Teller, P. (1991): "Substance, Relations, and Arguments About the Nature of Space–Time," *Philosophical Review*, 100, 363–97.

Weingard, R. (1979): "Some Philosophical Aspects of Black Holes," *Synthese*, 42, 191–219.

Zangari, M. (1994): "A New Twist in the Conventionality of Simultaneity Debate," *Philosophy of Science*, 61, 267–75.

Chapter 10

Interpreting Quantum Theories

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Introduction: Interpretation

The foundational investigation of quantum theories is inevitably specialized, but it ought not be exclusively so. Continuities of theme and approach should link it to the foundational investigation of other physical theories, to the general philosophy of science, to metaphysics and epistemology more broadly construed. The interpretation of quantum theories is the furnace in which these links are forged. To interpret a physical theory is to say what the world would be like, if the theory were true. A realist about a theory believes that theory to be true. Interpretation gives the realist's belief content, tells the constructive empiricist what he does ***not*** believe, and makes available to all parties the understanding of a theory constituted by a grasp of its truth conditions. Interpretation can promote theory development: Howard Stein offers the example of interpretive questions about the ether's state of motion in Maxwell theory, questions whose answers "revolutionized the theory and deepened our understanding of nature very considerably" (Stein, 1972, p. 423).

Having issued this apology for interpretation, this chapter surveys the interpretation of quantum theories. It chronicles past highlights (pp. 200–9); covers current work (pp. 209–17); and presents future directions (pp. 217–21). The remainder of this section sets the stage.

The Heisenberg–Born–Jordan matrix mechanics and Schrödinger's wave mechanics were twin formulations of quantum theory so fraternal it took von Neumann to pinpoint their relation. He called the structure they shared "Hilbert Space."[1] Pure quantum states are normed Hilbert space vectors; quantum observables are self-adjoint Hilbert space operators; once the Hamiltonian operator $\hat{H}$ is provided, Schrödinger's equation determines dynamical trajectories through state space. Because classical observables are functions from state space elements to the reals, a system's classical state fixes the values of all classical observables pertaining to it. In quantum mechanics (QM), this is not so. A state $|\psi\rangle$ does not in general fix the value of an observable $\hat{A}$; rather $|\psi\rangle$, via the Born Rule, determines a probability distribution over $\hat{A}$'s possible values. In its standard Hilbert space formulation, QM lacks what I'll call a *semantics*, an account of which observables have determinate values on a quantum system, and of what those values are or might be.

The quartet {state space, observables, dynamics, semantics} characterizes what is or can be true for a theory over time, and so constitutes an interpretation of a theory. Correlatively, degrees of freedom available to those engaged in interpretive projects include freedoms to propose and modify members of the quartet. One way to see the venerable debate about the nature of space(time) is as a debate about how best to tune classical theory's state space, observable set, and dynamics to one another. But interpretations can be – many interpretations of QM are – efforts in creative physics.

Bohr and Complementarity

Interpretive efforts can also retard creative physics, as Einstein feared Bohr's philosophy of complementarity would. The doctrine is too intricate to explicate in a short space, so I will settle for listing a few of its key elements, some of which persist in influence, by seeming to those working at present either to deserve explication, or to constitute exculpation.

Bohr denies that position and momentum can be simultaneously determinate on a quantum system. Position and momentum are what he styles ***complementary*** modes of description, "complementary in the sense that any given application of classical concepts precludes the simultaneous use of other classical concepts which in a different guise are equally necessary for the elucidation of the phenomenon" (1934, p. 10). For Bohr, it is as though reality were a stereoscopic image we were constrained to view one eye at a time. The doctrine originates in an insistence on the use of classical concepts, which Bohr couples to an operationalism governing their use. He observes that the experimental circumstances warranting the use of the momentum concept are incompatible with those warranting the use of the position concept. The upshot is the complementarity of position and momentum concepts, representatives of the complementary classes of kinematic (that is, spatio-temporal) and dynamic (that is, subject to conservation laws) concepts. Bohr takes the position-momentum uncertainty relations to express – and be explained by – this deeper principle of complementarity (1934, p. 57); see also Murdoch (1987, ch. 3).[2]

Bohr repeatedly emphasizes that the quantum of action is central to the doctrine. But the quantum of action seems to have gone missing from the foregoing reconstruction. One place it might lurk is a loophole through which a sort of counterfactual discourse might sneak. Having bolted our diaphragm to the table, we may

with Bohr's blessing speak of the position of an electron passing through our experimental arrangement. Could we also, and in defiance of complementarity, speak of its momentum, by appeal to experimental results we ***would have*** obtained, ***had we*** instead dangled our diaphragm from a spring balance? Not if the ***uncontrollable*** exchange of the quantum of action blocks such extrapolation. A disturbance theory of measurement fertilizes yet another root of complementarity.

Consider how Bohr's philosophy could interact with living physics. A physics community embracing the philosophy of complementarity would thereby abandon the project, declared inconceivable by the doctrine of complementarity, of "completing" QM by developing a theory which described the simultaneously determinate positions and momenta of systems. Einstein (Fine, 1986, p. 18) feared that "the Heisenberg-Bohr tranquilizing philosophy – or is it religion? – is so delicately contrived that, for the time being, it provides a gentle pillow for the true believer from which he cannot very easily be aroused." In 1935, with Podolsky and Rosen, he issued a wakeup call.

The Einstein-Podolsky-Rosen (EPR) Argument

Bohr denies that complementary magnitudes are simultaneously determinate. Einstein, Podolsky and Rosen (1935) argue that quantum statistics themselves imply that Bohr is wrong. Crucial to their argument is the "criterion of reality":

> If without in any way disturbing a system we can predict with certainty . . . the value of a physical quantity, then there exists an element of physical reality corresponding to this physical quantity (Einstein et al., 1935, p. 777).

(I take the consequent to be equivalent to "this physical quantity has a determinate value.") They argue that there are circumstances in which ***complementary*** observables satisfy the reality criterion. Their key move is to consider quantum states of ***composite*** systems instituting ***correlations*** between observables pertaining to component subsystems. Bohm (1951) reformulates the argument for a pair of electrons in the spin singlet state:

$$|\Psi\rangle_{\text{singlet}} = \frac{1}{\sqrt{2}}(|\rightarrow\rangle_I |\leftarrow\rangle_{II} - |\leftarrow\rangle_I |\rightarrow\rangle_{II}) \qquad (10.1)$$

Although (10.1) expresses $|\psi\rangle_{\text{singlet}}$ in terms of eigenstates $|\rightarrow\rangle$ and $|\leftarrow\rangle$ of the x-component of spin $\hat{\sigma}_x$, $|\psi\rangle_{\text{singlet}}$ assumes biorthogonal form for, and institutes perfect correlations between, $\hat{\sigma}_n$ eigenstates of the two systems for all n. Thus, $|\psi\rangle_{\text{singlet}}$ assigns Born rule probability 1 to the experimental result that the outcomes of $\hat{\sigma}_n$ measurements on systems I and II disagree. EPR consider a pair of electrons prepared in $|\psi\rangle_{\text{singlet}}$ and sent to laboratories remote from one another.

Measuring $\hat{\sigma}_x$ on system I affords the prediction, with certainty, that an $\hat{\sigma}_x$ measurement performed on system II will yield the opposite result. The remoteness of the laboratories ensures that a measurement on system I cannot in any way disturb system II – provided the universe is "local" in a way that renders distance an assurance of isolation. By the reality criterion, then, $\hat{\sigma}_x$ on system II is an element of reality. EPR could well have stopped here (Fine, 1986, ch. 3) musters evidence that Einstein wishes they had). They have shown that, for those who would withhold determinateness from the quantum realm, it is as though the spin measurement in first laboratory, instantaneously and at a distance, brings into being an element of reality in the second laboratory.

But EPR continue. We might rather have measured $\hat{\sigma}_y$ on system I. $|\psi\rangle_{singlet}$ anticorrelates $\hat{\sigma}_y$ eigenstates just as well as it anticorrelates $\hat{\sigma}_x$ eigenstates. By parity of reasoning, in this counterfactual situation, $\hat{\sigma}_y$ on system II would be an element of reality. EPR again appeal to locality to conclude from this that $\hat{\sigma}_y$ on system II *is* an element of reality – otherwise "the reality of $[\hat{\sigma}_x]$ and $[\hat{\sigma}_y]$ depend on the process of measurement carried out on the first system, which does not disturb the second system in any way. No reasonable definition of reality could be expected to permit this" (Einstein et al., 1935, p. 780). (Bohr's reply to EPR is to permit what they deem impermissible: the nonlocal dependence of system II's matters of fact on system I manipulations. There is "no question of a mechanical disturbance," Bohr writes, but there is "the question of an influence on the very conditions which define the possible types of predictions regarding the future behavior of the system" (Bohr, 1935, p. 699).) Because the correlations $|\psi\rangle_{singlet}$ institutes are thoroughgoing, if the EPR argument works, it works for every spin observable. Those convinced by the argument should undertake the project of "completing" QM, for instance, by devising a theory which attributes a determinate value to every element of reality established by the EPR gambit, a theory which moreover respects a "locality" requirement of the sort EPR exploit. (Those convinced *ab initio* that the project of completing QM is worth undertaking needn't be constrained by locality, or by reconstituting reality EPR element by EPR element.) One of John Bell's groundbreaking contributions to the foundations of QM was to bring "local" hidden variable theories (HVTs) in contact with empirical data.

Bell's Theorem and Other No-Go Results

Bell's theorem shows that local HVTs are committed to sets of statistical predictions known as Bell Inequalities. Insofar as there exist quantum states predicting the violation of the Inequalities, Bell's theorem sets up a crucial test of local HVTs vs. standard QM. Experiment upholds QM, violates the Inequalities, and falsifies local HVTs. The field is set for the game of experimental metaphysics. To play, show how to derive Bell Inequalities from a set of premises bearing philosophically

fraught names ("determinism," "completeness," "locality"). Observe that the experimental violation of the Inequalities reveals at least one of these premises to be false. Invoking priors of various sorts, single out leading suspects. The literature is vast; see Cushing and McMullin (1989) for a sample. In this section, I'll review a few of its defining moments, express a concern that locality is a red herring, and touch upon questions the violation of the Bell Inequalities raises about the nature of explanation.

The Bell inequalities

Like the EPR argument, Bell's (1964) theorem concerns distant correlations established by $|\psi\rangle_{singlet}$. In Bell's version of the experimental setup, the distant devices need not measure the same component of spin. Thus the generic outcome of a Bell correlation measurement is $(x,y|a,b)$ where $x, y \in \{+,-\}$ are the outcomes of measurements of spin components $\hat{\sigma}_a,\hat{\sigma}_b$ on particles I and II respectively. The Born rule probability $|\psi\rangle_{singlet}$ assigns $(x,y|a,b)$ is $\frac{1}{2}sin^2\theta ab/2$, where θ_{ab} is the angle between orientations a and b. Consider how a HVT might handle such probabilities. Let λ denote a complete set of parameters by which such a theory characterizes the state of a physical system; let Λ denote the full set of such states. Let $Pr_\lambda(x,y|a,b)$ be the probability the hidden state λ assigns the experimental result $(x,y|a,b)$. So-called deterministic HVTs countenance only probabilities of 1 or 0; stochastic HVTs countenance non-trivial probabilities. A quantum system has a hidden state $\lambda \in \Lambda$, we know not which; a normalized probability density $\rho(\lambda)$ over Λ encodes our ignorance. To obtain the empirical probability for a Bell-type measurement outcome, a HVT integrates, over the set Λ, the probabilities each λ assigns this outcome, weighted by the density $\rho(\lambda)$:

$$Pr(x,y|a,b) = \int_\Lambda Pr_\lambda(x,y|a,b)\rho(\lambda)d\lambda \tag{10.2}$$

To derive the Bell Inequalities, one imposes additional constraints on the HVT's probability assignment. Appealing broadly to intuitions about locality, Bell required the joint probability to *factorize* into probabilities for outcomes on each wing, which probabilities conditionalize only on settings proper to that wing:

$$Pr_\lambda(x,y|a,b) = Pr_\lambda(x|a) \times Pr_\lambda(y|b) \tag{10.3}$$

HVTs obedient to the factorization condition (3) obey the Inequality[3]

$$-1 \leq Pr(+,+|a,b) + Pr(+,+|a,b') + Pr(+,+|a',b') - Pr(+,+|a',b) \\ - Pr(+|a) - Pr(+|b) \leq 0 \tag{10.4}$$

This is a Bell Inequality. There are quadruples of orientations (a,a',b,b') – for instance $(\pi/3,\pi,0,2\pi/3)$ – for which standard QM predicts its violation. Upholding standard QM, experiment falsifies local HVTs.

For the purposes of probing locality, the factorization condition is blunt. In his 1983 dissertation – Jarrett (1986) provides a précis – John Jarrett sharpened it, by demonstrating its equivalence to the pair of conditions:

$$\Pr_\lambda(x|a,b) = \Pr_\lambda(x|a) \qquad \text{(Jarrett Locality)}$$

$$\Pr_\lambda(x|a,b,y) = \Pr_\lambda(x|a,b) \qquad \text{(Jarrett Completeness)}$$

The first expresses the desideratum that the outcome of a particle I measurement be independent of detector II's setting (ergo Shimony's (1984b) label: "parameter independence"). Jarrett equates it to a prohibition on superluminal signaling, which prohibition he supposes the special theory of relativity (STR) to issue. If Jarrett Locality fails, by changing the setting of her detector, a physicist in laboratory I can send instantaneously to laboratory II a signal in the form of altered measurement statistics (Shimony calls this "controllable non-locality" or action-at-a-distance). Jarrett Completeness expresses the desideratum that the outcome of a particle I measurement be independent of the outcome of a particle II measurement (ergo Shimony's label: "outcome independence"). Because the laboratory I physicist has no control over laboratory II outcomes, she can not exploit breakdowns in Jarrett locality to signal (Shimony call this "uncontrollable non-locality" or "passion-at-a-distance").

The violation of the Bell Inequalities implies that one of the assumptions generating them must be false. Having furnished his factorization of (10.3), and supposing our commitment to the special theory of relativity theory to commit us, at least morally, to Jarrett Locality, Jarrett fingers Completeness as the culprit (1986, p. 27). Setting $\lambda = |\psi\rangle$, standard quantum mechanics itself can be cast as a stochastic hidden variable theory violating completeness: $|\psi\rangle_{\text{singlet}}$ makes particle I probabilities sensitive to particle II outcomes. It appears that the quantum domain is ruled by passion-at-a-distance. Enlisting a Lewis-style counterfactual analysis of causation, Butterfield (1992) has argued that this violation of Jarrett completeness signals a causal connection between distant wings of the apparatus. Much care has been lavished on articulating relativistic locality constraints suited to this stochastic setting, so that the question of whether QM and the STR can "peacefully coexist" (Redhead, 1983) might be settled once and for all.

No-Go results without locality

I would advocate postponing the question. STR does not issue bans on superluminal causation. It does not address causation at all. It rather requires of that class of space–time theories formulated in Minkowski space–time that they be Lorentz-covariant.[4] Non-relativistic QM, which is not a space–time theory, is not subject to STR's requirements. So the question of whether STR and QM can peacefully coexist is ill-posed. Another question – can there be Lorentz-covariant quantum theories? – is well-posed. Quantum field theory (QFT) associates observables

$\hat{A}(\mathcal{D})$ with regions of space–time $\mathcal{D}$. The inhomogeneous Lorentz group[5] Λ is represented on the Hilbert space which is the common domain of these observables by a group of unitary operators $\hat{U}(\Lambda)$. QFT so formulated is Lorentz covariant iff the observables associated with the Lorentz transform $\Lambda\mathcal{D}$ of a region $\mathcal{D}$ is the corresponding unitary transform of the observables associated with $\mathcal{D}$:

$$\hat{A}(\Lambda\mathcal{D}) = \hat{U}(\Lambda)\hat{A}(\mathcal{D}) \qquad \text{(LC-QFT)}$$

That there are QFTs satisfying (LC-QFT) should settle the peaceful coexistence question. In the QFT context, bans on superluminal signal propagation are expressed by the microcausality requirement that operators associated with space-like separated regions commute (intuitively, it does not matter what order they act in). That this microcausality requirement is independent of the requirement of Lorentz covariance suggests that the folkloric connection between STR and the prohibition on superluminal signal propagation is only that.

Bell's theorem may be profitably analyzed without recourse to locality notions tenuously linked to STR. Fine (1982a,b) showed the (Clauser–Horne form of) the Bell Inequalities to be equivalent to

1 the existence of a deterministic HVT
2 the existence of joint distributions for all pairs and triples of observables
3 the existence of a stochastic HVT satisfying (10.3).

Intuitions about locality might motivate (3), but they are not directly implicated in either (1) or (2), which simply offer ambitious patterns of determinate value assignment. Indeed, a family of arguments originating with Bell (1966 – which he wrote *before* the 1964 Bell Inequalities paper) but refined by Kochen and Specker reveals that the project of assigning determinate values to sufficiently rich sets of observables is untenable, if the value assignment is subject to prima facie reasonable constraints.

Here's an informal sketch of Bell's version of the No-Go result; see Redhead (1987, ch. 5) for more details and references. Consider a project of determinate value assignment satisfying

(Spectrum) $\hat{O}$'s determinate value $[\hat{O}]$ is one of its eigenvalues

and

(FUNC) If $\hat{A} = f(\hat{B})$, then $[\hat{A}] = f([\hat{B}])$

In a Hilbert space of dimension three, any trio $\{\hat{P}_i\}$ of mutually orthogonal projection operators furnishes a resolution of the identity operator $\hat{I}$:

$$\hat{I} = \hat{P}_1 + \hat{P}_2 + \hat{P}_3 \qquad (10.5)$$

By the Spectrum rule $[\hat{I}] = 1$ and $[\hat{P}_i] \in \{0,1\}$. The $\{\hat{P}_i\}$ commute pairwise; there is therefore an operator of which each of them is a function. So the FUNC rule requires

$$[\hat{I}] = [\hat{P}_1] + [\hat{P}_2] + [\hat{P}_3] \tag{10.6}$$

Equations (10.5) and (10.6) together imply that for any trio of mutually orthogonal projectors, one of them will be assigned the value 1 while the other two will be assigned the value 0. This assignment induces a linear, normalized map from the set of projection operators on Hilbert space to the interval $[0,1]$ – indeed to the set $\{0,1\}$ containing only the endpoints of that interval. This map is also a probability measure over the closed subspaces of Hilbert space. According to Gleason's theorem, for Hilbert spaces of dimension three or greater, all such probability measures are continuous. But the map induced by the project of complete determinate value assignment is discontinuous – intuitively, as it sweeps through the set of projectors, it is going to have to leap from a projector it maps to 0 to a projector it maps to 1, without assigning intermediate projectors intermediate values. A HVT inducing such a map from Hilbert space operators to their determinate values is therefore inconsistent.

Bell needs infinitely many observables – the full set of projection operators on a three-dimensional Hilbert space – to generate the contradiction. Kochen and Specker showed that 117 projectors on a four-dimensional Hilbert space could not without contradiction be assigned determinate values obedient to the FUNC and Spectrum rules; Bell–Kochen–Specker type contradictions for ever smaller sets of observables have been emerging ever since.[6] Mermin's excellent presentation of Bell–Kochen–Specker results (1993) situates one version of the Bell Inequalities among them. The No-Go argument just sketched attributes $\hat{P}_1$ the same determinate value whether it's considered an element of the orthogonal triple $T = \{\hat{P}_1, \hat{P}_2, \hat{P}_3\}$ or an element of the *different* orthogonal triple $T' = \{\hat{P}_1, \hat{P}'_2, \hat{P}'_3\}$. It assigns a non-maximal observable a *non-contextual* value, that is, one not relativized to a particular eigenbasis of the observable. (The question of contextualizing does not arise for maximal observables, whose eigenbases are unique.) Contextualizing determinate value assignments, one can avert No-Go results, by without contradiction assigning $\hat{P}_1$ in the context of the basis T a value different from the one it's assigned in the context of the basis T'.

While such a move might seem shamefully *ad hoc*, it is precisely the move Bell makes after presenting his version of the No-Go result. The argument, he writes, "tacitly assumed that the measurement of an observable must yield the same value independently of what other measurements must be made simultaneously" (Bell, 1966, p. 451). To see what this vaguely Bohrian pronouncement has to do with contextualism, and to anticipate its connection with the Bell inequalities, consider the dramatically non-maximal composite system observable $\hat{I} \otimes \hat{\sigma}_x$. One way to select an eigenbasis from the myriad available for this observable is to specify a spin observable for particle one: for instance $\hat{\sigma}_x \otimes \hat{\sigma}_x$ has a unique eigenbasis which

is also an eigenbasis for $\hat{I} \otimes \hat{\sigma}_x$. To attribute $\hat{I} \otimes \hat{\sigma}_x$ a non-contextual value admitting faithful measurement is to assume that a $\hat{I} \otimes \hat{\sigma}_x$ measurement has the same outcome regardless of which particle I measurement is made. Seeing no reason to suppose that measurement outcomes are in general insensitive to measuring environments, Bell rejects the non-contextuality requirement.

Whether this "judo-like maneuver" (Shimony, 1984a) of invoking Bohr to protect ambitious plans of value assignment succeeds or not, it suggests a connection between Bell–Kochen–Specker arguments and the Bell Inequalities. The locality assumptions invoked in deriving the Inequalities are a species of a non-contextuality requirement. Mermin (1993) shows how to use locality-as-non-contextuality to convert an eight-dimensional Bell–Kochen–Specker result into one version of the Bell Inequalities. (What is lost in the translation is the state independence of the Bell–Kochen–Specker result; contradiction ensues in the converted case only for certain states.) I would regard this conversion as further evidence that to focus on locality is to distort the discussion. What precipitates No-Go results are overambitious plans of non-contextual determinate value assignment, whether the systems at issue are composite and spatially separated, or simple. Others (including Bell!) would say that it is only in the cases where locality motivates the requisite non-contextuality that the No-Go results have any bite.

Correlation and explanation

In articulating his principle of the common cause, Hans Reichenbach heeded twentieth-century revolutions in physics. Taking quantum mechanics to preclude deterministic causes and relativity to preclude non-local ones, he offered *common causes* as causes which act both locally and stochastically. Roughly, where A and B are events correlated in the sense that

$$\Pr(A \& B) \neq \Pr(A) \times \Pr(B)$$

their common cause C is an event in the overlap of their backwards lightcones rendering A and B probabilistically independent in the sense that

$$\Pr(A \& B|C) = \Pr(A|C) \times \Pr(B|C)$$

The principle of the common cause frames an influential and intuitively attractive account of explanation. Correlations – for instance, the correlations effected by $|\psi\rangle_{singlet}$ – are what require explanation; explanation proceeds by specifying a common cause for the correlated events. Straightforwardly applied to quantum correlations, the principle comes to grief. Articulated to regulate demands for explanation in the context of *statistical* theories, the principle, applied to the perfect (anti)correlations established by $|\psi\rangle_{singlet}$, is satisfied only by *deterministic* common causes, that is, Cs such that $\Pr(A|C), \Pr(B|C) \in \{0,1\}$ (van Fraassen, 1989). What's more, the assumption that there are common causes for correlations observed in the Bell experiments implies the Bell Inequalities (van Fraassen, 1989). Thus any theory satisfying Reichenbachian demands for explanation will be empirically false.

Explanatory activity adheres to standards: not all demands for explanation are legitimate; not all putative *explanans* are satisfactory. Philosophers of science would like to tell the difference. One way to tell the difference is, as it were, ahead of time, by articulating a template to which scientific explanations always and everywhere conform. Taking the common causal account of explanation as just such a template, Fine (1989) and van Fraassen (1989) present its quantum travails as evidence that essentialism about explanation is misplaced, that explanatory strategies arise within the various sciences variously. But for many, the feeling persists that QM's capacity to predict correlations falls dramatically short of a capacity to explain those correlations.

The Measurement Problem

These No-Go results can be read as fables whose moral is that we ought not be too ambitious in ascribing quantum observables determinate values. One way to moderate our ambition is to adopt the semantics typically announced by textbooks:

> [I]t is strictly legitimate to say that $\hat{O}$ has a value in a state $|\psi\rangle$ if and only if a measurement of $\hat{O}$ on this state is certain to yield a definite result – i.e. if and only if $|\psi\rangle$ coincides with an eigenvector of $\hat{O}$ (Gillespie, 1973, p. 61).

Although this ***eigenstate/eigenvalue link*** averts No-Go results, there is another debacle in store for it. A measurement is an interaction between an object system S and an apparatus R prepared in its ready state $|p_0\rangle$, ideally one that establishes a perfect correlation between eigenstates of the object observable $\hat{O}$ and pointer observable $\hat{P}$. If measurement is a quantum mechanical process, this correlation-establishing evolution should be Schrödinger evolution, and so implemented by a unitary operator $\hat{U}_M$:

$$\hat{U}_M(|o_i\rangle|p_0\rangle) = |o_i\rangle|p_i\rangle \tag{10.7}$$

The right-hand side of (10.7) is the post-measurement state, a state in which both the object and pointer observables have determinate values, according to textbook semantics; a state in which the pointer value reflects the value of the object observable. This consolidates the status of evolution driven by $\hat{U}_M$ as ***measurement*** evolution. But consider what happens when an object system initially in a superposition $\Sigma_i c_i|o_i\rangle$ of $\hat{O}$ eigenstates is subject to a measurement of the sort just

described. To obtain the post-measurement state of the composite system, apply $\hat{U}_M$ to the premeasurement state.

$$\hat{U}_M\left[\sum_i c_i|o_i\rangle|p_0\rangle\right] = \sum_i c_i\hat{U}_M(|o_i\rangle|p_0\rangle) = \sum_i c_i|o_i\rangle|p_i\rangle \tag{10.8}$$

(Use $\hat{U}_M$'s linearity to move from the first expression to the second, and (10.7) to move from the second to the third.) Unitary measurement leaves the object + apparatus system in the entangled state $\Sigma_i c_i|o_i\rangle|p_i\rangle$ which is ***not*** an eigenstate of the pointer observable $\hat{P}$. According to textbook semantics, then, the pointer observable has no determinate value, and the measurement has no outcome. (One version) of ***the measurement problem*** is that if measurement processes obey the laws of quantum dynamics, then measurements rarely have outcomes. Cautious enough to avoid No-Go results, textbook semantics are too cautious to accommodate the manifest and empirically central fact that experiments happen. If QM as interpreted by textbook semantics were true, we'd be unable to confirm it!

Recognizing this problem, von Neumann (1955 [1932]) responded by invoking the ***deus ex machina*** of measurement collapse, a sudden, irreversible, discontinuous change of the state of the measured system to an eigenstate of the observable measured. According to this (quite orthodox – many texts accord this "Collapse Postulate" axiomatic status) way of thinking, reconciling textbook semantics with the datum that there are empirical data requires suspending unitary dynamics in measurement contexts, and interpreting Born Rule probabilities as probabilities for collapse. Collapse is a Humean miracle, a violation of the law of nature expressed by the Schrödinger equation. If collapse and unitary evolution are to coexist in a single, consistent theory, situations subject to unitary evolution must be sharply and unambiguously distinguished from situations subject to collapse. And despite evocative appeals to such factors as the intrusion of consciousness or the necessarily macroscopic nature of the measuring apparatus, no one has managed to distinguish these situations clearly.

Contemporary Work

I now have on hand material sufficient to frame much recent philosophical work on QM. The challenge is to offer an interpretation of the theory which makes sense of measurement outcomes without running afoul of NoGo results. Such an interpretation will have to revise one or more of the following naive identifications, the set of which precipitates the measurement problem:

Quantum *states* are normed vectors $|\psi\rangle$ on a Hilbert space $\mathcal{H}$.
Quantum *observables* are self adjoint operators on $\mathcal{H}$.
Quantum *dynamics* is unitary Schrödinger dynamics.
Quantum *semantics* are given by the eigenstate/eigenvalue link.

The revisions that require the least new physics, are semantic revisions; revisions which retain the standard state space but reconfigure its dynamical trajectories are more radical; most radical of all are revisions to the fundamental state space and observable set of QM. A recurrent feature of interpretations of QM is that their conservative exteriors hide radical hearts.

Changing the dynamics: The GRW model

The GRW model of quantum processes (Ghirardi, Rimini and Weber, 1986) – see also Pearle (1989) – would avoid having to reconcile Schrödinger and non-Schrödinger evolution by dispensing with Schrödinger evolution. GRW offers in its stead a more general form of state evolution, to which Schrödinger evolution is nearly approximate. The GRW equation of motion for an isolated quantum system supplements the usual unitary term with a non-unitary term. The effect of this extra term is, rarely and at random, but with a uniform probability per second (10^{-15}), to multiply the system's configuration space state $|\psi(x) >$ by a Gaussian (bell curve) of width 10^{-7} meters, then normalize. The result of a hit by a Gaussian centered at $x = q$ is a wave function $|\psi_q(x) >$ localized about q. Given that a particle in the state $|\psi(x) >$ is hit by a Gaussian, the GRW dynamics set the probability that it's hit by a Gaussian centered at $x = q$ equal to the Born Rule probability $|\psi(q)|^2$ that a position measurement performed on a system in the state $|\psi(x) >$ has the outcome q.

Generally, when systems interact, their composite state becomes entangled. For instance, a purely unitary $\hat{\sigma}_x$ measurement coupling a pointer system containing N particles to an electron in initial state $c_+|\rightarrow> + c_-|\leftarrow>$ generates the post measurement state

$$c_+|\rightarrow>\bigotimes_{i=1}^{n}|\chi_+(x)\rangle_i + c_-|\leftarrow>\bigotimes_{i=1}^{n}|\chi_-(x)\rangle_i \qquad (10.9)$$

where $|\chi_\pm(x) > i$ represents the i^{th} particle in a pointer localized about $x = \pm L$. As the number of particles in the pointer grows, so too does the probability that one of them experiences a GRW collapse. The entanglement of (10.9) ensures that multiplying the state of any particle in the pointer by a Gaussian centered at $+L$ renders the second term on the right-hand side negligible, and so leaves the composite system localized about $+L$. Because our measuring apparatuses (generally) couple a macroscopic number of systems together, such a reduction is overwhelmingly likely to occur practically immediately upon the completion of measurement.

Thus, the GRW dynamics imply that the quantum states of individual systems will almost always Schrödinger evolve, while the quantum states of macroscopic measuring apparatuses are almost always highly localized. But this does not render GRW an unqualified success. It accounts only for measurement outcomes recorded in positions. However, it may not be that all measurement outcomes are so

recorded (Albert (1992, ch. 5) presents one which, *prima facie*, is not). And in addition to modifying the dynamics of the naive interpretation, GRW must modify its semantics, and perhaps even its observable set. For GRW reductions are not reductions to strictly localized states (that is, states $|\varphi(x)\rangle$ such that for some finite interval Δ, $\int_\Delta \varphi^*(x)\varphi(x)dx = 1$ – recall that there are no point-valued position eigenstates). Rather, they are reductions to states with infinite tails in configuration space. The *problem of tails* is that adhering strictly to the orthodox semantics motivating their pursuit of reduction, GRW cannot attribute even interval-valued determinate positions to even systems in post-reduction states such as $|\psi_q(x) >$. Relieving us of peculiar measurement dynamics, GRW does not supply our pointers with determinate positions.

A possible recourse is to liberalize eigenstate/eigenvalue semantics so that "System S in $|\psi(x)\rangle$ is localized in the interval Δ" is true iff

$$\int_\Delta \psi^*(x)\psi(x)\mathrm{d}x > 1 - \varepsilon$$

where $0 < \varepsilon < \frac{1}{2}$ (Albert and Loewer, 1996). Setting $\varepsilon = 0$ reinstitutes the eigenstate/eigenvalue link; setting $\varepsilon = \frac{1}{2}$ allows incompatible propositions (for instance, those associated with the projectors P_Δ and $I - P_\Delta$) to be true at once. Setting ε somewhere in between implies that a system in the state $|\psi(x)\rangle$ can be *localized in* Δ while a system in the state $|\psi'(x)\rangle$ is *localized in* Δ', where Δ and Δ' are disjoint, even though $|\psi(x)\rangle$ and $|\psi'(x)\rangle$ are not orthogonal. If so, GRW's *localized* observable is not a standard quantum mechanical one. For quantum observables, self-adjoint operators, are projection-valued measures which associate distinct eigenvalues of the observable with orthogonal subspaces of Hilbert space. Though Δ and Δ' are distinct values of GRW's *localized* observable, $|\psi(x)\rangle$ and $|\psi'(x)\rangle$ are not orthogonal, and *localized* is not a self-adjoint operator. By liberalizing textbook semantics, GRW makes the more radical interpretive move of revising QM's observable set.[7]

Changing the state space: The Bohm theory

On Bohm's causal interpretation – originating in Bohm (1952); see Cushing et al. (1996) for subsequent developments – all particles have determinate positions. Thus Bohm attributes a system of N particles of mass m moving in three dimensions a determinate *configuration* $Q \in \mathbb{R}^{3N}$ in a configuration space of their possible joint positions. The quantum wave function $\psi(x_1, \ldots, x_{3N})$ for the system can be expressed as a function over this configuration space. Manipulating the Schrödinger equation, and reasoning by analogy with other bits of physics, Bohm offers a set of velocity functions

$$\dot{x}_i = \upsilon((x_1, \ldots, x_{3N})) = \frac{1}{m}\mathrm{Im}\left(\frac{\nabla_i \psi(x_1, \ldots, x_{3N})}{\psi(x_1, \ldots, x_{3N})}\right)$$

which make each component of each particle's velocity depend both on the quantum state and on the configuration of the composite system.[8] Possessing at all times "precisely definable and continuously varying values of position and momentum" (Bohm, 1952, p. 373), a Bohmian particle follows a deterministic trajectory.[9]

Bohmian Mechanics is the ***guidance condition*** (the velocity functions $\dot{x}_i$) along with the requirement that $\psi(x_i)$ evolve in accordance with the Schrödinger equation. Given the appropriate initial conditions, an ensemble of particles following their Bohmian trajectories can reconstitute quantum statistics. If at some initial time t_0, the distribution of determinate positions among particles in an ensemble assigned $\psi(x_i,t_0)$ is described by the probability density $|\psi(x_i,t_0)|^2$, then at all later times, probability densities are well-behaved, and described by the appropriate Schrödinger developments of $\psi(x,t_0)$. Bohm's ***distribution postulate*** is that $|\psi(x_i)|^2$ does give the probability density.[10]

Bohm's interpretation does not assign noncontextual determinate values to observables other than position. Non-position observables it relegates to the realm of dispositions manifested in the post-measurement positions of pointer systems. These dispositions are contextual: whether a particle described by some superposition of spin eigenstates will wind up in a position indicating spin up or spin down depends not only on the initial position of the particle but also on how the measuring device is configured; Albert (1992, ch. 7) gives a simple illustration. And positions themselves are subject to manifestly non-local influences: the velocities of individual particles are functions of the configurations of the composite systems they comprise, so that (reverting to the EPR case) changes in particle II's position instantaneously alter particle I's velocity. Bohmians deem this non-locality benign. Maintaining that we can not predict or control particle positions, they argue that we can not harness the non-locality for signaling purposes.

Stingy, contextual, non local, the Bohm interpretation avoids No-Go results. It accounts for measurement outcomes recorded in particle positions. And it seems to its adherents "the most obvious," "most natural," and "simplest" (Dürr et al., 1996, pp. 21, 24) account of quantum phenomena – so much so that Cushing (1994) has suggested that had Bohm beaten Bohr to prominence, standard physics curricula would include Bohmian, rather than quantum, mechanics.

Pleas for the naturalness of Bohmian mechanics sometimes derive illicit support from glosses like the following (which Bell supplied for its ancestor, de Broglie's pilot wave theory):

> De Broglie showed in detail how the motion of a particle passing through just one of two holes in a screen could be influenced by waves propagating through both holes. And so influenced that the particle does not go to where the waves cancel out, but is attracted to where they cooperate. This idea seems to me so natural and simple, to resolve the wave-particle dilemma in such a clear and ordinary way, that it is a great mystery to me that it was so generally ignored. (Bell, 1987, p. 191)

On the picture enchanting Bell, the particle surfs the pilot wave through space. It is a picture that results when the Bohm theory's state space – the configuration space $\mathbb{R}^3$ on which the wave function of a single particle is defined – is identified with physical space. Such an identification threatens to confuse the representational with the concrete, and is anyway foiled once the theory considers systems composed of $N > 1$ particles, with wavefunctions $\psi(x_1, \ldots, x_{3N})$. Then surf's up not in $\mathbb{R}^3$ but in a $3N$-dimensional configuration space which it is not tempting to identify with physical space.

Natural or not, the Bohm theory is significant. By refusing to constitute matters of fact from determinate quantum observables, Bohm not only circumvents the usual No-Go results, but shows how they stack the deck against the non-quantum physicist by foisting upon her a quantum-theoretic space of possibilities.

Changing the semantics: Modal interpretations

Modal interpretations – see Kochen (1985), Healey (1989), Dieks (1989), van Fraassen (1991) and, for more recent work, Dieks and Vermaas (1998) – would resolve the Measurement Problem by maintaining the universality of Schrödinger evolution while revising the eigenvector/eigenvalue link. A stock example of a modal interpretation exploits the biorthogonal decomposition theorem, according to which any vector $|\psi\rangle^{SR}$ in the tensor product space $\mathcal{H}_S \otimes \mathcal{H}_R$ admits a decomposition of the form

$$|\Psi\rangle^{SR} = \sum_i c_i |a_i > |b_i >$$

where $\{c_i\}$ are complex coefficients, $\{|a_i>\}$ and $\{|b_i>\}$ are sets of orthogonal vectors on $\mathcal{H}_S$ and $\mathcal{H}_R$ respectively, and the summation index i does not exceed the dimensionality of the smaller factor space. If the set $\{|c_i|^2\}$ is non-degenerate, then this biorthogonal decomposition of $|\psi\rangle^{SR}$ is unique. Modal interpretations replace the orthodox eigenvector/eigenvalue link with the following semantic rule:

> If $|\Psi\rangle^{SR} = \Sigma_i c_i |a_i > |b_i$ is the unique biorthogonal decomposition of the state of a composite $S + R$ system, then subsystem S has a determinate value for each $\mathcal{H}_S$ observable with eigenbasis $\{|a_i>\}$, and subsystem R has a determinate value for each $\mathcal{H}_R$ observable with eigenbasis $\{|b_i>\}$. $|c_i|^2$ gives the probability that these observables' actual values are the eigenvalues associated with $|a_i > |b_i >$.

Consider the unitarily evolved post-measurement state

$$|\Psi>^{S+R} = \sum_i c_i |o_i > |p_i >$$

The eigenbasis of the pointer observable $\hat{P}$ conspires in its biorthogonal decomposition. By modal semantics, then, the pointer observable $\hat{P}$ is determinate on the apparatus system after measurement. Moreover, the probability that $\hat{P}$'s actual value is p_n is just the Born Rule probability. Thus would modal interpretations explain what textbook interpretations can not: how measurement interactions obedient to the laws of quantum dynamics issue determinate outcomes corroborating quantum statistical predictions.

Four problems for this stock modal interpretation are listed here:

(i) What to say when the biorthogonal decomposition is degenerate
In the extreme case where $\mathcal{H}_S$ and $\mathcal{H}_R$ are each of dimension $N > 2$ and

$$|\Psi\rangle^{SR} = \sum_i \frac{1}{\sqrt{N}} |a_i>|b_i>$$

the eigenbasis of every observable on the component systems conspires in some biorthogonal decomposition, and Kochen–Specker contradictions threaten.

(ii) What to say about the dynamics of possessed values
A viable option, one preserving the status of the modal interpretation as an interpretation that succeeds not by developing new physics but by adjusting semantics to existing physics is: nothing. Dickson (1998a) describes modal dynamics which are dramatically underdetermined by the requirement that they return single time probabilities conforming to the Born Rule, and discusses that underdetermination.

(iii) What to say about state preparation, the laboratory processes whereby we assign states to quantum systems
Modal interpretations cannot avail themselves of the standard account that measurement collapse leaves the prepared system in the eigenstate of the measured observable corresponding to the eigenvalue obtained. Perhaps modal interpretations can account for preparation by appeal to conditional probabilities: the "prepared" state is the one mimicking the post-preparation composite state's predictions for the prepared system, conditional on the "outcome" of the preparation – Wessels (1997) treats preparation along these lines. Adopting standard quantum expressions for conditional probabilities, modal interpretations can take this way with preparation at the cost, in certain settings, of violating the Markov consistency requirement that

$$\Pr(a|b) = \sum_i \Pr(a|c_i) \times \Pr(c_i|b)$$

where $\{c_i\}$ is an exhaustive set of mutually exclusive events intermediate between a and b. Using non-standard conditional probabilities, modal interpretations embark on value state dynamics, with the class of candidate dynamics narrowed to those that make sense of preparation.

(iv) What to make of non-ideal measurements (Albert, 1992, appendix)
These are measurements which fail to correlate eigenstates of the designated pointer observable with orthogonal states of the object system, so that the

pointer eigenbasis fails to furnish a biorthogonal decomposition of the post-measurement composite state. By the biorthogonal decomposition theorem, some apparatus eigenbasis will furnish a biorthogonal decomposition, and observables with this eigenbasis, not the pointer observable, are determinate after measurement, according to modal semantics. Perfectly error-free measurements confront modal intepretations with this problem, and there is a class of observables whose only error-free measurements are of this sort (Ruetsche, 1995).

Responses to (iv) (and also (i) and (iii)) appeal to *decoherence* processes – interactions between the pointer and its environment that tend to correlate distinct pointer eigenstates with nearly orthogonal states of the environment.[11] The suggestion is that decoherence carries post non-ideal measurement systems into states biorthogonally decomposed by apparatus observables close enough to the designated pointer observables that one needn't fret (Bacciagaluppi and Hemmo, 1996). Because decoherence is not perfect, this response leaves the modal interpretation with its own version of the problem of tails, a problem whose resolution might lie in the now-familiar maneuver of constituting matters of fact from something other than determinate quantum observables.

Relative state formulations

"Postulat[ing] that a wave function that obeys a linear wave equation everywhere and at all times supplies a complete mathematical model for every isolated physical system without exception" (Everett, 1983 [1957], p. 316), Hugh Everett's Relative State Formulation promises an interpretation according to which the quantum state description is complete and the quantum dynamics are universal. Although the entangled post-measurement state

$$|\Psi>^{S+R} = \sum_i c_i|o_i>|p_i>$$

associates no $\hat{O}$ ($\hat{P}$) eigenstates with the object (apparatus) simpliciter, it correlates $\hat{O}$ and $\hat{P}$ eigenstates with one another. This illustrates Everett's moral that "the state of one subsystem does not have an independent existence, but is fixed only by the state of the remaining subsystem" (1983 [1957], p. 316), so that "it is meaningless to ask the absolute state of a subsystem – one can only ask the state relative to a given state of the remainder of the system" (1983 [1957], p. 317). (Relatively speaking) when system S has determinate $\hat{O}$ value o_n, system R has determinate $\hat{P}$ value p_n, and "this correlation is what allows one to maintain the interpretation that a measurement has been performed" (1983 [1957], p. 320). Thus Everett purports to reconcile the uncollapsed composite state $|\psi>^{S+R}$ with determinate measurement outcomes.

But the terms of reconciliation are notoriously unclear. An option proposed by physicists but embraced by the science fiction community is that "the universe is constantly splitting into a stupendous number of branches, all resulting from the measurement-like interactions between its myriads of components" (DeWitt, 1970, p. 161); within each branch, the relative state of the pointer registers a determinate outcome. Criticisms of this version of Everett (Albert and Loewer, 1988) include that its profligate creation of new universes violates the conservation of mass/energy required by unitary evolution, and that it makes hash of quantum probabilities by rendering every outcome *certain* to occur along some branch. What's more, to disambiguate this version of Everett, its proponents must furnish an account of when splitting occurs, and into what branches. Such an account would serve also on the von Neumann collapse interpretation to distinguish systems subject to collapse from systems evolving unitarily, rendering that interpretation consistent, unambiguous, and free of suspect metaphysics.

More recent Everett-style interpretations have responded to the disambiguation problem in one of two broad ways. The more fanciful notes that it is, after all, only our determinate experiences which must be reconciled with universal unitary evolution, and so postulates "eigenstates of mentality" – brain states to which correspond mental states whose contents are determinate beliefs – as a preferred basis of relative states. Perhaps the most astonishing variation of this approach is the Many Minds interpretation, a radical dualism which invites us to

> Suppose that every sentient physical system there is is associated not with a single mind but rather with a *continuous infinity* of minds; and suppose (this is part of the proposal too) that the measure of the infinite subset of those minds which happen to be in some particular mental state at any particular time is equal to the square of the absolute value of the coefficient of the brain state associated with that mental state, in the wave function of the world at that particular time. (Albert, 1992, p. 130)

A more prosaic response (Griffiths, 1993; Hartle, 1990) to the disambiguation problem offers *consistent* (or *decoherent*) histories as the prefered basis of relative states. A time-indexed set of determinate observables generates a family of *histories* for a system; an individual history in the family assigns observables in the time-indexed set determinate values. Given the initial state of the system and the unitary operator governing its evolution, a generalized Born Rule assigns probabilities to such histories. A family of histories is said to be consistent if the probabilities so assigned do not "interfere" – roughly, they are Markov consistent. Thus histories in a consistent family admit multi-time probability assignments that constitute a tractable dynamics.

The rub is that while the initial state and the system Hamiltonian constrain which families of histories are consistent, they don't determine a unique family of consistent histories. So, while there may be a consistent family of histories declaring the pointer observable determinate at measurement's completion, there will

also be other consistent families which do not. What assures that a consistent family containing the pointer observable, and not one excluding it, corresponds to what actually occurs in the laboratory? Branding families of consistent histories which foil merger into a family satisfying the non-interference condition "complementary," Griffiths rejects this yen for reassurance on broadly Bohrian grounds: "A question of the form, 'Which of these really took place?' asked in terms of comparing two mutually incompatible histories, makes no sense quantum mechanically" (Griffiths 1993, p. 2204).

Like the perspectival metaphysics of the many worlds approach, this response is philosophically suspect. Yet Everett-style approaches are the preferred quantum framework for many working physicists. Rovelli (1997) sees in "relational QM" the seeds of a solution to the problem of time in quantum gravity; Hartle (1990) puts the consistent histories approach, and the tractable (if perspectival) dynamics it underwrites, to cosmological use. Meanwhile, interpretations of QM more philosophically respectable languish relatively unloved.[12] Saunders offers a stark diagnosis: "The disturbing feature of both the Bohm and GRW approaches is that they seem to require that we *redo* high energy physics" (Saunders, 1996, pp. 125–6). Requiring a preferred time foliation, both approaches fundamentally (if not phenomenologically) violate Lorentz and general covariance, and thus deprive physicists of a powerful criterion for winnowing down the set of acceptable theories. This should remind us at least that non-relativistic QM is not the only game in town – a lesson those working on the foundations of quantum theories have increasingly taken to heart.

Future Directions: Interpreting QFT

With apologies to those who have been working in the field for years – for a very recent review, see Huggett (2000) – I offer QFT – and quantum gravity, a theory about whose eventual shape QFT on curved spacetimes might hold a clue – as one future direction for the philosophy of quantum theories. Moving from the least to the most exotic space–time settings, this section sketches some issues that are kicked up by the pursuit of quantum theories in such settings.

Minkowski space–time

The proper setting for questions about "locality," QFT is also a provocative one. A striking example is the Reeh–Schlieder theorem, which states that where $\{A(O)\}$ is the set of observables the theory associates with an open bounded region of space–time O and $|0\rangle$ is the Minkowski vacuum state, $\{A(O)\,|0\rangle\}$ is *dense* in the theory's state space – that is, any state the theory recognizes can be approximated arbitrarily closely by acting on the vacuum by polynomial combinations of observables in $\{A(O)\}$. If it were appropriate to model events in the region O as applications of elements of $\{A(O)\}$ to the global vacuum state, this would mean that machinations in local regions could produce arbitrary approximations of arbitrary global states! The model is not apt, but its whiff of non-locality is. The Reeh–Schleider theorem implies that $|0\rangle$ is an eigenstate of no observable associated with a finite space–time region, which in turn implies that the vacuum spreads correlations far and wide. Redhead (1995) illuminates the Reeh–Schleider theorem by explicating analogies between how the vacuum stands to local algebras of observables and how the spin-singlet state stands to algebras of spin observables pertaining to the component systems. Clifton et al. (1998) show that states with $|0\rangle$'s feature that given any pair of space–time regions, any observable from one is correlated with some observable from the other, are dense; Butterfield (1994) discusses the capacity of such correlations to violate Bell-type inequalities (they can, even maximally). The nature and extent of such non-local features of QFT, as well as the theory's hospitability to causal talk, are topics of ongoing research.

To see how questions about the ontology of QFT, as well as its state space, arise, we need to go into a bit more detail. The canonical approach to quantization casts a classical theory in Hamiltonian form, then promotes its canonical observables q_k, p_k to symmetric operators $\hat{q}_k, \hat{p}_k$ obeying *canonical commutation relations* arising from the Poisson brackets of the classical theory. A classical *field* theory assigns a field configuration $\varphi(x)$ and a conjugate momentum density $\pi(x) \equiv \partial L/\partial\dot{\phi}$ (where L is the theory's Lagrangian density) to every point x of space–time; its quantization proceeds by finding operators $\hat{\varphi}(x)$ and $\hat{\pi}(x)$ obeying the relevant canonical commutation relations.[13] I will refer in what follows to a mathematically well-behaved exponential form of these commutation relations known as the Weyl relations, and call sets of operators satisfying them *representations* of the Weyl relations.

A simple classical field is the Klein–Gordon field $\varphi(x)$, which satisfies

$$(g^{ab}\nabla^a\nabla_a - m^2)\varphi(x) = 0$$

Its solutions can be Fourier-decomposed into uncoupled normal modes with angular frequency ω_k, and so the classical field can be modeled as an infinite collection of independent oscillators. The textbook route to quantization exploits this analogy by introducing creation and annihilation operators $\hat{a}_k^\dagger$ and $\hat{a}_k$ for field modes obeying

$$[\hat{a}_k, \hat{a}_{k'}] = 0 = [\hat{a}_k^\dagger, \hat{a}_{k'}^\dagger], [\hat{a}_k, \hat{a}_{k'}^\dagger] = \delta_{kk'}\hat{I} \qquad (10.10)$$

Formal expressions for operators $\hat{\varphi}(x)$ and $\hat{\pi}(x)$ satisfying the canonical commutation relations can be constructed from these. The resulting quantization is the free boson field; imposing anti-commutation relations in lieu of (10.10) yields the free fermion field.

The state $|0>$ such that $\hat{a}_k|0> = 0$ for all k is the lowest energy eigenstate of the quantum Hamiltonian for the free boson field. The state $(\hat{a}_k^\dagger)^n|0>$ is an eigenstate of the Hamiltonian with the same energy a system of n particles each with energy $\hbar\omega_k$ would have – provided the momenta and rest masses of these particles are given by standard relativistic expressions. Thus, the theory tempts a particle interpretation:

- the vacuum state $|0>$ is the no particle state
- the state $\hat{a}_k^\dagger|0>$ describes one particle of energy $\hbar\omega_k$. . . .
- $\hat{N}_k = \hat{a}_k^\dagger\hat{a}_k$ is the number operator for particles of type k
- $\hat{N} = \sum_j \hat{a}_j^\dagger\hat{a}_j$ is the total number operator

and so on. Countering this temptation in the first instance are some distinctly unparticulate features of the theory so interpreted (Teller, 1995, ch. 2). For one thing, the theory hosts states with indeterminate particle numbers. For another, even states which are eigenstates of the total number operator are constrained by (10.10) to be symmetric – that is to be unchanged under permutations of particle labels. Whether this, and their ensuing obedience to Bose-Einstein statistics, deprives bosons of the genidentity one might expect from particles has been a topic of lively debate, well-represented in Castellani (1998).

A prior challenge to the viability of particle interpretations has excited somewhat less interest among philosophers. Consider two quantum theories, each taking the form of a Hilbert space $\mathcal{H}$, and a collection of operators $\{\hat{O}_i\}$. When are these theories physically equivalent? A natural criterion of equivalence is that the theories recognize the same set of states, that is, probability distributions over eigenprojections of their observables. And a sufficient condition for this is that the theories be *unitarily equivalent* in the sense that there exists a unitary map $\hat{U}: \mathcal{H} \rightarrow \mathcal{H}'$ s.t. $\hat{U}^{-1}\hat{O}'_i\hat{U} = \hat{O}_i$ for all values of i, in which case the expectation values assigned observables $\{\hat{O}_i\}$ by any state $|\psi\rangle$ in the first theory are duplicated by those assigned $\{\hat{O}'_i\}$ by the state $\hat{U}|\psi\rangle$ in the second. If the observable set is rich enough, unitary equivalence is necessary as well. If physical equivalence is unitary equivalence, the quantization of a classical theory yields a unique quantum theory if and only if all representations of the relevant Weyl relations are unitarily equivalent. The Stone–Von Neumann theorem ensures that representations of Weyl relations expressing the quantization of a classical theory with a finite dimensional state space are unique upto unitary equivalence. But classical fields have infinitely many degrees of freedom. The Stone–Von Neumann theorem does not apply. Indeed, the Weyl relations encapsulating the quantization of classical Klein–Gordon theory admit continuously many inequivalent representations.

Let $\{\hat{a}_k, \hat{a}_k^\dagger\}$ be one quantization of some classical field theory, and $\{\hat{a}'_k, \hat{a}'^\dagger_k\}$ be another, unitarily inequivalent to the first. In general, the primed vacuum state will not be a state in the unprimed representation, nor will the primed total number operator be an operator there, and *mutatis mutandis*. One might say that associated with the unitarily inequivalent quantizations are *incommensurable* particle notions. Even granting that it is appropriate to run a particle interpretation of a quantization $\{\hat{a}_k, \hat{a}_k^\dagger\}$, one cannot run a sensible particle interpretation of QFT unless one can privilege as *physical* a unitary equivalence class of representations admitting particle interpretations – as Saunders (1988) discusses, not all do.

The default setting for a QFT is Minkowski space–time. And this furnishes a de facto criterion of privilege: physical representations respect the symmetries of the space–time, in the sense that their vacua are invariant under its full isometry group. Coupled with the requirement that physical representations admit only states of non-negative energy, this singles out a unitary equivalence class of representations. But this strategy for privilege breaks down in generic curved space–time settings, which do not supply the symmetries it requires.[14]

The *algebraic approach* to quantum theories grounds an entirely different response to unitarily inequivalent representations. The algebraic approach articulates the physical content of a theory in terms of an abstract algebra $\mathcal{A}$. Observables are elements of $\mathcal{A}$, and states are normed, positive linear functionals $\omega: \mathcal{A} \to \mathbb{C}$. The expectation value of an observable $A \in \mathcal{A}$ in state ω is simply $\omega(A)$. Abstract algebras can be realized in concrete Hilbert spaces. A map π from elements of the algebra to the set of bounded linear operators on a Hilbert space $\mathcal{H}$ furnishes a Hilbert space representation of the algebra. In particular, all Hilbert spaces carrying a representation of the Weyl relations are also representations of the abstract Weyl algebra. For the algebraist, "[t]he important thing here is that the observables form some algebra, and not the representation Hilbert space on which they act" (Segal, 1967, p. 128). Inequivalent representations need not puzzle him, for conceiving the state space of a quantum theory as the space of algebraic states, he has rendered unitary equivalence an inappropriate criterion of physical equivalence. (Early proponents of the algebraic approach concocted baldly operationalist motivations for alternative glosses on physical equivalence; see Summers (1998) for a review.) Nor need he trouble with particles, for particle notions are (at least prima facie) the parochial residues of concrete representations.

Standard quantum states are probability measures over closed subspaces of Hilbert space. The class of algebraic states is broader than the class of such probability measures. There are, for instance, algebraic states which can accomplish what no Hilbert space state can: the assignment of precise and punctal values to continuous observables (Clifton, 1999). Some would advocate restricting admissable algebraic states. A restriction that looks down the road to quantum gravity is the Hadamard condition, which requires admissable state to be states for which a prescription assigning the stress-energy tensor an expectation value succeeds. (Provocatively, in closed space–times such states form a unitary equivalence class (Wald, 1994, §4.6). Both mathematical and physical features of algebraic states merit, and are receiving, further attention, attention which should inform discussion about the state space, and maybe even the ontology, appropriate to QFT.

Curved space–time

Different notions of state demand different dynamical pictures. Hilbert space dynamics are implemented by unitary Hilbert space operators. Having jettisoned Hilbert spaces as essential to QFT, the algebraist has jettisoned as well this account of the theory's dynamics. In its stead, he implements quantum field dynamics by means of automorphisms of the abstract algebra $\mathcal{A}$ of observables (that is, structure-preserving maps from $\mathcal{A}$ to itself). A question of equipollence arises: is it the case that every dynamical evolution implementable by an automorphism on the abstract algebra is also implementable as a unitary evolution in some fixed Hilbert space? Algebraic evolution between Cauchy slices related by isometries can be implemented unitarily, but more general algebraic evolution can not be; see Arageorgis et al. (2001) for details. One moral we may draw from these results is that in exactly those space–times whose symmetries furnish principles by appeal to which a unitary equivalence class of representations might be privileged, dynamical automorphisms are unitarily implementable. In more general settings, the algebraic formulation is better suited to capturing the theory's dynamics.

Unitarity breaks down even more dramatically in the exotic space–time setting of an evaporating black hole. Hawking has argued that a pure to mixed state transition – the sort of transition von Neumann's collapse postulate asserts to happen on measurement – occurs in the course of black hole evaporation. Not only unitarity but also symmetries of time and pre/retrodiction are lost if Hawking is right. Belot et al. (1999) review reactions to the Hawking Information Loss Paradox; not the least of the many questions the Hawking paradox raises is how to pursue QFT in non-globally hyperbolic space–times.

Quantum gravity

The QFTs discussed so far are free field theories, whereas the QFTs brought into collision with data from particle accelerators are *interacting* field theories, whose empirical quantitites are calculated by perturbative expansions of the free field. The divergence of these expansions calls for the art and craft of renormalization, chronicled in Teller (1995, chs 6 and 7). Cushing (1988) argues that this (and every other!) feature of QFT raises not "foundational" but "methodological" issues. Insofar as methodological predilections are affected by foundational commitments and affect the shape of future theories, the two domains might not be so cleanly separable as Cushing suggests; ongoing work on Quantum Gravity is one place to look for their entanglement.

Notes

1 See Hughes (1989, pt. I) for an introduction. As space limitations prohibit even a rudimentary review, I attempt in what follows to minimize technical apparatus.

2 Such explanation has its limits. Consider $\hat{\sigma}_x$ and $\hat{\sigma}_y$, perpendicular components of

intrinsic angular momentum or spin. They obey uncertainty relations, and their measurement requires incompatible experimental apparatus. Yet spin is explicitly and infamously a quantum phenomenon, and $\hat{\sigma}_x$ and $\hat{\sigma}_y$ are individually conserved. They are not *classical* concepts precluding one another's application, nor occupants different sides of the kinematic/dynamic divide.

3 For a derivation, which is simply a matter of bringing a home truth about sums of 1s and −1s to bear upon probabilities the HVT assigns, see Redhead (1987, pp. 97–8). This form of Bell Inequality applies to both deterministic and stochastic HVTs. A particularly simple derivation of a Bell Inequality, due to Wigner, applies to deterministic HVTs requiring perfect (anti-) correlation; see Hughes (1989, pp. 170–2). One can derive Bell-type inequalities without the intermediary of hidden variables, provided one assumes that joint probability distributions for non-commuting observables are well-defined; see Redhead (1987, pp. 81–3) for this Stapp-Eberhard form of the inequalities.

4 Maudlin (1994) discusses STR's real and imagined implications, and the constraints they place on the interpretation of QM.

5 Conventionally denoted by the same letter as, but not to be confused with, the space of hidden variable states.

6 Bub (1997) gives a thorough review, and presents Bub and Clifton's trend-bucking "Go" theorem, which characterizes the *largest* set of observables that can without contradiction be attributed determinate values obedient to the Spectrum and FUNC rules.

7 So setting ε also has the repercussions that our discourse about localization features oddities reminiscent of discourse involving vague predicates. For instance, each of a pair of predicates ("is localized in Δ" and "is localized in Δ'") can be true of some system S without their conjunction ("is localized in $\Delta \cap \Delta'$") being true of S. Whether we can live with this is a topic of ongoing debate; see, for instance, Clifton and Monton (1999).

8 The analogies plied in Bohm's original presentation invoke a quantum potential with disquieting features; Dürr et al. (1996) attempt to eliminate this invocation by showing how the velocity function is suggested (if not implied) by symmetry considerations alone.

9 Vink (1993) extends Bohm's approach to assign every observable a determinate value (albeit a contextual one), and offer for those possessed values a generalization of Bohmian dynamics which is stochastic when the observables are discrete.

10 Valentini (1991) would like to unifiy the role of $\psi(x_i)$ in the Bohm theory by proving a "quantum H-theorem" according to which arbitrary initial distributions evolve under the influence of the Bohmian equations of motion to the distribution $|\psi(x_i)|^2$. This would render the distribution postulate otiose. Dickson (1998b, pp. 123–5) offers criticisms of Valentini's approach.

11 Zurek (1982) offers toy models of decoherence processes, as well as the claim that decoherence solves the measurement problem. An apparatus entangled not only with the object system but also with its environment is still entangled, and not a system to which eigenstate/eigenvalue semantics attribute determinate values. To respond to the measurement problem, decoherence proposals need to be accompanied by non-standard semantics. Modal semantics work admirably.

12 But see Huggett and Weingard (1994) for Bohmian approaches to QFTs, and Pearle (1992) for Lorentz-invariant quantum field version of continuous spontaneous localization.

13 Mathematical nicety demands that the quantum field be cast not (as the foregoing suggests) as a map from space–time points to operators, but as operator-valued distributions over space–time regions. Wald (1994) is an excellent introduction to this and other issues discussed in this section.

14 Notoriously, it even breaks down in a subset of Minkowski space–time. Positive energy states correspond to solutions to the Klein–Gordon equation that oscillate with purely positive frequency. States in the standard Minkowski representation are positive frequency with respect to time as measured by families of inertial observers. But restricting our attention to the right Rindler wedge of two-dimensional Minkowski space–time setting $c = 1$, this is the region where x is positive and $|x| < t$ – we can quantize the Klein-Gordon field by admitting solutions that oscillate with positive frequency with respect to time as measured by observers whose accelerations are constant, for Lorentz boost isometries generate a global time function for the Rindler wedge. The Rindler representation we thereby obtain has a natural particle interpretation – but the Rindler representation is unitarily inequivalent to the Minkowski representation! This is sometimes, and loosely, expressed as the Unruh effect: observers accelerating through the Minkowski vacuum "see" a thermal flux of particles (Wald, 1994, ch. 5).

References

Albert, D. (1992): *Quantum Mechanics and Experience*. Cambridge, Mass: Harvard.

Albert, D. and Loewer, B. (1988): "Interpreting the Many Worlds Interpretation," *Synthese*, 77, 195–213.

Albert, D. and Loewer, B. (1996): "Tails of Schrödinger's Cat," in R. Clifton (1996, pp. 81–92).

Arageorgis, A., Earman, J. and Ruetsche, L. (2001): "Weyling the Time Away: The Non-Unitary Implementability of Quantum Field Dynamics on Curved Space–time and the Algebraic Approach to Quantum Field Theory," *Studies in History and Philosophy of Modern Physics*, forthcoming.

Bacciagaluppi, G. and Hemmo, M. (1996): "Modal Interpretations, Decoherence and Measurements," *Studies in the History and Philosophy of Modern Physics*, 27B, 239–77.

Bell, J. (1964): "On the Einstein–Podolsky–Rosen Paradox," *Physics*, 1, 195–200. Reprinted in Wheeler and Zurek (1983, pp. 403–8).

Bell, J. (1966): "On the Problem of Hidden Variables in Quantum Mechanics," *Reviews of Modern Physics*, 38, 447–52. Reprinted in Wheeler and Zurek (1983, pp. 397–402).

Bell, J. (1987): *Speakable and Unspeakable in Quantum Mechanics*. Cambridge: Cambridge University Press.

Belot, G., Earman, J. and Ruetsche, L. (1999): "The Hawking Information Loss Paradox: The Anatomy of a Controversy," *British Journal for the Philosophy of Science*, 50, 189–229

Bohm, D. (1951): *Quantum Theory*. Englewood Cliffs, NJ: Prentice-Hall.

Bohm, D. (1952): "A Suggested Interpretation of the Quantum Theory in Terms of 'Hidden' Variables, I and II," *Physical Review*, 85, 166–93. Reprinted in Wheeler and Zurek (1983, pp. 367–96).

Bohr, N. (1934): *Atomic Theory and the Description of Nature*. Cambridge: Cambridge University Press.

Bohr, N. (1935): "Can Quantum-Mechanical Description of Reality be Considered Complete?" *Physical Review*, 48, 696–702. Reprinted in Wheeler and Zurek (1983, pp. 145–51).

Brown, H. and Harre, R. (eds.) (1988): *Philosophical Foundations of Quantum Field Theory*. Oxford: Clarendon Press.

Bub, J. (1997): *Interpreting the Quantum World*. Cambridge: Cambridge University Press.

Butterfield, J. (1992): "Bell's Theorem: What it Takes," *British Journal for the Philosophy of Science*, 58, 41–83.

Butterfield, J. (1994): "Vacuum Correlations and Outcome Dependence in Algebraic Quantum Field Theory," in D. M. Greenberger and A. Zeilinger (eds.), *Fundamental Problems in Quantum Theory, Annals of the New York Academy of Sciences*, 755, 768–85.

Castellani, E. (1998): *Interpreting Bodies: Classical and Quantum Objects in Modern Physics*. Princeton: Princeton University Press.

Clifton, R. (ed.) (1996): *Perspectives on Quantum Reality: Non-Relativistic, Relativistic, and Field-Theoretic*. Dordrecht: Kluwer.

Clifton, R. (1999): "Beables for Algebraic Quantum Field Theory," in J. Butterfield and C. Pagonis (eds.), *From Physics to Philosophy*, Cambridge: Cambridge University Press, 12–44.

Clifton, R. and Monton, B. (1999): "Losing Your Marbles in Wavefunction Collapse Theories," *British Journal for the Philosophy of Science*, 50, 697–717.

Clifton, R., Feldman, D., Halvorson, H. and Wilce, A. (1998): "Superentangled States," *Physical Review*, A58, 135–45.

Cushing, J. (1988): "Foundational Problems in and Methodological Lessons from Quantum Field Theory," in Brown and Harre (1988, pp. 25–39).

Cushing, J. T. (1994): *Quantum Mechanics: Historical Contingency and the Copenhagen Hegemony*. Chicago: University of Chicago Press.

Cushing, J. T. and McMullin, E. (eds.) (1989): *Philosophical Consequences of Quantum Theory*. Dordrecht: University of Notre Dame Press.

Cushing, J. T., Fine, A. and Goldstein, S. (1996): *Bohmian Mechanics: An Appraisal*. Dordrecht: Kluwer.

DeWitt, B. (1970): "Quantum Mechanics and Reality," *Physics Today*, 23(9); reprinted in B. DeWitt and N. Graham (eds.), *The Many-Worlds Interpretation of Quantum Mechanics*, Princeton: Princeton University Press, 155–67.

Dickson, W. M. (1998a): "On the Plurality of Dynamics: Transition Probabilities and Modal Interpretations," in R. Healey and G. Hellman (eds.), *Quantum Measurement: Beyond Paradox*, Minneapolis: University of Minnesota Press, 160–82.

Dickson, W. M. (1998b): *Quantum Chance and Nonlocality*. Cambridge: Cambridge University Press.

Dieks, D. (1989): "Quantum Mechanics without the Projection Postulate and its Realistic Interpretation," *Foundations of Physics*, 19, 1397–1423.

Dieks, D. and Vermaas, P. (eds.) (1998): *The Modal Interpretation of Quantum Mechanics* Dordrecht: Kluwer.

Dürr, D., Goldstein, S. and Zanghi, N. (1996): "Bohmian Mechanics as the Foundation of Quantum Mechanics," in Cushing et al. (1996, pp. 21–44).

Einstein, A., Podolsky, B. and Rosen, N. (1935): "Can Quantum Mechanical Description of Physical Reality be Considered Complete?" *Physical Review*, 47, 777–80. Reprinted in Wheeler and Zurek (1983, pp. 138–41).

Everett, H. (1957): "'Relative State' Formulation of Quantum Mechanics," *Review of Modern Physics*, 29, 454–62. Reprinted in Wheeler and Zurek (1983, pp. 315–23).

Fine, A. (1982a): "Joint Distributions, Quantum Correlations, and Commuting Observables," *Journal of Mathematical Physics*, 23, 1306–10.

Fine, A. (1982b): "Hidden Variables, Joint Probabilities, and the Bell Theorems," *Physical Review Letters*, 48, 291–95.

Fine, A. (1986), *The Shaky Game: Einstein, Realism and The Quantum Theory*. Chicago: University of Chicago Press.

Fine, A. (1989): "Do Correlations Need to be Explained?" in Cushing and McMullin (1989, pp. 175–94).

Ghirardi, G., Rimini, A. and Weber, T. (1986): "Unified Dynamics for Microscopic and Macroscopic Systems," *Physical Review*, D34, 470–84.

Gillespie, D. T. (1973): *A Quantum Mechanics Primer*. New York: International Textbook Company.

Griffiths, R. B. (1993): "The Consistency of Consistent Histories: A Reply to D'Espagnat," *Foundations of Physics*, 23, 2201–10.

Hartle, J. B. (1990): "The Quantum Mechanics of Cosmology," in S. Coleman, T. Piran and S. Weinberg (eds.), *Quantum Cosmology and Baby Universes*, Singapore: World Scientific, 65–151.

Healey, R. (1989): *The Philosophy of Quantum Mechanics: An Interactive Interpretation*. Cambridge: Cambridge University Press.

Huggett, N. (2000): "Philosophical Foundations of Quantum Field Theory," *British Journal for the Philosophy of Science*, 51, Supplement, 617–37.

Huggett, N. and Weingard, R. (1994): "Interpretations of Quantum Field Theory," *Philosophy of Science*, 61, 370–88.

Hughes, R. I. G. (1989): *The Structure and Interpretation of Quantum Mechanics*. Cambridge, MA: Harvard University Press.

Jarrett, J. (1986): "An Analysis of the Locality Assumption in the Bell Arguments," in L. M. Roth and A. Inomata (eds.), *Fundamental Questions in Quantum Mechanics*, London: Gordon and Breach, 21–7.

Kochen, S. (1985): "A New Interpretation of Quantum Mechanics," in P. Lahti and P. Mittelstaedt (eds.), *Symposium on the Foundations of Modern Physics*, Singapore: World Scientific, 151–69.

Maudlin, T. (1994): *Quantum Nonlocality and Relativity*. Oxford: Blackwell.

Mermin, D. (1993): "Hidden Variables and the Two Theorems of John Bell," *Reviews of Modern Physics*, 65, 803–15.

Murdoch, D. (1987): *Niels Bohr's Philosophy of Physics*. Cambridge: Cambridge University Press.

Pearle, P. (1989): "Combining Stochastic Dynamical State Vector Reduction with Spontaneously Localization," *Physical Review*, A39, 2277–89.

Pearle, P. (1992): "Relativistic Models of State Vector Reduction," in P. Cvitanovic, I. Percival and W. Wirzba (eds.), *Quantum Chaos–Quantum Measurement*, Dordrecht: Kluwer, 283–97.

Redhead, M. (1983): "Nonlocality and Peaceful Coexistence," in R. Swinburne (ed.), *Space, Time and Causality*, Dordrecht: Reidel, 151–89.

Redhead, M. (1987): *Incompleteness, Nonlocality, and Realism*. Oxford: Clarendon Press.

Redhead, M. (1995): "More Ado About Nothing," *Foundations of Physics*, 25, 123–37.

Rovelli, C. (1997): "Relational Quantum Mechanics," revised and updated as Los Alamos archive paper quant-ph/9609002. Originally published as "Relational Quantum Mechanics," *International Journal of Theoretical Physics*, 35, 1637–78.

Ruetsche, L. (1995): "Measurement Error and the Albert–Loewer Problem," *Foundations of Physics Letters*, 8, 331–48.

Saunders, S. (1988): "The Algebraic Approach to Quantum Field Theory," in Brown and Harre (1988, pp. 149–86).

Saunders, S. (1996): "Relativism," in Clifton (1996, pp. 125–42).

Segal, I. (1967): "Representations of the Canonical Commutation Relations," in F. Lurcat (ed.), *Cargese Lectures in Theoretical Physics*, New York: Gordon and Breach, 107–70.

Shimony, A. (1984a): "Controllable and Uncontrollable Non-Locality," in S. Kamefuchi, P. Fusaichi, and D. Drarmigan (eds.), *Foundations of Quantum Mechanics in Light of New Technology*, Tokyo: Physical Society of Japan, 225–30.

Shimony, A. (1984b): "Contextual Hidden Variable Theories and Bell's Inequalities," *British Journal for the Philosophy of Science*, 35, 25–45.

Stein, H. (1972): "On the Conceptual Structure of Quantum Mechanics," in R. A. Colodny (ed.), *Paradigms and Paradoxes: The Philosophical Challenge of the Quantum Domain*, Pittsburgh: University of Pittsburgh Press, 367–438.

Summers, S. J. (1998): "On the Stone–von Neumann Uniqueness Theorem and its Ramifications," to appear in M. Redei and M. Stolzner (eds.), *John von Neumann and the Foundations of Quantum Mechanics*, forthcoming.

Teller, P. (1995): *An Interpretive Introduction to Quantum Field Theory*. Princeton: Princeton University Press.

Valentini, A. (1991): "Signal Locality, Uncertainty and the Subquantum H-Theorem. I and II," *Physics Letters*, A156, 5–11 and 158, 1–8.

van Fraassen, B. (1989): "The Charybdis of Realism," in Cushing and McMullin (1989, pp. 97–113).

van Fraassen, B. (1991): *Quantum Mechanics*. Oxford: Oxford University Press.

Vink, J. (1993): "Quantum Mechanics in Terms of Discrete Beables," *Physical Review*, A48, 1808–18.

von Neumann, J. (1955 [1932]): *Mathematical Foundations of Quantum Mechanics*. Princeton: Princeton University Press. Translated from *Mathematische Grundlagen der Quantenmechanik* Berlin: Springer.

Wald, R. M. (1994): *Quantum Field Theory in Curved Space–Time and Black Hole Thermodynamics*. Chicago: University of Chicago Press.

Wessels, L. (1997): "The Preparation Problem in Quantum Mechanics," in J. Earman and J. D. Norton (eds.), *The Cosmos of Science*, Pittsburgh: University of Pittsburgh Press, 243–73.

Wheeler, J. and Zurek, W. (eds.) (1983): *Quantum Theory and Measurement*. Princeton: Princeton University Press.

Zurek, W. J. (1982): "Environment Induced Superselection Rules," *Physical Review*, D26: 1862–80.

Chapter 11

Evolution

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Introduction

It has become almost standard practice for philosophers of biology to bracket their writings with a pair of manifestoes. The first manifesto proclaims that the philosophy of science is not just about physics anymore. This is usually accompanied by an argument suggesting that a myopic focus on physics has led the philosophy of science to misrepresent the true nature of science. As we face a new millennium, the time has come to dispense with such proclamations. The philosophy of biology has come into its own and no longer needs to justify its existence. One look at the most recent Philosophy of Science Association conference should be enough to convince anyone of that fact: approximately one fourth of the presentations are in the philosophy of biology. The first manifesto has become manifest.

The second manifesto, on the other hand, often takes the form of a "call to arms" for philosophers to venture into fields of biology outside of evolutionary theory, such as ecology and molecular biology. That this call to arms has been at least partially successful is reflected in the inclusion of an essay in this volume on Developmental and Molecular Biology, distinct from the present essay on Evolution. Thus, I need not apologize, as many have done before me, for focusing exclusively on evolutionary theory. Yet, since many of the debates in the philosophy of evolution overlap and intertwine with those in the philosophy of developmental and molecular biology, it is not entirely possible to separate the issues. In fact, philosophers seldom use the phrase "philosophy of evolution." Philosophy of biology has often *meant* philosophy of evolution. However, perhaps it is time to be more explicit.

The philosophy of evolution considers issues that are both conceptual and empirical, and, consequently, it is practiced by philosophers and scientists alike. The following discussion will reflect that diversity. Some philosophy of evolution has looked to evolutionary theory to answer broader questions in the philosophy of science. For example, a recent volume explores epistemological issues through