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Critique of Putnam's Quantum Logic

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Putnam gives a strongly realist account of quantum logic. This has been criticised as suggesting a hidden variable interpretation for quantum mechanics. Friedman and Glymour have done this in the framework of noncontextual hidden variable theories, which, however, does not fully represent Putnam's ideas. Here Putnam's approach to quantum logic is understood in terms of contextual truth-value assignments. The concept of a measurement is discussed. It follows that in order to reproduce quantum mechanical predictions a kind of disturbance is necessary, which is then analyzed. Finally, it is shown that the Putnam approach does not escape proofs of nonlocality, and thus shares, indeed, the unwelcome features of a hidden variable theory.

1. INTRODUCTION

In many of Hilary Putnam's writings on quantum logic, beginning with his well-known Putnam (1968), and continuing with Putnam (1974) and Friedman and Putnam (1978), a strong form of realism was present, in which definite values were ascribed simultaneously to all observables of an individual system.[2] This is well illustrated by the following [adapted from Putnam (1968, pp. 184–185) and Putnam (1974, pp. 52–53)]. Consider an n-dimensional Hilbert space $\mathscr{H}$. Associated with a maximal (nondegenerate) observable X is a decomposition of the Hilbert space with respect to atomic propositions (corresponding to the eigenvectors of X):

$$\mathscr{H} = x_1 \vee x_2 \vee \cdots \vee x_n$$

Under any interpretation, this is a true proposition. Its truth is understood by Putnam as meaning that the observable X has indeed a value corre-

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[2]Putnam (1968) was reprinted in Putnam (1975); Putnam (1974) in Putnam (1976).

sponding to one of the x_i. As the reasoning is independent of the particular choice of X, Putnam concludes that the system possesses values for all nondegenerate observables (Putnam, 1968) or, in fact, a state vector relative to each nondegenerate observable (Putnam, 1974). He then interprets measurements as simply revealing those preexisting values.

Strongly realist positions in the interpretation of quantum mechanics typically run into difficulties with no-hidden-variables theorems. However, Putnam has denied that his quantum logical interpretation is a hidden variable theory (Putnam, 1974, pp. 49–50; see further Putnam, 1968, pp. 188, 191, 197). And also Stairs (1983), while not thinking that 'value-definiteness' can be convincingly upheld (he argues for a more modest form of realism), states that, if successfully defended, the approach would give us 'all the benefits of a hidden variable account and none of the grief' (Stairs, 1983, pp. 584–585). It is this claim we wish to prove mistaken.

In order to develop at all an analogy between Putnam's ideas and the situation where one has quantum mechanics and an underlying hidden variable theory, we must be able to make a distinction between the *logic*, which is quantal, and an underlying *interpretation*, which in some sense is classical, as done in Friedman and Glymour (1972) and in Redhead (1987, Chapter 7). The interpretation corresponds then to some hidden ontology. Putnam's original claim was that quantum logic is 'the' true logic (Putnam, 1968, p. 184; 1974, p. 53), and this rules out the possibility of a classical metalogic. However, Redhead (1987, pp. 166–167) has argued that in this case any 'strong realism' in Putnam's position would be nothing but a play with words. Also, Putnam has revised his position at least since Putnam (1981), and accepts now the idea of a classical interpretation as legitimate (Putnam, 1994). Granted that we have such a framework, it is then possible to exploit the tension between the quantal nature of the logic and the classicity of the interpretation in terms of the standard no-hidden-variables theorems. One version of this program has been carried out by Friedman and Glymour (1972). They consider a number of possible truth valuations on the logic, and show that the Kochen and Specker (1967) theorem places restrictions on them which contradict Putnam's aims. However, their strategy applies only if one interprets Putnam as advocating a *noncontextual* assignment of truth values, while Putnam's suggestions seem rather to fit the pattern of a *contextual* hidden variable theory [for this distinction, see Shimony (1984)]. Accordingly, in this paper we develop a contextual version of Friedman and Glymour's program, concentrating on the treatment of measurements and the problem of disturbance. Eventually, we show that Stairs' own proof of nonlocality for (contextual) hidden variable theories (Stairs, 1983, pp. 579–581) applies also in the case of Putnam's quantum logic. These features undermine the realist and individualist appeal of the Putnam approach: it may have the advantages of a hidden variable theory, but it shares also the grief.

2. PUTNAM AND CONTEXTUALITY

In the following we adopt the formalism of partial Boolean algebras (PBA), which Putnam turns to in Friedman and Putnam (1978). Instead of giving up distributivity, in a PBA one has a primitive two-place relation called compatibility, the logical connectives being defined only for pairs of compatible propositions. The essential literature, especially the papers by Specker and by Kochen and Specker, is contained in Hooker (1975). In this terminology, what Friedman and Glymour do is to show that there are fairly well-behaved mappings from a typical quantum mechanical PBA onto the Boolean algebra $\{T, F\}$, but that these cannot be homomorphisms. In fact, the Kochen–Specker result shows that these PBAs always have even a finite subalgebra which does not admit PBA-homomorphisms onto $\{T, F\}$. In the spin-1 example (Friedman and Glymour, 1972, p. 27) it is easy to see that for any truth valuation of the kind Friedman and Glymour are considering there will be a direction x such that all three propositions $\sigma_x = 1$, $\sigma_x = 0$, and $\sigma_x = -1$ are assigned the truth value F. In Friedman and Glymour's words, we cannot 'maintain that in every quantum system every observable has a precise value' (p. 27). But now the idea of a measurement which is both revealing truth values and nondisturbing breaks down: if we do not want to say that the 'measurement' might be unfaithful, telling us that, say, $\sigma_x = 1$, while leaving the system as it was in an '*FFF*' state, then at least in some cases 'it must be admitted that measurement effects a change in what is true of [...] the quantity measured' (Friedman and Glymour, 1972, p. 27). This we shall refer to as a *strong disturbance* by the measurement.

Our question now is: can Putnam really be understood as implying the existence of such truths valuations? And our answer is: no, he cannot. In Putnam (1974, p. 52–53) he suggests that a quantum mechanical system possesses many different state vectors, in fact, a state vector for each nondegenerate observable. One finds similar statements in Friedman and Putnam (1978, p. 309) and Putnam (1981, p. 211). He further says that only one of these state vectors can be known to pertain to the system at any given time (Putnam, 1974, p. 53). And assigning a state vector to the system seems to depend on the choice of some perspective (Friedman and Putnam, 1978, p. 314; Putnam, 1981, p. 211) [see also the equivalence of 'perspectivalism' and 'quantum logic' in Putnam (1994)]. Notice that each nondegenerate observable is uniquely associated with an orthonormal basis in the Hilbert space. This basis, seen as a set of propositions, generates a maximal Boolean subalgebra B of the PBA, and a truth valuation on B can be defined by specifying that exactly one of these propositions is true, i.e., by specifying a (state) vector ψ_B in the corresponding basis. From this point of view, it appears natural to read Putnam as suggesting that the interpretation of the logic is given not by one single truth valuation defined on the whole of the PBA, but by a collection of partial truth valuations defined on maximal Boolean subalgebras of the logic.

Borrowing the terminology from Redhead (1987, p. 164), we define a *Putnam state* to be a pair (B, ψ_B) where B is a maximal Boolean subalgebra of the PBA of propositions, and ψ_B is a Boolean truth valuation, i.e., a PBA-homomorphism

$$\psi_B: \quad B \to \{T, F\}$$

onto the two-point (partial) Boolean algebra $\{T, F\}$. We now consider Putnam (1974) as suggesting that the interpretation of the logic (the ontology of the system) is defined by the collection of all Putnam states, one for each maximal Boolean subalgebra of the algebra of propositions. Every proposition b possesses a truth value relative to each (B, ψ_B) with $b \in B$. These values need not necessarily match across different Boolean subalgebras, and this enables Putnam to escape a Kochen–Specker contradiction. Further, every observable, being associated via its spectral resolution with (at least) one maximal Boolean subalgebra, has (at least) one precise value. The measurement of an observable can now be understood as revealing the value of the observable determined by a Putnam state. If there is more than one such value (which will in general be the case for a nonmaximal observable), the measurement will involve selection of a maximal Boolean subalgebra, but in any case it will not strongly disturb the system. (Since truth values of propositions and values of observables are closely linked, we shall often mix loosely talk about the two.)

3. MEASUREMENTS AND DISTURBANCE

In Putnam's understanding, a measurement simply reveals the value of an observable which is already possessed by the system. Measurements do not create values nor induce mysterious state transitions in the system. Further, all possible values for measurement results preexist any act of measurement, and the choice of performing one measurement rather than another is merely a matter of perspective (Putnam, 1968, pp. 186–187; 1974, pp. 49–51). In the framework of Putnam states, the natural description of a measurement is that it consists in the assignment of a Putnam state to the system. This fulfills Putnam's basic *desideratum* that when we want to measure the value of an observable, there is a preexisting value which is revealed without being disturbed. However, if there is no link at, all between different Putnam states (and the fact that there is no link is what allows this account to escape the criticism of Friedman and Glymour), it is not clear that this notion of measurement will truly represent a quantum mechanical measurement. It seems that we should get no information about the results of measurements of other, incompatible observables (or even partially compatible ones, see below). If our idea of a measurement is that it reveals definite truth values of propositions (which under any reading is what Putnam requires), is it possible that it should give information also about outcomes of different measurements which we know empirically to obey the statistical predictions of quantum mechanics? Putnam is aware of this, and he argues using Gleason's theorem that once given information about the truth values of certain propositions, the structure of the logic forces upon us the choice of the corresponding quantum mechanical states as giving us probabilities (interpreted as epistemic) on the rest of the algebra (Friedman and Putnam, 1978, p. 313). There is a detailed discussion of probabilities in the paper by Friedman and Putnam (1978), precisely in the sense of probabilities conditional on the results of previous measurements. They treat in particular the problem of defining probabilities conditional upon a nonmaximal measurement, i.e., the question of the Lüders rule, and they claim that they can reproduce all the statistical predictions of quantum mechanics in a very natural way. Or better still: their aim is to show that the basic (nonprobabilistic) concepts of the Putnam interpretation are *sufficient* to derive, in a sense as logical necessities, also the statistical predictions of quantum mechanics. We shall discuss these claims (especially the problem of the Lüders rule) in a separate paper. Here we wish to analyze instead what the *necessary* consequences are, at the level of the truth value assignments, of requiring that the quantum mechanical predictions be reproduced.

We start by reasoning analogously to Friedman and Glymour. Single measurements or measurements of compatible observables do not present any problems. But if one tries to account even for the simplest predictions of quantum mechanics concerning successive measurements of incompatible observables, one is forced to admit that a measurement cannot be faithful as well as nondisturbing. Consider for example measuring the spin of a particle first in the x direction, then in the y direction, and again along the x axis. Quantum mechanics tells us that the two spin-x measurements will not always yield the same result. In order to accommodate this in the Putnam account, one has to admit either that the last measurement is unfaithful, but then it is not revealing any preexisting value at all, or else that the measurement of the spin in the y direction, although not disturbing the value that is being measured (no strong disturbance), does, in fact, disturb the value of the spin along the x axis! We shall call this a *weak disturbance* of the system. In the concrete case of a spin-1 particle we have

$$(\sigma_x = 1) \vee (\sigma_x = 0) \vee (\sigma_x = -1) = (\sigma_y = 1) \vee (\sigma_y = 0) \vee (\sigma_y = -1) = \mathbf{1}$$

where **1** is the unit of the algebra; thus both propositions are trivially true. According to Putnam's ontology, the particle has both a spin-x and a spin-y state. Let us say the respective Putnam states assign the value T to $\sigma_x = 1$ and $\sigma_y = 0$. If we perform an x measurement, this will reveal the (ψ_{σ_x}-)truth of $\sigma_x = 1$. If we subsequently perform a y measurement, this will yield that $\sigma_y = 0$ is (ψ_{σ_y}-)true. But now, for a further measurement of spin in the x direction, quantum mechanics predicts the following outcomes:

$$\mathrm{Prob}(\sigma_x = 1) = \tfrac{1}{2}$$

$$\mathrm{Prob}(\sigma_x = 0) = 0$$

$$\mathrm{Prob}(\sigma_x = -1) = \tfrac{1}{2}$$

We infer that during the spin-y measurement, the particle has a 50% chance of retaining the value 1 of spin-x it had prior to the spin-y measurement, and a 50% chance of flipping to the value -1. In fact, we can calculate these probabilities by the usual Born rule for quantum mechanical statistics.

One may legitimately ask if Putnam has not just swept the problem (the 'mysterious physical disturbance') under the carpet. Not only do we have a disturbance of the system, but while the Putnam states obviously deterministically fix the outcomes of measurements, in general the associated weak disturbance seems to be stochastic and to follow the same pattern as the projection postulate! Putnam might have a reply here. Actually, he does accept that in general measurements will disturb the values of magnitudes other than those being measured (Putnam, 1968, p. 186). But with a 'judo-like manoeuvre' [adapting a phrase from Shimony (1984)], he refers to Gleason's theorem, and turns the fact that the truth valuations change during the spin-y measurement into a *logical necessity* (Putnam, 1968, p. 186). As Putnam states it, the physics has to comply with the true logic, which is *quantum* logic (Putnam, 1974, p. 53). Values of other Putnam states than the one associated with the measurement are disturbed, but they are then ready for any further measurement we may want to perform. From this point of view weak disturbance, far from being an *ad hoc* device for recovering quantum mechanical predictions, can be interpreted as a direct consequence of the structure of the logic.

4. KOCHEN–SPECKER AND WEAK DISTURBANCE

Even if that is so, a further discussion of how weak disturbance manifests itself will show that the price one has to pay in order to reproduce the quantum mechanical predictions is high. Indeed, now that the phenomenon of weak disturbance provides a link between different Putnam states, Kochen and Specker will côme back with a vengeance.

Take a Putnam state (B, ψ_B). It defines values for observables associated with the subalgebra B and for no other observables. In ordinary quantum mechanics, instead, when the quantum state is ψ_B (seen as a vector), one can predict with probability 1 the value of just any observable A of which ψ_B is an eigenstate, irrespective of whether A commutes with the maximal observables associated with the subalgebra B or not. If we consider a spin-1 particle and use the Kochen–Specker theorem, we can find five directions in space, say, x, y, z, y', z', such that (x, y, z) and (x, y', z') are orthogonal triads, and such that if A and B are the maximal Boolean subalgebras generated by the propositions $\sigma_x = 0$, $\sigma_y = 0$, $\sigma_z = 0$ and $\sigma_x = 0$, $\sigma_{y'} = 0$, $\sigma_{z'} = 0$, respectively, then the Putnam states (A, ψ_A) and (B, ψ_B) assign the following truth values to the propositions:

$$\psi_A: \quad \begin{cases} \sigma_x = 0 \mapsto T \\ \sigma_y = 0 \mapsto F \\ \sigma_z = 0 \mapsto F \end{cases} \qquad \text{and} \qquad \psi_B: \quad \begin{cases} \sigma_x = 0 \mapsto F \\ \sigma_{y'} = 0 \mapsto T \\ \sigma_{z'} = 0 \mapsto F \end{cases}$$

This now is a 'perverse' situation in which we have two quantum mechanical observables sharing an eigenvector (partial compatibility), and two Putnam states assigning contradictory truth values to the corresponding proposition! Now consider a measurement associated with A. This selects the Putnam state (A, ψ_A), and will reveal that $\sigma_x = 0$ is T. Similarly, a B-measurement, which selects (B, ψ_B), will yield that $\sigma_x = 0$ is F. The predictions of quantum mechanics are that after a measurement of a maximal observable associated with A, measurement of one associated with B will necessarily yield that $\sigma_x = 0$ is T. In order to explain this in the Putnam framework, weak disturbance has to be invoked. But this time we infer that measurement of (A, ψ_A) has disturbed the truth values assigned according to (B, ψ_B) not merely stochastically, but with *certainty*. In this case, the disturbance is weak, by definition, only due to the fact that it is not disturbing what one is looking at; but indeed it is stronger than the disturbance induced by a quantum mechanical measurement, because it induces a transition from a homomorphism defined by one vector to a homomorphism corresponding to a vector *orthogonal* to the first one. And this is the case because the Putnam-state formalism is, in fact, more contextual than orthodox quantum mechanics.

5. VARIABLES UNDER THE CARPET

The Putnam interpretation provides contextual value assignments for all observables of a quantum system. These value assignments determine the results of measurements, and it is claimed that they are always distributed over an ensemble of systems such as to recover the quantum mechanical predictions. We thus have a theory which is at least formally equivalent to a deterministic contextual hidden variable theory. We cautiously say the equivalence is formal, because Putnam emphasizes that his approach is not in the spirit of a hidden variable interpretation (Putnam, 1974, pp. 49–50). As a matter of fact, hidden variables are physical variables which, if known, would permit a (more) complete description of the system than the one given by quantum mechanics. Putnam instead insists that thinking of giving any more complete representation of a quantum system is logically *inconsistent*. Nevertheless, even formal analogies with hidden variable theories are enough to derive unwelcome features. In particular, we expect the Putnam approach to be also more nonlocal than orthodox quantum mechanics. In fact, as a deterministic hidden variable theory, it should exhibit *parameter* dependence in an EPR setup, instead of *outcome* dependence (see Shimony, 1986, pp. 188ff, 191ff). We shall not spell this out in terms of Bell inequalities, but of an 'algebraic' proof of nonlocality, in the tradition inaugurated by Heywood and Redhead (1983) and by Stairs (1983), i.e., of a result involving a nonprobabilistic contradition with quantum mechanics if one postulates a suitable locality condition. Our proof is in fact the same as Stairs', but applied to the Putnam-state formalism.

In order to derive nonlocality, we use again weak disturbance, adapting the Kochen–Specker example to the case of a 2-particle system consisting of two spin-1 particles in the singlet state ψ_{sing}. The state can always be written as

$$\psi_{\text{sing}} = \frac{1}{\sqrt{3}}\left(|\sigma_x^1 = 1\rangle|\sigma_x^2 = -1\rangle - |\sigma_x^1 = 0\rangle|\sigma_x^2 = 0\rangle + |\sigma_x^1 = -1\rangle|\sigma_x^2 = 1\rangle\right)$$

irrespective of the direction x, and it is characterized by the strict anticorrelation of measured spin values for the two particles along any axis x. In this composite system, maximal Boolean subalgebras will be generated by sets of propositions of the form

$$\{\sigma_u^1 = 0 \wedge \sigma_{u'}^2 = 0 \mid u \in \{x, y, z\}, u' \in \{x', y', z'\}\}$$

with $\{\sigma_x^1 = 0, \sigma_y^1 = 0, \sigma_z^1 = 0\}$ and $\{\sigma_{x'}^2 = 0, \sigma_{y'}^2 = 0, \sigma_{z'}^2 = 0\}$, respectively, generating maximal Boolean subalgebras of the PBA of propositions for particles 1 and 2. Call these maximal Boolean subalgebras respectively A_1

and B_2, and the subalgebra of the composite system (which has a tensor product structure) $A_1 \otimes B_2$. The Putnam state $(A_1 \otimes B_2, \psi_{A_1 \otimes B_2})$ associated with $A_1 \otimes B_2$ of course defines truth values for the propositions in A_1 and the propositions in B_2. Indeed, we can identify, e.g., $\sigma^1_x = 0$ with

$$\sigma^1_x = 0 \wedge (\sigma^2_{x'} = 0 \vee \sigma^2_{y'} = 0 \vee \sigma^2_{z'} = 0)$$

which is a proposition of the combined system, and so on. In other words, for any A_1 and B_2,

$$(A_1, \psi_{A_1 \otimes B_2}|_{A_1}) \qquad \text{and} \qquad (B_2, \psi_{A_1 \otimes B_2}|_{B_2})$$

are well-defined Putnam states on the two particles, respectively. Given a particular subalgebra of one of the subsystems (say A_1), in general different Putnam states $(A_1 \otimes B_2, \psi_{A_1 \otimes B_2})$ and $(A_1 \otimes B'_2, \psi_{A_1 \otimes B'_2})$ of the combined system will not define a unique Putnam state for particle 1. This is the case because the two homomorphisms $\psi_{A_1 \otimes B_2}$ and $\psi_{A_1 \otimes B'_2}$ need not *a priori* match on A_1. In the following, however, we shall assume that for any A_1, B_2, and B'_2 as above,

$$\psi_{A_1 \otimes B_2}|_{A_1} = \psi_{A_1 \otimes B'_2}|_{A_1}$$

(and similarly with particles 1 and 2 interchanged). That is, we assume that for each subsystem (say 1) the Putnam states of the combined system define 'reduced' Putnam states of the form (A_1, ψ_{A_1}), where

$$\psi_{A_1} := \psi_{A_1 \otimes B_2}|_{A_1}$$

independently of the choice of B_2. This corresponds to Heywood and Redhead's (1983, p. 487) condition of ontological locality (OLOC). It is an essential assumption for the following argument, because it guarantees that one can assign well-defined truth values to propositions referring only to one particle, and thus perform genuinely *local* measurements on either of the two subsystems. Consider again the Putnam states (A, ψ_A) and (B, ψ_B) of the previous section, or better (A_1, ψ_{A_1}) and (B_1, ψ_{B_1}) referring to particle 1, and (B_2, ψ_{B_2}) accordingly for particle 2. We have

$$\psi_{A_1}: \quad \begin{cases} \sigma^1_x = 0 \mapsto T \\ \sigma^1_y = 0 \mapsto F \\ \sigma^1_z = 0 \mapsto F \end{cases} \qquad \text{and} \qquad \psi_{B_1}: \quad \begin{cases} \sigma^1_{x} = 0 \mapsto F \\ \sigma^1_{y'} = 0 \mapsto T \\ \sigma^1_{z'} = 0 \mapsto F \end{cases}$$

while (B_2, ψ_{B_2}) makes truth-value assignments for $\sigma^2_x = 0$, $\sigma^2_{y'} = 0$, and $\sigma^2_{z'} = 0$. In what follows we shall be forced to conclude that the disturbance we have described earlier now acts nonlocally.

If we perform a B_1-measurement, we reveal that $\sigma^1_{y'} = 0$ is T and $\sigma^1_x = 0$ is F. In order to respect the predictions of quantum mechanics, we have to say that a subsequently performed B_2- measurement will yield that $\sigma_{y'}^2 = 0$ is T and so that $\sigma_x^2 = 0$ is F. If we start with an A_1-measurement on particle 1 instead, we shall obtain the value T for $\sigma_x^1 = 0$. Because the system is in the singlet state, we conclude that a later B_2-measurement will now necessarily yield that $\sigma_x^2 = 0$ is T: what in the case of one particle was simply context dependence of values (truth value of $\sigma_x^1 = 0$ dependent on measurement of A_1 or B_1) has now become *parameter dependence* of the truth value of $\sigma_x^2 = 0$ on whether a measurement of A_1 or B_1 is performed on particle 1. If we assume, as we said before, that the Putnam state (B_2, ψ_{B_2}) is well defined, we have to admit that either the measurement of B_1 or the measurement of A_1 *nonlocally* disturbed the truth-value assignments of ψ_{B_2}. The only alternative we have is to deny that reduced Putnam states as above are well-defined. This will rid us of the disturbance at a distance, but it explains parameter dependence by introducing a full-blown nonlocal contextualism, in which truth values of propositions are in general only defined relative to the maximal Boolean subalgebras of the *global* PBA of propositions (referring to the composite system). Our conclusions are perfectly analogous to the ones of Heywood and Redhead. As a matter of fact, the Putnam approach is equivalent to the one they have in mind ['to de-Ockhamize QM *à la* van Fraassen,' (Heywood and Redhead, 1983, p. 485)]. Our assumption that local Putnam states are well-defined corresponds exactly to their ontological locality, and the nonlocal disturbance corresponds in Heywood and Redhead's (1983, p. 488) terminology to a violation of environmental locality (ELOC). Thus, a Heywood–Redhead–Stairs proof of nonlocality can be used to strengthen Stairs' own argument in Stairs (1983) against the strongly realist position in quantum logic.

6. CONCLUSIONS

A *shadow cabinet* theory: that would be a most fitting description for the Putnam interpretation. The results of any measurement we may choose to perform are already there, just as a shadow cabinet is ready for whenever it is voted in. But while politics at least can (perhaps) be described by classical logic, quantum mechanical systems cannot. This has the consequence that such an account of quantum logic has all distinctive properties of a hidden variable theory, in particular nonlocal effects at the level of the 'hidden variables.' On the other hand, the Putnam approach also has some quite appealing features: indeed, it provides a realist and individualist account of quantum systems. However, in our opinion, one of the major *desiderata* for a realist, individualist interpretation is that such a scheme may have additional explanatory power as compared with the Copenhagen interpretation, and in particular that it may open up the possibility for measurements to be analyzed as physical processes. We do not feel, however, that Putnam achieves this. While he justly criticises Copenhagen for having 'erected into an article of faith in the state of Denmark that there can be no [. . .] theory [of measurement]' (Putnam, 1968, p. 183), if in Putnam's account the disturbance induced by measurements is a logical necessity, then it is just as unanalyzable as Copenhagen's. One is left with a vaguely uncomfortable feeling that Putnam dissolves the problems by *fiat*, and that he might perhaps be proving too much.

If the Putnam interpretation does not seem satisfactory, this does not mean that quantum logic forbids an ontologically realist picture altogether. The example of Jauch and Piron (1969) and Stairs' (1983) arguments show that one can take the logic seriously as reflecting the ontology of the system, thus eliminating the tension between logic and ontology: propositions compatible with the state vector can be given truth values, while the others are considered merely as potential. Indeed, it is worth keeping in mind that there are alternative possibilities for trying to implement a realist, individualist program, and that Putnam's quantum logical interpretation is but one interpretation of quantum logic.

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