the form $\mathcal{Z}(\mathcal{C}_0)$ for some spherical fusion category \mathcal{C}_0 . A major goal of this chapter is to set up and prove Muger's theorem

We now define the Drinfeld center. The Drinfeld center is a somewhat direct categorification of the usual notion of center for finite groups. If G is a finite group, its center is defined as follows:

$$Z(G) = \{g \in G \mid gh = hg \ \forall h \in G\}.$$

The first guess at $Z(\mathcal{C})$ is thus

$$\mathcal{Z}(\mathcal{C}) = \{A \in \mathcal{C} \mid A \otimes B \cong B \otimes A \ \forall B \in \mathcal{C}\}.$$

This is almost correct, but not quite. The issue is that $\mathcal{Z}(\mathcal{C})$ is not quite a braided monoidal category yet. Even though $A \otimes B \cong B \otimes A$ for all $A, B \in \mathcal{Z}(\mathcal{C})$, we don't have a distinguished choice of isomorphism. A braided monoidal category requires a distinguished isomorphism $\beta_{A,B} : A \otimes B \xrightarrow{\sim} B \otimes A$. Moreover, these distinguished isomorphisms are required to satisfy the hexagon equations. Hence, we make a new definition of center which keeps track of the choice of isomorphism and enforces the hexagon equation along the way:

Proposition 5.17. *The Drinfeld center $\mathcal{Z}(\mathcal{C})$ of a monoidal category \mathcal{C} is a braided monoidal category defined as follows:*

- (Objects) Pairs $(A, \beta_{A,-})$, where $A \in \mathcal{C}$, and $\beta_{A,-}$ is a natural isomorphism of monoidal natural isomorphism between the two functors $A \otimes -$ and $- \otimes A$ from \mathcal{C} to \mathcal{C} , satisfying the additional condition that

$$\beta_{A,B \otimes C} = (\text{id}_B \otimes \beta_{A,C}) \circ (\beta_{A,B} \otimes \text{id}_C).$$

- (Morphisms) Given $(A, \beta_{A,-}), (B, \beta_{B,-}) \in \mathcal{Z}(\mathcal{C})$, $\text{Hom}_{\mathcal{Z}(\mathcal{C})}((A, \beta_{A,-}), (B, \beta_{B,-}))$ is the subspace of morphisms $f \in \text{Hom}_{\mathcal{C}}(A, B)$ such that for all $C \in \mathcal{C}$

$$(\text{id}_C \otimes f) \circ \beta_{A,C} = \beta_{B,C} \circ (f \otimes \text{id}_C).$$

- (*Tensor product*) Given $(A, \beta_{A,-}), (B, \beta_{B,-}) \in \mathcal{Z}(\mathcal{C})$, we define

$$(A, \beta_{A,-}) \otimes (B, \beta_{B,-}) = (A \otimes B, (\beta_{A,-} \otimes \text{id}_{\mathcal{C}}) \circ (\text{id}_{\mathcal{C}} \otimes \beta_{B,-})).$$

- (*Unit*) The element $(1, \rho \circ \lambda^{-1})$
- (*Braiding*) We define the braiding between two elements $(A, \beta_{A,-}), (B, \beta_{B,-}) \in \mathcal{Z}(\mathcal{C})$ to be $\beta_{A,B} = \beta_{A,B}$.

Inheriting associativity, unitors, and composition from \mathcal{C} , this gives $\mathcal{Z}(\mathcal{C})$ the structure of a braided monoidal category.

Proof. Since morphisms in $\mathcal{Z}(\mathcal{C})$ are a subspace of morphisms in \mathcal{C} , commutative diagrams don't change when going from \mathcal{C} to $\mathcal{Z}(\mathcal{C})$. Hence, the triangle and pentagon axioms for $\mathcal{Z}(\mathcal{C})$ follow immediately from the triangle and pentagon axioms on \mathcal{C} . One thing to be checked is that evaluation/co-evaluaiton satisfy the compatibility condition required to a morphism in $\mathcal{Z}(\mathcal{C})$, but this is straightforward. We remark on the hexagon identities. The condition imposed on $\beta_{A,B \otimes C}$ given is technically incorrect. To make the parentheses work in the braiding one has to add associators, and impose the longer condition

$$\beta_{A,B \otimes C} = \alpha_{C,A,B}^{-1} \circ (\text{id}_B \otimes \beta_{A,C}) \circ \alpha_{A,C,B} \circ (\beta_{A,B} \otimes \text{id}_C) \circ \alpha_{A,B,C}^{-1}.$$

This condition makes the second hexagon identity tautological. Similarly, the definition of tensor product given is not strictly correct - one must add the correct associator terms, making the first hexagon identity immediate. Lastly one must verify the half-braidings defined on the tensor unit/tensor product are actually half braidings, i.e., that they satisfy the hexagon condition. These follow from straightforward computations, which we leave as exercises. This completes the proof. □

5.7.2 Muger's theorem

[WORK: through Mueger's theorem. The exposition will greatly differ based on what the proof looks like, which I haven't done before.

The standard proof people use is completeley steeped in the language of module categories. The original proof uses the tube algebra, and is “elementary” in the sense that you don't need tube algebras but it is heineous, and I would very much like to avoid it. I'll figure out what to write hence once I've done the module category section and I've digested the proof. Maybe there's a way to de-module-ify it, but maybe not. Maybe I don't want to do that if it'll lose its essence.]

5.7.3 Discrete gauge theory as a quantum double and Morita equivalence

We saw in chapter [ref] that \mathbf{Vec}_G and $\text{Rep}(G)$ are both naturally spherical fusion categories. Thus, Muger's theorem tells us that $\mathcal{Z}(\mathbf{Vec}_G)$ and $\mathcal{Z}(\text{Rep}(G))$ are both modular categories. Hence, given a finite group G we have three different modular categories we can associate to it: $\mathfrak{D}(G)$, $\mathcal{Z}(\mathbf{Vec}_G)$, $\mathcal{Z}(\text{Rep}(G))$. The amazing fact is that these are all the same category:

Proposition 5.18. *Let G be a finite group. There are equivalences of modular categories $\mathfrak{D}(G) \cong \mathcal{Z}(\mathbf{Vec}_G) \cong \mathcal{Z}(\text{Rep}(G))$.*

Proof. .[WORK: do proof. Shouldn't be too hard.] \square

We now make a few comments about this theorem. The first is that it proves that $\mathfrak{D}(G)$ is a quantum double modular category. Secondly, it gives a second proof that $\mathfrak{D}(G)$ has a non-degenerate braiding, using Muger's theorem. Thirdly, it demonstrates the concept of *Morita equivalence*.

[WORK: introduce Morita equivalence. I know that there's some important basic facts to tell - I should include those.]

5.7.4 Factorizability and time reversal symmetry

Given a modular category \mathcal{C} , we can forget the braiding on \mathcal{C} and only remember its structure as a spherical fusion category. Hence, Muger's theorem tells us that $\mathcal{Z}(\mathcal{C})$ is canonically a modular category. It is a fantastic fact that in this case $\mathcal{Z}(\mathcal{C})$ can be explicitly computed in terms of \mathcal{C} . We describe this computation now.

[WORK: Define the time-reversal conjugate $\overline{\mathcal{C}}$. Setup the map $\mathcal{C} \boxtimes \overline{\mathcal{C}} \rightarrow \mathcal{Z}(\mathcal{C})$]

Proposition 5.19. *Let \mathcal{C} be a pre-modular category. The canonical map $\mathcal{C} \boxtimes \overline{\mathcal{C}} \rightarrow \mathcal{Z}(\mathcal{C})$ is an equivalence of categories if and only if \mathcal{C} is modular.*

Proof. .[WORK: do proof] \square

This theorem is fantastic because it not only computes $\mathcal{Z}(\mathcal{C})$ for every modular category \mathcal{C} , but also it gives an equivalent definition of modularity. This gives us our third definition of modularity. Namely a pre-modular category \mathcal{C} is modular if and only if its braidings are all non-degenerate, or equivalently if its S -matrix is non-degenerate, or equivalently if it is factorizable in the above sense.

5.7.5 Levin-Wen model

[WORK: work though the Levin-Wen model.

I think that this model is fantastic because it shows how all of the ideas of tensor category theory can manifest themselves extremely concretely on the level of gapped Hamiltonians. Namely, the coherence relations on the category theory side correspond exactly to the formulas needed to make terms in a Hamiltonian commute with one another. It would be nice if I could give a motivation for which the category which describes the Levin-Wen model is the Drinfeld center, though I've never seen that before.]

5.8 Unitarity

5.8.1 Characterization of unitarizability

[WORK: An early reference about unitary MTCs is [KJ96]. I should read this.]

As we've mentioned before, the algebraic theory of topological order is not modular category theory but *unitary* modular category theory because hom-spaces need Hilbert space structure to define valid quantum systems. We define unitary fusion categories and unitary modular categories in section [ref].

We've mostly ignored unitary fusion categories for the following reasons: their additional structure is a large semantic burden, and is in a way inessential. We mean this in the following sense:

1. Given a spherical fusion category, \mathcal{C} every unitary structure on \mathcal{C} is equivalent. That is, all the unitary fusion categories obtained by enriching \mathcal{C} with unitary structure are equivalent to one another as unitary fusion categories;
2. All braidings on unitary fusion categories are automatically unitary so there is no subtlety in passing to unitary modular categories;
3. A spherical fusion category admits a unitary structure if and only if its quantum dimensions are all positive.

Still, it is worthwhile to study unitary fusion categories. They are the correct algebraic structure to describe topological order, so understanding *why* points (1-3) above are true gives insight into unitarity in practice. There is also a characterization of unitary fusion categories in the skeleized perspective which is of much utility in practice. Namely, a fusion category admits a unitary structure if and only if its F -matrices and R -matrices can all be made unitary. We discuss this more and offer a proof in section [ref].

We now prove the first main result of the section:

Proposition 5.20. *Let \mathcal{C} be a spherical fusion category. There is a compatible unitary structure on \mathcal{C} if and only if the quantum dimensions d_A are positive for all simple objects $A \in \mathcal{C}$.*

Proof. Suppose first that \mathcal{C} admits a unitary structure. Then, for all simple objects $A \in \mathcal{C}$,

$$d_A = \text{tr}(\text{id}_A) = \langle \text{id}_A | \text{id}_A \rangle > 0$$

because the inner product is positive definite. Conversely, suppose that all of the quantum dimensions of \mathcal{C} are positive. We will suppose for simplicity that \mathcal{C} is skeletal, which is possible by proposition [ref]. Suppose that $f : A \rightarrow B$ is a morphism between objects $A, B \in \mathcal{C}$. We know that there are direct sum decompositions $A = \bigoplus_i A_i$ and $B = \bigoplus_j B_j$. Writing f as an element of $\bigoplus_{i,j} \text{Hom}(A_i, B_j)$, we find that f acts by some scalar $f_{i,j}$ on each homspace $\text{Hom}(A_i, B_j)$, where $f_{i,j} = 0$ if $A \not\cong B$. We define $f^\dagger : B \rightarrow A$ to be the map whose $\text{Hom}(B_j, A_i)$ is $f_{j,i}^\dagger = \overline{f_{i,j}}$.

We now prove that this defines a unitary structure on \mathcal{C} . Given any two morphism $f : A \rightarrow B$, we compute

$$\langle f | f \rangle = \text{tr}(f^\dagger \circ f) = \sum_{i,j} |f_{i,j}|^2 d_{A_i}.$$

Clearly, $\langle f | f \rangle = 0$ if and only if $f = 0$ because all of the quantum dimensions d_{A_i} are positive real numbers. Hence, $\langle \cdot | \cdot \rangle$ is positive definite.

It is clear that $(f^\dagger)^\dagger = f$, $(f \circ g)^\dagger = g^\dagger \circ f^\dagger$, and $(f \otimes g)^\dagger = f^\dagger \otimes g^\dagger$ for appropriate choices of f, g . It remains to show that $(\text{ev}_A^R)^\dagger = \text{coev}_A^L$ and $(\text{coev}_A^R)^\dagger = \text{ev}_A^L$.

[WORK: I don't know how to do this part of the proof. It's something to come back to.] \square

5.8.2 Unitary braidings

[WORK: Show that every braiding is automatically unitary. This is the content of [Gal14].]

5.8.3 Uniqueness of unitary structure

[WORK: Show that unitary structures are unique. This is the content of [Reu23b].]

5.8.4 Skeletonization of unitarity

[WORK: Show that a fusion category admits a unitary structure if and only if its F and R symbols can be made unitary in some basis. There are good references for this in the papers cited above.]

5.9 Number theory in modular categories

5.9.1 Techniques and first results

[WORK:

A good general reference about this stuff is [GS19]. To-read, for sure.

There is also interesting work in [MS12] which gives some counterexamples.

Another reference is [DHW13].]

5.9.2 Galois conjugation

[WORK: An early reference is [CG93].

Some other papers not to ignore in this area are [PSYZ23, BR22].]

5.9.3 Ocneanu rigidity

5.9.4 Rank-finiteness theorem

[WORK: Of course there is the original paper on the topic. However, there is also the generalization for G -crossed MTCs and fermionic MTCs - [JMNR21]. Should I bring this up now or later?]

5.9.5 Schauenberg-Ng theorem

.[WORK: Go through Schauenberg-Ng's original paper and understand the proof.

It seems on the face of it like it is a hard theorem. Certainly, it uses strongly the theory of the Drinfeld center as well as the modular representation. It is good to push this proof as far down as possible since it will use a lot of machinery. I think it can be boiled down to something elegant, though, if the machinery has been set up.]

.[WORK: I would quite like to show that fusion categories (or at least modular categories) have finitely automorphisms of the identity functor. The proof I know follows from generalities about the universal grading group - [GN08]. I wonder if this proof can be done completely without the use of grading. Then, the fact that the universal grading group can be interpreted in terms of grading can be put as an exercise in the "further structure" section!]

[WORK: this section is going to host a lot more theorems]

History and further reading:

Modular categories were born from conformal field theory in the late 1980s. In a series of papers, Moore and Seiberg analysed deeply the underlying content within conformal field theory to find what essential algebraic data lied within it [MS88, MS89]. They wrote out the axioms of this essential algebraic data in their subsequent notes on conformal field theory [MS90]. They used the name modular category to describe their data, as suggested by Igor Frenkel. This definition was then refined and re-introduced by Turaev [Tur92]. The first major application of modular categories was the Reshetikhin-Turaev construction [RT91, Tur10b]. Prior to this result nobody had succeed in constructing topological quantum field theories. In this way, modular categories and the Reshetikhin-Turaev construction completed Witten's programme of quantizing Chern-Simons theory.

By the early 2000s, the proposal of topological quantum computing was attracting a lot of interest in anyons and their algebraic properties. Seeing as topological order can be described by topological quantum field theory and topological quantum field theory is essentially equivalent to modular categories, it was understood that modular categories could be used to understand topological order. This latent description of anyons in terms of modular categories was made explicit in an appendix in the seminal 2006 paper of Kitaev [Kit06]. This approach to anyons in terms of modular categories was popularized by Wang's early monograph [Wan10]. This has since become the standard approach towards the algebraic theory of topological quantum information.

Exercises:

- 5.1. .[WORK: apply Verlinde formula to group-theoretical modular categories to recover classical theorem by Burnside]
- 5.2. .[WORK: show that irreducible G -graded G -reps are equivalent to irreducible reps of centralizers of conjugacy classes]
- 5.3. . [WORK: One of the pivotal axioms automatically holds in fusion categories. Namely, the one with the twists going the two different ways. The proof is very standard - reduce to the case A, B are both simple, reduce to case $A = B$, $f = \lambda \cdot \text{id}_A$ and hence commutes with everything. This would be a nice exercise.]
- 5.4. .[WORK: there's a notion of prime factorization of MTCs. If \mathcal{C} is an MTC, then it can be decomposed as a tensor product

$$\mathcal{C} = \mathcal{C}_1 \boxtimes \mathcal{C}_2 \dots \boxtimes \mathcal{C}_n,$$

where $\mathcal{C}_i \neq 0$ all have no proper non-degenerate braided fusion subcategories other than $\mathbf{Vec}_{\mathbb{C}}$. We call such MTCs *prime MTCs*. Suppose that \mathcal{C} has NO abelian anyons. Then, the decomposition is unique. However, if \mathcal{C} has abelian anyons then the decomposition can fail to be unique. In particular, the color code is equivalent to both bilayer toric code and bilayer 3-fermion (c.f. [KPEB18]). The factorization results I asserted come from the (largely ignored?) paper by Muger: [Müg02].

This could make for a very nice exercise. First, prove Muger's double centralizer theorem. Then, prove the factorization into prime MTCs. Then, if \mathcal{C} has no abelian anyon anyons, prove that the decomposition is unique. Then, prove that bilayer toric code is equivalent to bilayer 3-fermion (or maybe just leave that as a comment)!]

6 Further structure

6.1 Overview

6.1.1 Introduction

We've seen this general picture throughout the book:

[WORK: add triangle.]

In this chapter we will talk about aspects of the algebraic theory of topological quantum information *beyond* this basic model. In particular, we will talk about physical phenomena beyond plain topological order. These phenomena are not just a physical or mathematical curiosity. They are central to almost all proposals for topological quantum computation, including the majority of the proposals we will present in chapter [ref].

We will show by example how easy it is to start running into phenomena beyond plain topological order. Suppose we want to make a topological quantum computer on some chip in the lab. Realistically, this chip is going to be *flat*. It's difficult and unpractical to make chips bend into spheres or tori for all sorts of reasons. The chip is also going to be finite. Hence, the chip is going to have *boundary*. This boundary and the way it is manipulated have serious and subtle impacts on the computation being performed. It cannot be ignored. Up to now, we have given no indication about how boundaries are to be dealt with.

[WORK: add picture of chip with boundary.]

These boundaries are not described by modular categories - we need additional algebraic structures to define them. In general, whenever we have any phenomenon beyond plain topological order we will need to introduce new algebraic structures.

In this chapter, we will discuss three different classes of important physical phenomena:

1. **Boundaries and domain walls.** In this section, we will talk about how topological phases can interact with one another. If you have one topological phase next to another, they can be physically separated by a *domain wall*:

[WORK: add picture of domain wall.]

If tuned correctly, this domain wall can be chosen so that the composite system keeps its topological gap. That is, we can make a composite topological system with two different sub-phases! The theory of domain walls includes the theory of boundaries as a special case, because a boundary can be viewed as a domain wall between a topological phase and empty space (which is itself a trivial topological phase).

2. **Symmetry enriched topological order.** [WORK: describe SETs]
3. **Fermionic topological order.** [WORK: describe fermionic TO.]

One of the fantastic features of the above structures is that they are all mathematically based on the same ideas: *module categories* and *graded categories*. The basic theory of module categories and graded categories will take us very far. Hence even though physically we will be stopping-and-starting every section with new phenomena, mathematically we will see a very smooth narrative.

6.2 Boundaries and domain walls

6.2.1 Physical picture

In this section we will talk about boundaries and domain walls in topological systems. To begin, we'll talk about boundaries in topological phases. We have the following general picture:

[WORK: topologically ordered system with boundary. DON'T use the categorical " \mathcal{C} ". Use the words "topologically ordered system". Keep with this convention for the whole section. I'm not assuming that the reader knows any category theory.]

This topologically ordered system with boundary will be described by some Hamiltonian. Within the bulk of the system the Hamiltonian would be the same as if the topological system had no boundary. On the boundary, however, it will necessarily be different. A-priori, we could choose our boundary terms in the Hamiltonian to behave arbitrarily. This will cause a problem, however. It can close the topological gap!

The main point of a topologically ordered system is that it has an energy gap in its spectrum between its lowest energy eigenstate and its next lowest energy eigenstate:

[WORK: add picture of gap.]

Adding the wrong terms on the boundary could close the gap. Namely, the size of the gap could tend to zero as the system size tends to infinity. To preserve the gap, the boundary terms in the Hamiltonian need to have a very special form. This special form implies a quite rigid structure on the space of gapped boundaries.

[WORK: Talk about how gapped boundaries can be equivalent to one another, they can be stable or unstable. When we say boundary theory, we mean *stable gapped boundary* theory. State that there are finitely many equivalence classes of boundary theories for a given topological order.

I don't really know how this equivalence works and how stability works. This is something to learn.]

Within a boundary there can be *boundary defects*. These are points in which the boundary theory changes from one theory to another, as shown below:

[WORK: add basic picture of a boundary defect.]

Just like with boundaries, it is important that the terms of the Hamiltonian around the boundary defect are chosen so that the topological gap is maintained. There can be stable and unstable boundary defects. There is a natural equivalence relation on the space of stable boundary defects, under which there are finitely many equivalence classes.

One important feature of boundary defects is that they can be *fused*. This happens when two defects are brought very close to one another, as shown below:

[WORK: add a picture for the fusion of boundary defects.]

This fusion can be nondeterministic. That is, repeating the fusion of two boundary defects can yield different results. In this sense boundaries can host non-abelian defects. Importantly, we will find that abelian bulk theories can host non-abelian boundary defects! This is the heart of our proposal for quantum computing with $\mathfrak{D}(\mathbb{Z}_2)$ topological order. Instead of working with the abelian anyons, we work with the nonabelian boundary defects!

More generally, in this section we will also discuss *domain walls*. Domain walls are boundaries between two different phases. That is, a domain wall is a one-dimensional line separating two-dimensional topologically ordered systems, as shown below:

[WORK: add picture.]

Domain walls include boundaries as a special case. Empty space can be viewed as a topologically ordered system. Namely, it is the topologically ordered space with a unique

ground state - emptiness! This is the trivial topological theory. A domain wall can be chosen between any topological theory and the empty theory. This is a boundary!

Of course, domain walls should be chosen to maintain the topological gap. There is a notion of stable and unstable domain walls. There is a notion of equivalence of stable domain walls. There are finitely many equivalence classes of stable domain walls between any two phases.

Within a domain wall there can be a boundary defect, just like before, which can fuse:

[WORK: add picture of boundary defect in domain wall, alongside a picture for fusion of defects.]

Domain walls also have the added structure of a notion of fusion. Given three topological phases and successive domain walls between them, bringing the domain walls will result in a new domain wall:

[WORK: fusion of domain walls.]

It is important to observe that just like how every boundary can be interpreted as a domain wall, we can interpret every domain wall in terms of a boundary. This goes as follows. Given a domain wall between two phases, we can *fold* the domain wall to get a boundary of the resulting bilayer theory:

[WORK: add folding picture.]

This establishes a bijection,

$$\begin{pmatrix} \text{domain walls between} \\ \text{phase 1 and phase 2} \end{pmatrix} \xrightarrow{\sim} \begin{pmatrix} \text{boundaries of the bilayer phase} \\ \text{obtained by stacking phase 1} \\ \text{and a flipped copy of phase 2} \end{pmatrix}.$$

This bijection is known as the *folding trick*. [WORK: Should I attribute it to Wen? People at the conference at the Perimeter institute were doing that. Kitaev and Kong are an early reference for the trick.]

In this section, we will be discussing the theory of domain walls, boundary defects, and their fusion.

6.2.2 Boundaries and domain walls in the toric code

[WORK: Do boundaries and domain walls in the toric code. The original reference is [BK98]. There's also a good discussion in "Topological quantum". I suppose I'll stay away from the surface code literature because they all do it for Bravyi's $\pi/4$ twisted version.]

6.2.3 The algebraic theory of gapped boundaries

In this subsection, we will take all of the principles discussed in section [ref] and interpret them in terms of category theory. This will give us an algebraic theory of gapped boundaries.

[WORK: Motivate why gapped boundaries correspond to Lagrangian algebras. As you approach a gapped boundary, every anyon is *confined* (become a boundary excitation) or *condensed* (disappear at the boundary). A nice talk about this perspective is [Bur18].

I still don't quite get how you get from this to Lagrangian algebras. What does the multiplication map mean, physically? Why does it need to be commutative? Why does it need to be separable? Why does it need to be connected? What's with the condition on dimensions?

Also, there's this quote by Burnell which is very mysterious to me:

"In a conventional ordering transition, the order parameter results from condensing a bosonic excitation— for example, the ordering of a magnet can be viewed as the Bose condensation of spinons."

What does it mean?

What do boundary defects correspond to? If the boundary defects go from one boundary type to the same boundary type then the answer is that it is the idempotent completion of the quotient of the original category by the algebra object. This quotient procedure is supposed to describe condensation of anyons at the domain wall. What is the general definition of boundary defect? I can't find it in Cong-Cheng-Wang.

I'll talk to Zhenghan and he will explain everything to me.]

[WORK:

- A discussion of anyon condensation at boundaries as a physical mechanism for gapped boundaries - smoothly get to Lagrangian algebras without needing to go through module categories or the Levin-Wen model!
- Prove that every module category is equivalent to the category of modules over some algebra (this is Ostriks original paper);
- Classification of irreducible module categories in terms of Lagrangian algebras - this is proposition 4.8 of [DMNO13].
- Characterization of Lagrangian algebras in terms of algebras with nice properties - this is due to Cong-Cheng-Wang;
- Proof that an MTC is doubled iff it contains a Lagrangian algebra
- Characterization of Lagrangian algebras in $\mathcal{Z}(\mathbf{Vec}_G)$;

]

[WORK:

There's the very important issue of what the ground state degeneracy of systems with boundaries and boundary defects is. What is a basis for these states? How does braiding act on them? All that stuff. This is done well in Cong-Cheng-Wang.

Another part of this is giving a treatment of topological charge measurement. What are observables? What are the associated probabilities?]

6.2.4 Lagrangian algebras and algebra objects

.[WORK: Give a nice mathematical treatment of Lagrangian algebras and algebra objects.]

6.2.5 Module categories and boundaries in the Levin-Wen model

[WORK: This subsection needs to be totally re-done. This is now an *alternate* perspective on gapped boundaries. I'm including it because the theory of module categories is necessary for the next section. I'm also including it because its nice to have multiple perspectives on this sort of stuff. I'm also including it because the theory of modules is very powerful. This might be the only reasonable way to get at Muger's theorem and Lagrange's theorem for fusion categories.]

In this subsection, we will take all of the principles discussed in section [ref] and interpret them in terms of category theory.

To begin, we re-iterate that not every topological phase admits a gapped boundary. In fact, we have the following physical principle [KKO⁺22]:

$$\left(\begin{array}{c} \text{the topological order described} \\ \text{by the modular category } \mathcal{B} \\ \text{admits a gapped boundary} \end{array} \right) \iff \left(\begin{array}{c} \text{there is an equivalence of categories} \\ \mathcal{B} \cong \mathcal{Z}(\mathcal{C}) \text{ for some} \\ \text{spherical fusion category } \mathcal{C} \end{array} \right).$$

For this reason, we will restrict our attention to *doubled* topological order for the rest of the section. That is, modular categories \mathcal{B} which are equivalent to $\mathcal{Z}(\mathcal{C})$ for some spherical fusion category \mathcal{C} . We will simplify notion for doubled topological order. We will write

[WORK: bulk theory with letter \mathcal{C} in it = topological ordered described by $\mathcal{Z}(\mathcal{C})$.]

This now leads us to the big question: what, algebraically, are boundaries theories in the $\mathcal{Z}(\mathcal{C})$ topological order?

[WORK: The answer is \mathcal{C} -module categories. I am suspecting that to motivate this answer one needs to take some physical principles from the string-nets developed in the Levin-Wen model. I haven't written the Levin-Wen section yet, so this can wait until that.]

More generally, a domain wall between the phases described by $\mathcal{Z}(\mathcal{C})$ and $\mathcal{Z}(\mathcal{D})$ will be a $(\mathcal{C}, \mathcal{D})$ -bimodule category. [WORK: motivate this as well].

We rigorously define the above intuitive notions as follows.

Definition (Left module category). A *left module category* over a fusion category \mathcal{C} is the following data:

1. A \mathbb{C} -linear category \mathcal{M} ;
2. (Action) A functor $\otimes : \mathcal{C} \times \mathcal{M} \rightarrow \mathcal{M}$;
3. (Left associator) A natural isomorphism $m : ((-) \otimes (-)) \otimes (-) \xrightarrow{\sim} (-) \otimes ((-) \otimes (-))$ of functors $\mathcal{C} \times \mathcal{C} \times \mathcal{M} \rightarrow \mathcal{M}$;
4. (Unit) A natural isomorphism $\lambda_m : \mathbf{1} \otimes (-) \rightarrow (-)$ of functors $\mathcal{M} \rightarrow \mathcal{M}$.

Such that:

1. There is an equivalence $\mathcal{C} \cong \mathbf{Vec}_{\mathbb{C}}^n$ as \mathbb{C} -linear categories for some integer $n \geq 1$.
2. The functor \otimes is \mathbb{C} -linear.
3. For all $A, B, C \in \mathcal{C}$ and $M \in \mathcal{M}$ the diagram

$$\begin{array}{ccccc}
 & & ((A \otimes B) \otimes C) \otimes M & & \\
 & \swarrow \alpha_{A,B,C} \otimes \text{id}_M & & \searrow m_{A \otimes B, C, M} & \\
 (A \otimes (B \otimes C)) \otimes M & & & & (A \otimes B) \otimes (C \otimes M) \\
 \downarrow m_{A, B \otimes C, M} & & & & \downarrow m_{A, B, C \otimes M} \\
 A \otimes ((B \otimes C) \otimes M) & \xrightarrow{\text{id}_A \otimes m_{B, C, M}} & & & A \otimes (B \otimes (C \otimes M))
 \end{array}$$

commutes.

4. For all $A \in \mathcal{C}$ and $M \in \mathcal{M}$, the diagram

$$\begin{array}{ccc}
 (A \otimes \mathbf{1}) \otimes M & \xrightarrow{m_{A, \mathbf{1}, M}} & A \otimes (\mathbf{1} \otimes M) \\
 \downarrow \rho_A & \swarrow id_A \otimes \lambda_m & \\
 A \otimes M & &
 \end{array}$$

commutes.

□

The definition of a *right* module category is completely symmetric, with no surprises. The definition of a bimodule category comes from putting these definitions together:

Definition (Bimodule category). A $(\mathcal{C}, \mathcal{D})$ -bimodule category \mathcal{M} over fusion categories \mathcal{C}, \mathcal{D} is the following data:

1. The structure of a right \mathcal{C} -module on \mathcal{M} ;
2. The structure of a left \mathcal{D} -module on \mathcal{M} ;
3. (Middle associator) A natural isomorphism $b : ((-) \otimes (-)) \otimes (-) \xrightarrow{\sim} (-) \otimes ((-) \otimes (-))$ of functors $\mathcal{C} \times \mathcal{M} \times \mathcal{D} \rightarrow \mathcal{M}$;

Such that:

1. For all $A, B \in \mathcal{C}$, $M \in {}_{\mathcal{C}}\mathcal{M}_{\mathcal{D}}$, $A', B \in \mathcal{D}$, the diagrams

$$\begin{array}{ccc}
 & ((A \otimes B) \otimes M) \otimes C & \\
 m_{A,B,M}^L \swarrow & & \searrow b_{A \otimes B, M, C} \\
 (A \otimes (B \otimes M)) \otimes C & & (A \otimes B) \otimes (M \otimes C) \\
 \downarrow b_{A, B \otimes M, C} & & \downarrow m_{A, B, M \otimes C}^L \\
 A \otimes ((B \otimes M) \otimes C) & \xrightarrow{\text{id}_A \otimes b_{B, M, C}} & A \otimes (B \otimes (M \otimes C))
 \end{array}$$

and

$$\begin{array}{ccc}
 & ((A \otimes M) \otimes B) \otimes C & \\
 b_{A, M, B} \otimes \text{id}_C \swarrow & & \searrow m_{A \otimes M, B, C}^R \\
 (A \otimes (M \otimes B)) \otimes C & & (A \otimes M) \otimes (B \otimes C) \\
 \downarrow b_{A, M \otimes B, C} & & \downarrow b_{A, M, B \otimes C} \\
 A \otimes ((M \otimes B) \otimes C) & \xrightarrow{\text{id}_A \otimes m_{M, B, C}^R} & A \otimes (M \otimes (B \otimes C))
 \end{array}$$

commute.

□

Given two fusion categories \mathcal{C}, \mathcal{D} and a $(\mathcal{C}, \mathcal{D})$ -bimodule categories \mathcal{M} , we define \mathcal{M} to be a *simple* module if for all simple objects $N, M \in \mathcal{M}$ there exists some $A \in \mathcal{C}$ such that $\text{Hom}(N, A \otimes M) \neq 0$. This condition of irreducibility corresponds exactly the condition of stability of boundaries. Hence, we arrive at the following correspondance:

[WORK: boundary theories correspond to simple module categories;]

[WORK: domain walls correspond to simple bimodule categories.]

We now turn our attention towards boundary defects.

[WORK: motivate why boundary defects are bimodule functors. I think this one can be easily seen by passing boundary excitations through the defect and thus noticing that the defect must act like a functor.]

Formally, we get the following definition:

Definition (Left module functor). A *left module functor* between left \mathcal{C} -module categories \mathcal{M}, \mathcal{N} is the following data:

1. A functor $F : \mathcal{M} \rightarrow \mathcal{N}$;
2. A natural isomorphism $s : F((-) \otimes (-)) \xrightarrow{\sim} (-) \otimes (-)$ of functors $\mathcal{C} \times \mathcal{M} \rightarrow \mathcal{N}$.

Such that:

1. For all $A, B \in \mathcal{C}, M \in \mathcal{M}$, the diagram

$$\begin{array}{ccccc}
 & & F((A \otimes B) \otimes M) & & \\
 & F(m_{A,B,M}^{\mathcal{M}}) \nearrow & & \searrow s_{A \otimes B, M} & \\
 F(A \otimes (B \otimes M)) & & & & (A \otimes B) \otimes F(M) \\
 \downarrow s_{A,B \otimes M} & & & & \downarrow m_{A,B,F(M)}^{\mathcal{N}} \\
 A \otimes F(B \otimes M) & \xrightarrow{\text{id}_A \otimes s_{B,M}} & & & A \otimes (B \otimes F(M))
 \end{array}$$

commutes.

2. For all $M \in \mathcal{M}$ the diagram

$$\begin{array}{ccc}
 F(\mathbf{1} \otimes M) & \xrightarrow{s_{\mathbf{1},M}} & \mathbf{1} \otimes F(M) \\
 \downarrow F(\lambda_{m,M}^{\mathcal{M}}) \searrow & & \downarrow \lambda_{m,F(M)}^{\mathcal{N}} \\
 & & F(M)
 \end{array}$$

commutes.

□

A right module functor is defined in complete analogy to left module functor. A $(\mathcal{C}, \mathcal{D})$ -bimodule functor is defined to be a functor which is simultaneously equipped with the structure of a left and right module functor. We denote by $\text{Fun}_{\mathcal{C}}(\mathcal{M}, \mathcal{N})$, the space of \mathcal{C} -module functors from \mathcal{M} to \mathcal{N} , and we denote by $\text{Fun}_{\mathcal{C}|\mathcal{D}}(\mathcal{M}, \mathcal{N})$ the space bimodule functors.

To turn these spaces of module functors into categories, we will need to define an appropriate notion of natural transformation

Definition (Natural transformation of left module functors). A *natural transformation of left module functors* between left \mathcal{C} -module functors $F, G : \mathcal{M} \rightarrow \mathcal{N}$ is a natural transformation $\eta : F \rightarrow G$ such that for all $A \in \mathcal{C}, M \in \mathcal{M}$

$$\begin{array}{ccc}
F(A \otimes M) & \xrightarrow{\eta_{A \otimes M}} & G(A \otimes M) \\
s_{A,M}^F \downarrow & & s_{A,M}^G \downarrow \\
A \otimes F(M) & \xrightarrow{\text{id}_A \otimes \eta_M} & A \otimes G(M)
\end{array}$$

□

Just like in our general theory, we define an object in a functor category to be irreducible if it cannot be expressed as the biproduct of two nonzero elements. We arrive at the following correspondance:

[WORK: boundary defects correspond to simple bimodule functors.]

The fusion of boundary defects is clear. Given two boundary defects $F : \mathcal{M} \rightarrow \mathcal{M}'$ and $G : \mathcal{M}' \rightarrow \mathcal{M}''$ the composition $(G \circ F) : \mathcal{M} \rightarrow \mathcal{M}''$ gives another boundary defect. This is the correct physical choice because [WORK: depends on the physical motivation I gave for defects corresponding to functors].

The fusion of domain walls is a bit more subtle. It requires the notion of relative tensor product:

Definition (Relative Deligne tensor product). [WORK: give definition]

Physically, we can define fusion of domain walls \mathcal{M} and \mathcal{N} to be the relative Deligne tensor product $\mathcal{M} \boxtimes_{\mathcal{K}} \mathcal{N}$, where \mathcal{M} is a $(\mathcal{C}, \mathcal{K})$ -bimodule and \mathcal{N} is a $(\mathcal{K}, \mathcal{D})$ -bimodule. This is the correct definition because [WORK: give physical motivation].

All together, this gives us the following physics-math dictionary:

[WORK: add dictionary.]

As we have just seen, the algebraic theory of topological boundaries and domain walls is modules and bimodules. So, mathematically, we now take the time to set up the theory of modules and bimodules over fusion categories.

The first thing we are going to do is reinterpret bimodules in terms of modules. This will allow us to state all of our results in this section in terms of modules without loss of generality, because dealing with bimodules adds unnecessary notational difficulty. This is done via the folding trick. As mentioned before, we have a correspondance as follows:

[WORK: add folding trick. The flat plane should have the label $(\mathcal{C}, \mathcal{D})$ -bimodules and the folded plane shuld have the label $(\mathcal{C} \boxtimes \overline{\mathcal{D}})$ -modules.]

This is codified mathematically in the following proposition:

Proposition 6.1 (Folding trick). *Let \mathcal{C}, \mathcal{D} be fusion categories. Given a $(\mathcal{C}, \mathcal{D})$ -bimodule \mathcal{M} , define a left $(\mathcal{C} \boxtimes \overline{\mathcal{D}})$ -module structure on \mathcal{M} via the formula*

$$(A \boxtimes \overline{B}) \otimes M = (A \otimes M) \otimes B$$

for any $A \in \mathcal{C}$, $B \in \mathcal{D}$, $M \in \mathcal{M}$. This is a well-defined left $(\mathcal{C} \boxtimes \overline{\mathcal{D}})$ -module structure. This assignment induces a bijection between the set of $(\mathcal{C}, \mathcal{D})$ -bimodules and the set of left $(\mathcal{C} \boxtimes \overline{\mathcal{D}})$ -modules.

Proof. This is a straightforward and unenlightening exercise. □

Next, we take a closer look at simple modules. That is, module categories with no proper nontrivial submodule categories. Given two categories \mathcal{M} and \mathcal{N} , we can define their Cartesian product $\mathcal{M} \times \mathcal{N}$, whose objects are pairs (A, B) , $A \in \mathcal{M}$, $B \in \mathcal{N}$. To avoid confusion with the Deligne tensor product, we will use the notation \boxplus to denote the

Cartesian product in this context, and refer to it as the direct sum. This is consistent with our general principles of higher linear algebra (see section [ref]). If \mathcal{M} and \mathcal{N} are left \mathcal{C} -module categories, then $\mathcal{M} \boxplus \mathcal{N}$ can be equipped with the structure of a left \mathcal{C} -module via the action

$$A \otimes (M \boxplus N) = (A \otimes M) \boxplus (A \otimes N).$$

We seek to prove that every module is the direct sum of simple modules. This follows after a lemma.

Lemma 6.1. *Let \mathcal{C} be a fusion category. Let \mathcal{M} be a left \mathcal{C} -module. For all $A \in \mathcal{C}$ and $N, M \in \mathcal{M}$, we have that*

$$\text{Hom}_{\mathcal{M}}(N, A \otimes M) \cong \text{Hom}_{\mathcal{M}}(A^* \otimes N, M)$$

as \mathbb{C} -vector spaces.

Proof. The proof is exactly the same as of proposition [ref]. We opt for a proof using the graphical language of morphisms instead of the algebraic approach, because it is more lucid. We have not formally introduced the graphical language for module categories, but it is completely analogous to the graphical language for fusion categories. We define the map

$$\begin{array}{ccc} \text{Hom}_{\mathcal{M}}(N, A \otimes M) & \longrightarrow & \text{Hom}_{\mathcal{M}}(A^* \otimes N, M) \\ \begin{array}{c} \text{AM} \\ \square \\ \alpha \\ \text{N} \end{array} & \longmapsto & \begin{array}{c} M \\ \square \\ \alpha \\ A^* \text{ N} \end{array} \end{array}$$

Twisting the A^* the other way serves as an inverse, completing the proof. \square

Proposition 6.2. *Let \mathcal{C} be a fusion category. Every left \mathcal{C} -module category is equivalent to a direct sum of simple \mathcal{C} -module categories.*

Remark. The usual way of summarizing this proposition is by saying that the space of left \mathcal{C} -modules is *semisimple*.

Proof. Let \mathcal{M} be a left \mathcal{C} -module. Denote by $\mathcal{L}_{\mathcal{M}}$ the set of isomorphism classes of simple objects in \mathcal{M} . We define a equivalence relation \sim on $\mathcal{L}_{\mathcal{M}}$ by saying that $[N] \sim [M]$ if there exists some $A \in \mathcal{C}$ such that $\text{Hom}(N, A \otimes M) \neq 0$. This relation is reflexive since $\text{Hom}(M, 1 \otimes M) \neq 0$, it is transitive since if $\text{Hom}(M, A \otimes M') \neq 0$ and $\text{Hom}(M', B \otimes M'') \neq 0$, then $\text{Hom}(M, A \otimes B \otimes M'') \neq 0$ by composing morphisms. Seeing as $\mathcal{M} \cong \mathbf{Vec}_{\mathbb{C}}^n$ for some $n \geq 1$ as a category, we have that $\text{Hom}(N, M) \cong \text{Hom}(M, N)$ as vector spaces for all $N, M \in \mathcal{M}$. Hence, the equivalence relation is symmetric by lemma [ref].

For all $i \in \mathcal{L}_{\mathcal{M}} / \sim$, we define a full subcategory of \mathcal{M} via

$$\mathcal{M}_i = \left\{ M \in \mathcal{M} \mid M \cong \bigoplus_{[N] \in i} n_N \cdot N, \text{ for some integers } n_A \geq 0 \right\}.$$

Now, our equivalence relation guarantees that \mathcal{M}_i is closed under the action of \mathcal{C} for all $i \in \mathcal{L}_{\mathcal{M}} / \sim$. That is, it forms a left \mathcal{C} -submodule of \mathcal{M} . We claim that the map

$$\begin{aligned} \bigoplus_{i \in \mathcal{L}_{\mathcal{M}}/\sim} \mathcal{M}_i &\rightarrow \mathcal{M} \\ (M_i)_{i \in \mathcal{L}_{\mathcal{M}}/\sim} &\mapsto \bigoplus_{i \in \mathcal{L}_{\mathcal{M}}/\sim} M_i \end{aligned}$$

is an equivalence of categories. It is faithful by the general properties of the direct sum. It is full because $\text{Hom}(M_i, M_j) = 0$ for all $i \neq j$ since M_i, M_j are the direct sums of disjoint sets of simple objects, and hence every morphisms can be broken into its i -components for each $i \in \mathcal{L}_{\mathcal{M}}/\sim$. It is essentially surjective because every simple objects appears in the category \mathcal{M}_i for some $i \in \mathcal{L}_{\mathcal{M}}/\sim$. Hence, it is an equivalence of categories by proposition [ref] so our proof is complete. \square

With propositions [ref] and [ref] in hand, we can set up more directly the theory of fusion of domain walls. Let $\mathcal{C}, \mathcal{K}, \mathcal{D}$ be fusion categories. Let \mathcal{M} be a $(\mathcal{C}, \mathcal{K})$ -bimodule and let \mathcal{N} be a $(\mathcal{K}, \mathcal{D})$ bimodule. We have an equivalence

$$\mathcal{N} \boxtimes \mathcal{M} \cong \bigoplus_{[\mathcal{M}']} n_{\mathcal{M}'}^{\mathcal{N}, \mathcal{M}} \cdot \mathcal{M}',$$

where $[\mathcal{M}']$ runs over the set of equivalence classes of simple $(\mathcal{C}, \mathcal{D})$ -bimodules. This is completely analogous to the algebraic theory of fusion of anyons introduced in section [ref]. One of the major differenes is that we are now up a categorical level. Instead of fusing objects in a category, we are fusing the categories themselves.

We now turn our attention towards functors between module categories. When the source and target of our functors are the same, we have the following main result:

Proposition 6.3. *Let \mathcal{C} be a fusion category. Let \mathcal{M} be a left \mathcal{C} -module. Define*

$$\mathcal{C}_{\mathcal{M}}^{\vee} = \text{Fun}_{\mathcal{C}}(\mathcal{M}, \mathcal{M}).$$

We call $\mathcal{C}_{\mathcal{M}}^{\vee}$ the dual of \mathcal{C} with respect to \mathcal{M} . We define a monoidal structure on $\mathcal{C}_{\mathcal{M}}^{\vee}$ via the formula

$$F \otimes G = G \circ F.$$

Along with a natural choice of rigid structure, $\mathcal{C}_{\mathcal{M}}^{\vee}$ has the structure of a fusion category.

If \mathcal{C} is a spherical fusion category, then $\mathcal{C}_{\mathcal{M}}^{\vee}$ is naturally equipped with the structure of a spherical fusion category.

Proof. .[WORK: do proof. This is going to use some stuff from “on fusion categories”. Might need the same results for the next proposition as well.]

I’m in trouble the proof uses Yetter cohomology and other scary-looking tools. Is there a simpler proof? The hard part is showing that if \mathcal{N}, \mathcal{M} are module categories the space of module functors is equivalent to $\text{Vec}_{\mathbb{C}}^n$ for some $n \geq 0$. In particular, we need to show that there are finitely many isomorphism classes of simple module functors. How do I do it?! It’s a sort of generalized Ocneanu rigidity. \square

When the source and target are different, we have the follwing result:

Proposition 6.4. *Let \mathcal{C} be a fusion category. Let \mathcal{N}, \mathcal{M} be left \mathcal{C} -modules. The category $\text{Fun}_{\mathcal{C}}(\mathcal{N}, \mathcal{M})$ is naturally equipped with the structure of a left $\mathcal{C}_{\mathcal{N}}^{\vee}$ -module and a right $\mathcal{C}_{\mathcal{M}}^{\vee}$ -module.*

Proof. .[WORK: do proof. Again, the finiteness part is hard.] \square

Now, we verify the finiteness result that we claimed in section [ref]:

Proposition 6.5. *Let \mathcal{C} be a fusion category. There are finitely many equivalence classes of simple left \mathcal{C} -modules.*

Proof. .[WORK: do proof. In “On fusion categories” they say that this follows from the finiteness condition on functors, along with a decategorified version of the same result which follows from counting.] \square

The following result serves as the centerpiece of our study of module categories:

Theorem 6.1. *Let \mathcal{C} be a fusion category. Let \mathcal{M} be a left \mathcal{C} module. The map*

$$\begin{aligned} \text{can} : \mathcal{C} &\xrightarrow{\sim} (\mathcal{C}_{\mathcal{M}}^{\vee})_{\mathcal{M}}^{\vee} \\ A &\mapsto (M \mapsto A \otimes M) \end{aligned}$$

is an equivalence of categories. The $\mathcal{C}_{\mathcal{M}}^{\vee}$ -module functor structure on can_A is defined for all $A \in \mathcal{C}$, $M \in \mathcal{M}$, $F \in \mathcal{C}_{\mathcal{M}}^{\vee}$ by

$$s_{F,M}^{\text{can}_A} : \text{can}_A(F \otimes M) = A \otimes F(M) \xrightarrow{s_{A,M}^{-1}} F(A \otimes M) = F \otimes \text{can}_A(M).$$

If \mathcal{C} is a spherical fusion category, can is an equivalence of spherical fusion categories.

Remark. This result is an analogue of the *double centralizer theorem* from classical representation theory.

Proof. .[WORK: do proof. The proof in EGNO uses algebras and algebra structures. It would be nice to get around this.] \square

As a corollary, we get the following:

Corollary 6.1. *Let \mathcal{C} be a fusion category and let \mathcal{M} be a left \mathcal{C} -module. The module \mathcal{M} is simple as a right $\mathcal{C}_{\mathcal{M}}^{\vee}$ -module.*

Proof. .[WORK: do proof] \square

Of course, the above discussion guarantees a proper theory of fusion of boundary defects. Suppose that \mathcal{C}, \mathcal{D} are fusion categories, and $\mathcal{M}, \mathcal{M}', \mathcal{M}''$ are simple $(\mathcal{C}, \mathcal{D})$ -bimodules. Suppose that $F : \mathcal{M} \rightarrow \mathcal{M}'$ and $G : \mathcal{M}' \rightarrow \mathcal{M}''$ are simple $(\mathcal{C}, \mathcal{D})$ bimodule functors. Since $\text{Fun}_{\mathcal{C} \mid \mathcal{D}}(\mathcal{M}, \mathcal{M}'')$ is equivalent to $\text{Vec}_{\mathbb{C}}^n$ for some $n \geq 0$ as a \mathbb{C} -linear category, we have a decomposition

$$G \circ F \cong \bigoplus_{[H]} n_H^{F,G} \cdot H$$

in $\text{Fun}_{\mathcal{C} \mid \mathcal{D}}(\mathcal{M}, \mathcal{M}'')$ where $[H]$ runs over isomorphism classes of simple functors in $\text{Fun}_{\mathcal{C} \mid \mathcal{D}}(\mathcal{M}, \mathcal{M}'')$. Again, this is completely analogous to the fusion of anyons.

6.2.6 The Drinfeld center via bimodules

[WORK: The paper [KR08] has some interesting additional insights. Maybe if I took the time to understand it it would be relevant as a comment?]

So far we have been setting up a powerful theory of domain walls to study doubled topological order.

We already know one thing about doubled topological order - its anyons are described by the Drinfeld center $\mathcal{Z}(\mathcal{C})$ of the input spherical fusion category \mathcal{C} . These anyonic bulk excitations can interact with the boundaries, and produce non-trivial effects. This means that the category $\mathcal{Z}(\mathcal{C})$ necessarily has deep connections to the theory of boundaries, and thus mathematically will have a module-theoretic interpretation. This will also help answer the question of why the Drinfeld center $\mathcal{Z}(\mathcal{C})$ is the category which describes anyons in the Levin-Wen model.

To start, consider an anyon in the bulk. This anyon is a stable localized excitation. Above it and below it the theory is in its ground state. Hence, we can think of this anyon as an interface between the trivial boundary and itself! This is seen visually below:

[WORK: add picture of anyon as interface between trivial boundary and itself.]

Thus, from our general description of boundary defects, we should conclude the following:

$$(\text{anyon types}) = \left(\begin{array}{c} \text{boundary defects} \\ \text{between trivial theory and itself} \end{array} \right) = \left(\begin{array}{c} \text{simple objects in} \\ \text{Fun}_{\mathcal{C}|\mathcal{C}}(\mathcal{C}, \mathcal{C}) \end{array} \right).$$

However, we know from our first discussion of the Levin-Wen model that anyon types should correspond to simple objects in $\mathcal{Z}(\mathcal{C})$. Hence, we conclude that simple objects in $\mathcal{Z}(\mathcal{C})$ should correspond to simple objects in $\text{Fun}_{\mathcal{C}|\mathcal{C}}(\mathcal{C}, \mathcal{C})$. Of course, this equivalence should respect fusion, duality, braiding, and all the other algebraic structures of anyons. Thus, if our boundary theory is correct, it predicts an equivalence of categories $\mathcal{Z}(\mathcal{C}) \cong \text{Fun}_{\mathcal{C}|\mathcal{C}}(\mathcal{C}, \mathcal{C})$. This mathematical prediction is accurate, by the following proposition:

Proposition 6.6. *Let \mathcal{C} be a fusion category. The map*

$$\begin{aligned} \text{Fun}_{\mathcal{C}|\mathcal{C}}(\mathcal{C}, \mathcal{C}) &\xrightarrow{\sim} \mathcal{Z}(\mathcal{C}) \\ F &\mapsto (F(1), \beta_{F(1), -}) \end{aligned}$$

is an equivalence of categories, where $\beta_{F(1), -}$ is defined by the composition

$$\beta_{F(1), -} : F(\mathbf{1}) \otimes B \xrightarrow{s_{\mathbf{1}, B}^R} F(\mathbf{1} \otimes B) = F(B) = F(B \otimes \mathbf{1}) \xrightarrow{(s_{B, \mathbf{1}}^L)^{-1}} B \otimes F(\mathbf{1})$$

and s^L, s^R are the compatibility maps implicit in the left/right module functor structure of F .

If \mathcal{C} is a spherical category, then the map $\text{Fun}_{\mathcal{C}|\mathcal{C}}(\mathcal{C}, \mathcal{C}) \xrightarrow{\sim} \mathcal{Z}(\mathcal{C})$ induces an equivalence as pre-modular categories.

Proof. . [WORK: do proof] □

This equivalence of categories gives the proper interpretation for $\mathcal{Z}(\mathcal{C})$ in the theory of boundaries. If we wanted we could have taken it as a *definition* of $\mathcal{Z}(\mathcal{C})$. We now examine how anyons interact with boundaries to get more module-theoretic properties of the Drinfeld center.

Suppose \mathcal{C}, \mathcal{D} are spherical fusion categories, and ${}_{\mathcal{C}}\mathcal{M}_{\mathcal{D}}$ is a bimodule. If a $A \in \mathcal{C}$ is a simple object describing an anyon, we can imagine trying to push A from \mathcal{C} to \mathcal{D} through the boundary theory ${}_{\mathcal{C}}\mathcal{M}_{\mathcal{D}}$:

[WORK: add picture: first its on one side, then a bubble starts to form, then the bubble detaches. It leaves behind a dotted line with label ${}_{\mathcal{D}}\mathcal{M}_{\mathcal{C}}^{op} \boxtimes_{\mathcal{C}} {}_{\mathcal{C}}\mathcal{M}_{\mathcal{D}}$]

Suppose that the domain wall is *invertible*. This means that ${}_{\mathcal{D}}\mathcal{M}_{\mathcal{C}}^{op} \boxtimes_{\mathcal{C}} {}_{\mathcal{C}}\mathcal{M}_{\mathcal{D}} \cong \mathcal{D}$. Thus, the anyon can pass directly through the boundary, and the boundary can re-close without leaving a trace of the fact that the anyon went through.

This gives an assignment of anyons in the $\mathcal{Z}(\mathcal{C})$ to anyons in the $\mathcal{Z}(\mathcal{D})$ phase. This assignment is reversible, because anyons in the $\mathcal{Z}(\mathcal{D})$ phase can pass through the domain wall in the other direction. Thus, our boundary theory predicts that every invertible bimodule ${}_{\mathcal{C}}\mathcal{M}_{\mathcal{D}}$ will induce an equivalence of categories $F_{\mathcal{M}} : \mathcal{Z}(\mathcal{C}) \xrightarrow{\sim} \mathcal{Z}(\mathcal{D})$. We verify this prediction with a theorem. Before we can prove it, we need a lemma.

Lemma 6.2. *Let \mathcal{C} be a fusion category. The forgetful functor*

$$\begin{aligned} \mathcal{Z}(\mathcal{C}) &\rightarrow \mathcal{C} \\ (A, \beta_{A,-}) &\mapsto A \end{aligned}$$

is surjective onto the objects of \mathcal{C} .

Proof. .[WORK: this proof is steeped in the theory of module categories.] \square

Theorem 6.2. *Let \mathcal{C}, \mathcal{D} be spherical fusion categories. There is a bijection*

$$\left(\begin{array}{c} \text{invertible } (\mathcal{C}, \mathcal{D})\text{-bimodule} \\ \text{categories } \mathcal{M} \end{array} \right) \xleftrightarrow{\sim} \left(\begin{array}{c} \text{equivalences of pre-modular} \\ \text{categories } \mathcal{Z}(\mathcal{C}) \xrightarrow{\sim} \mathcal{Z}(\mathcal{D}) \end{array} \right),$$

defined by sending a $(\mathcal{C}, \mathcal{D})$ -bimodule category \mathcal{M} to the functor $F_{\mathcal{M}} : \mathcal{Z}(\mathcal{C}) \rightarrow \mathcal{Z}(\mathcal{D})$ defined by

$$F_{\mathcal{M}} : \mathcal{Z}(\mathcal{C}) \cong \text{Fun}_{\mathcal{C}|\mathcal{C}}(\mathcal{C}, \mathcal{C}) \rightarrow \text{Fun}_{\mathcal{D}|\mathcal{D}}(\mathcal{D}, \mathcal{D}) \cong \mathcal{Z}(\mathcal{D}).$$

$$G \mapsto \left(\mathcal{D} \rightarrow \mathcal{M}^{op} \boxtimes_{\mathcal{C}} \mathcal{C} \boxtimes_{\mathcal{C}} \mathcal{M} \xrightarrow{\text{id} \boxtimes G \boxtimes \text{id}} \mathcal{M}^{op} \boxtimes_{\mathcal{C}} \mathcal{C} \boxtimes_{\mathcal{C}} \mathcal{M} \rightarrow \mathcal{D} \right)$$

For all triples $\mathcal{C}, \mathcal{K}, \mathcal{D}$ of fusion categories, $(\mathcal{C}, \mathcal{K})$ -bimodules \mathcal{M} , and $(\mathcal{K}, \mathcal{D})$ -bimodules \mathcal{N} , we have that

$$F_{\mathcal{M} \boxtimes_{\mathcal{K}} \mathcal{N}} = F_{\mathcal{N}} \circ F_{\mathcal{M}}.$$

If \mathcal{C} and \mathcal{D} are spherical fusion categories, then the assignment $\mathcal{M} \mapsto F_{\mathcal{M}}$ induces a bijection.

[WORK: what is the correct statement? Do I have to reduce the class of bimodules I am considering?]

Remark. This theorem is originally due to Etingof-Nikshych-Ostrik [ENO10]. It was observed in the theory of gapped boundaries by Kitaev and Kong [KK12]. The proof presented here is original.

Proof. .[WORK: do proof. The first direct is making sure this is a well-defined map and filling in details. That's not very hard.

The big question is showing that this is a *bijection*. That means that we need an inverse to this map. The inverse is surprisingly cute and nice. The forgetful functor $\mathcal{Z}(\mathcal{C}) \rightarrow \mathcal{C}$ induces a $\mathcal{Z}(\mathcal{C})$ -bimodule structure on \mathcal{C} . Given any equivalence $F : \mathcal{Z}(\mathcal{C}) \xrightarrow{\sim} \mathcal{Z}(\mathcal{D})$, we can thus construct the tensor product

$$\mathcal{C} \boxtimes_{\mathcal{Z}(\mathcal{C}) \cong \mathcal{Z}(\mathcal{D})} \mathcal{D}.$$

This tensor product is a $(\mathcal{C}, \mathcal{D})$ bimodule. The fact that it is an invertible bimodule uses in a key way the fact that the forgetful functors $\mathcal{Z}(\mathcal{C}) \rightarrow \mathcal{C}$ and $\mathcal{Z}(\mathcal{D}) \rightarrow \mathcal{D}$ are surjective. It's not hard to show that these two constructions are inverses to one another. This gives the proof.] \square

As a corollary we find the following theorem, which was first proved by Kitaev and Muger and communicated by Etingof-Nikshych-Ostrik [ENO11]. We recall that two fusion categories \mathcal{C}, \mathcal{D} are called *Morita equivalent* if their Drinfeld centers are equivalent as braided fusion categories

[WORK: What is the correct statement of this theorem for spherical fusion categories?]

Corollary 6.2. *Two fusion categories \mathcal{C}, \mathcal{D} are Morita equivalent if and only if there exists an invertible bimodule category between them.*

Proof. This is clear from Theorem [ref]. \square

6.2.7 Muger's theorem and Lagrange's theorem for fusion categories

[WORK: There's some results about Frobenius-Perron dimensions which will play a key role in the proof. Can I use quantum dimension instead? Probably not because I'm not assuming spherical.]

[WORK: Out of all this we get Lagrange's theorem for fusion categories, that given a fusion \mathcal{C} with full subcategory \mathcal{C} , the ratio $\text{FPdim}(\mathcal{D})/\text{FPdim}(\mathcal{C})$ is an algebraic integer. Do I care? Should I include this? Maybe I should have a section with a proof of Muger's theorem plus a proof of Lagrange's theorem for fusion categories.]

6.2.8 Skeletonization of gapped boundaries

[WORK: To be explicitly workable it is nice to reinterpret gapped boundaries in the skeletal language just like how one does with modular categories. In this case the symbols are M3j and M6j symbols, as discussed in Cong-Cheng-Wang. Not sure if I want to do this, though.]

6.3 Symmetry enriched topological order

6.3.1 Physical picture

[WORK: Two additional good references to read for this stuff are [AMF16, AFM20]]

[WORK:

This section should give an elementary physical picture of what symmetry enriched topological order. The basic idea is that a topological gap might be closable, but maybe closing it has to break some symmetry! This is easy to explain and motivate.

To talk about the next part, people need some idea about how phase transitions correspond to symmetry breaking. Moreover, it would be nice if it was clear how symmetry breaking corresponds to condensing quasiparticles. Burnell has a very nice quote about this which is mysterious to me:

“In a conventional ordering transition, the order parameter results from condensing a bosonic excitation— for example, the ordering of a magnet can be viewed as the Bose condensation of spinons.”

I think that getting the example of the magnet phase transition across would be great. If people can understand the magnet, they can understand everything. I don’t understand the magnet yet.

]

[WORK: I would like to include the Curie principle - the symmetry of the causes are to be found in the effects! The symmetries $\text{Aut}^\otimes(\mathcal{C})$ are necessarily interesting objects. Why? Curie principle! The curie principle in this case manifests in the possible defects which can appear. While \mathcal{C} and $\mathcal{C} \otimes \overline{\mathcal{C}}$ might *look* like very similar categories, $\mathcal{C} \otimes \overline{\mathcal{C}}$ may have way more symmetries than \mathcal{C} and hence might look like a better candidate for TQC in this guise

The original reference for the Curie principle is [Cur94]. The theory of magnetic space groups gives a quite nice example.]

6.3.2 Haldane spin chain

[WORK:

This section should include the easiest example possible of an SPT phase, or at least one that maximizes simplicity and relevance. I’m sure Kitaev has a nice model. People online say that the Haldane phase of the Heisenberg spin chain is easy but I haven’t seen it before so I can’t be sure.

]

6.3.3 Twist defects in the toric code

[WORK:

Talk about the toric code. Introduce the twist defects, and show how they work. This gives the main example of an SET phase with its defects nice and visible.]

6.3.4 The algebraic theory of SET phases

In this section, we will introduce the algebraic theory of symmetry enriched topological phases.

[WORK:

Why categorical G -actions? Why G -graded categories? It would be great if I could work up to the picture

$$\mathcal{C} <-- \text{confinement}, \text{defectification} --> \mathcal{C}_G^\times$$

If I could even get in the fact that there are obstructions that would be awesome. Of course, I can’t really prove much. The hope is that I can state everything trivially, and not have to make any big claims...]

[WORK: Another good thing to note is that there are generalizations of the notion of symmetry in this context! In particular we have the work of Cui-Zini-Wang [CZW19]. It’s not clear that group actions are the only symmetries which could protect defects.]

[WORK:

A really important theorem in this space is EGNO's observation that Aut and Pic are isomorphic. Does this come up on its own already, or do I need to add this in? I'm hazy now on the physical significance of this isomorphism, and on what categorical level it holds.

Zhenghan made a point to me that noninvertible defects break the theorem (they aren't in the picard group), but the philosophy is still there. Namely, the category (appropriately defined) $\text{Pic}(\mathcal{B}_G^\times)$ should be isomorphic to some endomorphism group. The philosophy holds (though I am hazy on it), but there is no known theorem. This is something to look into!]

Definition (Categorical group action). A *categorical group action* of a finite group G on a fusion category \mathcal{C} is the following data:

1. A function $\rho : G \rightarrow \text{Aut}^\otimes(\mathcal{C})$;
2. A natural isomorphism of functions of fusion categories

$$\eta_{g,h} : \rho(g) \circ \rho(h) \xrightarrow{\sim} \rho(gh)$$

for all $g, h \in G$.

Such that:

1. The square

$$\begin{array}{ccc} \rho(g_0) \circ \rho(g_1) \circ \rho(g_2) & \xrightarrow{\eta_{g_0,g_1}} & \rho(g_0g_1) \circ \rho(g_2) \\ \downarrow \eta_{g_1,g_2} & & \downarrow \eta_{g_0g_1,g_2} \\ \rho(g_0) \circ \rho(g_1g_2) & \xrightarrow{\eta_{g_0,g_1g_2}} & \rho(g_0g_1g_2) \end{array}$$

commutes for all $g_0, g_1, g_2 \in G$.

2. There is an equality of functors $\rho(e) = \text{id}_{\mathcal{C}}$, where $e \in G$ is the identity element³.

□

We can similarly define a categorical group action on a spherical fusion category and modular category. Namely, the group of automorphisms is replaced by the group of automorphisms which respect the additional structure, and the natural isomorphisms between those functors are required to respect that extra structure as well. In this case we allow ourselves extra adjectives to emphasize the additional structure. Namely, we refer to a *braided* categorical group action of a group on a braided fusion category. We can thus give the following physical principle:

$$(\text{Symmetry enriched topological phases}) \iff \left(\begin{array}{c} \text{triples } (\mathcal{C}, G, \rho), \\ \text{where } \mathcal{C} \text{ is a modular category, } G \text{ is a finite group} \\ \text{and } \rho \text{ is a braided categorical group action} \end{array} \right).$$

Algebraically, the physical phenomenon of *symmetry breaking* is extremely clear. It is the act of taking a symmetry enriched phase (\mathcal{C}, G, ρ) and passing some subgroup $H \leq G$, yield

³The condition is added for simplicity of notation down the line - the axioms already imply that there is an equivalence $\rho(e) \cong \text{id}_{\mathcal{C}}$.

$(\mathcal{C}, H, \rho|_H)$. Passing all the way to the trivial group yields the ordinary topological phase \mathcal{C} .

We now move on to discussing symmetry fractionalization.

[WORK: motivate the answer physically.]

It's quite important that NOT every symmetry can be fractionalized. This is where the gauging anomaly shows up. Talk about this.]

Definition (G -graded fusion category). A G -graded fusion category over a finite group G is the following data:

1. A fusion category \mathcal{C} ;
2. A full subcategory \mathcal{C}_g of \mathcal{C} for each $g \in G$.

Such that:

1. The functor

$$\begin{aligned} \bigoplus_{g \in G} \mathcal{C}_g &\xrightarrow{\sim} \mathcal{C} \\ (A_g)_{g \in G} &\mapsto \bigoplus_{g \in G} A_g \end{aligned}$$

is an equivalence of categories;

2. The tensor product on \mathcal{C} restricts to a map $\otimes : \mathcal{C}_g \times \mathcal{C}_h \rightarrow \mathcal{C}_{gh}$;
3. (Faithfulness⁴) The subcategories \mathcal{C}_g are nonzero for all $g \in G$.

□

Definition (G -crossed braided fusion category). A G -crossed braided fusion category over a finite group G is the following data:

1. A G -graded fusion category \mathcal{C}_G^\times with g -component \mathcal{C}_g ;
2. A categorical group action ρ of G on \mathcal{C}_G^\times ;
3. A natural isomorphism

$$\beta_{g,A,B} : A \otimes B \rightarrow \rho(g)(B) \otimes A$$

for all $A \in \mathcal{C}_g$, $B \in \mathcal{C}_h^\times$, $g \in G$.

Such that:

[WORK: add compatibility. Namely, heptagon.]

Also there's another big one:

$$\rho(g)(\mathcal{C}_h) \subseteq \mathcal{C}_{ghg^{-1}}.$$

]

□

⁴This condition is not included in the definition by many authors.

Finally, we arrive at the definition of a G -crossed modular category:

Definition (G -crossed modular category). A G -crossed modular category over a finite group G is the following data:

1. A G -crossed braided fusion category \mathcal{C}_G^\times ;
2. A spherical structure on \mathcal{C}_G^\times ;

Such that:

1. [WORK: the spherical structure should be compatible with the G -crossed braided structure].
2. \mathcal{C}_e is a modular category, where $e \in G$ is the identity element. [WORK: say this nicer. Should I introduce better notation for group elements?]

□

We observe from our definitions that the identity graded component \mathcal{C}_e of any G -crossed modular category is itself a modular category. The assignment $\mathcal{C}_G^\times \mapsto \mathcal{C} = \mathcal{C}_e$ corresponds to the physical transition from a phase where defects exist and can be freely manipulated, and a phase in which no free defects can be made. This is known physically as *confinement*. This gives us the following picture:

$$\begin{array}{ccc} \mathcal{C} & \xrightleftharpoons[\text{confinement}]{\text{fractionalization}} & \mathcal{C}_G^\times \\ \uparrow & & \uparrow \\ \text{modular category} & & \text{G-crossed} \\ \text{with braided} & & \text{modular category} \\ \text{categorical } G\text{-action} & & \end{array}$$

as mentioned, it is not always possible to fractionalize symmetries. We will discuss this in more detail in section [ref].

6.3.5 Gauging symmetries

[WORK: This is the section where I talk about gauging. Not sure how much I would have mentioned it before. First off, it would be nice to talk about gauging from a physical perspective. The end game is to state the maps

$$\mathcal{C}_G^\times \dashleftarrow \text{condensation, gauging} \dashrightarrow (\mathcal{C}_G^\times)^G$$

properly. General discussion gauge symmetry breaking phase transition would be nice to talk about. If I could state everything clearly mathematically that would be great.]

Proposition 6.7. Let \mathcal{C} be a fusion category with a categorical G action. Define a category \mathcal{C}^G as follows. Its objects are pairs (A, γ) , where $A \in \mathcal{C}$ is an object and

$$\gamma_{g,A} : \rho(g)(A) \xrightarrow{\sim} A$$

is an isomorphism for all $g \in G$, $A \in \mathcal{C}$ such that

[WORK: add diagram]

commutes. The morphisms in \mathcal{C}^G are morphisms in \mathcal{C} such that

[WORK: add diagram]

commutes. [WORK: define the rest of the structure]. This endows \mathcal{C}^G with the structure of a fusion category. We call \mathcal{C}^G the equivariantization of G .

[WORK: the space of possible maps γ making $(n \cdot \mathbf{1}, \gamma)$ an element of \mathcal{C}^G is equivalent to the space of n -dimensional representations of G . To make this canonical, however, one needs to choose a choice of trivialization of $\rho(g)(\mathbf{1})$ for all $g \in G$. That is, an isomorphism $\rho(g)(\mathbf{1}) \cong \mathbf{1}$.

I figured it out. I added the axiom that $\rho(e)$ is the identity. We thus have the distinguished map

$$\begin{aligned} & \rho(g)(\mathbf{1}) \\ & \rightarrow \rho(g)(\mathbf{1}) \otimes \rho(e)(\mathbf{1}) \\ & \rightarrow \rho(g)(\mathbf{1}) \otimes \rho(g)\rho(g^{-1})(\mathbf{1}) \\ & \rightarrow \rho(g)(\mathbf{1} \otimes \rho(g^{-1})(\mathbf{1})) \\ & \rightarrow \rho(g)(\rho(g^{-1})(\mathbf{1})) \\ & \rightarrow \rho(e)(\mathbf{1}) \\ & \rightarrow \mathbf{1} \end{aligned}$$

which fixes the issue of canonical trivialization. Should I add this as a footnote maybe? Feels overly technical.]

[WORK: figure out how to pass this lemma through once there's extra structure added onto \mathcal{C} .]

Proof. .[WORK: do proof] □

In light of the above proposition, we give the following definition:

Definition (Internal G -symmetry). An *internal G -symmetry* of a fusion category \mathcal{C} is a fully faithful functor $\text{Rep}(G) \rightarrow \mathcal{C}$.

[WORK: I'm realizing that I haven't done much with representation categories up to here. Should I talk about them more? Say some generalities about representation theory?]

Proposition [ref] tells us that every fusion category with a categorical G -action induces a fusion category with an internal G -symmetry. [WORK: give motivation for why this is gauging].

We now discuss *condensation*. As discussed above, condensation should serve as an inverse to gauging. Thus, it should take as an input a category with an internal G -symmetry and output a category with a categorical G -action. We define this procedure as follows:

Proposition 6.8. .[WORK: define equivariantization and assert that it behaves well]

Proof. .[WORK: do proof] □

We now verify that these two constructions are indeed inverses to one another:

Proposition 6.9. . [WORK: define maps

$$\mathcal{C} \rightarrow (\mathcal{C}^G)_G$$

and

$$\mathcal{C} \rightarrow (\mathcal{C}_G)^G.$$

State that these are equivalences of categories. Furthermore, assert that these equivalences of categories go through when we add all of our extra structure back on.

The principle here is that the data of a G -grading is the same as the data of an internal G -symmetry. This is the “main principle” of Bruguieres and Muger [Bru00, Müg00]. A good exposition of it is in [DGNO10].]

Thus, combining the results of this section with the ones from the previous section we get the following picture:

[WORK: add picture; gauging, condensing, defectifying, confining]

6.3.6 Obstruction theory, anomalies, and zesting

[WORK: this section is devoted to mathematical results which are neccecary from the first two sections. This might have to be broken up for readability]

[WORK: A lot of the stuff in this section is going to be about graded categories. There needs to be a good discussion of equivariantization and deequivariantization. I suppose there also needs to be some discussion of Tannakian subcategories. Lots of stuff to do - this section could be quite mathematically intricate. I will NOT prove the main obstruction theory theorem from ENO. It would be nice to have a proper statement of it though... Maybe the modern theory of zesting is relevant for a good statement?]

The theory of zesting seems to be pretty good. There are several papers which have made major progress on it: it was formally introduced in [DGP⁺21]. The paper [DKP21] explains how zesting gives MTCs with the same modular data - this could be a good exercise for the end of the section. The paper I care about is [DGP⁺24], which allows me to go from solution-to-solution in obstruction theory, realizing the torsor action.]

6.3.7 G -crossed modularity

[WORK:

There’s nice stuff to be said about the modular representation of a G -crossed MTC. After the original Barkeshli paper, the follow up paper [BB23] included a lot of the nice results.

There is a G -crossed verlinde formula - [DM19]. Seeing as it is new it is not very well known, and could easily be overlooked. I’m sure the proof isn’t that bad, and it would be really good to have included. It feels important to me.]

[WORK: Here are some tricks I found useful a few months ago when I was doing this

1. Given any $A \in \mathcal{C}_g$, all other elments of \mathcal{C}_g an be obtained by fusion with an element of \mathcal{C}_e ;
2. $\text{Rank}(\mathcal{C}_g) = \# \text{ fixed points of } \mathcal{C}_e \text{ under action of } g$

is this a good section to include them in?]

6.3.8 Skeletonization of SET phases

.[WORK: I want to list data for SET phases soon, so skeletonization is important. Talk about it. Set it up right. Do it.]

6.4 Fermionic topological order

6.4.1 Physical picture

6.4.2 Kitaev spin chain

[WORK: talk about Kitaev spin chain. Emphasize that bulk defects are fermions! These fermions can interact with the boundary Majorana modes.]

6.4.3 The algebraic theory of fermionic topological order

.[WORK: give the algebraic theory. This should be very easy to set up based on the work I've done in the previous sections.

I should talk about 16 fold way too. Prove that if a gauging exists then there are 16 solutions. State the 16 fold way theorem, but don't prove it. The proof is [JFR24] and it is very deep. Talking to Theo, he emphasizes that the proof is highly nontrivial and is not very physically motivated.]

6.4.4 The 16-fold way on the trivial fermionic theory

.[WORK: Talk about Kitaev's 16-fold way on the trivial fermionic theory.

This means showing that the obstruction vanishes, and then constructing all 16 theories, and saying the relevant things about them.]

6.4.5 Modular representations of fermionic topological order

[WORK:

One easy thing to do is the modular representation of fermionic topological order. It just has a small subtlety (you have to pass to a congruence subgroup for spin structure reasons) but then everything works out! This is shown here: [BRWZ18].

Maybe there's more to say in this section? Talk about Verlinde formula and other hallmarks of modular categories, but just insofar as they have special applications to fermionic topological order.]

6.4.6 Super-groups and Deligne's theorem

[WORK:

A very important theorem in this area is Deligne's theorem for fusion categories: Every symmetric fusion category is equivalent to the category of representations of a finite group or super-group. The original proof is in [Del02], but that's way too general and I'm sure there are better/more modern proofs for my case.

Depending on how important this theorem is, I might have to move this up. In Theo's paper [JFR24] he talks about Deligne's theorem + his obstruction results gives a complete picture for modularization and gauging away centers. This is an interesting perspective which I would like to include.]

History and further reading:
(Boundaries and domain walls)

The first major work on boundaries in conformal field theory was Cardy's seminal 1989 paper [Car89]. This work was then reinterpreted in terms of category theory by Fuchs and Schweigert in 2001 [FS01], who introduced the idea of interpreting boundaries in terms of module-categories and Frobenius algebras. This theory was then generalized by Frohlich, Fuchs, Runkel, and Schweigert to include a categorical description of domain walls as well [FFRS05].

In parallel, Bravyi and Kitaev (along with similar work by Freedman and Meyer [FM01]) studied the toric code with boundary and discovered that it had two different boundary types [BK98]. In 2008 Bombin and Martin-Delgado defined a version of the Kitaev quantum double model with boundary based on any finite group [BMD08, BMD11]. In Kitaev and Kong's 2012 paper [KK12], they generalized this to an arbitrary doubled topological order using the Levin-Wen model. They brought the well-established theory of boundaries in conformal field theory to topological order. The full theory of boundaries and domain walls in topological order was then established in 2016 by Cong, Cheng, and Wang [CCW16].

The canonical reference for the theory of module categories is Etingof-Gelaki-Nikshych-Ostrik [EGNO16]. This reference collected a lot of the earlier literature references, such as [Ost03, ENO05, ENO11].

History and further reading:
(Symmetry enriched topological order)

Many of the most famous examples of topological quantum systems are not topologically ordered systems, but symmetry protected topological systems. The first example of symmetry protected topological order was the Haldane phase of the odd-integer spin chain [Hal83], which is protected by $SO(3)$ spin rotation symmetry. Another famous example is that of the topological insulator, which is protected by $U(1)$ phase symmetry and time-reversal symmetry [KM05, HK10]. The existence of these important examples begs for a description in the category-theoretic language.

This mathematical language, however, had to wait for a while. There are lots of subtleties involved, and people were still working out the theory for ordinary topological order. The notion of G -crossed modular category was introduced in Turaev's giant work on homotopy quantum field theory [Tur10a]. Homotopy quantum field theory was not intended as a physically relevant work - it was mostly a mathematical curiosity, which was worked on because it seemed natural. The theory of G -crossed fusion categories was then applied to the process of orbifolding in conformal field theory [KJ04], though most of the seminal work such as Etingof-Nikshych-Ostrik's application of obstruction theory to G -crossed categories was still being done as a purely mathematical exercise.

It was not until the work of Barkeshli-Bonderson-Cheng-Wang that G -crossed modular categories were introduced as the correct formalism for symmetry enriched topological order [BBCW19]. This bridging of the mathematical literature and the physical literature led to a huge amount of progress in the area.

History and further reading:
(Fermionic topological order)

There is a very strong link between fermions and \mathbb{Z}_2 -gradings. This link was made popular by the discovery of *supersymmetry* in string theory, which means passing to a \mathbb{Z}_2 -graded geometry [Var04]. The fact that topological order whose fundamental degrees of freedom are fermionic need special attention was brought to light by Gu, Wang, and Wen, in a series of foundational papers [GWW14, GW14, GWW15]. The formalism of super-modular categories was introduced in 2017 [BGH⁺17]. This paper brought together a lot of the surrounding literature into a coherent physical narrative. The central conjecture of this narrative was the 16-fold way, which was established in 2024 by Johnson-Freyd and Reutter [JFR24].

Exercises:

- 6.1. .[WORK: show that there is a canonical bijection between structures of a left \mathcal{C} -module on \mathcal{M} and \mathbb{C} -linear monoidal functors

$$\mathcal{M} \rightarrow \text{End}(\mathcal{C}).$$

This is a categorical version of the fact that modules and representations are in bijection.]

- 6.2. . [WORK: people love the Tambara-Yamagami theorem, and consider it to be very deep. One of the original motivation of the algebraic theory of SET phases is that it gives a quick proof! It would thus be nice to have the Tambara-Yamagmi result here as an exercise. [TY98].]
- 6.3. .[WORK: categorical group actions can be reinterpreted in terms of functors. Namely, a category group action is a monoidal functor from the categorified version of the group. The categorified version of the group is defined so that it has an object for each group element, its only morphisms are the identity, and its monoidal structure is given by the group law.]
- 6.4. . [WORK: There is the principle of *modularization*. Here, we take the center (which will be some symmetric fusion category) and “gauge it away”. The fact that this is possible comes from our setup of symmetry enriched toplogical order and fermionic topological order. I don’t quite see how it is physically relevant. It seems best left here as an exercise.]
- 6.5. .[WORK: I think that a SFC is abelian if and only if its Drinfeld center is. This would be a nice exercise to include.]

7 Topological quantum computation

7.1 Overview

7.1.1 Introduction

In this section we will discuss topological quantum computing, the concept of making a computer based on topological quantum systems. We recall now how this fits into the overall framework of this book:

[WORK: add diagram]

We recall some general principles and motivations for topological quantum computing, most of which were outlined in Chapter [ref]. Seeing as we will be making repeated use of the term, we abbreviate topological quantum computing to *TQC*.

The most important idea in the subject is that TQC is inherently *fault tolerant*. Noise in the environment of a quantum computer, when properly controlled, can be made small in magnitude and local in effect. By the definition of a topological system, the information in the topological computer is invariant under small local changes. Hence, the information remains invariant under noise and the computation can proceed as intended without errors. If there are errors, which is always possible with some small probability, topological quantum systems typically have mechanisms whereby the experimenter can remove the error and restore the information how it was. This is the general picture for fault tolerance in TQC. We will make this picture more precise as we give examples of methods for TQC.

We additionally recall that TQC splits into two major branches. The first is the method of finding physical materials that naturally exhibit topological quantum behavior. These systems can then be used to make a computer. The second approach is to simulate a topological quantum system on a quantum computer. This simulation is used to inherit the fault-tolerant properties of the topological quantum systems on the original quantum computer. So long as the simulation is efficient and local noise on the physical system corresponds to local noise on the simulated system, this method works as described. This gives the following diagram for TQC:

[WORK: add diagram - its already used somewhere else!]

In this chapter, we will talk about lots of different approaches to TQC. Some of them are naturally amenable to the approach of topological quantum materials, and some of them are naturally amenable to the approach of topological quantum error correction. We will flag these differences and the status of experimental progress as we go along.

Before moving on with our discussion of TQC, it is good to be aware of the limitations of the algebraic approach.

1. Introducing the algebraic theory is a lot of overhead for not very many examples. Overwhelmingly, proposals for TQC center around just a few algebraic models. Topological quantum error correction is mostly centered around the toric code (see section [ref]), and topological quantum materials are mostly centered around Majorana fermions (see section [ref]). The vast majority of algebraic models have no serious proposals for TQC associated with them. It is for this reason that much of the literature is focused on working out the details of small models and examples, instead of the development of general theory which is largely useless in this lens.

2. The algebraic structures fail to capture a lot of important details about proposals for TQC. It only captures the high-level information flow, and none of the microscopic features. For example, a breakthrough in the field of topological quantum error correction come with the introduction of *color codes* in 2006 [BMD06]. These color codes have very nice properties, and have been an important player in the field of TQC. However, algebraically the color code is equivalent to bilayer toric code:

$$(\text{color code}) \cong (\text{toric code}) \boxtimes (\text{toric code}).$$

The entirety of the novelty of the color code comes in its specific choice of Hamiltonian and microscopic details - there is no new algebra involved.

3. There are many topological quantum systems whose algebraic theory is beyond description in terms of modular categories. These phases require increasingly intricate algebraic descriptions, past what is described in this book. [WORK: I'm really talking about *fractons* here. They were introduced in [VHF15] and have garnered a lot of interest since. What else is there?]

All this is not to say that the algebraic theory of topological quantum information is useless. It has been an important guide in the subject, and has provided footing and motivation for the continued development of TQC. Large-scale fault-tolerant quantum computation is one of the defining technological challenges of the 21st century. It seems very likely that topological methods will be part of its realization!

7.1.2 Universality

An important concept for understanding TQC methodology is *universality*. To illustrate this concept we begin with an example.

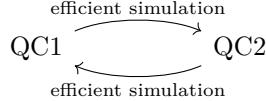
In 1994, Peter Shor developed an efficient quantum algorithm for factoring integers [Sho94]. This algorithm is important because much of modern cryptography is based on the hardness of factoring integers and similar problems. Hence, an efficient factoring algorithm could jeopardize internet security.

However, we must pose ourselves the question: what does it mean, really, for Shor to have found an efficient quantum factoring algorithm? A quantum computer is a computer whose information processing is fundamentally described by quantum mechanics. A priori there are lots of different quantum computers one could make. Which one did Shor find a factoring algorithm for? Maybe when we finally make a quantum computer it will be the type of quantum computer which cannot run Shor's algorithm so internet security will be safe.

The key in understanding this situation is the concept of universality. There are certainly quantum computers which cannot run Shor's algorithm. For instance, if a quantum computer has only a finite number of degrees of freedom to store information then it certainly cannot run large cases of Shor's algorithm because it can't even record the input! Even if a quantum computer can store arbitrarily large inputs, that doesn't mean it will have the capabilities to run Shor's algorithm because it might just not have a physical method for running the algorithm.

However, every *sufficiently powerful* quantum computer can run Shor's algorithm. Moreover, every algorithm you can run efficiently can run efficiently on one quantum computer can then be run on every other sufficiently powerful quantum computer.

Here is the explication for this phenominon. Computers can be viewed as simulation devices. Quantum computers are simultation devices which can simulate quantum systems. Suppose that you are given two quantum computers QC1 and QC2. If QC1 is sufficiently powerful, it should be able to efficienlty simulate the behavior of QC2. If QC2 is sufficiently powerful, it should be able to efficiently simulate the behavior of QC1, as illustrate below:



This gives an easy way to turn any efficient algorithm on QC1 into an efficient algorithm on QC2. First you simulate QC1, and then you run the algorithm on QC1! This means that any problem solved on one of the computers can be efficiently solved on the other. In this way these two computers are *computationally equivalent*. The non-trivial insight is that every sufficiently powerful quantum computer is able to efficiently simulate every other sufficiently powerful quantum computer. These powerful quantum computers which can simulate every other computer are known as *universal* quantum computers. The existence of universal quantum computers is the heart of universality. What Shor did was make a factoring algorithm which could be run on any universal quantum computer.

This sort of universality has been known for a long time. It was first proposed by pioneers of computation Alan Turing and Alonzo Church [Tur39, Cop97]:

Church-Turing thesis: “All sufficiently powerful computational models yield efficiently intersimulable classes - there is one theory of computation”.

Of course, this thesis does not account for the possibility of quantum computation. Classical and quantum computation seem to be turly distinct theories of computation, violating Church-Turing’s intention. This leads to a modern revised version of the Church-Turing thesis:

Revised Church-Turing thesis: “All sufficiently powerful classical computational models yield efficiently intersimulable classes - there is one theory of classical computation”.

A quantum version of the Church-Turing thesis was introduced in an early review article on topological quantum computation by Michael Freedman, Alexei Kitaev, Michael Larsen, and Zhenghan Wang [FKLW03]. It goes as follows:

Freedman-Church-Turing thesis: “All sufficiently powerful computational models which add the resources of quantum mechanics to classical computation yield efficiently intersimulable classes - there is one theory of quantum computation”.

The goal, then, is to make a *universal* topological quantum computer. In a sense this makes designing a scheme for topological quantum computation difficult. It gives constraints, and forces a certain amount of computation power. In another sense, it is liberating. The existence of universal quantum computation means that once we have implemendent a certain amount of computational power into our proposal, we are done. There is no need to search for clever ways to add more power because or system is already universal, and hence finding more techniques in unneccesary. It gives an end goal to building a quantum computer - a bell to ring when we are done.

Formally, a universal quantum computer will be one which can approximate any unitary transformation. This means that for every $n \geq 0$, the a universal quantum computer can be prepared in such a way that its information is stored in a Hilbert space V of dimension greater or equal to n . The space of possible computations on the computer should form a dense subgroup of the projective unitary group $PU(V)$.

One important question is whether a computer which is universal in the sense above can *efficiently* simulate any other computer is an important question. This is generically true by to Solovay-Kitaev theorem [Kit97, NC10]. However, a finer discussion of these points and other notions in computational complexity is beyond the scope of this book and is unnecessary for understanding the topics at hand.

[WORK: there is the general correlation between computational power, difficulty of implementation, and non-abelian flavor. Give the nice table and talk about it.]

[WORK: boson fermion - easy to simulate. Needs to exploit the weird, complicated (non-abelian) nature of braiding. Hence needs to be sufficiently non-abelian, hence the above picture.]

7.2 Computation with Fibonacci anyons

7.2.1 Methodology

[WORK: I'm realizing that I don't know enough about Fibonacci anyons to write this. What's the deal with the "golden chain"? How does the relationship with $SU(2)$ work again? What's the history? What does the density of braid group reps say, exactly, and how does the proof go?]

7.2.2 The Jones invariant

[WORK: here we can connect things to the Jones invariant. The first step is the definition via the Kauffman bracket. The key part is to observe that the Skein relation can be enforced physically by finding an anyon (i.e. Fibonacci) which satisfies the Skein relation as operators on a Hilbert space! Such a Skein relation *must* exist, and hence this also explains in a fundamental way why there is a good Skein relation which gives a knot invariant. Also motivates quantum topology in a way, so maybe say a few words about quantum topology? A good reference is [AA11] which gives a self-contained elementary proof of BQP completeness of Jones invariant.]

7.2.3 Proof of universality

[WORK:

Prove that the braid group representations are dense in the appropriate sense.

There's lots of different ways of doing this of course, coming from the work of different authors. One fun and informative way of doing this could be to use *Kliuchnikov's exact synthesis* method. This method is due to [KBS14]. Important follow up works are [KBRY15] and [KY15]. This method is nice because it demonstrates the power of number theory and deep mathematical ideas in topological quantum information. This is a textbook about the algebraic theory, after all!

]

7.3 Computation with doubles of finite groups

.[WORK: need to read Mochon's papers and Zhenghan's clarifications again to write this section. Maybe have two sections, one for non-solvable and one for solvable non-nilpotent?]

.[WORK: Maybe add a little word about the idea of having \mathbb{Z}_2 bulk and then interfacing with S_3 islands?]

.[WORK: using gapped boundaries. This is the surface code.]

.[WORK: maybe bring up universal TQC with gapped boundaries in $\mathfrak{D}(\mathbb{Z}_3)$ and projective charge measurement somehow? It would be nice to include somewhere. Maybe have an intro about how hard you have to try to get universal TQC with different groups, and how you can get universal TQC even with abelian groups if you try hard enough.]

]

7.4 Computation with the toric code

[WORK: A good (older) reference which I've never read is [DKLP02]]

7.5 Computation with Ising anyons

.[WORK: universal TQC with Ising twsit defects]

.[WORK: There is a general picture of computing using mapping class group representations on high genus surfaces. This is discussed in its most general form in [BW18]]

.[WORK: I would like this section to include a general discussion of genons. This includes the fact that the obstruction to gauging vanishes and that they actually behave like genons as we would want them to. The relevant papers are [GJ19, Del19, BJQ13]]

7.6 Computation with Majorana zero modes

.[WORK: Theory of Majoranas, as \mathbb{Z}_2 -crossed extensions of sVec.]

.[WORK: This section requires a real discussion of physics. There are three key systems to discuss.

1. $\nu = 5/2$ FQH. This system is described by a supermodular category, up to the typical caveat that it is only quasi-topological order and not pure topological order so some phases might not be protected. This supermodular category has Ising as a subcategory. However, the simple object which makes the nonabelian anyon in the Ising MTC is *not* “fundamental” in the system. It is composite, made out of two physically creatable anyons. In this sense $\nu = 5/2$ doesn’t really have fundamental Ising anyons, only composite ones.
2. The ends of nanowires. Kitaev has his famous paper introducing this idea. These are Majorana zero modes in their purest form. This is NOT ising. It is a \mathbb{Z}_2 -crossed extension of sVec which is algebraically essentially the same as Ising, but the distinction is that some of the phases which are well-defined in Ising are unphysical in the \mathbb{Z}_2 -crossed extension. (original paper is [Kit01], and then [FK11] has big improvements)
3. Superconductor/topological insulator heterostructures. If you have a sample of topological insulator and you make its boundary conditions oscillate between magnet and superconductor you get Majoranas at the interface between those boundaries. The online course on topology in condensed matter has a good section on this, and there is a

lot of literature on the subject. Algebraically, this should be the \mathbb{Z}_2 -crossed extension of sVec as well. This could be a good reference: [SA19].

Pointing out the key subtle differences between these models is of utmost importance. There should be sections summarizing each experiment and describing its algebraic theory. Another thing to know about is intrinsically topological superconductors, c.f. [BNQ13].]

[WORK: This could become a lot of work. It is very relevant to physicists (perhaps the most relevant part of this book), but unnecessary and cumbersome for mathematical thinkers. Maybe this should be its own chapter?]

History and further reading:

The idea of topological quantum computing was first introduced by 1997 by Kitaev and Freedman [Kit03, Fre98]. Soon, Freedman, Kitaev, Wang, and Larsen wrote a review article about topological quantum computing which formally started the field in 2002 [FKLW03]. In these early years, these authors and others introduced a number of techniques for universal topological quantum computation [FLW02, Moc03, BK05]. From here, the goal of research became the task of achieving universal topological quantum computation in the simplest possible experimental setup.

In the world of quantum materials, this has mostly taken the form of hunting for *Majorana bound states*. Majorana bound states are topological quasiparticles which are bound to defects in materials. Some theories suggest that these Majorana bound states could be braided in a fashion which allows for topological quantum computing. Algebraically, Majorana bound states behave as [WORK: what do they behave as?]. Theorists have engineered increasingly simple materials which are predicted to host Majoranas [FK08, STL⁺10, Ali10]. Braiding Majorana bound states does not allow for universal topological quantum computation, so most proposals for Majorana quantum computing include some non-topological gates.

In the world of quantum error correction, the search for simple experimental setups has centered around the surface code. The surface code on its own is not universal, and requires a single extra gate to be made universal. There have been a large number of proposals for how to do this final extra gate, which are more or less feasible depending on the architecture of the underlying quantum computer [BK05, BMD11, Bom15].

There are many good references for topological quantum computing. From the perspective of materials, there are several excellent review articles by Freedman, Nayak, Das Sarma, and others [NSS⁺08, SFN15]. From the perspective of topological quantum error correcting codes, the best approach to learn more is to delve into the general theory of quantum error correction. A good place to start is the chapter in Nielsen-Chuang [NC10]. After this there are several review articles [Ter15, Got97].

Exercises:

7.1. .[WORK: make exercises]

A Odds and ends

[WORK: I'm not sure where to do this, but I'd like to make a little comment about non-semisimple modular categories. There are inherent limits to the applications of MTCs to quantum topology, as shown in these papers [Reu23a, Dav11]. This motives going beyond semisimplicity.

A nice paper about this from Zhenghan's perspective is [CKWZ24].

The canonical reference is [CDGG21]. The summary is that “the only physical thing is the derived category”. Makes me look at derived categories differently. Another important paper in this area is [Shi19].

]

A.1 Topological quantum field theories

[WORK: re-do now that this has been moved]

In this chapter we will be exploring topological quantum field theory, a particular way of mathematically formalizing topological order. We recall below how this fits into the general framework of this book:

[WORK: add pretty picture.]

Topological quantum field theory is not our primary approach for understanding topological order - we will mainly be performing our analysis of topological order using modular tensor categories. For this reason, the present chapter is essentially auxillary to the rest of the book. This chapter is primarily included for the following reasons:

1. To help give a comprehensive picture of the algebraic theory of topological quantum information. Much of the work in the field is written in the language of topological quantum field theory. Not knowing topological quantum field theory can make existing in the world of topological quantum information more painful than it needs to be.
2. Topological quantum field theory makes some important ideas clear which are opaque in the language of modular tensor categories. For instance, the proper way of interpreting the modular representation of a modular tensor category is to use concepts of topological quantum field theory.
3. In contrast to modular tensor categories, the definition of a topological quantum field theory is very concise. This means that the fact that topological quantum field theories and modular tensor categories are equivalent is a strong indication that the definition of modular tensor category is well-chosen. As we will discuss in Chapter [ref], the definition of modular tensor category was explicitly chosen to be the way it so that they would be equivalent to topological quantum field theories.

For brevity, we will use the acronym *TQFT* to abbreviate Topological Quantum Field Theory. We will now move on to describing the big idea of TQFT. We will do this by starting with an abstract topological order \mathcal{C} . Of course, we haven't defined what an abstract topological order is yet. The point is that we will imagine that \mathcal{C} is some family of gapped Hamiltonians which has topologically protected ground spaces. All the different gapped Hamiltonians should be related in the sense that they have the same underlying algebraic theory, though we haven't described what that theory is yet. That underlying

algebraic theory will exactly be a TQFT. Because they are the only examples we have given, we will always imagine that \mathcal{C} is the Kitaev quantum double model based on some finite group G , or even more specifically the toric code.

The first step in defining TQFT is to think about \mathcal{C} is a machine which takes in topological spaces and spits out quantum systems, by taking the space and putting the topological order \mathcal{C} on it. For instance, if \mathcal{C} is the Kitaev quantum double model based on a finite group G and the input space is a torus, then the corresponding system is the Hamiltonian $H = \sum_v (1 - A_v) + \sum_p (1 - B_p)$ defined by choosing a lattice structure on the torus and collecting flatness and gauge invariance conditions, as pictured below:

[WORK: add picture.]

The Hilbert space $\mathcal{N} = \bigotimes_{\text{edges}} \mathbb{C}[G]$ of this system clearly depends on the choice of lattice on the torus. However, as demonstrated in proposition [ref] the ground states of \mathcal{N} are $\mathcal{C} = \mathbb{C}[\text{Hom}(\pi_1(T^2, v), G)/(\text{conjugation})]$. Since the fundamental group is a topological invariant, we see thus that this Hilbert space does not depend on our choice of lattice - its dimension is fixed by the topology of the torus.

More generally, this is what we should expect when putting a topological order on space. The excited states will depend on the details of the gapped Hamiltonian we choose, but the ground states are a topological invariant of the space. The fact that the ground states are topologically invariant is the *defining feature* of topological order. Hence, the topological order \mathcal{C} gives a well-defined assignment from topological spaces to quantum systems.

Not every topological space can host topological order, however. We recall that our definition of topological order required physical space to be *two dimensional*. Of course, there can be global curvature like in the torus. What's important is that locally the system is flat. Hence, the topological spaces on which we can apply our topological order are subspaces of \mathbb{R}^3 which locally look like \mathbb{R}^2 . Topological spaces of this type are called *surfaces*. The most important examples are the g -holed surfaces, for any $g \geq 0$, called Σ_g :

[WORK: add Σ_g picture]

Hence, we find that every topological \mathcal{C} induces an assignment

[WORK: add formula - surfaces get assigned to the Hilbert space of ground states of \mathcal{C} on that surface.]

This is the general picture for TQFT. A TQFT is an assignment from surfaces to vector spaces, with extra restrictions which are required to get a reasonable theory. This sort of approach to quantum field theory can be generalized beyond TQFT. In these generalized cases, however, the assignment won't take surfaces as inputs. Instead, it will take surfaces with detailed geometric structure which encodes the fact that the resulting quantum system will depend on distances and local geometric information. This approach is feasible in some cases but is quite technical [Seg88]. Typically in non-topological cases people opt for other techniques.

A.2 Quasitriangular weak Hopf algebras

[WORK:

Weak Hopf algebras were introduced in ???. A good early source about them is [Nik04].

Weak Hopf algebras are relevant to the algebraic theory of topological quantum information because the representation category of a weak Hopf algebra is a fusion category. Adding more structure to the weak Hopf algebra gets you all the way up to modular categories. This is Tannaka duality in action. The reference for tannaka duality for modular categories is [Pfe09].

They are also intimately linked to the theory of module categories. This was first established in [Ost03], and then was shown much more explicitly in [KK12].]

A.3 Quantum groups

A.4 Subfactors

A.5 Vertex operator algebras

[WORK:

The connection between vertex operator algebras and topological order comes through conformal field theory. VOAs are at their heart tools for conformal field theory. Of course, since algebraically conformal field theory and topological field theories are so similar, this means that well behaved VOAs describe topological order.

This was first proved in the landmark paper of Huang [Hua05]. Of course, there are versions for G -crossed and fermionic theories - [Hua21, CGH23].

One very nice thing to be aware of is the work of Nikita Sopenko. He is able to prepare topologically ordered states using vertex operator algebras, thus realizing the implicit program in the topological order interpretation of Huang's work [Sop23].

A big thing in all of this is the Kazhdan-Lusztig correspondence, which I do not understand very well. A great reference seems to be [Tan20].

]

B Anyon data

B.1 Low-rank modular categories

.[WORK: list of all low-rank MTCs.
The relevant papers are [RSW09, BNRW16].]

B.2 Abelian modular categories

.[WORK: classification of abelian MTCs, give data for a lot of examples]

B.3 Group-theoretical modular categories

.[WORK: give theorems to characterize all of the data for group-theoretical MTCs, give generously many examples]

B.4 Miscellaneous examples

.[WORK: miscellaneous high-rank non-abelian non-group theoretical categories of interest. Probably Haagerup and E6 subfactors [HRW08]. Maybe $SU(2)$ quantum group MTCs for various roots of unity would be nice too.]

.[WORK: I would also like to add some fermionic modular categories and symmetry enriched modular categories. Where should I put those?]

[WORK:

List of things to add:

- A big “physics-math dictionary” which allows the translation of everything;
- A table of notation;

]

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