Lorentzian quantum reality: postulates and toy models

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We describe postulates for a novel realist version of relativistic quantum theory or quantum field theory in Minkowski space and other background space–times, and illustrate their application with toy models.

1. Introduction

In previous papers [1,2], we described ideas aimed at defining realist and Lorentz covariant versions of relativistic quantum theory or quantum field theory in Minkowski space. The definitions also extend to other background space–times with appropriate asymptotic properties in the far future. This paper introduces new ideas into this programme, including a new and different proposal for the definition of beables.

The realist version of quantum theory proposed here is defined by postulates that involve only expressions obtained from the predictions that standard quantum theory would make for various specified measurements on the unitarily evolved quantum state. This makes it relatively easy to illustrate its implications in toy models. The proposal relies on the existence of particles, or wave packets, or field perturbations, that travel at light speed. In this sense, it is intrinsically relativistic. It is motivated by the fact that quantum electrodynamics suggests it is sensible—albeit not rigorously justifiable—in some circumstances to think of photons as individual particles that do indeed travel at essentially light speed in the vacuum. Most approaches to unifying quantum theory and gravity suggest similar intuitions about gravitons, although of course the direct evidence for quantum gravity is thin.

In a certain sense, these light-speed particles (packets, field perturbations) play the role of an environment, in that we infer the actual state of a physical system in a
given region of space–time from the final state of some of these particles. However, we do not need to postulate any formal separation between ‘system’ and ‘environment’ particles. Our postulates treat all types of particle, and all subsystems of the universe, on an equal footing.

The proposal is most simply illustrated in models that include photons travelling along light ray segments along with slower-moving systems made up of massive particles, with particle number and photon number both conserved. Although models of this type can be good approximations in the appropriate regime, they are not rigorously describable by relativistic quantum field theory, as particle numbers are not conserved in relativistic quantum field theory and as anyway we have no rigorous formulation of non-trivial relativistic field theories. Nor are they rigorously described by non-relativistic quantum mechanics, as we need both that photon wave packets travel at the speed of light and that massive particle wave packets travel at slower than light speed. Nonetheless, they incorporate standard intuitions about the behaviour of an essentially non-relativistic physical system interacting with photons in the environment.

Also, because the proposal requires distinguishing between future data that could arise inside or outside the future light cone of a given point in space–time, it does not behave well in the non-relativistic limit. However, one can try to define non-relativistic postulates that produce approximately similar pictures of reality, even if they rely on somewhat different assumptions. In a future paper we will consider some postulates of this type, which are also of some independent interest.

2. The reality problem for relativistic quantum theory

As noted in [2], as relativistic quantum field theory is not itself presently based on a rigorous foundation, we cannot hope to give a fully rigorous definition of a realist version of that theory. However, we can aim to give a conceptually coherent description of reality within relativistic quantum field theory, even if it involves quantities that presently have no rigorous mathematical definition. We can also illustrate the implications of our postulates in toy models that incorporate intuitions from relativistic quantum field theory but do not assume the full theory. These are the main aims of this paper. We put to one side here the precise proposals made in [2] in favour of the new simpler formulation given below.

We now suppose that the initial state $|\psi_0\rangle$ is given on some spacelike hypersurface $S_0$, and that some relativistic unitary evolution law is given. The Tomonaga–Schwinger formalism allows us to define formally the evolved state $|\psi_S\rangle$ on any hypersurface $S$ in the future of $S_0$ via a unitary operator $U_{S_0 S}$. These future hypersurfaces $S$ play the same role here as in [2]. That is, we initially give definitions for some fixed hypersurface $S$ in the far future of $S_0$, but have in mind ultimately taking the limit in which $S$ tends to future infinity. For this asymptotic procedure to work, we have to assume that the relevant limits are well defined. We examine this assumption in the toy models we consider.

For any given hypersurface $S$ in the future of the initial hypersurface $S_0$, we consider the effect of joint measurements of the local mass-energy density operators $T_S(x) = T_{\mu\nu}(x)n_\mu n_\nu$ carried out at each point $x \in S$, where $n_\mu$ is the forward-pointing timelike unit four-vector orthogonal to the tangent plane of $S$ at $x$. This gives us a probability distribution on possible mass-energy distributions $t_S(x)$ on $S$. In a universe in which physics starts on $S_0$ and ends on $S$, our picture of reality is that one $t_S(x)$ is randomly selected, from the Born rule probability distribution. In other words, there is a randomly selected final boundary condition on $S$, which is defined mathematically in the same way that it would be if $T_S(x)$ were actually measured on $S$. However, we treat this as simply as a mathematical algorithm. We do not suppose that a physical measurement actually takes place on $S$, or anywhere else. Our aim, instead, is to give a mathematical description of reality applicable to closed quantum systems, for which there are no external observers able to carry out measurements.

To give a description of reality between $S_0$ and $S$, we use the initial state on $S_0$, the randomly chosen final outcome data $t_S(x)$ on $S$, and the unitary evolution law arising from the quantum
dynamics. We use these data to calculate generalized expectation values of the stress-energy tensor \((T_{\mu\nu}(y))\), at each point \(y\) between \(S_0\) and \(S\). However, these expectation values are not obtained in the familiar way by simply considering a unitary evolution of the initial state to some spacelike hypersurface through \(y\). Instead, we use the following recipe.

\[\text{(a) Description of reality using generalized stress-energy tensor expectation values as beables}\]

We wish to define a generalized expectation value for the stress-energy tensor at a point \(y\) between \(S_0\) and \(S\), using post-selected final data \(t_S(x)\) on \(S\). More precisely—and this is an important new ingredient compared with our previous proposals—we will use the post-selected data \(t_S(x)\) for all points \(x \in S\) outside the future light cone of \(y\), and only for those points.

In words, our recipe is to take the expectation value for \(T_{\mu\nu}(y)\) given that the initial state was \(|\psi_0\rangle\) on \(S_0\), conditioned on the measurement outcomes for \(T_S(x)\) being \(t_S(x)\) for \(x\) outside the future light cone of \(y\). So, for any given point \(y\), our calculation ignores the outcomes \(t_S(x)\) for \(x\) inside the future light cone of \(y\).

Mathematically, we can calculate such expressions as follows. For any point \(y\) between \(S_0\) and \(S\), define the effective future boundary \(\Lambda^1(y)\) of \(y\) in our model to be \(\Lambda_1(y) \cup S^1(y)\). Here, \(\Lambda_1(y)\) is the set of points in the lightlike future of \(y\) that are either in the past of \(S\) or \(S\) itself, and \(S^1(y)\) is the set of points in \(S\) outside the future light cone of \(y\). Let \((S_i(y))\) be a sequence of smooth spacelike hypersurfaces that include \(y\) and include all points \(x \in S^1(y)\) such that the spacelike separation \(d(x, z) \geq \epsilon_i\) for all \(z \in \Lambda_1(y) \cap S\). Suppose that \(\epsilon_i \to 0\) as \(i \to \infty\), so that

\[
\lim_{i \to \infty} S_i(y) = \Lambda^1(y).
\]

In words, the \(S_i(y)\) are spacelike hypersurfaces that include almost all of the part of \(S\) outside the future light cone of \(y\) and include \(y\), and that tend to \(\Lambda^1(y)\) as \(i \to \infty\).

Now for any of the \(S_i(y)\), we can consider the Born rule probability distribution of outcomes of joint measurements of \(t_S(x)\) (for all \(x \in S \cap S_i(y)\)) and of \(T_{\mu\nu}(y)\). These are calculated in the standard way, taking the initial state \(|\psi_0\rangle\) on \(S_0\), unitarily evolving to \(S_i(y)\), and applying the measurement postulate there. This gives us a joint probability density function \(P(t_S(x), t_{\mu\nu}(y); x \in S \cap S_i(y))\). (The notation here means that we consider the joint probability for values of \(t_S\) for all \(x \in \text{the relevant set as well as for the value of } t_{\mu\nu}\text{ for the single value } y\).) By taking the limit as \(i \to \infty\), we obtain a joint probability density function \(P(t_S(x), t_{\mu\nu}(y); x \in S^1(y))\). From this, we can calculate conditional probabilities and conditional expectations for \(t_{\mu\nu}(y)\), conditioned on any set of outcomes for \(t_S(x)\) (for \(x \in S^1(y)\)), in the standard way.

Our mathematical description of reality, in a hypothetical world in which physics takes place only between \(S_0\) and \(S\) and in which the outcomes \(t_S(x)\) were randomly selected, is then given by the set of conditional expectations \(\{t_{\mu\nu}(y)\}\), for each \(y\) between \(S_0\) and \(S\), calculated as above. We stress that the calculations for the beables \(t_{\mu\nu}(y)\) at each point \(y\) all use the same final outcome data \(t_S(x)\). However, different subsets of these data are used in these calculations: for each \(y\), the relevant subset is \(\{t_S(x); x \in S^1(y)\}\).

We then consider the asymptotic limit in which \(S\) tends to the infinite future of \(S_0\). Suppose that \(S_1\) is some fixed hypersurface in the future of \(S_0\). Let \(C_{S_1}\) be any coarse-grained subsets of the sets of continuous tensor functions \(\{t_{\mu\nu}(x); x \in R^4, S_0 < x < S_1\}\), where the notation \(S_0 < x < S_1\) means that \(x\) lies in the future of some point in \(S_0\) and the past of some point in \(S_1\). Let

\[
\text{Prob}_S(C_{S_1})
\]

be the probability that the configuration \(\{t_{\mu\nu}(x); x \in R^4, S_0 < x < S_1\}\) belongs to \(C_{S_1}\), given our constructed probability density function on the set of possible functions \(T(x): S \to R\). Then, assuming that

\[
\text{Prob}_\infty(C_{S_1}) = \lim_{S \to \infty} (\text{Prob}_S(C_{S_1}))
\]
exists, we define this to be the probability that reality between \(S_0\) and \(S_1\) is described by time-evolving mass distributions belonging to \(C_{S_1}\). This completes this version of our proposed description of reality. The limiting values of \(t_{\mu\nu}(x)\) define a tensor field distribution on space–time, which defines the beables in our description. The limiting probability distribution defines the probability distribution on configurations of beables (here, real tensor functions) in space–time.

### 3. Toy model illustrations

To illustrate the implications of these rules, we consider a toy version of ‘semi-relativistic’ quantum theory, in which a non-relativistic system interacts with a small number of ‘photons’. We treat the photons as following lightlike path segments. We model their interactions with the system as bounces, which alter the trajectory of the photon. For simplicity, we neglect the effect of these interactions on the non-relativistic system, and also neglect its wave function spread and self-interaction, so that in isolation its Hamiltonian \(H_{\text{sys}} = 0\). We simplify further by working in one spatial dimension, and we take \(c = 1\).

**(a) Model 1**

We suppose that the initial state of the system is a superposition of two separate localized states, 
\[
\psi_{\text{sys}}^0 = a\psi_1^{\text{sys}} + b\psi_2^{\text{sys}}.
\]
Here, \(|a|^2 + |b|^2 = 1\) and the \(\psi_i^{\text{sys}}\) are states localized around the points \(x = x_i\), with \(x_2 > x_1\). For example, the \(\psi_i^{\text{sys}}\) could be taken to be Gaussians (but recall that we are neglecting changes in their width over time). We take \(|x_1 - x_2|\) to be large compared with the regions over which the wave functions are non-negligible. We thus have a crude model of a superposition of two well-separated beams, or of a macroscopic object in a superposition of two macroscopically separated states.

We suppose that the environment consists of a single photon, initially unentangled with the system. It is initially propagating rightwards from the direction \(x = -\infty\), so that in the absence of any interaction it would reach \(x = x_1\) at \(t = t_1\) and \(x = x_2\) at \(t = t_2 = t_1 + (x_2 - x_1)\).

We take the photon–system interaction to have the effect of instantaneously reversing the photon’s direction of travel, while leaving the system unaffected. (As noted above, we neglect the effect on the system: this violates conservation of momentum but simplifies the overall picture.)

Thus, for \(t < t_1\), the state of the photon-system combination in our model is
\[
\delta(X - x_1 - t + t_1)(a\psi_1^{\text{sys}}(Y) + b\psi_2^{\text{sys}}(Y)),
\]
where \(X, Y\) are the position coordinates for the photon and system, respectively. For \(t_1 < t < t_2\), the state is
\[
\delta(X - x_1 + t - t_1)(a\psi_1^{\text{sys}}(Y)) + \delta(X - x_1 - t + t_1)(b\psi_2^{\text{sys}}(Y)).
\]
For \(t > t_2\), the state is
\[
\delta(X - x_1 + t - t_1)(a\psi_1^{\text{sys}}(Y)) + \delta(X - x_2 + t - t_2)(b\psi_2^{\text{sys}}(Y)).
\]

The possible outcomes of a (fictitious) stress-energy measurement at a late time \(t = T \gg t_2\) are thus either finding the photon heading along the first ray \(X = x_1 + t_1 - t\) and the system localized in the support of \(\psi_1^{\text{sys}}\), or finding the photon heading along the second ray \(X = x_2 + t_2 - t\) and the system localized in the support of \(\psi_2^{\text{sys}}\).

Suppose, for example, we consider a real world defined by the first outcome. Our rules for constructing the system’s beables imply that, for \(t < 2t_1 - t_2\), and for \(x = x_1\) or \(x_2\), we condition on none of these outcomes, as all of them correspond to observations within the future light cone. Up to this time, then, the mass density beables for the system are distributed according to \(|\psi^{\text{sys}}(Y)|^2\), with a proportion \(|a|^2\) localized around \(Y = x_1\) and a proportion \(|b|^2\) localized around \(Y = x_2\).

For \(2t_1 - t_2 < t < t_1\), the observation of the photon on the first ray is outside the future light cone of the component of the system localized at \(x_2\), but not outside the future light cone of the
component localized at \( x_1 \). This gives us mass density beables distributing a proportion \(|a|^2\) of the total system mass around \( x_1 \), but zero mass density beables around \( x_2 \).

For \( t > t_1 \), the observation is outside the future light cone of both localized components of the system. This gives us mass density beables distributing the full system mass around \( x_1 \), and zero around \( x_2 \).

In other words, in the picture given by the beables, the system is a combination of two mass clouds with appropriate Born rule weights initially, and ‘collapses’ to a single cloud containing the full mass after \( t > t_1 \).

In the frame defined by our coordinates, this picture has the seemingly peculiar property that, for \( 2t_1 - t_2 < t < t_1 \), the object is, so to speak, only ‘partly present’, with only a proportion \(|a|^2\) of its mass represented by the beables. The point here is that the ‘realization’ of the object is determined in this case by the reflected light ray corresponding to the photon wave function

\[
\delta(X - x_1 + t - t_1).
\]

Along lines parallel to this light ray, the system collapses from a combination of two mass clouds to a single cloud of full mass, without any intervening interval in which ‘the mass is only partly present’. In general, we expect our rules to define ‘collapses’ in a way that will depend on the relationship between the frame considered and the possible interactions between the particles. However, with realistic interaction rules and initial states, it will not generally define ‘collapses’ to be instantaneous and need not necessarily associate the same frame with all collapses.

Note that, in this model, so long as \( T \gg t_2 \), its precise value makes no difference to the conclusion set out so far. To this extent, our asymptotic condition holds.

To complete the story in this model, we should characterize the beables associated with the photon. For \( t < t_1 \), this is unproblematic: the photon follows the incoming light ray \( \delta(X - x_1 - t + t_1) \). It also follows from our rules that, given a final detection of the photon along the outgoing light ray \( \delta(X - x_1 + t - t_1) \), the second possible outgoing light ray \( \delta(X - x_2 + t - t_2) \) is empty of beables, as the observation of the first light ray takes place outside the future light cone of points on the second. It would be tempting, then, to say that the beable description implies that the photon follows the first outgoing light ray. However, our rules imply a discontinuity precisely along this light ray. The only relevant measurement outcome (until \( t \) close to \( T \)) is the observation of the photon at the end of the light ray, and this is lightlike rather than spacelike separated from points along the ray. Strictly speaking, given that we consider the limit of spacelike surfaces approaching the ray, this observation never enters into the calculation, so that the beables describe only \(|a|^2\) of the photon energy-momentum following this ray. The fact that other final state observations (of the system, and of the absence of the photon from the second possible outgoing light ray) become spacelike separated for \( t \) close to \( T \) does not rescue this part of the picture, as we ultimately take the limit as \( T \to \infty \). We tentatively interpret this feature as an artefact, due to the parsimony of this model, as if many photons were interacting with the system, from both directions, we would expect some final measurement outcomes to be spacelike separated from all possible outgoing rays. This point is illustrated in our next model.

(b) Model 2

To test these intuitions a little further, we consider a symmetric two photon version of the above model. The initial system wave function is as before. Now, though, there are two incoming photons, one from the left, which arrives at \( x_1 \) at time \( t_1 \), and the other from the right, which arrives at \( x_2 \) at the same time \( t_1 \). As before, when interacting with the system, each photon bounces, instantaneously reversing its direction of travel, while the system is unaffected.

Thus, for \( t < t_1 \), the state of the photon–system combination is

\[
\delta(X_1 - x_1 - t + t_1)\delta(X_2 - x_2 + t - t_1) (a\psi_1^{\text{sys}}(Y) + b\psi_2^{\text{sys}}(Y)),
\]

where \( X_1, X_2 \) are the position coordinates for the two photons and \( Y \) is the coordinate of the system.
Suppose, again, that we find a final measurement outcome consistent with the two photons bouncing off the system localized at $x_1$. That is, at late time $T$, we find one photon along the intersection of $t = T$ with the ray $\delta(x_1 - x_2 + t - t_1)$, a second at the intersection of $t = T$ with the ray $\delta(x_2 - x_2 - t + t_2)$, and the system localized at $x_1$.

Now there are four relevant outgoing rays: the two at whose ends we find the two photons, the ray $\delta(x_1 - x_2 + t - t_2)$, at whose end we do not find photon 1, and the ray $\delta(x_2 - x_2 - t + t_1)$, at whose end we do not find photon 2. These last two rays correspond to the reflection of a photon from the system when localized at $x_2$.

These extra rays change the picture. For the system, our rules now give beables that distribute its mass density in the expected proportions $|a|^2$ and $|b|^2$ up to time $t = t_1 - (x_2 - x_1)$. At this point, the beables at both $x_1$ and $x_2$ ‘collapse’, and from then onwards all the beable mass density is localized at $x_1$, with none at $x_2$.

As for the photons, the beables corresponding to their incoming states show them along the incoming rays, as expected, until $t = t_1$. Because the measurement outcome corresponding to the absence of a photon along the outgoing ray $\delta(x_2 - x_2 - t + t_1)$ is outside the future light cone of the outgoing ray $\delta(x_1 - x_1 + t - t_1)$, and vice versa, the beables describe a full photon along the latter ray and none along the former. The outcomes for these two rays also ensure that the beables describe a full photon along the ray $\delta(x_2 - x_1 - t + t_2)$ and none along $\delta(x_1 - x_2 + t - t_2)$.

The entire beable picture for system and photons is thus consistent (if possibly initially a little challenging to the intuition). According to the beables, the system collapses to a mass density cloud localized around $x = x_1$ at time $t = t_1 - x_2 + x_1$. The photons follow definite paths, propagating until they hit this localized cloud, and then bouncing to reverse their courses.

The symmetry of this example ensures that this picture would be simply reflected if we obtained the other set of possible measurement outcomes at late time $T$, with the system beables now localizing around $x = x_2$ and the photons now bouncing off this localized cloud.

4. Discussion

We have presented a new proposal for a solution to the Lorentzian quantum reality problem. The proposal has several features in common with earlier suggestions [1,2]. But it has a significant new feature: beables at any point in space–time are defined using only final outcome data from outside that point’s future light cone. We are particularly interested in the beables associated with massive quasi-classical systems—that is, physical systems whose coarse-grained mass densities follow approximately classical equations of motion most of the time, with stochastic corrections arising (inter alia) from interactions with microscopic quantum systems [3]. The intuition here is that such a quasi-classical system generally leaves effective records of its location outside its future light cone, arising from its possible interactions with photons and other massless particles. This means that we could infer beables characterizing its quasi-classical behaviour at any given time from a (hypothetical) measurement at late times, even if we restrict attention to measurement data outside its future light cone, as these data include records arising from recent interactions with photons. Following this prescription allows us to calculate a beable picture by an algorithm that involves only calculating the outcome probabilities for various measurements in standard (unitary) relativistic quantum theory.

Our proposal was set in the context of Minkowski space–time, but can be simply extended to other background space–times that have well-defined asymptotic futures. We do not believe it is true in all cosmological models that quasi-classical systems leave effective records outside their future light cone. For example, models with a final ‘big crunch’ do not have this property. In this sense, our proposal is contingent on both the large-scale structure of space–time and the particular quantum dynamics defining the quantum evolution law within that space–time. However, while future cosmological observations that modify the presently preferred cosmological models could possibly change the picture, we do not presently see any compelling reason to believe that it is not viable in our own universe. At present, it seems that photons and (if they exist) gravitons and other massless particles interacting with quasi-classical structures in our cosmos.
can and apparently do escape to future infinity, without further interactions that significantly alter their direction or slow their progress, at an appreciable rate. Indeed, the observed accelerating expansion of the universe seems (if it continues indefinitely) to guarantee this will always be the case.

We have illustrated our proposal in toy models and shown that it gives sensible and intuitively appealing descriptions of quantum reality in these models. The models considered here are admittedly extremely simplified. They are set in one space and one time dimension. While the choice of one space dimension makes the geometry simple, it introduces an unrealistic feature. If we considered many massive systems localized in different places, and many incoming photons, and also maintained the simplified interaction rule that photons always bounce from massive systems, then photons would always remain in the regions to the left or right of all the massive systems, or else trapped between two systems. This is an artefact of the $1 + 1$ dimensional geometry and of the unrealistic assumption that interactions between photons and matter produce a bounce with probability one. In a $3 + 1$ dimensional geometry in which the angular distribution of massive objects has large gaps and in which photons are incoming from all directions (a more realistic model of the present state of our own cosmos), and with realistic interactions, photons scattering from a massive object can and often will escape to future infinity without further scatterings that significantly alter their space–time path.

Our toy models also neglect Schrödinger dynamics for massive particles or objects, setting $H = 0$ for these subsystems. They treat photons and particles as effectively pointlike, rather than allowing for some spread in their wavepackets, and assume a simplified interaction in which the photons undergo perfect instantaneous reflection without causing any reaction on the massive reflecting particles. We believe they nonetheless illustrate the key intuitions well and indicate that more realistic $3 + 1$ dimensional models that allow for spreading wavepackets and imperfect interaction with reaction would produce beables giving the same qualitative description of reality. Of course, demonstrations in more sophisticated models would be useful: we aim to produce such demonstrations in future works.

Our models also necessarily rely on a hybrid ‘semi-relativistic’ treatment of photons travelling at light speed interacting with particles or objects for which relativistic effects are negligible. This treatment can certainly be significantly improved, although it will be hard to provide a fundamentally correct and rigorous treatment, given the difficulties in modelling any realistic quasi-classical physics within quantum field theory as presently understood. Nonetheless, the consistency of quantum field theory with Minkowski causality, reflected in the fact that propagators fall off rapidly outside the future light cone and in the existence of the Tomonaga–Schwinger formalism, suggests to us that the strategy of defining beables at a point by conditioning on final outcomes outside its future light-cone is natural in quantum field theory. It suggests too that the qualitative features of the beables descriptions of reality in our toy models should accurately reflect the pictures that should emerge from a fully relativistic field-theoretic treatment.

Mass density beable ontologies were first proposed for non-relativistic collapse models by Pearle & Squires ([4]; see also [5]). An extension of these ontologies to relativistic collapse models using constructions previously defined in [6,7] was proposed in [8]. These proposals apply to generalizations of quantum theory rather than to quantum theory itself.

By contrast, the solution proposed here requires no alteration to quantum dynamics. Our models predict ‘collapses’ in the sense that the beables are initially consistent with a superposition of macroscopically distinct states, and later with a single component of the superposition. However, these ‘collapses’ have no direct effect on the underlying dynamics: rather, they are all inferred from the final outcome state. In particular, unlike non-relativistic or relativistic dynamical collapse models [8–11], our models predict no violation of conservation of energy.

Our proposal is also consistent with the standard understanding of Bell experiments. Without producing a detailed toy model here, it is easy to give the basic intuitions. First, it predicts that the beables associated with quasi-classical measuring devices in two spacelike separated wings of a Bell experiment will correspond to definite measurement outcomes after the experiment,
as the final state will include ‘photon’ states characterizing one of the possible outcomes on each wing. Second, it predicts that neither the beables in the entangled pair of particles being measured, nor those in the measuring devices, are correlated with the eventual outcomes before the measurements. No ‘photons’ significantly interact with the entangled particles before the measurements, and the ‘photon’ interactions with the measuring devices are the same whatever the eventual outcomes. Third, it predicts the same correlations among measurement outcomes as those predicted by standard quantum theory, as the dynamics are those of standard quantum theory, and its predictions for measurement outcomes are inferred from the final state records of those outcomes, whose probability distribution is obtained from standard quantum dynamics and the Born rule.

Our solution also seems to fit naturally within the existing framework for relativistic quantum field theory and to be consistent with standard intuitions about relativistic quantum field theories. Of course, making it mathematically rigorous nonetheless seems impossible without solving the problem of rigorously defining realistic relativistic quantum field theories in $3+1$-dimensional Minkowski space. Perhaps for the moment, we should be content to have a coherent conceptual account of reality defined at the same (low) level of rigour as our understanding of relativistic quantum field theory. In this context as elsewhere, the programme of rigorizing existing relativistic quantum field theories may not necessarily be the most scientifically fruitful way forward: other options include investigating quantum theories of gravity or other potentially deeper underlying theories.

The proposal presented here admits many variations, including the various interesting possibilities discussed in Kent [2] (adapted to allow for the new treatment here of outcome data, distinguishing data from inside and outside the future lightcone). Intriguingly, it also suggests natural ways to generalize quantum theory, following [2,12,13]. For, given a beable formulation—a description of quantum reality—consistent with standard quantum dynamics, we can define modified theories by altering the rules for assigning probabilities to beable configurations. For example, we can introduce weight factors that depend directly on the beable configuration, and use these to rescale the probabilities for beable configurations that are defined by standard quantum theory. Even more general possibilities are also available.

More extended comments and discussions of possibilities in these directions, which apply—and indeed, apply with more force—to the present proposal can be found in [2]. The fact that our beables take the form of a generalized stress-energy tensor also suggests a possible new path to considering the coupling of quantum theory and gravity, while avoiding the problems afflicting standard semi-classical gravity [2].

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