

Everettian chance in no uncertain terms

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Abstract

The current landscape of views on the role of chance in the Everett interpretation is rocky. Everettians (Wallace 2012, Sebens and Carroll 2018, McQueen and Vaidman 2019) agree that chance values should be derived using principles governing uncertain or partial belief, but cannot agree on how. Critics (Baker 2007, Dawid and Thébault 2015, Mandolesi 2019) maintain that any such approach is circular. We smooth the landscape by shifting focus from what Everettians take to be uncertain to what they should think is certain: namely, the conditions under which branches are isolated. Our approach to isolation resolves the main tensions among the different Everettian chance derivations while clarifying how they avoid circularity.

Contents

1	Introduction	2
2	Most attitudes towards ‘chance’ tacitly invoke certainty	3
3	Certainty of a system’s isolation can resolve the reference class problem	6
4	But ‘isolation’ is ambiguous	10
5	Probing Everettian ‘isolation’ with theory sectors	15
5.1	Orthodoxy’s sectors for microstates and macrostates	17
5.1.1	Orthodoxy’s microstate sectors	17
5.1.2	Orthodoxy’s macrostate sectors	19
5.1.3	Symmetries in orthodox macrostates	21
5.2	Challenges to orthodox isolation	22
5.3	Revising the sectors to accommodate non-repeatability	23
5.4	Revising the sectors for Everettians	24

5.5	Challenges to Everettian isolation	29
5.5.1	Isolation does not require a measure of chance	30
5.5.2	Isolation might be challenged by how decoherence spreads	31
6	Sectors unite and strengthen Born rule derivations	32
6.1	Wallace’s two-branch, equal-weight case	33
6.2	Sebens and Carroll’s two-branch, equal-weight case	38
6.3	McQueen and Vaidman’s three-branch, equal-weight case	43
7	Discussion	45

1 Introduction

The current landscape of views on chance in Everettian quantum mechanics is a curious one. On the one side, longstanding critics of many-worlds theories maintain that probability is needed to make sense of the machinery that Everettians use to derive chance values, resulting in circularity (Baker 2007, Dawid and Thébault 2015, Mandolesi 2019). On the other side, Everettians seem to agree that chance values should be derived in terms of agents’ uncertain or partial beliefs—but they cannot agree on how.¹

Perhaps the most famous of these uncertainty-based approaches is the decision-theoretic program explored by Deutsch (1999) and Wallace (2012): they purport to prove that a rational Everettian agent must order their preferences over acts in a way that recovers the Born rule. Sebens and Carroll (2018) take issue with one of the principles of rationality in Wallace’s approach and instead aim to derive Everettian probabilities from principles governing self-locating uncertainty. McQueen and Vaidman (2019) offer yet another self-locating uncertainty approach, taking issue with Sebens and Carroll’s metaphysical view of branching. Notably, all three approaches claim that symmetries of quantum states are central to their arguments, but none attempt to characterize the symmetries at play.²

We suggest a shift in focus from what these Everettians take to be uncertain to what they take to be certain. In fact, we can fruitfully understand these Everettians as holding a tacit commitment to a functional link between chance and certainty—namely, that chances supervene on the totality of physical information relevant to the relative frequencies. We flesh out this link with an *ontic separability principle*, which says, roughly, that chance

¹While this is arguably the majority view, not all parties agree. For instance, Wilhelm (2022, 2023) gives an account of Everettian probability, which he calls the ‘Centered Everett Interpretation’, that takes the Born rule to prescribe the objective chances that a particular centered proposition is true. Wilhelm’s view seems to suggest that self-locating uncertainty is derived from the objective chances, rather than the other way around—a moral that we endorse. However, Romagosa (2024) identifies a tension between Wilhelm’s approach and many Everettians’ desire not to place special emphasis on observers (among other potential issues). While this dialectic warrants further attention, we set it aside for our present purposes; our approach avoids using centered propositions altogether, thereby hewing closer to Everettian orthodoxy. We thank an anonymous referee for raising this point.

²In this article, we focus on the approach of Wallace (2012) over that of Deutsch (1999) simply because doing so suffices to make a series of (we hope fruitful and interesting) comparisons with Sebens and Carroll (2018) and McQueen and Vaidman (2019) (in e.g. §6) without overly complicating our narrative; doing so also gives our article parity with Chua and Chen (2025).

assignments should supervene on, and respect the appropriate symmetries of, the states of isolated subsystems. We sketch a thin functional analysis of isolation that allows Everettians to identify branches as isolated subsystems with only a partial assignment of trivial chance values (i.e., chance values of either 0 or 1) and without a principle of locality. This analysis, in turn, provides a characterization of the symmetries of branches relevant to chance in non-relativistic quantum mechanics.

We argue that chance’s functional link with certainty does most of the explanatory work in deriving specific, non-trivial chance values: to wit, it does so by characterizing the precise sense in which symmetric, isolated states should yield equivalent chance assignments. This characterization yields a thin core approach to deriving chance in Everettian quantum mechanics that smooths the current rocky landscape. Each of the decision-theoretic and self-locating uncertainty programs discussed above can adopt this core approach while (a) changing very little of the structure of their arguments and (b) addressing the main criticisms they have levied at each other. Moreover, since our notion of a branch’s ‘isolation’ explicitly requires an assignment of 0 and 1 chance values but nothing more, the core approach clarifies the sense in which derivations of full chance measures avoid the usual charges of circularity.

Our plan for the article is as follows. In §2, we consider the various functional roles of chance which have been identified in the literature to date, and add to this mix the certainty link, which will play a central role in this article going forward. In §3, we show how fleshing out the certainty link with our ontic separability principle can resolve versions of the reference class problem faced by objectivist and subjectivist analyses of chance alike. In §4, we offer a thin functional analysis of ‘isolated subsystems’ to which our separability principle refers: namely, that (a) such systems should instantiate a recursive structure, such as Wallace’s (2022a) ‘theory sectors’, and (b) one should be able to study these ‘isolated subsystems’ in a laboratory. Then, in §5 we use these functional criteria to argue that Everettians should adopt two separate notions of isolation: one for microstates and one for macrostates. We show that while our notion of isolation for macrostates relies on a partial interpretation of the Hilbert space norm using chance-adjacent concepts, it does not depend on any prior complete assignment of chance values. Having done this, we demonstrate in §6 how our ontic separability principle specifies a core approach to deriving complete chance assignments. We then outline how each of the above-mentioned Born rule derivations deviates from the core, as well as how the core resolves the primary tensions among them. In §7 we wrap up.

2 Most attitudes towards ‘chance’ tacitly invoke certainty

Much cross-talk in philosophical literature stems from the difference between two key philosophical methodologies for confronting troubling terms: namely, *conceptual analysis* and *conceptual engineering* (Cappelen 2018). The former aims to give a precise definition of the term under consideration that yields maximal fidelity with all its uses in everyday talk. The latter aims to create a novel definition for the term that better suits a set of desired uses. Whereas the former is primarily descriptive, the latter is primarily normative.³

³For a discussion of this distinction in a very different context in the philosophy of physics, see Krátký and Read (2024).

Conceptual analysis and conceptual engineering, so construed, are meant to identify endpoints on a spectrum of methodologies. For example, Quine’s version of conceptual analysis (which we favor) falls somewhere in between. Quine, drawing on Carnap (1947, §2), rejects the sort of analysis that aims to ‘expose hidden meanings’; instead, he thinks analysis should ‘fix on the particular functions of the unclear expression that make it worth troubling about’ (2013, p. 238). One might summarize the pursuit of a Quinean analysis of ‘chance’ as an attempt to answer the following question: what would be a happy thing for ‘chance’ to mean, given roughly how we use it today?

For our attempt at answering this question, we follow Quine in fixing on the *functions* that ‘chance’ plays or might play in daily discourse. This approach already puts us in the good company of philosophers of probability such as Hájek (2019), and Everettians such as Wallace (2012) and Saunders (2010), all of whom aim to give ‘chance’ a functional definition. Given our present focus on chance in Everettian contexts, we will begin our search for a functional definition with the functional links that Everettians have deemed most important. All the chance functionalists we survey agree that the following links are *descriptive* of our use of the term ‘chance’. It varies whether any given chance functionalist treats any given link *normatively*, as a principle of desired use.

To begin, then: Saunders (2010), following Papineau (1996), helpfully identifies the following two functional roles which one might demand be satisfied by a notion of objective chance:

- (C1) **The inferential link:** The chance of an event is measured (roughly) by (actual)⁴ relative frequencies of that event; and
- (C2) **The credential link:** All else being equal, one’s subjective degree of belief or credence in an event ought to equal the chance of that event.

Brown (2011, p. 7), Wallace (2012, §4.4), and Read (2018, p. 137) all endorse that these links (or slight variations of them) are at least descriptive of chance talk.⁵

In addition to the above two functional roles for chance, Saunders (2010, p. 181) goes on to identify another:

⁴Here, our parenthetical ‘actual’ is intended to emphasize a contrast with hypothetical relative frequencies, recalling the standard distinction between actual versus hypothetical frequentism; see, e.g., Wallace (2012, ch. 4).

⁵We should be clear at the outset that while these links are popular, key parties to these discussions have starkly different attitudes towards the use of the term ‘chance’. For example, while Wallace (2012) takes Lewis’s (1980) principal principle to realize the functional roles of chance and thus yield *bona fide* objective chances, Deutsch (2016) does not make this move. Others, such as Greaves (2007), are willing to endorse much of Wallace’s interpretation of an Everettian decision-theoretic strategy while reserving the term ‘chance’ for descriptions of outcomes in single-world theories. Likewise, Vaidman (2012, p. 299) states that ‘there is no “probability” for an outcome of a quantum experiment in the usual sense’, and one could pursue a similar reading of Sebens and Carroll (2018). One might quite reasonably worry that it is inappropriate to put the word ‘chance’ in the mouths of these authors. To clarify, then, our focus in this article is on the various strategies for ‘deriving the Born rule’ proposed by Deutsch (1999) and Wallace (2012), Sebens and Carroll (2018) and McQueen and Vaidman (2019), which we propose can all be fruitfully understood as fleshing out a thin functionalist approach to ‘chance’—albeit not the exact functionalist approach proposed by Wallace (2012), as we discuss in the remainder of this section. We thank an anonymous referee for drawing our attention to this issue.

(C3) **The uncertainty link:** Chance events, prior to their occurrence, are uncertain.

Clearly this link is descriptively adequate, but should we add it to a list of normative requirements? Given that many-worlds theories specify that all possible measurement outcomes are actual, one might be wary of the normative upshot of uncertainty. One can read Greaves (2007), for example, as building an account of Everettian chance with the first two links alone. Still, Saunders and Wallace (2008) have taken pains to show how the Everettian *can* recover this link in a normative sense, if they so desire.

One immediate question to ask of the uncertainty link—whether one treats it as descriptive or normative—is this: should it have anything in particular to do with the *measure* of chance? To be sure, we might like to recover the lived experience of observers—but it seems a stretch to think that uncertainty about outcomes *per se* would be enough to inform us about the right choice of chance function.

Now, one way in which one might leverage the uncertainty link in order to get a bead on a specific chance measure is via the ‘principle of indifference’, which states that in the absence of any relevant evidence, agents should distribute their degrees of belief equally among all the possible outcomes under consideration—see Eva (2019) and references therein. However, even availing oneself of the principle of indifference is unlikely to fix a unique such chance measure, absent further constraints. This might be illustrated by familiar problems with applying the principle of indifference to an infinite state space. Consider the classic example of the hypothetical box factory due to van Fraassen (1989, p. 303), in which the principle of indifference is unable to discriminate between (e.g.) length- and area-preserving probability measures.⁶ To paraphrase, imagine that you are told that the factory produces cubic boxes with sides of any length from 4 to 5 centimeters and given no further information. What are your best estimates for the face area and the side length of a given box? If you apply the principle of indifference to side length, you get the uniform measure over the interval [4, 5] (in units of cm); if you apply the principle to face area, you get the uniform measure over [16, 25] (in units of cm²). The problem is that these choices yield estimates of 4.5 cm for the box’s side length and 20.5 cm² for the area of its faces, and *no* cubic box can have both of these measurements. To get consistent answers, only one of the two probability measures can be uniform—and the principle of indifference seems to offer little guidance as to which one it should be.

Ideally, one would seek more information about the box factory to resolve the paradox. Suppose, for example, that one learns that the factory starts with an imprecise rod-cutting machine, then nearly perfectly creates eleven copies of that initial rod to piece together a cubic frame. This additional information suggests that the chance value for a finished box’s parameter to lie within a given range should supervene on—i.e., be a function of—facts about how the rod-cutter operates. Along this line of reasoning, if one indeed strives for the recovery of a *specific* chance measure, then the following additional functional role might serve one better in making good on that:

(C4) **The certainty link:** Chance values supervene on the totality of physical information

⁶Van Fraassen credits Bertrand (1889) with the insight driving his example. Keynes (1978, pp. 48–49) gives a very similar version of the thought experiment (using volume and density rather than length and area), which he credits to von Kries (1886). Note that Adlam (2025) has also recently invoked Bertrand’s paradox to call into question subjective uncertainty based approaches to fixing an objective chance measure.

that is relevant to the (actual) relative frequencies of the event.

In particular, an intuitive application of the certainty link to our reasoning about the box factory suggests that there is some sense in which the rod-cutting process is *isolated*, allowing us to ignore any other details about the construction process.

In at least this case, the certainty might prove necessary to derive a specific measure of chance, since the uncertainty link does not suffice on its own. But in fact there is nothing to stop us having both (C3) and (C4)! We can recover lived experience (via (C3), as articulated above) *and* potentially obtain a unique chance measure (via (C4)); it's just that the explanation of the former has little to do with the explanation of the latter.

The promise of a solution to the box factory paradox provides some initial evidence that the certainty link latches onto something important in our usual 'chance' talk. In the next section, we adduce more evidence for this claim by arguing that both objectivists and subjectivists about chance can make good on that promise—and in the process, resolve their respective versions of the reference class problem.

3 Certainty of a system's isolation can resolve the reference class problem

We claim that the link with certainty—that chance values supervene on the totality of information relevant to the actual frequencies—plays an important role in our usual 'chance' talk. To defend this claim, we now show how several prominent objective and subjective analyses of chance can flesh it out to resolve their respective versions of the reference class problem, using the box factory paradox as a guiding example.

First, for objective analyses, consider various versions of frequentism. One can view frequentism as an approach to chance that begins by 'filling in' the inferential link with an analytic definition, one that identifies a chance value with with a (possibly limiting) relative frequency of a given event in an actual or hypothetical *reference class* of events of a certain type. For example, a finite frequentist might fill in the inferential link (C1) by defining a chance function, ch , as follows:

$$ch(X | R) = \frac{n(X)}{N}, \tag{1}$$

where X is an attribute and R is a finite reference class with N elements, and $n(X)$ gives the number of occurrences of X in R .⁷ Selecting a reference class becomes a problem if it is unclear what types of events are most relevant—and so it is unclear how to specify R . Suppose, for example, that one wants to assess the chance that their grandmother Sarah, a serial smoker, will contract lung cancer. One can do so by comparing the frequency of lung cancer cases in the total population to the frequency of cases among only folks like Sarah. However, it is not immediately clear which factors making someone 'like Sarah' ought to

⁷This particular definition of finite frequentism is due to Hájek (1997). Throughout this paper, for ease of exposition, we allow the arguments for the chance and credence functions to be attributes, propositions, or events. However, nothing in our argument should hinge on this permissive approach. For example, the reader should feel free to rewrite a definition using attributes in terms of propositions if they prefer.

be included in one’s assessment.⁸ This problem remains if we allow ourselves to consider hypothetical agents, expanding our reference class to be infinite; thus, the problem plagues hypothetical frequentism, too.⁹

The reference class problem does not exclusively affect these frequentist analyses of probability, however. Hájek (2007) argues cogently that it is a problem for subjectivists too (among others). He suggests that most subjectivist analyses of chance start with the rough idea that you ought to conform your degrees of confidence or belief to those of some better-informed expert, real or hypothetical. In other words, these approaches roughly suggest that one fills in the credential link (C2) with the following schema:

$$c(X \mid c_{\text{expert}}(X) = x) = x, \tag{2}$$

where c is your credence function and c_{expert} is some expert’s credence function. Hájek then notes that there are various legitimate-seeming ways of picking the right expert that might not all agree. For example, he considers that c_{expert} might be a Lewisian objective chance function, the logical probability given some total body of evidence, or a function derived from some physical theory T .¹⁰ Nothing stops these functions from conflicting with each other. Indeed, the box factory paradox provides a concrete example, as a subjectivist about chance might variously try to set c_{expert} by applying the principle of indifference to side length, face area, or box volume. As Hájek puts it, ‘You can’t serve all your masters at once, so you have to play favorites. But who trumps whom, and which trumps which?’ (2007, p. 597).

For the purposes of assigning chances to events that we model with some physical theory T , we assume that T trumps everything else. We view this assumption as enforcing self-consistency: if we think, for instance, that quantum theory is good enough to exhaustively describe a system’s observables, it had better be good enough to supply our expert credence or chance function about values of those observables, too. This pragmatic choice does not solve Hájek’s various reference class problems once and for all, as one is always free to reject that physical theory. However, the choice provides a happy way of resolving the issue once we have put our theorist hats on; i.e., once we are acting as if the world behaves as our theory describes.¹¹

Of course, wearing this hat, one might easily resolve any reference class problem by taking one’s physical theory T to stipulate a chance measure by fiat. How, instead, might a physical theory T explain why a given measure of chance is the right one? For a first pass, let’s return to the box factory, and in particular, to our intuition that more information about the construction process might help resolve the paradox. To construct T , let U describe the total process of constructing a single box, and let S describe the operation of the imperfect

⁸Venn (1876) presents an early version of this example.

⁹We will focus on definitions of chance that directly specify real-valued chance functions, such as hypothetical frequentism and the subjectivist ‘expert’ schema (2), below. However, see Saunders (2021, 2024) for a discrete Everettian approach that uses a version of finite frequentism. In this approach, real-valued probabilities emerge from appropriate limits.

¹⁰Moreover, while Lewis’s (1980) principal principle is one way of making the expert schema (2) precise with a chance–credence norm, it is not the only way. See Pettigrew (2012) for an overview of various alternative proposals and the key ways in which they differ.

¹¹This move is inspired by Arthur Fine’s (1986) characterization of Einstein’s approach to scientific realism, but note that no particular stance on realism *per se* is required for this paper’s argument.

rod-cutter; moreover, let T contain our previous assertion that S is isolated from the rest of U . Now, (C4) implies that the side length s fixed by the rod-cutter is the only physical parameter of the finished box that could be relevant to the chances. Should T now appeal to the principle of indifference and set the chances for ranges of the box's properties to agree with the uniform measure over side lengths? Perhaps not; at this stage in our theory-crafting, nothing physical seems to stop the imperfect rod-cutter from, say, producing many more rods closer to 4 cm than 5 cm. Thus, the uniform measure over $[4, 5]$ might *not* be measured, even roughly, by the actual relative frequencies of side lengths, violating (C1). In other words, to get a chance function satisfying our functional links, we might need to pay attention to the physical parameters of the rod-cutter, too.

So, for a second pass, let's add a bunch more information about the rod-cutter. Say the initial rod is fed into the cutter via a conveyor belt, where one end of the rod is flush with the left edge of the belt, and the other end hangs off the belt's right side. The rod-cutter consists of a series of sharp guillotines, one every centimeter, hovering over the right edge of the belt. Each guillotine is $\sqrt{2}$ cm long and set at a 45-degree angle from the edge, and the belt moves rods along at a rate of 1 cm per second, such that the length of an initial rod depends on the time t at which the guillotines drop. We have picked these dimensions so that if the initial rod ends up with a length of 4 cm, then if we had dropped the guillotines .5 seconds later, the rod would have been 4.5 cm long. In general, letting $s \in [0, 1]$ be the additional length of the rod past 4 cm, adjusting the drop time by i seconds corresponds to addition by i modulus 1 to s . Next, let us suppose that conveyor belt is divided into cells 1 cm long that line up with the guillotines at $t = 0$, each containing a rod, such that frequencies of positions of the rod within a cell are well-modeled by some unknown probability density function. Whatever this function is, it is identical to the density $\rho_{t=0}(s)$ giving the chance values for rod lengths when the guillotines drop at $t = 0$. In turn, let us say that T specifies a *state* of S by some distribution or other $\rho_t(s)$, such that there is one state for each t . At the very least, we have now arrived at a theory whose states fix the chance values.¹²

Unfortunately, even *this* much information is not enough to pin down a unique measure of chance! Granted, we can place some heavy constraints on the structure of states of S : in particular, we know that changing the time t at which the guillotines drop simply moves the probability density $\rho_t(s)$ up or down the unit interval (with wrap-around). A bit more formally, any time translation of t to $t' = t + i$ imposes a translation of s to $s' = s + i \bmod 1$ and a transformation of the state $\rho_t = \rho$ to $\rho_{t'} = \rho'$ satisfying

$$\rho'(s) = \rho(s'). \tag{3}$$

However, ρ can still specify any distribution over the unit interval. If T is truly the best we can do (and we still have our theorist hats on), we might now consider giving up on trying to solve the reference class problem, as there is no reason to suppose any ρ does a better job of modeling the relative frequencies than any other.¹³

In some cases, however, we can do better. In particular, for our third and final pass at T , we add the following crucial fact: shifting the time at which the guillotines drop for each run

¹²The authors thank Rami Jreige and Karim Thébault for assistance in the construction of this example.

¹³In this case, a subjectivist might use an imprecise probability space to formally structure their beliefs rather than the expert schema (2). See Bradley (2019) for an introduction to these spaces, and see Chandler (2014) for an extended argument that the box factory paradox motivates their adoption.

of the rod-cutter does not make an observable difference to the relative frequencies of side lengths. Formally, this fact asserts that any two states of S related by a translation are, in fact, identical; i.e.,

$$\rho' = \rho. \quad (4)$$

Combining (3) and (4), we obtain

$$\rho'(s) = \rho(s') = \rho(s), \quad (5)$$

and this last equality suffices to prove that $\rho(s)$ is uniquely specified by the uniform measure over the unit interval.¹⁴

To put these statements in terms of a more traditionally-defined chance function, consider arbitrary measurable subsets X of the unit interval, representing (roughly) ranges of side lengths at a given time. Adjusting the time that the guillotines drop by i seconds shifts any such range to $X' = X + i \bmod 1$, where addition is performed element-wise. With minor abuse of notation, let S and S' specify the states ρ and ρ' , respectively. Now, the structural constraint (3) amounts to the following claim:

$$ch(X | S) = ch(X' | S'); \quad (6)$$

likewise, the identity of states (4) becomes the claim $S = S'$, which, in turn, implies $ch(X | S) = ch(X' | S)$.

Let us draw out three important features of our third pass at T that led us to a specific chance function. First, as in our initial pass, T establishes a sense in which S is *isolated* from the rest of U ; thus, we should be able to assess the chance value for a range of side lengths X without considering any other part of the construction process. Second, as in our second pass, T establishes enough structure to define a *chance-preserving map* between states of S : namely, arbitrary translations of a range of side lengths X move a chance distribution up or down the unit interval. Finally, from our last pass, T asserts that states related by the chance-preserving map are *identical*, allowing that map to impose constraints on the measure assigned by a single state. We abstract each of these constraints into a general tool that we can use to fill in the certainty link (C4). We call this tool the *ontic separability principle*, or **OSP**:¹⁵

OSP: Suppose that a physical theory T specifies the totality of physical information that is relevant to the (actual) relative frequencies of an event X in a system U . Moreover, suppose that X occurs in an isolated subsystem S of U , and that a map from X in S to X' in S' preserves its relative frequency. Then T should assign chance values to X in S that are independent of the environment $U \setminus S$ and preserved by this map:

$$ch(X | U) = ch(X | S) = ch(X' | S'). \quad (7)$$

¹⁴By gluing the endpoints of the unit interval together, the proof of this fact becomes the same as showing that the translation invariance of a field defined on the unit circle S^1 implies that that field is isotropic, on which see, e.g., Marinucci and Peccati (2011, p. 13).

¹⁵Our ontic separability principle is inspired by, and modeled after, Sebens and Carroll's (2018) *epistemic separability principle*, or **ESP**. We detail the key differences between these two principles in §6.2.

Note that when $S = S'$, equation (7) implies $ch(X | S) = ch(X' | S)$; thus, **OSP** tacitly contains our identity constraint, in addition to the constraints imposed by isolation and chance-preserving maps.

Informally, **OSP** says that the chances of outcomes in an isolated subsystem should (a) supervene on the states of just that subsystem and (b) respect transformations of these states that make no difference to the subsystem's relative frequencies. This principle can be used in tandem with either of the objective and subjective analyses of chance discussed above. On the one hand, a hypothetical frequentist might want to ensure that they pick an infinite reference class that defines a function ch that satisfies **OSP**. On the other hand, a subjectivist might want to pick an expert credence function c_{expert} that is equal to a chance function ch satisfying **OSP**. In either case, **OSP** constrains the range of admissible options. Certainty of a system's isolation will not always solve these theorists' reference class problems, as we saw in our first two passes at constructing a physical theory T for the box factory. Nonetheless, we take the preceding to show that when we find what our theory T has to say about isolation compelling, then if T also imposes strong enough symmetries for **OSP** to pin down a unique measure, that measure is an attractive candidate for a chance function on objectivist and subjectivist approaches alike.

Of course, we have yet to characterize what it means for a physical theory T 's notion of isolation to be compelling, in general! For the box factory, we simply stipulate that the right sort of isolation is achieved. In the next section, we aim to do better.

4 But 'isolation' is ambiguous

In the previous section, we argued that both subjectivist and objectivist approaches to chance can, in special cases, compellingly resolve their respective versions of the reference class problem by appealing to a physical theory. In particular, they can use such a theory to fill in the certainty link (C4) with an appeal to the symmetries of isolated subsystems, yielding **OSP**. The notion of an 'isolated system' in physics, however, is a tricky one to pin down. Our strategy in this section is to do for 'isolated system' what we did for 'chance' in Section 2: namely, to give it a thin, functional analysis based on long-standing desiderata among physicists. Our analysis identifies two key functional links for 'isolation' talk: roughly, that 'isolated subsystems' instantiate a recursive structure, and that we can study them in a laboratory.

To motivate these links, we start with Einstein's (1948) classic concern about how the non-locality of quantum mechanics might challenge the possibility of empirical science:

The following idea characterizes the relative independence of objects far apart in space (A and B): external influence on A has no direct influence on B ; this is known as the 'principle of contiguity', which is used consistently only in the field theory. If this axiom were to be completely abolished, the idea of the existence of (quasi-)enclosed systems, and thereby the postulation of laws which can be checked empirically in the accepted sense, would become impossible. (Translation by Max Born, Einstein et al. 1971, p. 171)

Here, Einstein argues that something like a principle of local action is at least extremely useful

for making sense of the idea that some systems can be effectively closed off from external influence. However, as Howard (1985) notes, Einstein is less strongly committed to local action than to separability: it is ‘the existence of (quasi)-enclosed systems’, however justified, that achieves the possibility of empirical physics. Prospects seem dim for non-relativistic quantum mechanics to justify this separability principle on the back of local action; it is not a field theory, for one. Still, past physicists seem to have had a coherent notion of empirically testing non-local theories like Newtonian gravity and Coulombic electrostatics; it would be strange if one could not construct an analysis of ‘(quasi)-enclosed systems’ that did not depend on locality. Indeed, for Einstein, it only seems crucial that the physical laws—local or otherwise—can be studied in some controllable setting.

Wallace’s (2022a, 2022b) notion of *subsystem recursivity* provides a useful tool for sharpening how isolation relates to testability. Wallace does not aim to analyze ‘isolation’ *per se*; on the contrary, he is forward about treating it as a term of art in his analysis. Nonetheless, his discussion of a theory’s recursive ‘sectors’ is quite helpful for filling out the structural features that a theory needs for it to be ‘checked empirically in the accepted sense’. Very roughly, a *sector* of theory identifies a system whose kinematics and dynamics can be fully specified without reference to any other system. As long as one is justified in claiming that they have captured a sector of theory with a system in their laboratory, one can make inferences about that theory’s applicability writ large (i.e., for larger and larger subsystems, possibly with some fuzzy cutoff, but with the flexibility for this to in principle include the entire universe). The alignment of theory sectors and isolated subsystems, then, would seem to get very close to the thin notion of empirical testability that Einstein is after.

Wallace calls a theory *subsystem recursive* just in case it possesses this alignment. He fleshes out the idea of subsystem recursivity by illustrating how it arises in Coulombic electrostatics:

[Suppose] we now consider a dynamical model in an N -particle sector [of electrostatics], in which there is some sub-region R and some subset of $M < N$ particles which remain (over some time period) well inside R , while the other $N - M$ particles remain outside R ; and suppose that as we approach the [boundary] of R , [the electric potential] tends to some spatially-constant value [...]. This could happen, for instance, because the other $N - M$ particles are very far from R , or because positive and negative charges outside R approximately balance so as to exert little net Coulomb force within R . Then to a good approximation, we can study the dynamics of the subsystem of those M particles autonomously from the other $N - M$: they will behave as if those particles were not present and as if the M particles were interacting among themselves according to the appropriate M -particle sector dynamics. If we want to study the dynamics of that subsystem of particles, in other words, under these assumptions and if we are sufficiently tolerant of small errors, we can replace the full N -particle sector with the M -particle sector. In this sense, the latter can be interpreted as modelling an isolated subsystem of the former.

[...] In fact, Coulombic electrostatics has these two features:

- (i) Any sector of the theory can be interpreted as an idealized description of an

isolated subsystem of a sector of the theory;

- (ii) An isolated subsystem of any sector of the theory can be described, in idealization, as a sector of the theory.

I will call a theory with these two features *subsystem-recursive*: these are theories where any model can be interpreted in the first instance as modelling a dynamically isolated subsystem under certain idealizations about its environment and where, if we want to remove those idealizations, we can embed the model in a model of a larger system within the same theory—and where that larger system in turn is interpretable in the first instance as a subsystem of a still-larger system, with no assumption that we need to understand the whole setup in terms of an *ur*-system that describes the whole universe. (2022a, p. 242)

While Wallace treats ‘isolation’ as a term of art, he gives a formal definition of ‘sector’ by defining structure-preserving restriction maps that take a larger theory sector to a smaller one. For our purposes, we extract from this account three qualitative criteria that a restriction to a subsystem must satisfy for that subsystem to count as a sector: namely, the restriction must preserve the theory’s rules for *kinematics*, *dynamics*, and *symmetries*. To flesh out the kinematic part for electrostatics, we might describe the trajectories of N particles at a given time slice with vectors in the N -wise Cartesian product of \mathbb{R}^3 ; in this case, we are always free to project down onto M of these trajectories to describe those particles and those alone. We ensure that these trajectories obey the Coulomb force law by stipulating that the potential does not vary across the boundary of the spatial region containing these M particles, as Wallace describes above. Time reversal, spatial translations, and rotations applied to the state space of just the M particles preserve solutions to the Coulomb dynamics just as well as they do when applied to the larger system. All such solutions will adequately describe the M particles on their own, yielding a precise sense in which they behave as though the others ‘were not present’. Thus, the M -particle subsystem satisfies the three qualitative criteria and is therefore a sector of electrostatics. Moreover, in this case, we have good grounds to identify such sectors with *isolated* subsystems. The constant electric potential of the boundary (idealized to be at infinity) means, in particular, that the electric field vanishes there, and we can measure when and where there is a non-zero electric field. So, Coulombic electrostatics thereby satisfies features (i) and (ii), coming out as a subsystem recursive theory.¹⁶

In contrast with Coulombic electrostatics, Wallace cites Newtonian gravity as an example of a theory where the alignment between sectors and isolated subsystems breaks down.¹⁷ In this theory, the dynamics of point-source masses indexed by J is given by

$$\frac{d^2}{dt^2}x_J^i = \nabla_i\Phi(x_J^i, t), \tag{8}$$

¹⁶By contrast—and as Wallace (2022a, §6) points out explicitly—Barbour–Bertotti theory is not a subsystem recursive theory; this reveals that (at least in some cases) there is a tension between the desiderata of (i) subsystem recursivity versus (ii) having a suitably relational/‘Machian’ theory.

¹⁷Here, we recapitulate the relevant points from Wallace (2022a, §6); readers seeking more details should consult that section of his article.

where Φ is a scalar potential field that satisfies the Poisson equation,

$$\nabla^2\Phi(x, t) = -4\pi G \sum_J \delta(x - x_J)m_J. \quad (9)$$

We restrict our attention to Φ that are linear at the system boundary. Idealizing the boundary to be at $|x| = \infty$, we can write this condition as

$$\lim_{|x| \rightarrow \infty} \Phi(x, t) = \Phi_0(t) + a_i(t)x^i(t), \quad (10)$$

where the acceleration term $a_i(t)$ refers to the components of a time-dependent vector. In the case where $a_i(t) = 0$ for all t , the boundary condition is constant, and this theory recovers the usual inverse-squared force law of Newtonian gravity. We can then define sectors of the theory in much the same way that we did for Coulombic electrostatics. Just as the electric field disappears at the boundary of a Coulombic sector, the gravitational field vanishes at the boundary of a Newtonian sector. However, we now have a good reason to think that the resulting sectors are too strict to characterize a desirable notion of isolation: no such sector can describe a laboratory on Earth’s surface. We need a non-zero $a_i(t)$ at the laboratory boundary to account for the Earth’s gravitational field.

To fix this issue, we can simply change the definition of ‘Newtonian sector’ to contain dynamics with the full range of boundary conditions in (10). However, Wallace notes that one might also change the definition of ‘isolation’ to identify isolated subsystems with Newtonian sectors. So long as we know the specific point mass distribution generating a potential with non-zero $a_i(t)$ at a given subsystem’s boundary, we can always rewrite that subsystem’s dynamics within an appropriate sector. Considering only laboratories on the surface of the Earth, this issue does not appear to pose an immediate threat. However, Einstein’s elevator thought experiment challenges its general tenability. When one feels a greater push towards the elevator’s floor, there is no means of discriminating from within the elevator whether the cause of this effect is due to the Earth acquiring more mass or the elevator’s upward acceleration. This point challenges Newtonian sectors’ ‘upwards’ recursion, towards larger and larger systems. Progressing up this ladder, whenever we reach a cosmological system with a non-vanishing gravitational field at the boundary, Newton’s theory does not license us to stop at that rung—even though the particular completion of that system within a Newtonian sector might end up being arbitrary.

Contrast this situation with what happens when we revise electrostatics to include potentials with linear boundary conditions of the form (10). This revised theory of electrostatics has sectors with empirical differences: namely, some sectors will have vanishing electric fields at their boundary, and some will not. However, this move does not challenge the isolation of these sectors from other subsystems. The dynamics can still be studied in isolation. In other words, generalizing the form of the dynamics allows us to specify a sense in which the empirical differences between the sectors do not matter for the physics.

From these case studies, we can draw a general moral. When sectors impose conditions on the dynamics that we can check empirically but which are hard to control from within the system, we have cause to generalize those dynamics—which, in turn, disrupts that theory’s notion of isolation. The converse of this moral is that tenable notions of isolation can allow

for empirical differences between isolated subsystems by specifying a precise sense in which those empirical differences do not affect the dynamics.

We capture this moral as the interplay of two functional links the term ‘isolated subsystem’ ought to satisfy: a link to a recursive structure on the one hand and a link to empirical control on the other.

- (I1) **The recursion link.** Isolated subsystems instantiate a recursive structure (like that of a theory sector).
- (I2) **The control link.** We can study an isolated subsystem in a laboratory.

The recursion link (I1) sharply defines various senses in which a given subsystem might be isolated, and the control link (I2) pushes one to pick a sense of isolation that is useful for empirical practice.

One might want to be agnostic about the strength of the control link in the same way that one might be agnostic about the strength of the inferential link about chance. Subjectivists about chance often stipulate that any sequence of outcomes is possible for an event with any chance value. Thus, the actual outcomes of any finite sequence of chance events might end up being wildly misleading about the true value of the chance. Usually, these subjectivists invoke a version of the law of large numbers to argue that one’s credence in encountering such a case should be very low, thus recovering the inferential link (C1) approximately.¹⁸ Similarly, one might demand more or less rigor for a given laboratory setup to satisfy the control link (I2). For example, one might not demand that an isolated Coulombic system’s electric field is always exactly zero at the boundary, instead tolerating some small, random fluctuations.

In addition to filling out the recursion link (I1), Wallace’s subsystem structures provide a helpful way of characterizing the chance-preserving maps in **OSP**. Roughly, Wallace says that a property fixed by the state of a subsystem is *intrinsic* with respect to a given symmetry if it has the same value for any two states related by that symmetry (2022a, p. 245). For example, the relative distances of particles in a subsystem of Coulombic electrostatics are intrinsic for spatial translations of states of that subsystem, as well as for rotations of those states.¹⁹ Recall that the physical theory T in **OSP** maps each state S to a *function*, $ch(X|S)$, rather than to a scalar value like a relative distance. There is some flexibility in how one might extend Wallace’s intrinsicity condition to this case, but one natural way to do so is to view ch as a function of states that are pairs (X, S) , i.e., pairs of outcomes X for our original states S and those states themselves. In that case, the condition $ch(X | S) = ch(X' | S')$ in **OSP** asserts that chance is an intrinsic property of the states (X, S) in precisely Wallace’s sense. This intrinsicity condition amounts to saying that symmetric states S and S' yield the same chance function *up to a relabeling of the outcomes*. For example, in the box factory theory, if we obtain S' from moving S forward by .2 cm, then S gives the same chances as S' if we view it as using the label ‘4.1 cm’ to refer to a rod 4.3 cm long, and likewise for any other length specification. Henceforth, we will stick to using ‘states’ to talk about S rather

¹⁸For a precise formulation of this argument, see Greaves and Myrvold (2010).

¹⁹Note that this approach differs from the standard contrast in metaphysics between ‘intrinsic’ and ‘relational’ properties, as detailed, e.g., by Marshall and Weatherson (2023). On Wallace’s construal, relational properties (in the standard metaphysics sense) can be intrinsic (in his sense).

than pairs (X, S) , while continuing to refer to the preservation of chance up to a relabeling of outcomes as **OSP**'s intrinsicity condition.

Our strategy for deriving Everettian chance, then, runs as follows: first, we pin down appropriate theory sectors for Everettian quantum mechanics; next, we apply **OSP**'s intrinsicity condition to the appropriate symmetries of those sectors; and last, we look for symmetric states that end up being identical. We turn now to the first task, for which we use Newtonian and Coulombic case studies as templates.

5 Probing Everettian ‘isolation’ with theory sectors

In the previous section, we identified two core functional roles played by talk of ‘isolated subsystems’: (I1) isolated subsystems should instantiate a recursive structure, like that of Wallace’s (2022a) theory sectors, and (I2) agents should be able to study them in laboratories. ‘Isolated subsystems’ fulfill these roles in theories with nonlocal force laws, like Coulombic electrostatics and Newtonian gravity. That bodes well for the prospects of applying this functional analysis to ‘isolated subsystems’ in non-relativistic quantum mechanics.

However, there is a fly in the ointment: a theory like Newtonian gravity employs a single dynamical rule, and orthodox non-relativistic quantum mechanics infamously has two. According to orthodoxy, unitary dynamics describes self-evolution, projective dynamics describes measurement, and never the twain shall meet. In contrast, a traditional selling point of Everettian approaches is that they treat all dynamics as unitary. Thus, at first blush, one might justifiably think that orthodoxy will frustrate attempts to define a theory sector, while Everettian approaches will accommodate sectors similar to those in Newton’s theory.

In fact, we find that both of these impressions are false. Orthodoxy naturally accommodates two sorts of theory sector, one for each of its dynamical rules—and prospects for Everettians getting away with just one sort of sector seem dim. On the orthodox approach, one can helpfully use von Neumann’s measurement scheme to relegate projective dynamics to interactions with macroscopic objects. Similarly, Everettians rely on projections to specify the sense in which their macrostates—branches—are dynamically isolated. Thus, whether or not one is an Everettian, we recommend adopting two separate sectors and, accordingly, two separate notions of ‘isolation’ for non-relativistic quantum mechanics: one for microstates and one for macrostates.

The notion of isolation for Everettian macrostates amounts to what Franklin (2024) calls a ‘screening-off criterion’: namely, the emergence of a particular macrostate from a microstate amounts to the approximate dynamical isolation of one branch from any other branch. Macrostates include indicators of the outcomes of measuring quantum systems, such as a sharp position state of a measuring device’s pointer—a ‘pointer state’, in Zurek’s (1981, 1982) popular terminology. In short, pointer states correspond to measurement outcomes, the chance events that we care about. Thus, we use branching dynamics to fill in the recursion link (I1) and specify the notion of isolation relevant to Everettian chance.

The control link (I2) guides the precise subsystem structure that branching dynamics should instantiate. In particular, these dynamics pose two challenges to our empirical control of macroscopic subsystems. On the one hand, we have a repeatability problem: system eigenstates may fail to strongly couple to pointer states, resulting in non-repeatable

measurements. On the other hand, we have a preferred basis problem: system-environment interactions may fail to select well-localized pointer states, resulting in macroscopic states of affairs that are highly nonclassical. The repeatability problem is analogous to the case of the Newtonian who learns of Einstein’s elevator scenarios; it gives us reason to generalize our macrostate sectors.

What about the preferred basis problem? Several authors use it to charge Everettians who seek to derive chance values with circularity or incoherence (Baker 2007, Dawid and Thébaud 2015, Mandolesi 2019). For example, responding specifically to Wallace’s (2012) decision-theoretic derivation, Mandolesi writes the following:

Without the Born rule, Wallace’s solution of the preferred basis problem does not work as expected. Branches might be nothing like our world, lacking complex structures or behaving erratically. This compromises the whole decision-theoretic approach, which depends on narratives where agents exist, their actions have the expected consequences, and rationality is possible. (2019, p. 49)

In response to such challenges, Franklin (2024, p. 299) aims to confirm the emergence of classical worlds in Everettian quantum mechanics on the basis of phenomena that one can derive from decoherence without reference to probabilities. Our strategy is slightly different. We argue that the Everettians can use a partial measure of chance—one identifying all and only branches with a Hilbert space norm near 0 to have a chance near 0—to establish a robust sense in which their worlds are isolated, even if those worlds end up being non-classical. With this notion of isolation in hand, Everettians can then non-circularly derive that the full Born rule measure gives the chances of those worlds.

Note that the Everettian who allows for the theoretical possibility of deviant branches is strongly analogous to the electrostatics theorist who adopts dynamics with linear boundary conditions. The latter cannot use their weakened notion of ‘isolation’ to specify whether the electric field vanishes at an isolated subsystem’s boundary, even though they could easily check it empirically. Likewise, the Everettian cannot use our preferred notion of ‘isolation’ to guarantee that a branch is quasi-classical, even though they could easily observe it to be. Quite plausibly, these two setups latch onto different modal intuitions. Constant electric fields at system boundaries should be common, and non-classical macrostates should be either impossible or rare enough to ignore. However, this difference does not change the fact that both theorists fulfill our two thin functional criteria for isolated subsystems.

Building up to this point requires some technical exposition, which we present in the body of this section. We start in §5.1 by illustrating how microstate and macrostate sectors arise in a conservative approach to quantum mechanics; namely, the Schrödinger picture augmented with von Neumann’s account of measurement (§5.1.1–§5.1.2). This subsection concludes with a brief discussion of how symmetries that derive chance values already appear in the orthodox approach (§5.1.3). Then, §5.2 unpacks how the repeatability and preferred basis problems challenge the orthodox notions of isolation. §5.3 describes how the repeatability problem motivates quantum measurement theory’s generalization of orthodox sectors. §5.4 fleshes out how Everettians can make do with sectors nearly as strict as orthodoxy’s, despite the challenges posed by the preferred basis problem. Last, §5.5 details how these sectors address concerns about circularity and the role of locality in previous Everettian chance derivations.

5.1 Orthodoxy's sectors for microstates and macrostates

We start with a conservative approach—namely, the Schrödinger picture as presented in many textbooks on non-relativistic quantum mechanics, augmented with von Neumann's account of measurement. Infamously, this picture posits two different dynamical rules: one for a system's self-evolution, and one for measurements of that system. This immediately poses a challenge for applying Wallace's notion of subsystem structure, which tacitly assumes that a theory has just one fundamental dynamical rule. Our strategy is to propose a distinct notion of sector for each rule, associating the former with microstates and the latter with macrostates.

The kinematics of both sector types are the same, so we start there. On the Schrödinger picture, one identifies the state of a system at a given time with a vector in a complex Hilbert space. That state assigns values to observable properties via its projections \hat{P}_i onto certain orthogonal subspaces of the Hilbert space—subspaces which can be equivalently described as the spectra of Hermitian operators, in line with standard practice. However, there is some arbitrariness regarding which properties we assign to which subspaces. Intuitively, if we act on the Hilbert space's vectors in a way that changes none of their relative lengths or angles, then a difference in how they project onto subspaces should amount to a difference in our choice of labels for those subspaces. Handily, unitary maps encode such transformations, making them prime candidates for chance-preserving dynamical symmetries. In what follows, we take a *transformation* by a unitary \hat{U} to map vectors Ψ to $\Psi' = \hat{U}\Psi$ and operators \hat{A} to $\hat{A}' = \hat{U}\hat{A}\hat{U}^{-1}$.

5.1.1 Orthodoxy's microstate sectors

The Schrödinger equation specifies the dynamics of microstates. Anticipating our desire to define a subsystem structure, let us define these dynamics on a system S , its environment E , and a system SE containing both. We represent the state of SE at time t with a vector $\Psi(t) \in \mathcal{H}_{SE}$, where $\mathcal{H}_{SE} = \mathcal{H}_S \otimes \mathcal{H}_E$, the tensor product of Hilbert spaces describing the degrees of freedom of the system and the environment, respectively. $\Psi(t)$ must satisfy

$$i\hbar\partial_t\Psi(t) = \hat{H}\Psi(t), \quad (11)$$

for some self-adjoint operator \hat{H} , the Hamiltonian. The Hamiltonian defines a family of unitary operators $\hat{U}_{\hat{H}}(t, t_0) = e^{-i\hat{H}(t-t_0)/\hbar}$, and

$$\Psi(t) = \hat{U}_{\hat{H}}(t, t_0)\Psi(t_0) \quad (12)$$

solves (11). Note, too, that a unitary transformation \hat{U} of the underlying Hilbert space preserves the dynamical rule (11), in the following sense. Letting $\Psi'(t_0) = \hat{U}\Psi(t_0)$ and $\hat{H}' = \hat{U}\hat{H}\hat{U}^{-1}$,

$$\Psi'(t) = \hat{U}_{\hat{H}'}(t, t_0)\Psi'(t_0) \quad (13)$$

is a solution of

$$i\hbar\partial_t\Psi'(t) = \hat{H}'\Psi'(t) \quad (14)$$

corresponding to the old dynamics in our newly-labeled space. Hence, unitary operators encode dynamical symmetries of the total system SE .^{20,21}

Microstate dynamics and their symmetries behave well under a natural notion of restriction. Whenever we can express the kinematics of a total system as a simple tensor product, we can select any component in that product to recover a well-defined state that obeys the same dynamical rule with the same symmetries. To illustrate, let us say that S and E are kinematically separable at time t when the state of SE is separable in the usual sense; that is,

$$\Psi(t) = \Psi_S(t) \otimes \Psi_E(t). \quad (15)$$

Similarly, let us say the dynamics are separable when the Hamiltonian is separable in the usual sense, i.e.,

$$\hat{H} = \hat{H}_S \otimes \hat{H}_E. \quad (16)$$

The separability of the Hamiltonian implies that the unitaries it generates are separable, too. Thus, we can summarize the kinematic and dynamic separability of S and E with the following equation:

$$\hat{U}_{\hat{H}}(t, t_0)\Psi(t_0) = \hat{U}_{\hat{H}_S}(t, t_0)\Psi_S(t_0) \otimes \hat{U}_{\hat{H}_E}(t, t_0)\Psi_E(t_0). \quad (17)$$

The first component of this tensor product is a solution of equation (11) defined on \mathcal{H}_S , the Hilbert space of the subsystem S alone. Let us introduce the notation $(\cdot)|_S$ to denote restriction to the S component of a simple tensor product state—so, e.g., $\Psi|_S(t_0) = \Psi_S(t_0)$ and $\hat{H}|_S = \hat{H}_S$. By the same reasoning as before, unitary maps \hat{U}_S on \mathcal{H}_S identify symmetries of the restricted dynamics. We say that the restriction to S thus yields a new microstate sector of the theory. Likewise, with light abuse of notation, we can specify extensions $(\cdot)^{SE}$ of a microstate sector as a family of maps, each undoing the action of a given restriction from SE to S or E .

²⁰This is not to say that unitaries encode *all* dynamical symmetries! The time symmetry of Schrödinger dynamics, for example, is best represented by an anti-unitary operator (Sakurai and Napolitano, 2010, p. 291). For the Born-rule derivations we consider, it suffices that every unitary map encodes a symmetry, but not vice versa.

²¹An anonymous referee has expressed concern that this line (that unitary operators encode dynamical symmetries) trivializes the notion of a dynamical symmetry. We do not agree, although here we are confronted with deep waters which it would be wise not to enter without due caution. In brief: what we have in the above is the analog of what in the philosophy of spacetime are known as ‘local spacetime theories’, following the nomenclature of Earman and Norton (1987). Any such local spacetime theory has equations of motion invariant under arbitrary diffeomorphisms, which would seem to trivialize the notion of diffeomorphism invariance and simultaneously to diminish the extent to which general relativity is (supposedly) ‘special’ by comparison with prior theories of space and time (since those latter theories, e.g. Newtonian gravity, can also be formulated as local spacetime theories). We disagree (in both cases) that the relevant notion of a symmetry is thereby trivialized, for following the lead of Pooley (2017) it is still possible to make sense of theories with restricted symmetry groups (i.e., theories which respectively do not have unitary transformations or arbitrary spacetime diffeomorphisms as symmetries). But here is not the place to go deeper into this issue; see Read (2023, ch. 3) for further discussion.

5.1.2 Orthodoxy’s macrostate sectors

To pin down an appropriate notion of a macrostate sector, we consider orthodoxy’s approach to measurement. Although many textbook expositions of the Schrödinger picture skip it, von Neumann’s account of measurement fruitfully allows us to relegate projection dynamics to interactions between microscopic systems and macroscopic measuring devices. To illustrate the idea driving this account, consider a spin- $\frac{1}{2}$ system S shot through a (macroscopic) Stern–Gerlach measurement apparatus A measuring spin in the z -direction. We idealize our model of A such that \mathcal{H}_A is two-dimensional, just like \mathcal{H}_S . We suppose that SA is initially separable in the tensor-product sense described above. Then, we describe the measurement with the action of a unitary map $\hat{U}(t_1, t_0)$ that satisfies

$$(a|\uparrow\rangle + b|\downarrow\rangle)|A\rangle \xrightarrow{\hat{U}(t_1, t_0)} a|\uparrow\rangle|A_\uparrow\rangle + b|\downarrow\rangle|A_\downarrow\rangle, \quad (18)$$

where $|A_\uparrow\rangle$ and $|A_\downarrow\rangle$ are orthogonal (macroscopic) pointer states representing localized spots near the top and bottom of our observation screen, respectively. This particular Stern–Gerlach measurement is destructive, but we can imagine a non-destructive version of the experiment which prepares states $a|\uparrow\rangle|A_\uparrow\rangle$ by removing the top half of our screen. In this case, (18) describes a *pre-measurement*, i.e., a unitary interaction of the system and the apparatus before a projection. Note well that (18) falls entirely within the purview of a microstate sector; its dynamics can, in principle, be recovered by some Hamiltonian and equation (11). However, at time t_1 , measurement occurs, modeled by projecting onto $a|\uparrow\rangle|A_\uparrow\rangle$.

To describe realistic measurements, von Neumann de-idealizes the above schema in two ways that are important for our purposes. First, since many observables of interest, like position, lack discrete spectra, he generalizes (18) to apply to S and A defined on Hilbert spaces of arbitrary (countable) dimension. The dynamics of this generalization still consist of unitary evolution of a vector state of SA , followed by projecting it onto a particular vector state in a given orthonormal basis. Second, he notes that the significant difference in scale between the system S and the apparatus A implies that, for example, a range of different Planck-scale-differentiable positions of the apparatus’s pointer correspond to one macroscopically observable position. To this end, von Neumann introduces ‘macroscopic projections’ onto subspaces containing a range of apparatus vector states (1955, p. 407). All macroscopic projections possible at a given time commute with each other, such that they correspond to one or another event in an exhaustive set of mutually exclusive outcomes. In this spirit, when talking about macrostate sectors, we will use the term ‘microstate’ to refer to the vector state of SA and ‘macrostate’ to refer to a subspace X in which it lies, such that \hat{X} denotes a macroscopic projection onto this subspace.

Let us introduce a bit of notation to more easily describe the subsystem structure of macrostate sectors. We index projections by outcomes in subscripts and time in superscripts, so $\hat{P}_i^{t_m}$ denotes projection onto a microstate of SA lying in the macrostate $X_i^{t_m}$ at time t_m . In a similar manner, let us rewrite a microstate $\Psi(t_0)$ as Ψ^{t_0} , and its unitary microstate dynamics $\hat{U}(t_1, t_0)$ as $\hat{U}^{t_1 \leftarrow t_0}$. The (unnormalized) state after measurement is then given by

$$\Psi^{t_1} = \hat{P}_i^{t_1} \hat{U}^{t_1 \leftarrow t_0} \Psi^{t_0}. \quad (19)$$

These measurement dynamics straightforwardly preserve the unitary symmetries of microstate

sectors. For an arbitrary unitary transformation by \hat{U} , we have

$$\begin{aligned}
(\Psi^{t_1})' &:= \hat{U}\Psi^{t_1} = \hat{U}\hat{P}_i^{t_1}\hat{U}^{t_1\leftarrow t_0}\Psi^{t_0} \\
&= \hat{U}\hat{P}_i^{t_1}\left(\hat{U}^{-1}\hat{U}\right)\hat{U}^{t_1\leftarrow t_0}\left(\hat{U}^{-1}\hat{U}\right)\Psi^{t_0} \\
&= \left(\hat{P}_i^{t_1}\right)'\left(\hat{U}^{t_1\leftarrow t_0}\right)'(\Psi^{t_0})',
\end{aligned} \tag{20}$$

as expected. We can also iterate non-destructive measurements in a chain. For example, say we measure a new pointer observable at time t_2 , perhaps with a different apparatus B , and observe the outcome $X_j^{t_2}$. Recycling our notation for extensions and restrictions of microstates, let $\Psi^{t_1}|^{SAB} = \Psi^{t_1} \otimes |B\rangle$, where $|B\rangle$ is the ready state of B at t_1 . The evolution of the system is then given by

$$\Psi^{t_2} = \hat{P}_j^{t_2}\hat{U}^{t_2\leftarrow t_1}\left(\hat{P}_i^{t_1}\hat{U}^{t_1\leftarrow t_0}\Psi^{t_0}\right)\Big|^{SAB}. \tag{21}$$

Moreover, a specification of the times of all measurement events of interest, $\{t_m\}$, yields a set $\{X_i^{t_m}\}$ of time-indexed macrostates corresponding to a particular record of observable outcomes; i.e., precisely what Gell-Mann and Hartle (1993) call a *history* in the context of decoherence theory.

As such, we schematize a macrostate sector as a microstate Ψ following von Neumann measurement dynamics with projections at times in a sequence $(t_1, \dots, t_m, \dots, t_n)$, yielding a history $\{X_i^{t_m}\}$ of macrostates. We equip these sectors with a notion of extension via concatenation and restriction to substrings—that is, a contiguous sequence of characters from the original sequence without rearranging or changing them. (Note that substrings are less general than subsequences, which can omit intermediate elements of the substring.) So, for instance, for a sequence (t_2, t_3, t_4) , concatenations include (t_1, t_2, t_3, t_4) and (t_2, t_3, t_4, t_5) , and restrictions include (t_2, t_3) . Let us suppose that extension via concatenation is precisely the process of extending the initial sequence (19) to the final sequence (21), albeit generalized to initial and final sequences of arbitrary finite length. Likewise, let us take a restriction to a substring to be the dual operation undoing one such extension. Note that the projection postulate requires us to consider only substrings and not more general subsequences, as the ‘unused’ superposition terms in equation (18) need to be dismissed to guarantee that the future dynamics work as expected. Crucially, restrictions allow us to dismiss systems and apparatuses that are unused in the sector under consideration. For example, suppose the non-destructive Stern–Gerlach measurement of (18) indeed yields the state $\Psi^{t_1} = a|\uparrow\rangle|A_\uparrow\rangle$ on SA . Then, letting $\Psi^{t_1}|_S = a|\uparrow\rangle$ and $(\Psi^{t_1}|_S)|^{SB} = a|\uparrow\rangle|B\rangle$, we have

$$\Psi^{t_2}|_{SB} = \left(\hat{P}_j^{t_2}\hat{U}^{t_2\leftarrow t_1}\right)\Big|_{SB}(\Psi^{t_1}|_S)|^{SB}, \tag{22}$$

defining a macrostate sector for just apparatus B ’s measurement of S from the preparation at t_1 to the outcome at t_2 . Thus, there is a sense in which macrostate sectors robustly isolate systems in time: we are licensed to fully ignore past measurements of an isolated system when we wish to measure it anew. Indeed, it is straightforward to check that the extension (21) and the restriction (22) each retain unitary maps as dynamical symmetries of the sort demonstrated by equation (20).

5.1.3 Symmetries in orthodox macrostates

We have already noted that macrostate sectors preserve the unitary symmetries of microstates. Still, another sort of symmetry is important for our purposes: namely, different choices of coarse-graining a given set of microstates into macrostates, which we call *repartitionings*.

To describe repartitionings, we consider the de-idealization of the spin measurement (18) where we take the Hilbert space of the apparatus to be of infinite dimension. With minor abuse of notation, let S denote the microstate of SA at time t_1 , after pre-measurement; measurement will leave the system in either $S_1 = a|\uparrow\rangle|A_\uparrow\rangle$ or $S_2 = b|\downarrow\rangle|A_\downarrow\rangle$. Our specific choice of macrostates X_1 and X_2 containing S_1 and S_2 , respectively, is in some sense arbitrary. We can motivate this choice with facts about our scale relative to S , but a range of choices will work equally well; indeed, this arbitrariness reflects our insensitivity to differences in the finest-grained, microscopic degrees of freedom. So, consider a different choice of macrostates $X_{1'}$ and $X_{2'}$ containing S_1 and S_2 , respectively, such that $X_{1'} \neq X_1$ and $X_{2'} \neq X_2$. Since these new macrostates amount to a different way of labeling the outcomes of measuring S , they should have the same relative frequencies as the originals. In general, for a partition of macrostates $\{X_i\}$ coarse-graining the set of microstates $\{S_i\}$ that can result from measuring S , we say that $\{X_{i'}\}$ is a *repartitioning* of $\{X_i\}$ if $S_i \in X_i, X_{i'}$ for each i . Repartitionings should preserve chance values: that is, by **OSP**,

$$ch(X_i | S) = ch(X_{i'} | S). \quad (23)$$

Moreover, they clearly do not change the dynamics (19).

To set the stage for our use of symmetries to derive chance values in the Everett approach, it is also instructive to look at a concrete example of a unitary symmetry in a macrostate sector. Let $\hat{S} \in U(2)$ be a unitary map that flips the labels of ‘up’ and ‘down’ spin- $\frac{1}{2}$ states,

$$a|\uparrow\rangle + b|\downarrow\rangle \xrightarrow{\hat{S}} a|\downarrow\rangle + b|\uparrow\rangle. \quad (24)$$

Now, apply the unitary transformation $\hat{S} \otimes \hat{I}_A$ to the dynamics described by equation (18):

$$(a|\downarrow\rangle + b|\uparrow\rangle)|A'\rangle \xrightarrow{\hat{U}'(t_1, t_0)} a|\downarrow\rangle|A_\uparrow\rangle + b|\uparrow\rangle|A_\downarrow\rangle. \quad (25)$$

Using our notation from the repartitioning described above, equation (25) leaves SA in the microstate S' after pre-measurement, and projection will yield either S'_1 in the macrostate X'_1 or S'_2 in X'_2 . Note that the underlying symmetry of the microstates carries over to the macrostates. In other words, the freedom of choice in how to label the ‘up’ and ‘down’ microstates manifests in the arbitrariness of choosing between X_1, X_2 and X'_1, X'_2 to label the macrostates. Thus, we expect this symmetry to preserve chance values, and **OSP** implies

$$ch(X_i | S) = ch(X'_i | S') \quad (26)$$

for $i \in \{1, 2\}$.

Note well that, as of yet, we have not introduced any chance or probability rule into the orthodox approach. The Born rule figures nowhere in our schemas for microstate and macrostate sectors. Moreover, the symmetries outlined above suffice to derive that rule rather than stipulate it (as we will show in §6).

In sum, the orthodox approach to non-relativistic quantum mechanics has two separate but compatible subsystem structures, one for microstates and the other for macrostates. Macrostate sectors are specified by a sequence of microstate sectors, each of which might contain further microstate sectors as subsystems. In turn, orthodoxy yields candidates for two notions of isolation. We say that a subsystem is an isolated microstate if it is a microstate sector, and likewise for isolated macrostates.

5.2 Challenges to orthodox isolation

How defensible are orthodoxy’s notions of isolation? We consider two key challenges stemming from the same concern: that system observables might well fail to couple strongly with the pointer states of our macroscopic apparatuses. To illustrate, recall our simple example of a von Neumann measurement of a spin- $\frac{1}{2}$ system, (18). Let us now suppose that the unitary dynamics couples the definite spin states to apparatus states A_+ and A_- that do not align with the pointer states A_\uparrow and A_\downarrow :

$$(a|\uparrow\rangle + b|\downarrow\rangle)|A\rangle \xrightarrow{\hat{U}(t_1, t_0)} a|\uparrow\rangle|A_+\rangle + b|\downarrow\rangle|A_-\rangle, \quad (27)$$

where $\langle A_\pm|A_\uparrow|A_\pm\rangle$ is far from both 0 and 1. In this case, A_+ is a Schrödinger’s cat-like superposition of the upper and lower portions of our observation screen; thus, we would not expect de-idealizing the pointer states to yield coarse-grained projections that include A_+ and A_- in their ranges. Following Busch et al. (1995, §I.1.2), one nice way to motivate this generalization is to define the center-of-mass wavepacket coupled with each spin eigenstate, and then note that these wavepackets might fail to be strongly separated and localized.

The first issue is that a projection onto a localized apparatus region will fail to specify a state for the system. In particular, the non-destructive Stern–Gerlach measurement described in the previous section would now pick out the final state

$$\Psi(t_1) = (a\langle A_\uparrow|A_+\rangle|\uparrow\rangle + b\langle A_\uparrow|A_-\rangle|\downarrow\rangle)|A_\uparrow\rangle, \quad (28)$$

which, in general, is no longer a simple tensor product of states in S and A . Quantum measurement theory addresses this problem by targeting the fact that orthodox macrostates require measurements to be repeatable. For example, if one records the pointer state $|A_\uparrow\rangle$ in (18), the system S is left in the state $|\uparrow\rangle$; thus, another spin- z measurement performed immediately after the first should not change the state of the system in any way. Since measurement theory weakens this assumption to define more general system states, we call this first issue *the problem of non-repeatability*. This problem identifies a sense in which orthodox sectors are too strict: they do not allow for a notion of ‘isolation’ that accommodates systems we can reasonably expect to study in a lab.

The second issue is that nothing seems to stop us from projecting onto the apparatus state A_+ or A_- to finish the second step of von Neumann’s scheme. In short, von Neumann’s scheme does not explain why the projection postulate applies to some bases rather than others. We refer to this issue as *the preferred basis problem*.²² This problem shows a sense in

²²There are subtle but important differences between our use of the term ‘preferred basis problem’ and several uses of this term in the literature. To start, one often finds the term used to describe a problem in

which orthodox sectors might be too liberal: they allow for ‘isolated systems’ that we would never expect to observe.

Since quantum measurement theory’s answer to the problem of non-repeatability provides an exemplary case of sector-revision, we will illustrate it briefly in the next subsection. The preferred basis problem is trickier, but will serve as the backbone for our approach to revising orthodoxy’s sectors for Everettians.

5.3 Revising the sectors to accommodate non-repeatability

Our key notion of ‘isolation’ for branches comes from how Everettians revise orthodox sectors to resolve the preferred basis problem. Before turning to that revision, we briefly illustrate how quantum measurement theory revises orthodox sectors to address the repeatability problem, which provides a useful point of comparison.

First, recall that when an apparatus’s pointer states fail to strongly couple to a system’s eigenbasis, as in (27), the final state yielded by von Neumann’s scheme will, in general, fail to be a simple tensor product of S and A . Previously, we relied on this separability to justify the isolation of a macrostate over time. In other words, separability gave us the license to ignore the previous measurement interactions of a system when considering new ones. If a theory were to fail to recover this characteristic, then that would certainly threaten its empirical tenability.

Fortunately, a minor modification will suffice. Supposing that the *dynamics* of S and A remain separable in the conservative sense, we can describe S as evolving on its own by using the partial trace operation to suppress the apparatus’ degrees of freedom:

$$\hat{\rho}_S(t_1) = \text{Tr}_A (|\Psi(t_1)\rangle\langle\Psi(t_1)|). \quad (29)$$

The relevant generalization of the conservative approach’s Schrödinger dynamics is given by the standard Liouville–von Neumann equation,

$$i\hbar\partial_t\hat{\rho}_S = \left[\hat{H}_S, \hat{\rho}_S \right], \quad (30)$$

Everettian quantum mechanics *per se*, not an issue in von Neumann’s measurement scheme. For example, Wallace writes that the term ‘comes from asking the following question: if superpositions are supposed to represent distinct worlds, with respect to which basis are these superpositions defined?’ (2012, p. 39). The critiques of Baker (2007), Dawid and Thébault (2015), and Mandolesi (2019) mainly target Wallace, and they chiefly have his construal of the preferred basis problem in mind. However, Wallace ultimately resists using the term, suggesting that it emphasizes the question, ‘which basis?’, over the more important question, ‘why *any* basis?’ (2012, p. 40). Schlosshauer’s (2007) approach helpfully supplements this suggestion. Schlosshauer uses the term ‘the problem of the preferred basis’ to refer to the ‘why *any* basis?’ question, and he calls the ‘which basis?’ question ‘the problem of the non-observability of interference’, in reference to the fact that interference patterns indicative of non-localized superpositions disappear at larger length scales.

Unlike Schlosshauer, our use of the term ‘preferred basis problem’ refers to the ‘which basis?’ question, which we diagnose as applying equally pressingly to both von Neumann’s scheme and Everettian quantum mechanics. We agree with Wallace that the ‘why *any* basis?’ question is particularly important for Everettian quantum mechanics, where there is no dynamical projection postulate that literally collapses the state. Notably, Schlosshauer (2007, §2.5.2) argues for a sense in which the ‘why *any* basis?’ question identifies an explanatory deficit in von Neumann’s scheme, which we do not address in the main text; we explain why in fn. 24. We thank an anonymous referee for raising issues related these points.

which recovers (11) in the special case that $\hat{\rho}_S$ is pure. Thus, so long as the total system follows the separable Hamiltonian $\hat{H} = \hat{H}_S \otimes \hat{H}_A$, we can take S to identify a microstate sector of the theory. The system S , after all, now behaves precisely as if the system A were not present, insofar as it obeys the dynamical rule (30). However, it is no longer the case that specifying states on S and A suffices to uniquely specify a state on SA that restricts to them. We have to keep in mind that to piece together the kinematics of the total state, we need more information than just the kinematics (i.e., the reduced density matrices) of the subsystems on their own.²³ One can revise macrostate sectors to accommodate these generalized microstates by rewriting the equations (19) through (22) with the density operator formalism for states. So, for example, the macrostate dynamics of the generalized measurement in (27) is as follows:

$$(a|\uparrow\rangle + b|\downarrow\rangle) |A\rangle\langle A| (a^*\langle\uparrow| + b^*\langle\downarrow|) \xrightarrow{\hat{U}(t_1, t_0)} |a|^2|\uparrow\rangle\langle\uparrow| |A_+\rangle\langle A_+| + |b|^2|\downarrow\rangle\langle\downarrow| |A_-\rangle\langle A_-|. \quad (31)$$

For ease of exposition, we avoid unpacking additional examples here.

Moreover, in the next subsection, we will return to mainly using vector states with Schrödinger dynamics for our schema of Everettian sectors. The notation of this approach is arguably more compact and familiar, and we do not need to exploit measurement theory’s resolution of the repeatability problem to address Everettians’ approach to the preferred basis problem. However, nothing stands in the way of combining the two approaches, as we will demonstrate with an example along the way.

5.4 Revising the sectors for Everettians

At first glance, Everettians might seem to make the von Neumann scheme’s preferred basis problem worse rather than better. In general, the final state after a measurement’s unitary dynamics can be expressed as many different superpositions; by removing dynamical collapse, it is not immediately clear why we should interpret the terms of *any* of these superpositions as representing different macroscopic states of affairs. To borrow Wallace’s terminology, the Everettian not only faces the ‘which basis?’ question we put to von Neumann’s scheme, but also a ‘why *any* basis?’ question (2012, p. 40).²⁴

²³One can describe subsystems with non-separable dynamics, in addition to non-separable kinematics, by using completely positive, trace-preserving (CPTP) maps. By Stinespring’s (1955) dilation theorem, CPTP maps generically arise from the restriction of the unitary evolution of a vector state in a larger Hilbert space. For brevity, we avoid further discussion of non-separable dynamics here; see, e.g., Bertlmann and Friis (2023) for more details.

²⁴As flagged in fn. 22, Schlosshauer (2007, p. 54), following Zurek (1981), takes von Neumann’s scheme to face the ‘why *any* basis?’ question, too. In short, these authors suggest that the basis chosen by the projection postulate ought to be motivated by the unitary dynamics of the measurement interaction—and so if multiple superpositions can express the state after this interaction, then the projection basis is underdetermined here, too. However, as Schlosshauer (2007, §2.15) also notes, due to Schmidt’s (1907) theorem, any decomposition $|\Psi\rangle = \sum_i c_i |S_i\rangle |A_i\rangle$ is unique provided that the $|A_i\rangle$ are orthogonal, each c_i is real, $\sum_i c_i^2 = 1$, and all c_i are different from each other. Thus, if the initial state of S can be, say, any superposition of z -spin eigenstates, then the projection basis of the von Neumann measurement (18) must be $\{A_\uparrow, A_\downarrow\}$ for a host of these initial states. That fact implies, in turn, that any choice of basis other than $\{A_\uparrow, A_\downarrow\}$ would have to vary based on the initial state—an undesirable feature if we think that the unitary dynamics themselves are responsible for the pointer basis. All of which is to say: contra Schlosshauer and Zurek, it is not clear that von Neumann’s

Decoherence theory gives compelling answers to both questions, allowing us to sketch an account of Everettian theory sectors that at least takes significant steps towards resolving the von Neumann scheme’s preferred basis problem. In short, the Everettian replaces the dynamical collapse of conservative macrostates with the dynamical ‘isolation’ of quasi-classical branches—where the relevant notion of ‘isolation’ is different from that for Everettian microstates.

One might reasonably expect a full answer to the ‘which basis?’ question to explain what makes branches ‘quasi-classical’, but formally characterizing this notion has proven tricky. Without such a characterization, we cannot easily use theory sectors to identify all and only quasi-classical branches as isolated macroscopic subsystems. However, it seems that the Everettian can simply drop the quasi-classical requirement to recover a thin but robust notion of isolation that any branch must satisfy. We argue that this thin notion of isolation needs only invoke a partial measure of chance values, allowing us to non-circularly use it to derive a full chance measure.

In short, decoherence theory answers the ‘why *any* basis?’ question by positing a particular basis along which a system is especially resistant to getting entangled with its environment; in other words, that basis is preferred by the dynamics of the system-environment interaction. Zurek (2003) calls this phenomenon environmentally-induced superselection, or *einselection* in short. To answer the ‘which basis?’ question, decoherence theory stipulates that some system-environment interactions are particularly natural or typical, and these dynamics end up preferring bases that resemble the classical states of our macro-world. Landsman nicely sums up the interplay of these two answers: ‘The claim, then, is that einselected states are often classical, or at least that classical states [...] are classical precisely because they are robust against coupling to the environment’ (2006, p. 518).

To unpack these answers, recall that the standard story of how decoherence gives rise to branching begins with a single orthodox microstate sector: we let $\Psi(t_0)$ be the initial state of SE , a subsystem and its environment. Now, however, we stipulate that the dynamics must be modeled with a three-part Hamiltonian:

$$\hat{H} = \hat{H}_S + \hat{H}_{\text{int}} + \hat{H}_E, \tag{32}$$

where \hat{H}_S and \hat{H}_E are the self-Hamiltonians of the system S and the environment E (respectively), and \hat{H}_{int} describes the interaction of the two. This interaction couples the eigenstates of some observable on S with a corresponding set of approximately orthogonal—and thereby, roughly, ‘independent’—environment states. Thus, the environment states play an analogous role to the eigenstates of the pointer observable in the von Neumann measurement scheme. During a branching event, \hat{H}_{int} dominates, very quickly extending any coherence in S over the designated eigenstates to coherence over the environment states. Then the self-Hamiltonians take over again, preserving this ‘isolation’ of each coupled system and environment state for all future times.

scheme lacks a dynamical justification for projecting onto a specific basis, even if we restrict our attention to the unitary part of the scheme. One can debate the merits of this justification, as it lacks decoherence theory’s appeal to robustness across many different system-environment interactions. However, this debate is not our focus; thus, in the main text, we have chosen to restrict our attention to how the ‘why *any* basis?’ question arises for Everettians, and how they address it.

Schlosshauer (2007) illustrates the basic idea with a toy spin- $\frac{1}{2}$ system. Assuming that the system and environment are initially separable, we have (with slight abuse of notation)

$$(a|\uparrow\rangle + b|\downarrow\rangle)|E\rangle \xrightarrow{\hat{H}_{\text{int}}} a|\uparrow\rangle|E_{\uparrow}\rangle + b|\downarrow\rangle|E_{\downarrow}\rangle, \quad (33)$$

where $\langle E_{\uparrow}|E_{\downarrow}\rangle \approx 0$. We model the environment as a collection of systems not dissimilar to our system of interest—in this case, a collection of N additional spin systems. We take \hat{H}_{int} to couple the system and the environmental spins in the following way:

$$\hat{H}_{\text{int}} = \frac{1}{2}\hat{\sigma}_z \otimes \sum_{i=1}^N g_i \hat{\sigma}_z^{(i)}, \quad (34)$$

where $\hat{\sigma}_z$ and $\hat{\sigma}_z^{(i)}$ are the Pauli z -spin operators of the system and the i th environmental spin, respectively, and g_i is the strength of their coupling. Now, we no longer simply stipulate that the measurement-like evolution of (33) holds between a system and an apparatus. Rather, we derive this evolution from the interactions of many systems of the same kind. Moreover, if (34) describes typical interactions of spin degrees of freedom, including those present in measurement devices, then it will be difficult to couple spin superpositions to pointer states of those devices. To summarize this idea, we can couple a measuring apparatus A to our system S and say that we expect the following dynamics:

$$(a|\uparrow\rangle + b|\downarrow\rangle)|A\rangle|E\rangle \xrightarrow{\hat{H}_{\text{int}}} a|\uparrow\rangle|A_{\uparrow}\rangle|E_{\uparrow}\rangle + b|\downarrow\rangle|A_{\downarrow}\rangle|E_{\downarrow}\rangle. \quad (35)$$

In other words, we expect that apparatus states encoding definite pointer readings will couple to system states preferred by decoherence. This much suffices to illustrate einselection's answer to the 'why *any* basis?' question. To the extent that it can motivate appropriate system-environment dynamics \hat{H}_{int} in general, i.e., beyond specific cases like the spin system studied here, it can answer the 'which basis?' question, too. Our task now is to investigate the extent to which we might formalize these answers, to build them into new macrostate sectors for Everettians.

For ease of exposition and compactness of notation, we will not define Everettian sectors that address the problem of non-repeatability, but one can do so in principle; to illustrate, we describe a scenario where the system's spin states are not the ones favored by decoherence. Recall that in the previous subsection, we motivated the coupling of the system's spin eigenstates to macroscopic superpositions of pointer states in the following way: first, we identified the center-of-mass wavefunction of the electron coupled to each of its spin eigenstates, and then we noted that these wavefunctions might fail to be well-separated and localized. Now, consider an environment E of air particles; when these air particles scatter off a heavy object such as the apparatus's imaging screen, they strongly couple to its localized center-of-mass wavefunctions. So, if the center-of-mass wavefunction coupled with a spin eigenstate is not localized before reaching the screen, we expect evolution roughly like

$$(a|\uparrow\rangle + b|\downarrow\rangle)|A\rangle|E\rangle \xrightarrow{\hat{H}_{\text{int}}} (c|\uparrow\rangle + d|\downarrow\rangle)|A_{\uparrow}\rangle|E_{\uparrow}\rangle + (e|\uparrow\rangle + f|\downarrow\rangle)|A_{\downarrow}\rangle|E_{\downarrow}\rangle. \quad (36)$$

Here, the coupling of the apparatus' pointer states with approximately orthogonal environment states explains why we observe localized measurement records. However, the system's

observables of interest—the spin observables—fail to couple strongly to the localized pointer states. At this juncture, we could rewrite (33) in terms of density operators before defining our Everettian sectors. Instead, we will henceforth mostly stick to vector states and Schrödinger dynamics (as flagged in §5.3).

As such, we simply take our Everettian microstates to be the same as orthodoxy’s. We also model Everettian macrostates after the orthodox ones, but we remove the dynamical projections. Everettians view the seeming appearance of dynamical projection as the result of choosing to ignore certain microscopic degrees of freedom, which we capture by projecting the microstate onto an appropriate macrostate. In particular, let $|E_i^{t_1}\rangle \in \mathcal{H}_{SE}$ denote the environmental state coupled to the i th pointer observable at time t_1 , and let $\hat{X}_i^{t_1}$ be the macroscopic projection onto its corresponding macrostate, $X_i^{t_1}$. Then,

$$\Psi_i^{t_1} = \hat{X}_i^{t_1} \hat{U}^{t_1 \leftarrow t_0} \Psi^{t_0}. \quad (37)$$

defines the state of the i th branch of SE at time t_1 . Note that this prescription makes the state of a branch relative to our choice of coarse-graining for the macrostates, as we expect. Now, suppose that another branching event occurs at a later time $t_2 > t_1$. Let $\Psi_{ij}^{t_1 t_2}$ denote the branch state recovering pointer states i at t_1 and j at t_2 , respectively; that is,

$$\Psi_{ij}^{t_1 t_2} = \hat{X}_j^{t_2} \hat{U}^{t_2 \leftarrow t_1} \hat{X}_i^{t_1} \hat{U}^{t_1 \leftarrow t_0} \Psi^{t_0}. \quad (38)$$

The ‘isolation’ of branch states refers to the requirement that each branch state nearly determines its own future trajectory. If we idealize, we can take this determination to be complete. This move yields the *branching criterion*, which states that only one prior branch can contribute to the state of any future branch; i.e.,

$$|\Psi_{ij}^{t_1 t_2}| \neq 0, |\Psi_{i'j}^{t_1 t_2}| \neq 0 \iff i = i', \quad (39)$$

where $|\cdot|$ denotes the norm in the Hilbert space. Strictly satisfying (39) requires idealizing the environmental record states to be *exactly* orthogonal—bringing them much closer to playing the role of pointer states in a von Neumann measurement. Indeed, under this idealization, different choices of coarse-graining make no difference to the dynamics; i.e., in (37), $\{\hat{X}_i^{t_1}\}$ is now a partitioning of orthogonal summands in $\hat{U}^{t_1 \leftarrow t_0} \Psi^{t_0}$, and any repartitioning of it yields the same $\Psi_i^{t_1}$.

The branching criterion provides a formal characterization of einselection’s answer to the ‘why *any* basis?’ question. However, it does not address the ‘which basis?’ question. Dowker and Kent (1996, §3.2) illustrate this issue dramatically by showing that the branching criterion on its own allows for desirable pointer states in one branch and superpositions of those states in another. Following Gell-Mann and Hartle (1993, 1997, 2007), Everettians like Wallace (2012) opt to treat the branching criterion as a necessary, but not sufficient, condition for (perhaps idealized) descriptions of quasi-classical histories.²⁵

²⁵We take our formal statement of the branching criterion from Wallace (2012). Gell-Mann and Hartle (1993) and Dowker and Kent (1996) in fact discuss a slightly different condition on histories, called *medium decoherence*, which relates to the branching criterion in the following way: every branching history is medium decoherent, and every medium decoherent history is a coarse-graining of a branching history (Wallace, 2012, p. 93). Gell-Mann and Hartle claim that medium decoherence is ‘a natural condition to impose’, but they stress that it ‘is not, by itself, sufficient to single out a quasiclassical domain’ (1993, p. 3353).

How, then, should Everettians answer the ‘which basis?’ question? On the one hand, there seem to be several promising avenues towards giving quasi-classicality a formal characterization. Landsman (2006), for example, suggests combining decoherence with one of several formal frameworks used to sharpen discussions of the large-action (i.e., $\hbar \rightarrow 0$) limit. Gell-Mann and Hartle (1993, 1995) instead aim to recover quasi-classicality from an entropy-like measure on a coarse-graining of a branching structure, ideally recovering large-relative-to- \hbar actions as a consequence—although they explicitly assume the Born rule to do so. On the other hand, Everettians can reasonably forgo such formality. They might lean on plausibility arguments for suitable interaction Hamiltonians, or they might be happy to assess quasi-classicality on a case-by-case basis. Wallace seems to take the latter route, suggesting that quasi-classicality is ‘not the kind of criterion which can be formalized as a new law of physics’ but is rather a ‘criterion for emergent structure’ (2012, p. 99).

Giving quasi-classicality a formal definition would help us define theory sectors that pick out all and only quasi-classical branches as isolated microstates. By incorporating an answer to the ‘which basis?’ question, these macrostate sectors would end up being stricter than orthodoxy’s. Instead of developing these sorts of sectors, we take an approach closer to Wallace’s attitude. Namely, we schematize sectors for branches that are roughly as strict as orthodoxy’s sectors for macrostates by utilizing the branching criterion’s formal answer to the ‘why *any* basis?’ question. We view these sectors similarly to Coulombic sectors with linear boundary conditions for the electric potential: they miss differences that we can empirically discriminate (i.e., quasi-classicality), but they suffice to establish a robust and useful sense in which subsystems (i.e., branches) are isolated.

Given our aim to use branch sectors to derive a chance measure, we take extra care to flag when and where we use probabilistic concepts to schematize them. Namely, we start by partially interpreting the Hilbert space norm, stipulating that very close to 0 means negligible or nearly impossible. To be explicit, we assume the following correspondence between Hilbert space norms and chance measures on branches:

$$|\Psi_{ij}^{t_1 t_2}| \approx 0 \iff ch(\Psi_{ij}^{t_1 t_2}) \approx 0. \quad (40)$$

Recall that by the inferential link (C1), a chance measure approximately equal to zero means the relative frequencies of the event should be nearly zero. When we idealize a branching structure to satisfy the branching criterion (39) exactly, the approximate equalities in (40) become exact. After this idealization, (C1) does not imply that the measure-zero branches are strictly impossible; however, it does imply that they are negligible when measuring branches’ relative frequencies.²⁶ Thus, at each time t_1, t_2 , and so on, (C1) motivates us to ignore all the measure-zero branches.

How close to 0 in Hilbert space norm is close enough for this idealization to apply? At the very least, the answer should be somehow related to our choice of coarse-graining for the macrostates. We propose that further elaboration of this answer should be guided by one’s

²⁶This partial interpretation arguably amounts to a pragmatic reading of *Cournot’s principle*, i.e., Cournot’s claim that ‘the physically impossible event is therefore the one that has infinitely small probability’ (1984, p. 78). See Shafer (2023) for various readings and uses of this principle in many different approaches to probability, including both objective and subjective approaches. Steeger (2022) refers to something close to this partial interpretation as a ‘structural link’ (specifically, ‘the normalization link’), emphasizing how it relates Everettian subsystems to probabilistic subsystems.

answer to the ‘which basis?’ question. On a Wallace-style approach, one might decline to give a formal criterion for ‘close enough’. If we instead follow Landsman’s recommendation to couple decoherence with a formalization of the $\hbar \rightarrow 0$ limit, we might couple the largeness of a system’s action relative to \hbar to our tolerance of deviations away from a Hilbert space norm of 0. On at least these two approaches, applying the branching criterion does not require anything probabilistic above and beyond the partial interpretation stipulated by (40).

With either of these approaches in hand, we schematize a *branch sector* as a microstate Ψ evolving unitarily, a sequence of times $(t_1, \dots, t_m, \dots, t_n)$, and a history $\{X_i^{t_m}\}$ of macrostates for which Ψ ’s dynamics satisfy the branching criterion (39) exactly. In this case, the approximate equalities in the partial interpretation (40) are exact, and at each time t_m , the set of possible macrostates is an exhaustive set of mutually exclusive events. We equip these sectors with notions of extension and restriction between sequences and their subsequences. Note that the dynamics of branch sectors, (38), have a similar form as the dynamics of orthodox macrostate sectors, (21)—both are iterated sequences of unitary and projection operators indexed by the sequence of times t_m . The key structural difference afforded by the branch sectors is the more liberal notion of extension and restriction—to subsequences, rather than substrings. Basically, since the branching criterion ensures future macrostates are (to a very good approximation) unaffected by past ones, the observer is allowed to forget any given branching event in a sequence. As with orthodox macrostates, both unitary transformations and repartitionings are chance-preserving symmetries of branches.

In sum, Everettians do not change microstates; they can stick with the orthodox ones, or they can use the generalization from quantum measurement theory.²⁷ Their innovation lies in replacing orthodox macrostates with branches. We use branch sectors to characterize the subsystem structure imposed by the branching criterion, which differs only subtly from that of an orthodox macrostate sector. This much suffices to define a robust notion of ‘isolation’ of branches: to wit, we can say that a branch is isolated precisely when it can be described as a branch sector.

5.5 Challenges to Everettian isolation

What challenges does this new Everettian notion of isolation face? On the one hand, one might worry that using a partial interpretation of the Hilbert space norm to define it builds too much probability into the notion to use it to derive a full chance measure. In §5.5.1, we clarify the scope of our forthcoming chance derivation to defuse these sorts of circularity concerns. On the other hand, one might worry that our model of decoherence is too idealized, insofar as it fails to consider how decoherence might spread in space as well as time. In §5.5.2, we argue that while this issue might make an important difference in relativistic domains, it is safe to ignore it in non-relativistic ones.

²⁷If the Everettian does choose to generalize their microstates, they will also have to rewrite branch sectors using the density matrix formalism analogously to equation (31). See Chua and Chen (2025) for more illustrations of this type of re-expression.

5.5.1 Isolation does not require a measure of chance

We start with Mandolesi’s claim that removing a probabilistic interpretation from the quantum formalism scuttles Wallace’s solution of the preferred basis problem. Mandolesi worries that on non-quasi-classical branches, agents might not exist, undermining Wallace’s decision-theoretic approach. In our setup, if agents failed to exist, we would be unable to satisfy any of our functional criteria for either chance or isolation. However, on branches where agents do exist, we can satisfy each criterion—and that suffices to make sense of ‘chance’ and ‘isolation’ talk when and where it occurs. One might justifiably want their theory to do more explanatory work; i.e., one might want a quantum theory that explains why branches are quasi-classical, answering the ‘which basis?’ question and explaining why agents exist. However, it is not clear that we *require* an answer to this question to make good sense of both chance and isolation; answering the ‘why *any* basis?’ question seems to suffice.

Others have worried that answering the ‘why *any* basis?’ question requires reference to probability in a way that undermines Everettians’ chance derivations. Specifically, Baker (2007) and Dawid and Thébault (2015) argue that without a probabilistic interpretation of the Hilbert space norm, idealizing away branches with a norm near 0 lacks physical motivation. Dawid and Thébault (2025, p. 21) have sharpened their concern, clarifying that their target is an Everettian chance derivation that starts with an ‘entirely non-probabilistic account’ of the emergence of branches. They recommend, instead, starting with a partially-interpreted measure and recovering a full probability measure in the large-action limit. Their approach starts with partially-interpreted *quasi-probability measures*, which weaken Kolmogorovian probability measures by giving up the requirements of positivity and countable additivity.

Our approach is similar to Dawid and Thébault’s, except that we partially interpret the Hilbert space norm instead of a quasi-probability measure. This partial interpretation assumes our functional analysis of chance. In particular, we use the inferential link (C1) to interpret near-norm-0 as negligible. In other words, we do not seek to derive or justify our analysis of chance from any part of the Everettian theory; rather, we incorporate it into our statement of that theory. We view Everettian chance derivations as strictly analogous to our proposed resolution of the box factory paradox: they are a powerful way of pinning down a *specific* chance measure—not a way of figuring out what ‘chance’ means. Our approach thereby incorporates at least the moral of Dawid and Thébault’s critique. Moreover, if we use the large-action limit as a guide to the application of our measure-zero idealization, then our approach becomes even more strongly analogous to theirs.

In sum, Everettians can use the machinery of decoherence to recover a notion of isolation that requires neither a full measure of chance nor an answer to the ‘which basis?’ question. To wit, when paired with a partial interpretation of the Hilbert space norm, the branching criterion makes good sense of isolated macrostates that might be quasi-classical or not, just as a generalized Coulombic electrostatics makes good sense of isolated subsystems that might have vanishing electric fields at their boundaries or not.

5.5.2 Isolation might be challenged by how decoherence spreads

However, the Everettian is not fully off the hook. While the possibility that branches might be non-classical does not trouble their isolation, another possibility might: namely, that decoherence might spread in space as well as time. One might reasonably question whether branching is adequately modeled as occurring everywhere in space simultaneously.²⁸ This modeling assumption is tacit in the dynamics (34), since the coupling constants g_i are not functions of space or time; coupling occurs immediately at a given strength, regardless of how distant the i th spin system is from the system of interest. This modeling assumption persists in most applications of decoherence, including a de-idealization of (34) that Cucchietti et al. (2005) propose applying to nuclear magnetic resonance. In contrast, Wallace (2012, §8.5–8.6) sketches an explicitly local conception of branching, where decoherence spreads in space and time. As a first pass at unpacking this conception, we might de-idealize the decoherence interaction in (35) to occur in two steps, as follows:

$$\begin{aligned} (a|\uparrow\rangle + b|\downarrow\rangle)|A\rangle|E\rangle &\xrightarrow{\hat{H}_1} a|\uparrow\rangle|A\rangle|E_\uparrow\rangle + b|\downarrow\rangle|A\rangle|E_\downarrow\rangle \\ &\xrightarrow{\hat{H}_2} a|\uparrow\rangle|A_\uparrow\rangle|E_\uparrow\rangle + b|\downarrow\rangle|A_\downarrow\rangle|E_\downarrow\rangle. \end{aligned} \tag{41}$$

Here, \hat{H}_1 describes the decoherence of the spin system S with its environment E , after which \hat{H}_2 decoheres the composite system SE with the pointer states of the apparatus A . This establishes how decoherence might spread over time; if we have reason to believe that S , E , and A are each well-localized, it might also describe a sense in which decoherence spreads throughout space. Blackshaw et al. (2024) create a toy model that runs with this idea, de-idealizing the dynamics of (34) to demonstrate how decoherence spreading might work for a collection of spin systems.

Whether such de-idealizations will make an empirical difference to future physics is an interesting open question. One can imagine a robust model of spreading decoherence underpinning a covariant formulation of measurement histories. Such an approach would avoid adopting a preferred frame’s time coordinate, as we tacitly do in our descriptions of both conservative and Everettian macrostates. However, it seems likely that such developments would take us beyond the purview of *non-relativistic* quantum mechanics. So, as far as the Everettian dealing with the non-relativistic domain is concerned, branch sectors provide a thin, robust, and stable characterization of the sense in which macrostates are isolated subsystems that does not rely on locality. This isolation suffices to fill in quantum chances’ link with certainty, and so we take it to suffice to derive the values of those chances.

With that said, we turn to how our functional analysis of isolation with theory sectors begets a thin, core approach to deriving Everettian chances from symmetry. We show how this core approach helps ease disagreements among Everettians and defuse salient criticisms from non-Everettians.

²⁸Cf. our discussion of local versus global branching in the next section. Note, though, that when we question whether branching is adequately modeled as occurring everywhere in space simultaneously, we are not questioning Sebens and Carroll’s (2018) ‘global branching’ view; the latter is saying that *even if* we can make sense of spreading decoherence, the universe splits everywhere when it splits somewhere.

6 Sectors unite and strengthen Born rule derivations

With the above work on sectors and isolation in hand, we are now in a position to turn to different approaches to probability in the Everett interpretation. The three uncertainty-based approaches to quantum probability we consider here—the Deutsch–Wallace theorem (Deutsch 1999, Wallace 2010, 2012), Sebens and Carroll’s (2018) derivation, and McQueen and Vaidman’s (2019) modification of the latter—all claim that their derivations have something to do with symmetry. Wallace writes that his proof is driven by the fact that ‘in Everettian quantum mechanics not just the laws, but the actual microstate of the system are invariant under a symmetry transformation, as could not be the case if only one outcome was to occur’ (2012, p. 151). Sebens and Carroll do not explicitly distinguish between one- and many-outcome symmetries, but still require that transformations of subsystems are ‘symmetries of the dynamical laws’, i.e., symmetries of unitary evolution maps (2018, p. 48). Similarly, McQueen and Vaidman (2019) base their derivation explicitly on systems with rotational symmetry.

None of these authors attend to the difference between microstate and macrostate symmetries illustrated in the previous section. This elision might appear benign at first glance, as unitary transformation maps instantiate both sorts of symmetries. However, we take the elision to create avoidable conflict among Everettians and to obscure the core of the chance derivation. As such, our strategy in the present section is to leverage the work we have already undertaken in order to illustrate a core approach to deriving Everettian chances with macrostate symmetries—one which exists despite the different motivations and starting points of the above-mentioned authors—and then illustrate how this approach resolves these tensions in the literature.

In short, the core approach applies the ontic separability principle, **OSP**, to the chance-preserving symmetries of macrostate sectors: namely, repartitionings and unitary transformations. Since each of these symmetries amounts to a relabeling of the macrostates, we do not expect them to alter the relative frequencies. Thus, we use **OSP** to impose the intrinsicity constraint:

$$ch(X_i | S) = ch(X_{i'} | S) = ch(X'_i | S'), \quad (42)$$

where $X_{i'}$ and X'_i denote the actions of a repartitioning and a unitary transformation, respectively. Equation (42) combines our conclusions (23) and (26) from our study of macrostate symmetries in §5.1.3; it also corresponds to step (3) in our box factory derivation. Just as we did in that argument, we now look for pairs of states S and S' that are identical to derive specific chance values. This approach applies equally well to both orthodoxy and Everett; in the latter case, we simply specify that macrostates are branches.

Let us illustrate how the core approach works for an equal-weight, two-outcome case, returning to the symmetry we illustrated for orthodox macrostate sectors in §5.1.3. There, we applied a symmetry map flipping the labels of ‘up’ and ‘down’ spin- $\frac{1}{2}$ states, given by equation (24). This time, let’s include an environment E interacting with the spin system S before the latter interacts with our Stern–Gerlach device A :

$$(a|\uparrow\rangle + b|\downarrow\rangle)|A\rangle|E\rangle \xrightarrow{\hat{U}(t_1, t_0)} a|\uparrow\rangle|A_\uparrow\rangle|E_\uparrow\rangle + b|\downarrow\rangle|A_\downarrow\rangle|E_\downarrow\rangle. \quad (43)$$

Let t_1 be either the time at which either dynamical projection occurs or decoherence finishes,

such that the right-hand-side of equation (43) specifies a microstate S at that time with macrostates X_1 and X_2 containing the first and second summands of the superposition, respectively. Next, we apply a unitary transformation to (43) that obtains the following:

$$(a|\downarrow\rangle + b|\uparrow\rangle)|A'\rangle|E'\rangle \xrightarrow{\hat{U}'(t_1, t_0)} a|\downarrow\rangle|A_\downarrow\rangle|E_\downarrow\rangle + b|\uparrow\rangle|A_\uparrow\rangle|E_\uparrow\rangle. \quad (44)$$

The right-hand-side superposition gives the microstate S' , and the macrostates X'_1 and X'_2 contain its first and second summands, respectively. Finally, in the case that $a = b$, we have

$$S' = S. \quad (45)$$

In this case, the new macrostates are a repartitioning of the originals; i.e., there is a repartitioning such that $X'_{1'} = X_2$ and $X'_{2'} = X_1$. Note that (45) corresponds to step (4) in our box factory derivation. Combining (42) and (45), the analogs of (3) and (4), we obtain the analog of (5):

$$ch(X_1 | S) = ch(X'_{1'} | S') = ch(X'_{2'} | S) = ch(X_2 | S). \quad (46)$$

Since the first and last terms are equal, the rules of probability then require that they each equal one-half. To make this story Everett-specific, one stipulates that equations (43) and (44) are descriptions of branching events, rather than von Neumann pre-measurements soon to be followed by a dynamical projection.

Each of the three symmetry-based approaches described above deviates only slightly from this core approach. In the following subsections, we illustrate how each approach deviates in turn, and we argue that returning to the core unites them while addressing approach-specific critiques from Everettians and non-Everettians alike. For each approach, we explicitly discuss only a finite-outcome, equal-weight case like the one above; such cases suffice to exhibit the key conceptual differences, and the remainders of the proofs are mostly mechanical.

6.1 Wallace's two-branch, equal-weight case

In short, Wallace's approach to deriving Everettian chances deviates from the core approach by its lack of an intrinsicity condition. That is to say, roughly, that he allows for

$$ch(X_1 | S_1) = ch(X_2 | S_2) \quad (47)$$

whenever $ch(X_1 | S_1) = ch(X_2 | S_1)$ and $S_1 = S_2$, but he does not identify chance values when these macrostates and microstates are merely related by symmetry (i.e., when $X'_1 = X_2$ and $S'_1 = S_2$). Thus, constraints on chance must arise from identity alone, without the aid of symmetry. Among other things, this approach might help Wallace differentiate chance in a many-worlds theory from chance in a single-world theory: the latter could never derive chance values from the identity requirement alone, the thought roughly goes, as different outcomes seem to correspond to different states by fiat. However, we argue that this move ends up requiring dynamics that threaten the subsystem recursivity of Everettian macrostates.

To show how, we consider the two-branch, equal-weight case in Wallace's (2012) decision-theoretic Everettian derivation of the Born rule. The full derivation utilizes a total of ten

axioms: four ‘richness’ axioms that, roughly, enrich the state of a branch sector to closely mirror a Bayesian decision problem; and six ‘rationality’ axioms that, roughly, operationalize the role played by symmetries in the core approach. With this picture, Wallace refers to unitary dynamics as ‘acts’ (denoted U) available to agents in certain states, and he uses the term ‘reward’ to refer to sets of macrostates that yield the same payout for an agent. By operationalizing dynamics with acts and macrostates with rewards, Wallace builds into his Born-rule derivation a decision-theoretic justification of the probability axioms in the style of Savage (1972). The details of this portion of the proof do not matter for the conceptual difference with the core approach, so we will provide a rough translation of the key axioms into chance-theoretic terms along the way.

Two of the axioms, **ordering** and **diachronic consistency**, are key parts of standard expected utility (EU) theory that Wallace notes are ‘as relevant in the classical as in the quantum context’ (2012, p. 167). In either context, EU theory specifies a preference relation $U \succeq_\psi V$ (read as ‘the agent in the state ψ prefers U to V ’) for any two acts U and V available to an agent at a given time. **Ordering** specifies that the preference relation \succeq_ψ is a total ordering; hence, the agent always has a preference between any two acts, or else judges them to be equally good options. **Diachronic consistency** asserts that agents take a sophisticated approach to sequential reasoning, in the following sense: if, at $t = 0$, an agent prefers U followed by V_1 at t to U followed by V_2 at t , then if they perform U , they will still prefer V_1 to V_2 at t . Roughly speaking, the agent’s preferences do not change over time. Decision theorists working in the classical domain have criticized both of these axioms.²⁹ For now, we take them as writ, as they are not directly relevant to the difference from our core.

One of the remaining axioms (‘branching availability’) roughly ensures that branches occur with all possible weights for all possible outcomes. An additional three axioms (‘reward availability’, ‘problem continuity’, and ‘solution continuity’) ensure the state space is rich enough for Savage’s representation theorem to recover real-valued probability functions. Since these four axioms do not feature explicitly in the equal-weight case, we omit them from our exposition.

That leaves us with the following four axioms:

Erasure. ‘Given a pair of states $\psi \in E$ and $\phi \in F$, where E and F are macrostates in the same reward, there is an act U available at E and an act V available at F such that $U\psi = V\phi$.’ (Wallace 2012, p. 166)

Microstate indifference. ‘An agent doesn’t care what the microstate is provided it’s within a particular macrostate.’ (Wallace 2012, p. 170)

Branching indifference. ‘An agent doesn’t care about branching per se: if a certain operation leaves his future selves in N different macrostates but doesn’t change any of their rewards, he is indifferent as to whether or not the operation is performed.’ (Wallace 2012, p. 170)

State supervenience. An agent’s preferences between acts depend only on what physical state they actually leave his branch in: that is, if $U_1\psi_1 = U_2\psi_2$ and $V_1\psi_1 = V_2\psi_2$, then an agent who prefers U_1 to V_1 given that the initial state is ψ_1

²⁹For an overview of their criticisms, see Steele and Stefánsson (2025).

should also prefer U_2 to V_2 given that the initial state is ψ_2 .³⁰

To clarify how this last axiom meets up with our approach, first suppose that the acts U_i lead to macrostates X_1 and X_2 with rewards r_1 and r_2 , respectively, where the agent prefers r_1 over r_2 . Additionally, suppose that the acts V_i yield these same macrostates, but with the rewards swapped. Intuitively, given this setup, if an agent starting in ψ_1 is indifferent between U_1 and V_1 , then they should think that X_1 is just as likely as X_2 ; in fact, the equality of these outcomes' probability values is a consequence of Savage's representation theorem. Likewise, they should think X_1 and X_2 are equally likely if they start in ψ_2 and are indifferent between U_2 and V_2 . Representing the final states as $U_1\psi_1 = U_2\psi_2 = S_1$ and $V_1\psi_1 = V_2\psi_2 = S_2$, we have $ch(X_1 | S_1) = ch(X_2 | S_1)$ and $ch(X_1 | S_2) = ch(X_2 | S_2)$. In the case that $S_1 = S_2$, these equalities imply equation (47).³¹

Indeed, Wallace's two-branch, equal-weight case more or less rehearses this reasoning in reverse, for the special case that $S_1 = S_2$, to obtain an agent's indifference between two acts:

For a simple case, suppose we have two acts (A and B, say): in each, a system is prepared in a linear superposition $\alpha|+\rangle + \beta|-\rangle$ and then measured in the $\{|+\rangle, |-\rangle\}$ basis. On act A, a reward is then given if the result is '+'; on B, the same reward is given on instead '-'. The resultant states are

$$\begin{aligned} \text{A: } & \alpha|+\rangle \otimes |\text{reward}\rangle + \beta|-\rangle \otimes |\text{no reward}\rangle; \\ \text{B: } & \alpha|+\rangle \otimes |\text{no reward}\rangle + \beta|-\rangle \otimes |\text{reward}\rangle. \end{aligned} \tag{48}$$

By **erasure**, there will exist acts available to the agent's future self in the reward branch (for both A and B) which erase the result of what was measured, leaving only the reward. Performing these transformations, and the equivalent erasures in the no-reward branch, leaves

$$\begin{aligned} \text{A-plus-erasure: } & \alpha|0\rangle \otimes |\text{reward}\rangle + \beta|0'\rangle \otimes |\text{no reward}\rangle; \\ \text{B-plus-erasure: } & \beta|0\rangle \otimes |\text{reward}\rangle + \alpha|0'\rangle \otimes |\text{no reward}\rangle. \end{aligned} \tag{49}$$

Now, by **branching indifference**, the agent's future selves are indifferent to whether this erasure is or is not performed. (Branching indifference is needed because we have no guarantee that erasures are nonbranching; if we did, **microstate indifference** would suffice.) So by **diachronic consistency**, the original agent is indifferent between A and A-plus-erasure, and between B and B-plus-erasure.

But now: if $\alpha = \beta$, then A-plus-erasure and B-plus-erasure leave the system in the same quantum state. So by **state supervenience**, the agent is indifferent between them. Since we know from **ordering** that preferences are transitive, the agent must also be indifferent between A and B. (2012, pp. 172–173, emphasis added)

³⁰This statement of the state supervenience condition is identical to Wallace's (2012, p. 170), except that we have relabelled the pair U, U' as U_1, U_2 , the pair V, V' as V_1, V_2 , and the pair ψ, ψ' as ψ_1, ψ_2 , since we reserve the prime notation for denoting the result of mapping by a symmetry.

³¹The authors thank Catrin Campbell-Moore for assistance with this translation.

To start, let S_1 refer to the microstate of A-plus-erasure, and let the macrostates X_1 and X_2 stand in for the the summands of that state in (49), in the order in which they appear; in the same manner, let S_2 refer to B-plus-erasure’s microstate and let Y_1 and Y_2 stand in its summands. Note that the assertion $\alpha = \beta$ establishes the equalities $X_1 = Y_1$, $X_2 = Y_2$, and $S_1 = S_2$, as needed for the chance-theoretic analog of Wallace’s state supervenience condition, (47), to apply.

Next, note that branching indifference appears in this derivation with something of an apology. It is needed to deal with erasures, which are in turn needed to ensure that the microstates in equation (49) are identical (and not merely symmetric) as required by state supervenience. So, consider what happens when we modify state supervenience to include an intrinsicity condition that tracks the chance-preserving symmetries of branch sectors:

State intrinsicity. An agent’s preferences between acts depend only on what physical state they actually leave his branch in *up to unitary transformations and repartitionings that preserve the reward structure*: that is, if for X_i any macrostate, X'_i and $X_{i'}$ have the same rewards as X_i , and if $(U_1\psi_1)' = U_2\psi_2$ and $(V_1\psi_1)' = V_2\psi_2$, then an agent who prefers U_1 to V_1 given that the initial state is ψ_1 should also prefer U_2 to V_2 given that the initial state is ψ_2 .

Note that if we trivialize the symmetries, such that the unitary map $(\cdot)'$ acts as the identity and $X_{i'} = X_i$, **state intrinsicity** is just **state supervenience**. However, we can apply the former condition directly to equation (48), letting $(A\psi)'$ and $(B\psi)'$ be the results of swapping ‘+’ with ‘−’. **State intrinsicity** implies that if the agent prefers A to B at ψ , then they should prefer A' to B' at ψ' . In the case that $\alpha = \beta$, $(A\psi)' = B\psi$ and $(B\psi)' = A\psi$, so the agent who prefers A to B at ψ should also prefer B to A at ψ ; thus, by **ordering**, they should be indifferent between the two acts.

In short, augmenting supervenience with intrinsicity removes the need for erasures. While it also removes the need for microstate and branching indifference, there is a sense in which the spirit of these axioms is preserved: namely, the reward-preserving symmetries of the intrinsicity condition identify branch sectors as the isolated subsystems that matter for chances. Thus, removing the **erasure** axiom from Wallace’s argument essentially returns us to the core approach, albeit with an explicitly epistemic way of deriving the rules of probability built in.

Moreover, on closer inspection, the availability of erasure acts seems hard to justify. Erasures are meant to reflect the fact that a system’s microstate makes no difference to an agent as long as the agent gets the same reward. However, erasures do not enter the model as features of an agent’s ignorance, but rather as dynamical acts that the agent can perform. Wallace flags this tension in his motivation of erasure:

Erasure is slightly more complicated. It effectively guarantees that an agent can just forget any facts about his situation that don’t concern things he cares about (i.e. by definition: that don’t concern where in the reward space he is). In thinking about it, it helps to assume that any reward space has an ‘erasure subspace’ available (whose states correspond to the agent throwing the preparation system away after receiving the payoff but without recording the actual result of the measurement, say). An ‘erasure act’ is then an act which takes the quantum

state of the agent’s branch into the erasure subspace; the agent is (by construction) indifferent to performing any erasure act, and since he lacks the fine control to know which act he is performing, all erasures should be counted as available if any are. (2012, p. 167)

The scare quotes around ‘erasure subspace’ seem to convey the difficulty of physically motivating the reification of an agent’s ignorance of the microstate into a feature of the space of dynamical possibilities. Likewise, Wallace’s description of agents ‘throwing the preparation system away’ seems only tenuously related to the acts defined in the erasure axiom, which requires two agents to map their full microstates to one and the same state.

Even worse, as Mandolesi (2019, p. 42) argues, erasures stand in strong tension with the linearity of unitary dynamics. The gist of the issue runs as follows: erasures are pairs of acts that take two states in orthogonal branches to the same state; thus, they change the angle between those two states (to 0 degrees, to make them identical), and so they cannot be combined into a single act that is unitary. However, it seems unnatural to forbid such a combination of acts if they are each allowed individually. More precisely: let $\psi \in M$ and $\phi \in N$, where M and N are orthogonal macrostates in the same reward. By the erasure axiom, there are acts U and V available at M and N such that $U\psi = V\phi$. Now, note that M and N are sub-macrostates of $M \vee N$. In the superposition microstate $\psi + \phi$, it seems like the two versions of an agent on branches M and N should be able to perform U and V at the same time, yielding an act $W = U \vee V$ available at $M \vee N$. However, that cannot be the case. If it were, then we would have $W\psi = W\phi$, and so W would fail to be unitary.

Our functional analysis of isolation enables us to sharpen Mandolesi’s point: some erasures seem to lack natural extensions, threatening macrostate sectors’ subsystem recursivity. In other words, we diagnose the feature of erasures to which Mandolesi points as unnatural because it goes against natural desiderata for a theory sector. Recursivity is not broken outright, as U and V might each have different extensions to $M \vee N$ if we require the two versions of the branching agent to perform their acts at different times. However, this requirement makes the dynamics available to one version depend sensitively on the dynamics available to the other—which stands in tension with our desire to think of the two versions of the branching agent as isolated from each other.

With these issues in mind, let us turn to comments that Wallace offers that might motivate skepticism about an intrinsicity condition like ours:

If the dynamical laws are stochastic—as in a dynamical-collapse theory, for instance—then they already embed information about what the probabilities are, and it would seem question-begging to use the symmetries of those laws as a derivation of the probabilities. If instead we just leave the stochastic part of the dynamics unspecified (if, say, we consider quantum mechanics with the Schrödinger evolution interrupted periodically by an unspecified collapse mechanism), we leave open the possibility that the stochastic processes break the symmetries of the deterministic laws. (2012, p. 147)

We take our application of **OSP** to our box factory theory to defuse these concerns. In that case, there is no unique primitive distribution, and we stipulate a sense in which the theory is complete: namely, that it is empirically adequate for describing relative frequencies

in a given domain. Of course, the box factory theory had no dynamics whatsoever, but our application of **OSP** to orthodox quantum mechanics illustrates that these two features arise in that theory, too.³² We conclude that the supposed victory of many worlds over one is dubious; moreover, it seems more important to retain a natural subsystem structure than to secure that victory. Thus, Wallace seems better off augmenting supervenience with intrinsicity to forgo erasures, thereby adopting (an appropriately epistemicized version of) the core approach.

As a bonus, the core approach de-emphasizes the importance of the diachronic consistency axiom. Motivated by duplication thought experiments, Sebens and Carroll (2018) reject that this axiom states a universally-applicable principle of rationality. These considerations, combined with the aforementioned criticisms of classical decision theorists, might motivate a revision of Wallace’s approach that simply defends the applicability of diachronic consistency to the Everettian case. We will not pursue this proposal here, however; instead, we simply note that moving to the core might help Wallace to accommodate Sebens and Carroll’s criticism.

6.2 Sebens and Carroll’s two-branch, equal-weight case

Sebens and Carroll’s (2018) approach starts from intuitions very close to ours: they posit a separability principle to supplement cases of uncertain reasoning where the principle of indifference proves indecisive. Indeed, we have modeled our ontic separability principle, **OSP**, after their ‘epistemic separability principle’, **ESP**, for ease of comparison. Unlike **OSP**, **ESP** applies specifically to cases of self-locating uncertainty across multiple versions of an agent:

ESP: Suppose that the universe U contains within it a set of subsystems S such that every agent in an internally qualitatively identical state to agent A is located in some subsystem that is an element of S . The probability that A ought to assign to being located in a particular subsystem $X \in S$ given that they are in U is identical in any possible universe which also contains subsystems S in the same exact states (and does not contain any copies of the agent in an internally qualitatively identical state that are not located in S):

$$c(X | U) = c(X | S). \tag{50}$$

(Sebens and Carroll, 2018, p. 40)

Here, $c(X | U)$ means the self-locating probability of being in subsystem X given that one inhabits the universe U . Note that the terminology of S being a ‘set of subsystems’ might mislead one in identifying the role separability plays; to avoid confusion, one can think of S

³²Steeger (2022) also points to a sense in which these features are satisfied by a version of pilot-wave theory. However, note that March (2025) levies a number of cogent criticisms against the particular symmetry argument for Bohmian chances presented in that work. Here, we have aimed to address two of March’s key points: namely, (a) that Steeger (2022) conflates operational and functional analyses of chance, and (b) that Lewisian chances for Bohmian mechanics must necessarily be trivial, as Lewis denies that deterministic physical theories can yield genuine chances. In the present work, we take (C1)–(C4) to stipulate a functional analysis of chance, and we do not commit to any sense in which this analysis is operational. Moreover, we do not require a subjectivist to flesh out the expert credence scheme (2) with Lewis’s chance-credence norm, and we acknowledge that they might not be able to use our chance derivations if they do so.

as a state description common to a set of agents in a multiverse, one that marks all such agents as ‘internally qualitatively identical’. It is physical information in U above and beyond S that is meant to be redundant, as far as probability is concerned. As such, S in the above definition roughly recovers the sense in which we say S is ‘isolated subsystem’ in **OSP**. As Sebens and Carroll state, the ‘gist’ of **ESP** is this: ‘The credence one should assign to being any one of several observers having identical experiences is independent of the state of the environment’ (p. 40).

Helpfully for our purposes, Sebens and Carroll continue to sharpen the sense in which their subsystems S are isolated beyond the statement of **ESP** itself. After introducing the principle, they offer the following clarification:

The essential idea is that a subsystem is a part of the larger system that can be considered as a physical system in its own right. Slightly more formally, we imagine that the overall state of a system can be decomposed into the states of various subsystems, so that two constraints are satisfied: (i) the state of each subsystem, perhaps with some additional information about how the subsystems are connected, can be used to uniquely reconstruct the original state; and (ii) the information contained within each subsystem’s state is enough to specify its immediate dynamical evolution, as long as the other subsystems are not influencing it. (Sebens and Carroll, 2018, p. 41)

Note that (i) and (ii) cover very similar ground to subsystem recursivity. Thus, S in Sebens and Carroll’s **ESP** fulfills our functional criterion (II) for isolation.

However, Sebens and Carroll adopt the subsystem structure of microstates to define the sense in which S is isolated, rather than that of macrostates. In particular, they use the density matrix generalization of microstate sectors that we reviewed in Section 5.3, a stance they clarify in an appendix:

In general, the reduced density matrix for the composite system AB , $\hat{\rho}_{AB}$, cannot be constructed uniquely from the separated reduced density matrices for A and B . The matrix also encodes facts about the entanglement between A and B . But as these are just facts about how A and B are connected, condition (i) is satisfied. To see that condition (ii) is met, suppose that, at least for a time, subsystem A is isolated from everything else, E . Let \hat{U}_t be the unitary operator that gives the time evolution of the total state. Since A and E are non-interacting, $\hat{U}_t = \hat{U}_A \otimes \hat{U}_E$. The time evolution of $\hat{\rho}_A$ is then given by $\hat{U}_A \hat{\rho}_A \hat{U}_A^\dagger$ (\hat{U}_E is irrelevant). (Sebens and Carroll, 2018, p. 67)

As such, Sebens and Carroll take the following principle, **ESP-QM**, to be the realization of **ESP** in the quantum-mechanical context:

ESP-QM: Suppose that an experiment has just measured observable \hat{O} of system S and registered some eigenvalue O_i on each branch of the wavefunction. The probability that agent A ought to assign to the detector D having registered O_i in their branch when the universal wavefunction is Ψ , $c(O_i | \Psi)$, only depends on the reduced density matrix of A and D , $\hat{\rho}_{AD}$:

$$c(O_i | \Psi) = c(O_i | \hat{\rho}_{AD}). \quad (51)$$

(Sebens and Carroll, 2018, pp. 42-43)

In their criticism of the **ESP** derivation, Dawid and Friederich (2020, p. 716) note saliently that the reduced density matrix formalism might fail to track whether branching has occurred. In our setup, this issue amounts to the fact that Sebens and Carroll focus on microstate rather than macrostate sectors.

Moreover, neither **ESP** nor **ESP-QM** includes an intrinsicity condition; there is no analog of the condition $ch(X | S) = ch(X' | S')$ in **OSP**. Thus, Sebens and Carroll end up tacitly adopting Wallace’s condition that only identical states can constrain chance values. That is to say, in the present case, that only identical reduced density matrices can yield the same chances, and not merely symmetric ones.

The focus on microstate sectors gives rise to two curious moves in Sebens and Carroll’s chance derivation. First, they adopt a global view of branching. This view amounts to a particular way of interpreting the two steps of decoherence in the de-idealized dynamics (41), where the system S first decoheres with the environment E and then SE decoheres with the measuring device A . The global view says that there are two measuring devices in two universes once the first decoherence step occurs—that is, even when the device is still in its ready state. Sebens and Carroll further illustrate the view as follows:

[...] branching happens throughout the whole wave function whenever it happens anywhere. When the universal wave function splits into multiple distinct and effectively non-interacting parts, the entire world splits—along with every object and agent in it. (2018, p. 34)

As such, even when decoherence has not yet reached Alice in their equal-weight, two-outcome case, they can still infer that there are two copies of Alice uncertain about their location in the multiverse, because branching has occurred elsewhere.

Second, they derive chance values for just those instances in which decoherence has occurred for a system but not yet reached Alice. They present their equal-weight, two-branch case as follows:

Alice measures the z -spin of a single particle in the x -spin up state. One display (D1) will show the result of the experiment. If the spin is up, a second display (D2) will show \heartsuit . If it is down, a \diamond will appear on the second display. Alice is not immediately affected by the result; in particular, she is for a time unaware of the experiment’s outcome. The wave function of Alice, the detectors, the particle, and the environment (the rest of the universe) evolve from

$$|\Psi_0\rangle = |R_0\rangle_A |R\rangle_{D1} |R\rangle_{D2} |\uparrow_x\rangle |E_R\rangle \quad (52)$$

to

$$|\Psi_1\rangle = \frac{1}{\sqrt{2}} |R\rangle_A |\uparrow\rangle_{D1} |\heartsuit\rangle_{D2} |\uparrow_z\rangle |E_{\uparrow\heartsuit}\rangle + \frac{1}{\sqrt{2}} |R\rangle_A |\downarrow\rangle_{D1} |\diamond\rangle_{D2} |\downarrow_z\rangle |E_{\downarrow\diamond}\rangle. \quad (53)$$

To use **ESP-QM** to demonstrate that $P(\uparrow | \Psi_1) = P(\downarrow | \Psi_1) = \frac{1}{2}$, we will need to also consider an alternate scenario where the computer (part of the environment)

is programmed differently so that \heartsuit displays if down is measured and \diamondsuit displays if up. Then the post-measurement, pre-observation wave function would be:

$$|\Psi_2\rangle = \frac{1}{\sqrt{2}}|R\rangle_A|\uparrow\rangle_{D1}|\diamondsuit\rangle_{D2}|\uparrow_z\rangle|E_{\uparrow\diamondsuit}\rangle + \frac{1}{\sqrt{2}}|R\rangle_A|\downarrow\rangle_{D1}|\heartsuit\rangle_{D2}|\downarrow_z\rangle|E_{\downarrow\heartsuit}\rangle. \quad (54)$$

Step 1: Focus first on Alice and D1. The [(Alice + Detector 1)] reduced density matrices for Ψ_1 and Ψ_2 are the same,

$$\hat{\rho}_{AD1}(\Psi_1) = \hat{\rho}_{AD1}(\Psi_2) = \frac{1}{2}|R\rangle_A|\uparrow\rangle_{D1}\langle\uparrow|_{D1}\langle R|_A + \frac{1}{2}|R\rangle_A|\downarrow\rangle_{D1}\langle\downarrow|_{D1}\langle R|_A. \quad (55)$$

ESP-QM requires that the probabilities Alice assigns to the possible spin results be the same in these two universes as they have the same [(Observer + Detector 1)] reduced density matrix,

$$P(\downarrow | \Psi_1) = P(\downarrow | \Psi_2). \quad (56)$$

Step 2: If we ask what probability Alice should assign to the display being \heartsuit , we need to consider the reduced density matrix generated by tracing over $D1$, the spin of the particle, and the environment. Ψ_1 and Ψ_2 agree on $\hat{\rho}_{AD2}$. By **ESP-QM**, the probabilities assigned to \heartsuit must be equal,

$$P(\heartsuit | \Psi_1) = P(\heartsuit | \Psi_2). \quad (57)$$

Step 3: Next, note that the \heartsuit -branches ‘just are’ the \uparrow -branches in Ψ_1 and the \heartsuit -branches just are the \downarrow -branches in Ψ_2 . Thus Alice is in the \heartsuit -branch of $\hat{\rho}_{AD2}(\Psi_1)$ if and only if she is in the \uparrow -branch of $\hat{\rho}_{AD1}(\Psi_1)$. Similarly, she is in the \heartsuit -branch of $\hat{\rho}_{AD2}(\Psi_2)$ if and only if she is in the \downarrow -branch of $\hat{\rho}_{AD1}(\Psi_2)$. Therefore, Alice must assign

$$\begin{aligned} P(\uparrow | \Psi_1) &= P(\heartsuit | \Psi_1) \\ P(\downarrow | \Psi_2) &= P(\heartsuit | \Psi_2). \end{aligned} \quad (58)$$

Step 4: Putting together the results in Equations (56)–(58), we see that the probability of being on a \uparrow/\heartsuit -branch must be the same as that for being on a \downarrow/\diamondsuit -branch: $P(\uparrow | \Psi_1) = P(\downarrow | \Psi_1)$. So, the unique rational degrees of belief in the first scenario consider each branch to be equiprobable. (Sebens and Carroll, 2018, pp. 43–45)

McQueen and Vaidman (2019, p. 22) argue convincingly that this derivation applies *only* to cases where decoherence has not yet reached Alice—and so it seems that Alice’s credences can only be legitimate in the split-second before decoherence reaches her. Explicitly, they note that the application of **ESP-QM** in Step 2 becomes unsound once we fast-forward to the end of the decoherence process, after which Alice has split into orthogonal states. So, labeling Alice’s post-branching states R_\uparrow and R_\downarrow , we may write the fully decohered Ψ_1 and Ψ_2 as

$$|\Psi_1\rangle = \frac{1}{\sqrt{2}}|R_\uparrow\rangle_A|\uparrow\rangle_{D1}|\heartsuit\rangle_{D2}|\uparrow_z\rangle|E_{\uparrow\heartsuit}\rangle + \frac{1}{\sqrt{2}}|R_\downarrow\rangle_A|\downarrow\rangle_{D1}|\diamondsuit\rangle_{D2}|\downarrow_z\rangle|E_{\downarrow\diamondsuit}\rangle \quad (59)$$

and

$$|\Psi_2\rangle = \frac{1}{\sqrt{2}}|R_\uparrow\rangle_A|\uparrow\rangle_{D1}|\diamond\rangle_{D2}|\uparrow_z\rangle|E_{\uparrow\diamond}\rangle + \frac{1}{\sqrt{2}}|R_\downarrow\rangle_A|\downarrow\rangle_{D1}|\heartsuit\rangle_{D2}|\downarrow_z\rangle|E_{\downarrow\heartsuit}\rangle, \quad (60)$$

respectively. But now,

$$\hat{\rho}_{AD2}(\Psi_1) = \frac{1}{2}|R_\uparrow\rangle_A|\heartsuit\rangle_{D2}\langle\heartsuit|_{D2}\langle R_\uparrow|_A + \frac{1}{2}|R_\downarrow\rangle_A|\diamond\rangle_{D2}\langle\diamond|_{D2}\langle R_\downarrow|_A \quad (61)$$

and

$$\hat{\rho}_{AD2}(\Psi_2) = \frac{1}{2}|R_\downarrow\rangle_A|\heartsuit\rangle_{D2}\langle\heartsuit|_{D2}\langle R_\downarrow|_A + \frac{1}{2}|R_\uparrow\rangle_A|\diamond\rangle_{D2}\langle\diamond|_{D2}\langle R_\uparrow|_A, \quad (62)$$

and so $\hat{\rho}_{AD2}(\Psi_1) \neq \hat{\rho}_{AD2}(\Psi_2)$.

Of course, $\hat{\rho}_{AD2}(\Psi_1)$ and $\hat{\rho}_{AD2}(\Psi_2)$ specify isolated and symmetric Everettian macrostates, in our sense of these terms; note that (61) and (62) take the form of the right-hand side of (31). Thus, if Sebens and Carroll were to modify **ESP-QM** to track symmetry in macrostate sectors, they would essentially recover our core approach. This move would accommodate McQueen and Vaidman’s desire to treat fully decohered branches while altering very little of the structure of the proof. We believe that Sebens and Carroll should accommodate this desire, as their cases of interest are macroscopic states of affairs—i.e., primarily, events corresponding to measurement outcomes.

McQueen and Vaidman also reject the global conception of branching on metaphysical grounds. They present an alternative metaphysics, which they illustrate as follows:³³

[...] when you see [a notebook in which has been recorded the result of a Stern–Gerlach experiment] on your desk and ask, ‘what is the probability that I am in the world with the spin- \uparrow result?’, you ask a meaningless question. For prior to having branched, there is only one ‘you’ that is in *both* the world with the spin- \uparrow result and the world with the spin- \downarrow result. (McQueen and Vaidman, 2019, p. 16)

On one sensible reading of this passage, McQueen and Vaidman seek to apply a Lewisian ‘overlap’ view to a local de-idealization of decoherence of the sort sketched in Section 5.5, such that the number of objects or selves or whatever else in a given spatial region depends on whether decoherence has reached it.³⁴

The global and local branching pictures are certainly incompatible, but we think that each seems plausible. More to the point: neither view makes a difference to the subsystem structure of Everettian macrostates for non-relativistic quantum mechanics, which isolate branches only after the interaction Hamiltonian stops dominating the dynamics of decoherence. In other words, branching could occur everywhere at once or spread in space and time, and both options yield the same isolated macrostates once decoherence has finished. Thus, if Sebens and Carroll were to adopt our core approach, they could (a) generalize their derivation to apply to agents for longer than artificially short time scales *and* (b) retain their global view of branching.

³³See also Vaidman (2020).

³⁴As they say in a footnote: ‘We do not consider the self as a “history” in a sense which leads to a “divergent” instead of an overlapping picture’ (McQueen and Vaidman, 2019, fn. 8).

6.3 McQueen and Vaidman’s three-branch, equal-weight case

McQueen and Vaidman claim to derive the Born rule with just two principles: *symmetry*, which says that ‘if the physical situation has a particular symmetry, whatever will happen should respect this symmetry’; and *local supervenience*, which says that ‘whatever happens in region A depends only on the quantum description of this region and its immediate vicinity’ (2019, p. 17). They do not offer formal characterizations of these principles, but it is clear that their local conception of branching underpins both.

As such, McQueen and Vaidman’s finite-branch, equal-weight case focuses on a case with spatially distant, sharply-localized wavepackets:

Consider a particle in a superposition of three very far apart well localized wave packets centered at positions labeled A , B , and C . These three positions lie on the circumference of a perfect circle such that the particle has three-fold rotational symmetry. That is, the system appears the same when it is rotated by one third of a full turn about the circle’s center. Three identical detectors, distributed over those same positions, which also exhibit three-fold rotational symmetry, simultaneously measure the presence of the particle.

The quantum mechanical description is as follows. The initial state of the particle is:

$$(1/\sqrt{3})(|a\rangle + |b\rangle + |c\rangle) = \frac{1}{\sqrt{3}}(|1\rangle_{\mathbf{a}}|0\rangle_{\mathbf{b}}|0\rangle_{\mathbf{c}} + |0\rangle_{\mathbf{a}}|1\rangle_{\mathbf{b}}|0\rangle_{\mathbf{c}} + |0\rangle_{\mathbf{a}}|0\rangle_{\mathbf{b}}|1\rangle_{\mathbf{c}}), \quad (63)$$

where the right-hand-side is given in Fock representation, to help clarify the role that locality will play in our argument. The interaction process is described by

$$(1/\sqrt{3})(|1\rangle_{\mathbf{a}}|0\rangle_{\mathbf{b}}|0\rangle_{\mathbf{c}} + |0\rangle_{\mathbf{a}}|1\rangle_{\mathbf{b}}|0\rangle_{\mathbf{c}} + |0\rangle_{\mathbf{a}}|0\rangle_{\mathbf{b}}|1\rangle_{\mathbf{c}})|R\rangle_{\mathbf{A}}|R\rangle_{\mathbf{B}}|R\rangle_{\mathbf{C}} \quad (64)$$

$$\begin{aligned} &\rightarrow (1/\sqrt{3})(|1\rangle_{\mathbf{a}}|0\rangle_{\mathbf{b}}|0\rangle_{\mathbf{c}}|\checkmark\rangle_{\mathbf{A}}|R\rangle_{\mathbf{B}}|R\rangle_{\mathbf{C}} + \\ &\quad |0\rangle_{\mathbf{a}}|1\rangle_{\mathbf{b}}|0\rangle_{\mathbf{c}}|R\rangle_{\mathbf{A}}|\checkmark\rangle_{\mathbf{B}}|R\rangle_{\mathbf{C}} + \\ &\quad |0\rangle_{\mathbf{a}}|0\rangle_{\mathbf{b}}|1\rangle_{\mathbf{c}}|R\rangle_{\mathbf{A}}|R\rangle_{\mathbf{B}}|\checkmark\rangle_{\mathbf{C}}) \end{aligned} \quad (65)$$

$$\rightarrow \begin{cases} |1\rangle_{\mathbf{a}}|0\rangle_{\mathbf{b}}|0\rangle_{\mathbf{c}}|\checkmark\rangle_{\mathbf{A}}|R\rangle_{\mathbf{B}}|R\rangle_{\mathbf{C}} & \text{prob}_A \\ |0\rangle_{\mathbf{a}}|1\rangle_{\mathbf{b}}|0\rangle_{\mathbf{c}}|R\rangle_{\mathbf{A}}|\checkmark\rangle_{\mathbf{B}}|R\rangle_{\mathbf{C}} & \text{prob}_B \\ |0\rangle_{\mathbf{a}}|0\rangle_{\mathbf{b}}|1\rangle_{\mathbf{c}}|R\rangle_{\mathbf{A}}|R\rangle_{\mathbf{B}}|\checkmark\rangle_{\mathbf{C}} & \text{prob}_C \end{cases} \quad (66)$$

where signify the ready and clicked states of the detector in A , etc. The symmetry between A , B and C is explicit before measurement interaction (64). It is still explicit during the measurement unitary evolution (65). The symmetry principle then also requires that it will be manifested in the mixture (66) after the measurement. Therefore, $\text{prob}_A = \text{prob}_B = \text{prob}_C = 1/3$. (McQueen and Vaidman, 2019, p. 17)

Curiously, the fact that the symmetry is achieved by rotation of spatially distant states does not appear to play any special role in the derivation. In the end, the symmetries linking the three states are special cases of unitary transformations of branch states, as evidenced by (65). In particular, as we covered in our exposition of the core approach, there will be pairs

of environmental and system unitary symmetries that translate between any pair of the three branch states in equation (66) regardless of whether these states represent spatially distant systems.

Moreover, focusing on spatially distant, well-localized states seems to hamstring the scope of the derivation. Indeed, in order to achieve the generality of the core approach, McQueen and Vaidman are forced to tie arbitrary spatially-localized measurements to this rather specialized class of states:

The proof above considers a quantum state of a particle placed in different locations, and it is not applicable directly to a quantum state of a local system like a spin. But our proof is a consistency check: Born rule violation is inconsistent with relativity theory. Born rule violation for a local measurement of a spin will also lead to inconsistency. Quantum theory considers various legitimate methods of measurements. In particular, we can arrange a unitary swap mechanism between the spin state and the spatial state of the particle:

$$(\alpha|\uparrow\rangle + \beta|\downarrow\rangle)|ready\rangle \rightarrow |\uparrow\rangle(\alpha|a\rangle + \beta|b\rangle). \quad (67)$$

After this procedure we make a measurement of where the particle is. We have proved that this final measurement should comply with the Born rule, and the consistency of quantum theory tells us that direct local measurements of the spin should follow the Born rule too. (McQueen and Vaidman, 2019, p. 18)

Note that it might well be the case that a system’s spin degrees of freedom are isolated from its spatial degrees of freedom (in the microstate sense). This possibility licenses us to, for example, study the spin-chain decoherence model (34) without reference to any other system properties. Here, however, McQueen and Vaidman seem to render the symmetries of spin-only systems reliant on hidden degrees of freedom, threatening these systems’ sectorhood.

Charitably, McQueen and Vaidman aim to give a first pass at how a chance derivation should go in a relativistic generalization of the Everettian’s non-relativistic theory. They note briefly that their proof is meant to apply to ‘the relativistic generalization of the Schrödinger equation’ (p. 19), but they do not specify which generalization they have in mind (among, e.g., the Dirac equation, the Klein–Gordon equation, and so on). We think this project is a worthy one, although we would press McQueen and Vaidman to (a) provide an explicit, localized de-idealization of decoherence (of the sort sketched in §5.5.2) and (b) precisely characterize the relevant dynamical symmetries in the resulting picture. Moreover, it would be quite strange if such a project were to *invalidate* the core approach; one would naturally expect the latter to emerge from the low-velocity limit of the former. In such a case, McQueen and Vaidman should have no issue with decision-theorists or **ESPs** who adopt the core.

7 Discussion

We have argued that the apparent tensions among Everettian approaches to chance using uncertain or partial beliefs dissolve once we shift our attention to what these Everettians think is certain. They need to be certain of a given subsystem’s isolation to address the reference class problem associated with their preferred analysis of chance, and we offer a thin functional analysis of isolation in terms of (I1) instantiating a recursive subsystem structure and (I2) enabling empirical study. Getting clear on the subsystem structures that Everettians need enables us to unite decision-theoretic and self-locating uncertainty approaches under a thin common core that takes chance to be an intrinsic property of isolated branches. We use no notion of locality to describe how branches become isolated, resolving concerns about the need for either local or global branching; nor do we use any full measure of chance, defusing worries about circularity. Moreover, we formulate the common core as a specific application of a more general ontic separability principle, **OSP**. By showing how one can also use **OSP** to resolve van Fraassen’s box factory paradox, we illustrate a sense in which the logic of Everettian chance derivations might generalize to a broad range of physical theories.

To review, the key differences between our **OSP** derivation and the previous Born rule derivations of Wallace (2012), Sebens and Carroll (2018), and McQueen and Vaidman (2019) run as follows:

- Wallace’s (2012) derivation relies on *erasures*, acts which take subsystems in orthogonal branches to the same state. Erasures threaten Everettian branches’ subsystem recursivity by making acts available to one version of a branching agent sensitive to the acts available to another. Our common core avoids using erasures and retains a natural subsystem structure for branches.
- Sebens and Carroll’s (2018) derivation relies on *global branching*, a particular interpretation of a de-idealization of decoherence that models the process as spreading in space. Our common core relies only on the subsystem structure of the standard (non-spreading) decoherence models, so it accommodates, but does not require, global branching.
- McQueen and Vaidman’s (2019) derivation relies on a principle of *local supervenience* and invokes a relativistic generalization of the Schrödinger equation. We state our common core directly in terms of the kinematics, dynamics, and symmetries of non-relativistic quantum mechanics, requiring no principle of locality. Instead, we expect the common core to emerge from a relativistic approach in a suitable low-velocity limit.

For the latter two derivations, our revisions cost little: the authors can retain their respective local and global views of branching while agreeing that the common core captures the emergence of full chance measures in the non-relativistic domain. The situation for Wallace is more complicated. Recall that the need for erasures in Wallace’s proof arises from requiring that only identical (and not merely symmetric) quantum states constrain the values of chances. By emphasizing identity over symmetry, Wallace might recover a sense in which Everettian quantum mechanics better accounts for chance values than single-world theories. In our assessment, this apparent victory is not worth the threat it poses to branches’ subsystem recursivity. Thus, we conclude that our common core attractively unifies the key

commitments of the Everettians surveyed here while diagnosing why their differences have remained controversial.

It is not all wine and roses for Everettians, however! We ended up needing to specify two different subsystem structures corresponding to two different notions of isolation: one for microstates and one for macrostates. In one sense, this move is anathema to the Everettian ethos, which aims to unify all of quantum theory under a single dynamics. In another sense, the move amounts to a natural incorporation of lessons from statistical mechanics into an Everettian worldview. The extent to which this latter perspective might take precedence over the former, however, is a matter for future work.

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