

Foundations of quantum mechanics: decoherence and interpretation

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In this paper we review Castagnino's contributions to the foundations of quantum mechanics. First, we recall his work on quantum decoherence in closed systems, and the proposal of a general framework for decoherence from which the phenomenon acquires a conceptually clear meaning. Then, we introduce his contribution to the hard field of the interpretation of quantum mechanics: the modal-Hamiltonian interpretation solves many of the interpretive problems of the theory, and manifests its physical relevance in its application to many traditional models of the practice of physics. In the third part of this work we describe the ontological picture of the quantum world that emerges from the modal-Hamiltonian interpretation, stressing the philosophical step toward a deep understanding of the reference of the theory.

I. INTRODUCTION

Anybody who has been close to Prof. Mario Castagnino, even for a short time, knows that he is an ever-eager spirit: the many different subjects treated in this issue are a clear manifestation of the wide panoply of interests that have moved him during his long academic life. Nevertheless, the present article has a peculiarity with respect to the rest of the papers of the issue: Castagnino should be one of the authors of this work. In fact, since ten years ago he has been actively engaged with the foundations and the philosophy of physics, leading an always increasing research group to which we belong. In this field we have obtained relevant results with a remarkable repercussion.

As Castagnino uses to say, he is a senior physicist but a baby philosopher. However, this fact was not an obstacle to his eager spirit, which has been involved in the foundations of so many different subjects that cannot be addressed in a single article. In the present paper we will confine our attention to Castagnino's contributions to the foundations of quantum

mechanics (QM), in order to review his main results in this area. First, we will recall his work on quantum decoherence in closed systems, and the proposal of a general framework for decoherence from which the phenomenon acquires a conceptually clear meaning. Then, we will introduce his contribution to the hard field of the interpretation of QM: our modal-Hamiltonian interpretation solves many of the interpretive problems of the theory, and manifests its physical relevance in its application to many traditional models used in the practice of physics. In the third part of this work we will describe the ontological picture of the quantum world that emerges from our interpretation; here we will stress our philosophical step toward a deep understanding of the reference of the theory, a move not usual in the contemporary discussions about the interpretation of QM. Finally, we will briefly recall Castagnino's contributions in other areas of the foundations and the philosophy of physics.

II. FOUNDATIONS OF QUANTUM DECOHERENCE

More than a decade ago Castagnino developed, with Roberto Laura, a formalism that explains the limit reached by expectation values in closed quantum systems with continuous spectrum,[1]–[6] and begun to conceive that formalism in terms of decoherence. When, some years later, those works were reanalyzed in the context of our research group, we acknowledged the conceptual relevance and the fruitful perspectives of that work. So, the original proposal was further elaborated from a conceptual viewpoint, and presented in several meetings and papers.[7]–[16] In particular, we were invited by Prof. Fred Kronz, from the University of Texas at Austin, to discuss that new view, and he suggested the name ‘self-induced decoherence’ (SID) in contrast with the orthodox ‘environment-induced decoherence’ (EID) approach.[17]·[18]

In those first works, we presented SID as different from EID, that is, as the way in which decoherence manifests itself in closed systems. However, shortly after we realized that both approaches can be subsumed under a *General Theoretical Framework for Decoherence* (GTFD), which encompasses decoherence in open and closed systems.[19]–[21] According to this framework, decoherence is just a particular case of the general phenomenon of irreversibility in QM.[22]·[23] Since the quantum state $\rho(t)$ follows a unitary evolution, it cannot reach a final equilibrium state for $t \rightarrow \infty$. Therefore, if we want to explain the emergence of non-unitary irreversible evolutions, we must split the whole space \mathcal{O} of all possible ob-

servables into a relevant subspace $\mathcal{O}_R \subset \mathcal{O}$ and an irrelevant subspace. With this strategy we restrict the maximal information about the system: the expectation values $\langle O_R \rangle_{\rho(t)}$ of the observables $O_R \in \mathcal{O}_R$ express that relevant information. Of course, the decision about which observables are to be considered as relevant depends on the particular purposes in each situation; but without this restriction, irreversible evolutions cannot be described. In fact, the different approaches to decoherence always select a set \mathcal{O}_R of relevant observables in terms of which the time behavior of the system is described: gross observables in van Kampen,[24] macroscopic observables of the apparatus in Daneri *et al.*,[25] observables of the open system in EID,[26]–[18] relevant observables in Omnés.[29]:[30]

Once the essential role played by the selection of the relevant observables is clearly understood, decoherence can be explained in three general steps:

1. **First step:** The set \mathcal{O}_R of relevant observables is defined.
2. **Second step:** The expectation value $\langle O_R \rangle_{\rho(t)}$, for any $O_R \in \mathcal{O}_R$, is obtained. This step can be formulated in two different but equivalent ways:
 - $\langle O_R \rangle_{\rho(t)}$ is computed as the expectation value of O_R in the unitarily evolving state $\rho(t)$.
 - A coarse-grained state $\rho_G(t)$ is defined by $\langle O_R \rangle_{\rho(t)} = \langle O_R \rangle_{\rho_G(t)}$ for any $O_R \in \mathcal{O}_R$, and its non-unitary evolution (governed by a master equation) is computed.
3. **Third step:** It is proved that $\langle O_R \rangle_{\rho(t)} = \langle O_R \rangle_{\rho_G(t)}$ reaches a final equilibrium value $\langle O_R \rangle_{\rho_*}$:

$$\lim_{t \rightarrow \infty} \langle O_R \rangle_{\rho(t)} = \lim_{t \rightarrow \infty} \langle O_R \rangle_{\rho_G(t)} = \langle O_R \rangle_{\rho_*} \quad (1)$$

where the final equilibrium state ρ_* is obviously diagonal in its own eigenbasis, which turns out to be the final pointer basis. But the unitarily evolving quantum state $\rho(t)$ of the whole system *has only a weak limit*:

$$W - \lim_{t \rightarrow \infty} \rho(t) = \rho_* \quad (2)$$

This weak limit means that, although the off-diagonal terms of $\rho(t)$ never vanish through the unitary evolution, the system decoheres *from an observational point of view*, that is, from the viewpoint given by any relevant observable $O_R \in \mathcal{O}_R$.

This GTFD allows us to face the conceptual challenges that the EID approach still has to face. One of them comes from the fact that, since the environment may be “external” or “internal”, the EID approach offers no general criterion to decide where to place the “cut” between system and environment. Zurek considers this fact as a shortcoming of his proposal: *“In particular, one issue which has been often taken for granted is looming big, as a foundation of the whole decoherence program. It is the question of what are the ‘systems’ which play such a crucial role in all the discussions of the emergent classicality.”* [31] In order to address this problem, the first step is to realize that the EID relevant observables of the closed system U are those corresponding to the open system S :

$$O_R = O_S \otimes \mathbb{I}_E \in \mathcal{O}_R \subset \mathcal{O} \quad (3)$$

where $O_S \in \mathcal{O}_S$ of S and \mathbb{I}_E is the identity operator in \mathcal{O}_E of E . The reduced density operator $\rho_S(t)$ of S is defined by tracing over the environmental degrees of freedom,

$$\rho_S(t) = Tr_E \rho(t) \quad (4)$$

The EID approach studies the time-evolution of $\rho_S(t)$ governed by an effective master equation; it proves that, under certain definite conditions, $\rho_S(t)$ converges to a stable state ρ_{S^*} : $\rho_S(t) \rightarrow \rho_{S^*}$. But we also know that the expectation value of any $O_R \in \mathcal{O}_R$ in the state $\rho(t)$ of U can be computed as

$$\langle O_R \rangle_{\rho(t)} = Tr (\rho(t)(O_S \otimes \mathbb{I}_E)) = Tr (\rho_S(t) O_S) = \langle O_S \rangle_{\rho_S(t)} \quad (5)$$

Therefore, the convergence of $\rho_S(t)$ to ρ_{S^*} implies the convergence of the expectation values:

$$\langle O_R \rangle_{\rho(t)} = \langle O_S \rangle_{\rho_S(t)} \rightarrow \langle O_S \rangle_{\rho_{S^*}} = \langle O_R \rangle_{\rho_*} \quad (6)$$

where ρ_* is a final diagonal state of U , such that $\rho_{S^*} = Tr_E \rho_*$.

From this new conceptual perspective, we have studied the well-known spin-bath model: a closed system $U = P \cup P_1 \cup \dots \cup P_N = P \cup (\cup_{i=1}^N P_i)$, where (i) P is a spin-1/2 particle represented in the Hilbert space \mathcal{H}_P , and (ii) each P_i is a spin-1/2 particle represented in its Hilbert space \mathcal{H}_i . The Hilbert space of the composite system U is, then,

$$\mathcal{H} = \mathcal{H}_P \otimes \left(\bigotimes_{i=1}^N \mathcal{H}_i \right) \quad (7)$$

If the self-Hamiltonians H_P of P and H_i of P_i are taken to be zero, and there is no interaction among the P_i , then the total Hamiltonian H of the composite system U is given by the interaction between the particle P and each particle P_i . [32] By contrast to the usual presentations, we have studied different decompositions of the whole closed system U into a relevant part and its environment. [33]

a. Decomposition 1: A large environment that produces decoherence. In the typical situation studied by the EID approach, the open system S is the particle P , and the remaining particles P_i play the role of the environment E : $S = P$ and $E = \cup_{i=1}^N P_i$. This decomposition results in the system decoherence when the number of particles in the bath is very large.

b. Decomposition 2: A large environment with no decoherence We can conceive different ways of splitting the whole closed system U . For instance, we can decide to observe a particular particle P_j of what was previously considered the environment, and to consider the remaining particles as the new environment, in such a way that $S = P_j$ and $E = P \cup (\cup_{i=1, i \neq j}^N P_i)$. This decomposition results in that the system does not decohere.

c. Decomposition 3: A small environment that produces decoherence It may be the case that the measuring arrangement “observes” a subset of the particles of the environment, e.g., the p first particles P_j . In this case, the system of interest is composed by p particles, $S = \cup_{i=1}^p P_i$, and the environment is composed by all the remaining particles, $E = P \cup (\cup_{i=p+1}^N P_i)$. This decomposition results in the system decoherence when the number p is very large.

We have also studied a generalization of the spin-bath model, where a whole closed system was split into an open many-spin system and its environment. [34] In this case we studied different partitions of the whole system and identified those for which the selected system does not decohere. As stressed in that work, this might help us to define clusters of particles that can be used to store q-bits.

The results obtained in both cases allowed us to argue that Zurek’s “looming big” problem is actually a pseudo-problem, which is simply dissolved by the fact that the split of a closed quantum system into an open subsystem and its environment is just a way of selecting a particular space of relevant observables of the whole closed system. But since there are many different spaces of relevant observables depending on the observational viewpoint adopted, the same closed system can be decomposed in many different ways: each decomposition represents a decision about which degrees of freedom are relevant and which can be disregarded

in each case. And since there is no privileged or “essential” decomposition, there is no need of an unequivocal criterion for deciding where to place the cut between “the” open system and “the” environment. Summing up, decoherence is a phenomenon *relative* to the relevant observables selected in each particular case. The only essential physical fact is that, among all the observational viewpoints that may be adopted to study a quantum system, some of them determine subspaces of relevant observables for which the system decoheres.

Another conceptual difficulty of the EID approach relies on its definition of the pointer basis. This basis is clearly characterized in measurements situations, where the self-Hamiltonian of the system can be neglected and the evolution is completely dominated by the interaction Hamiltonian. In those cases, the pointer basis is given by the eigenstates of the interaction Hamiltonian.[35] However, there are two further regimes, differing in the relative strength of the system’s self-Hamiltonian and the interaction Hamiltonian, where the pointer basis lacks a general definition.[36] Our present research is directed to the search of a general and precise definition of the pointer basis of decoherence.

III. MODAL-HAMILTONIAN INTERPRETATION OF QUANTUM MECHANICS

Our work on decoherence from a closed-system perspective taught us that the decomposition of the total Hamiltonian has to be studied in detail in each case, in order to know whether the system of interest resulting from the partition decoheres or not under the action of its self-Hamiltonian and the interaction Hamiltonian. Once we acknowledged the central role played by the Hamiltonian in decoherence, the natural further step was to ask ourselves whether it plays the same central role in interpretation. This question led us to formulate our modal-Hamiltonian interpretation (MHI) of QM,[37]–[39] which belongs to the modal family:[40] it is a realist, non-collapse interpretation, according to which the quantum state describes the possible properties of a system but not its actual properties. Here we will only recall its main interpretative postulates.

The first step is to identify the systems that populate the quantum world. By adopting an algebraic perspective, a quantum system is defined as:

Systems postulate (SP): *A quantum system \mathcal{S} is represented by a pair (\mathcal{O}, H) such that (i) \mathcal{O} is a space of self-adjoint operators on a Hilbert space \mathcal{H} , representing the ob-*

servables of the system, (ii) $H \in \mathcal{O}$ is the time-independent Hamiltonian of the system \mathcal{S} , and (iii) if $\rho_0 \in \mathcal{O}'$ (where \mathcal{O}' is the dual space of \mathcal{O}) is the initial state of \mathcal{S} , it evolves according to the Schrödinger equation in its von Neumann version.

Of course, any quantum system can be partitioned in many ways; however, not any partition will lead to parts which are, in turn, quantum systems.[41][42] On this basis, a composite system is defined as:

Composite systems postulate (CSP): *A quantum system represented by $\mathcal{S} : (\mathcal{O}, H)$, with initial state $\rho_0 \in \mathcal{O}'$, is composite when it can be partitioned into two quantum systems $\mathcal{S}^1 : (\mathcal{O}^1, H^1)$ and $\mathcal{S}^2 : (\mathcal{O}^2, H^2)$ such that (i) $\mathcal{O} = \mathcal{O}^1 \otimes \mathcal{O}^2$, and (ii) $H = H^1 \otimes I^2 + I^1 \otimes H^2$, (where I^1 and I^2 are the identity operators in the corresponding tensor product spaces). In this case, the initial states of \mathcal{S}^1 and \mathcal{S}^2 are obtained as the partial traces $\rho_0^1 = Tr_{(2)}\rho_0$ and $\rho_0^2 = Tr_{(1)}\rho_0$; we say that \mathcal{S}^1 and \mathcal{S}^2 are subsystems of the composite system, $\mathcal{S} = \mathcal{S}^1 \cup \mathcal{S}^2$. If the system is not composite, it is elemental.*

Since the contextuality of QM, as implied by the Kochen-Specker theorem,[43] prevents us from consistently assigning actual values to all the observables of a quantum system in a given state, the second step is to identify the *preferred context*, that is, the set of the actual-valued observables of the system. Whereas the different rules of actual-value ascription proposed by previous modal interpretations rely on mathematical properties of the theory, our MHI places an element with a clear physical meaning, the Hamiltonian, at the heart of its rule:

Actualization rule (AR): *Given an elemental quantum system represented by $\mathcal{S} : (\mathcal{O}, H)$, the actual-valued observables of \mathcal{S} are H and all the observables commuting with H and having, at least, the same symmetries as H .*

This preferred context where actualization occurs is independent of time: the actual-valued observables always commute with the Hamiltonian and, therefore, they are constants of motion of the system. In other words, the observables that receive actual values are the same during all the “life” of the quantum system as such –precisely, as a closed system–: there is no need of accounting for the dynamics of the actual properties of the quantum system as in other modal interpretations.[44]

The fact that the Hamiltonian always belongs to the preferred context agrees with the many physical cases where the energy has definite value. The MHI has been applied to several well-known physical situations (hydrogen atom, Zeeman effect, fine struc-

ture, etc.), leading to results consistent with experimental evidence.[37] Moreover, it has proved to be effective for solving the measurement problem, both in its ideal and its non-ideal versions,[37] solving the deep challenges that non-ideal measurements pose to other modal interpretations.[45]·[46] In particular, the MHI distinguishes between reliable and non-reliable non-ideal measurements.[37] Furthermore, in spite of the fact that MHI applies to closed systems, we have proved its compatibility with EID.[47]·[48]

Once the MHI was clearly formulated, our further question was whether it satisfies the Galilean invariance of the theory. In fact, any continuous transformation admits two interpretations. Under the active interpretation, the transformation corresponds to a change from one system to another –transformed– system; under the passive interpretation, the transformation consists in a change of the viewpoint –reference frame– from which the system is described.[49] Nevertheless, in both cases the validity of a group of symmetry transformations expresses the fact that the identity and the behavior of the system are not altered by the application of the transformations: in the active interpretation language, the original and the transformed systems are equivalent; in the passive interpretation language, the original and the transformed reference frames are equivalent. Then, any realist interpretation should agree with that physical fact: the rule of actual-value ascription should select a set of actual-valued observables that remains unaltered under the transformations. Since the Casimir operators of the central-extended Galilei group are invariant under all the transformations of the group, one can reasonably expect that those Casimir operators belong to the preferred context.

As we have seen, the preferred context selected by AR only depends on the Hamiltonian of the system. Then, the requirement of invariance of the preferred context under the Galilei transformations is directly fulfilled when the Hamiltonian is invariant, that is, in the case of time-displacement, space-displacement and space-rotation:

$$H' = e^{iH\tau} H e^{-iH\tau} = H \quad (\text{since } [H, H] = 0) \quad (8)$$

$$H' = e^{iP_i r_i} H e^{-iP_i r_i} = H \quad (\text{since } [P_i, H] = 0) \quad (9)$$

$$H' = e^{iJ_i \theta_i} H e^{-iJ_i \theta_i} = H \quad (\text{since } [J_i, H] = 0) \quad (10)$$

However, it is not clear that the requirement completely holds, since the Hamiltonian is not invariant under Galilei-boosts. In fact, under a Galilei-boost corresponding to a velocity u_x ,

H changes as

$$H' = e^{iK_x^{(G)}u_x} H e^{-iK_x^{(G)}u_x} \neq H \quad (\text{since } [K_x^{(G)}, H] = iP_x \neq 0) \quad (11)$$

Nevertheless, when space is homogeneous and isotropic, a Galilei-boost only introduces a change in the subsystem that carries the kinetic energy of translation: the internal energy W remains unaltered under the transformation. This should not sound surprising to the extent that W –multiplied by the scalar mass m – is a Casimir operator of the central-extended Galilei group. On this basis, we can reformulate AR in an explicit Galilei-invariant form in terms of the Casimir operators of the central-extended group:

Actualization rule' (AR'): *Given a quantum system free from external fields and represented by $\mathcal{S} : (\mathcal{O}, H)$, its actual-valued observables are the observables C_i^G represented by the Casimir operators of the central-extended Galilei group in the corresponding irreducible representation, and all the observables commuting with the C_i^G and having, at least, the same symmetries as the C_i^G .*

Since the observables C_i^G –in the reference frame of the center of mass– are M , mW and m^2S^2 , this new version AR' is in agreement with the original AR when applied to a system free from external fields:[50]–[52]

- The actual-valuedness of M and S^2 , postulated by AR', follows from AR: these observables commute with H and do not break its symmetries because, in non-relativistic QM, both are multiples of the identity in any irreducible representation.
- The actual-valuedness of W might seem to be in conflict with AR because W is not the Hamiltonian: whereas W is Galilei-invariant, H changes under the action of a Galilei-boost. However, this is not a real obstacle because a Galilei-boost transformation only introduces a change in the subsystem that carries the kinetic energy of translation, which can be considered a mere shift in an energy defined up to a constant.[50]·[52]

Summing up, the application of AR' leads to reasonable results, since the actual-valued observables turn out to be invariant and, therefore, objective magnitudes. The assumption of a strong link between invariance and objectivity is rooted in a natural idea: what is objective should not depend on the particular perspective used for the description; or, in group-theoretical terms, what is objective according to a theory is what is invariant under

the symmetry group of the theory. This idea is not new: it was widely discussed in the context of special and general relativity with respect to the ontological status of space and time,[53] and since then it reappeared in several works. [54]–[59] From this perspective, AR says that the observables that acquire actual values are those representing objective magnitudes. On the other hand, from any realist viewpoint, the fact that certain observables acquire an actual value is an objective fact in the behavior of the system; therefore, the set of actual-valued observables selected by a realist interpretation must be also Galilean-invariant. But the Galilean-invariant observables are always functions of the Casimir operators of the Galilean group. As a consequence, one is led to the conclusion that any realist interpretation that intends to preserve the objectivity of actualization may not stand very far from the modal-Hamiltonian interpretation.

When AR is expressed in simple group terms, one can expect that it can be extrapolated to any quantum theory endowed with a symmetry group. In particular, the actual-valued observables of a system in quantum field theory would be those represented by the Casimir operators of the Poincaré group and of the internal symmetry group. On this basis, in a recent paper we presented an alternative version of the non-relativistic limit of the centrally extended Poincaré group and its consequences for interpretive problems.[60]

As it is well known, the Galilei group can be recovered from the Poincaré group by means of Inönü-Wigner contraction.[61] It is therefore natural to ask whether such a situation can be generalized to the central-extended Galilei group, which is the relevant group in QM. However, since the Poincaré group does not admit nontrivial central extensions,[62] we have to define a generalized Inönü-Wigner contraction from a trivial extension of the Poincaré group whose generators are H , P_i , J_i and K_{P_i} (where the last ones are the Lorentz boosts). With this purpose, we extend the group trivially, i.e., in such a way that all the generators of Poincaré group commute with a trivial central charge M . The basis of the resulting new algebra $I^M SO(1, 3) = ISO(1, 3) \times \langle M \rangle$ is $\{H, P_i, J_i, K_{P_i}, M\}$. Then, we perform the following change of the generators basis $\overline{H} = H - M$. In the new basis $\{\overline{H}, P_i, J_i, K_{P_i}, M\}$ all the commutators of the Poincaré group remain the same, with the only exception of

$$[P_i, K_{P_j}] = -i\delta_{ij}H = -i\delta_{ij}(\overline{H} + M) \quad (12)$$

The contraction is determined by the rescaling transformations (in the basis

$\{\overline{H}, P_i, J_i, K_{P_i}, M\}$) defined by

$$J'_i = J_i, \quad P'_i = \varepsilon P_i, \quad K'_{P_i} = \varepsilon K_{P_i}, \quad \overline{H}' = \overline{H}, \quad M' = \varepsilon^2 M \quad (13)$$

The space isotropy remains unchanged by this rescaling transformation and

$$\left[P'_i, K'_{P_j} \right] = -i\delta_{ij}(\varepsilon^2 \overline{H}' + M') \quad \text{so} \quad \lim_{\varepsilon \rightarrow 0} \left[P'_i, K'_{P_j} \right] = -i\delta_{ij} M' \quad (14)$$

Therefore, it turns out to be clear that the contracted algebra is isomorphic to the extension of the Galilei algebra. On the basis of this result, we have also proved that the Casimir operators of the trivially extended Poincaré group contract naturally to the Casimir operators of the extended Galilei group.[60]

Summing up, when AR is expressed in its explicit Galilei-invariant form AR', it leads to a physically reasonable result: the actual-valued observables are those represented by the Casimir operators of the mass central-extended Galilei group. The natural strategy is to extrapolate the interpretation to the relativistic realm by replacing the Galilei group with the Poincaré group. But when one takes into account that the relevant group of non-relativistic QM is not the Galilei group but its central extension, the mere replacement of the relevant group is not sufficient: one has to show also that the actual-valued observables in the relativistic and the non-relativistic cases are related through the adequate limit. As a consequence, the Poincaré group has to be trivially extended, in order to show that the limit between the corresponding Casimir operators holds, and this result counts in favor of the proposed extrapolation of our MHI to non-relativistic QM. Furthermore, this result is physically reasonable because mass and spin are properties supposed to be always possessed by any elemental particle,[63] and they are two of the properties that contribute to the classification of elemental particles. At present we are working on a further extrapolation of the MHI to the standard model.

IV. THE ONTOLOGICAL PICTURE OF THE QUANTUM WORLD

In general, the discussions about modal interpretations are concerned with the traditional problems, as the measurement problem and the no-go theorems. But these are not the only relevant issues: one should not forget the ontological question about the structure of the world referred to by QM.

All modal interpretations rely on a common assumption: QM does not describe what is the case, but rather what may be the case. The problem of the nature of possibility is as old as philosophy itself. Since Aristotle’s time to nowadays, however, two general conceptions can be identified. On the one hand, *actualism* reduces possibility to actuality. This was the position of Diodorus Cronus, who defined “*the possible as that which either is or will be*”. [64] This view survived up to 20th century; for instance, for Russell ‘possible’ means ‘sometimes’, whereas ‘necessary’ means ‘always’.[65] On the other hand, *possibilism* conceives possibility as an ontologically irreducible feature of reality. From this perspective, the stoic Crissipus defined possible as “*that which is not prevented by anything from happening even if it does not happen*”. [66] In present day metaphysics, the debate actualism-possibilism is still alive. For the actualists, the adjective ‘actual’ is redundant: non-actual possible items (objects, properties, facts, etc.) do not exist. According to the possibilists, on the contrary, not every possible item is an actual item: possible items—*possibilia*—constitute a basic ontological category.[67]

For our MHI, probabilities measure ontological propensities, which embody a possibilist, non-actualist possibility: a possible fact does not need to become actual to be real. This possibility is defined by the postulates of QM and is not reducible to actuality. This means that reality spreads out in two realms, the realm of possibility and the realm of actuality. In Aristotelian terms, being can be said in different ways: as possible being or as actual being, and none of them is reducible to the other. Moreover, the *ontological structure of the realm of possibility* is embodied in the definition of the elemental quantum system $\mathcal{S} : (\mathcal{O}, H)$, with its initial state ρ_0 : (i) the space of observables \mathcal{O} identifies all the *possible type-properties* (observables) with their corresponding *possible case-properties* (eigenvalues), and (ii) the initial state ρ_0 codifies the measures of the *propensities to actualization* of all the possible case-properties at the initial time, propensities that evolve deterministically according to the Schrödinger equation.

The fact that propensities belong to the realm of possibility does not mean that they do not have physical consequences in the realm of actuality. On the contrary, propensities produce definite effects on actual reality even if they never become actual. An interesting manifestation of such effectiveness is the case of the so-called “*non-interacting experiments*”, [68] [69] where non-actualized possibilities can be used in practice, for instance, to test bombs without exploding them.[70] This shows that possibility is a way in which reality

manifests itself, a way independent of and not less real than actuality.

One of the main areas of controversy in contemporary metaphysics is the problem of the nature of individual objects: is an individual a substratum supporting properties or a mere “bundle” of properties?[71] The idea of a substratum acting as a bearer of properties has pervaded the history of philosophy: it is present under different forms in Aristotle’s “primary substance”, in Locke’s doctrine of “substance in general” or in Leibniz’s monads. Nevertheless, many philosophers belonging to the empiricist tradition, from Hume to Russell, Ayer and Goodman, have considered the posit of a characterless substratum as a metaphysical abuse. As a consequence, they adopted some version of the *bundle theory*, according to which an individual is nothing but a bundle of properties: properties have metaphysical priority over individuals and, therefore, they are the fundamental items of the ontology.

In the Hilbert space formalism, states have logical priority over observables since observables apply to states. This logical priority favors the picture of an ontology of substances and properties, with the traditional priority of substances over properties. Our MHI, on the contrary, is based on the algebraic formalism, where the basic elements are observables and states are functionals over the space of observables. Then, the MHI favors the bundle theory, that is, an ontology of properties, where the category of substance is absent.

According to the traditional versions of the bundle theory, an individual is the convergence of certain case-properties, under the assumption that all the type-properties are determined in the actual realm. For instance, a particular billiard ball is the convergence of a definite value of position, a definite shape, say round, a definite color, say white, etc. Then, the properties taken into account are always actual properties: bundle theories identify individuals with bundles of actual properties. In QM, on the contrary, the Kochen-Specker theorem prevents the assignment of case-properties (eigenvalues) to all the type-properties (observables) of the system in a non-contradictory manner. Therefore, the classical idea of a bundle of actual properties does not work for the quantum ontology.

If, from the perspective of the MHI, the quantum world unfolds into two irreducible realms, the realm of possibility has to be taken into account when deciding what kind of properties constitutes the quantum bundle. Since the quantum system is identified by its space of observables, which represent possible properties, an individual quantum system turns out to be a *bundle of possible properties*: it inhabits the realm of possibility, which is as real as the realm of actuality.

This interpretation of quantum individual systems has the advantage of being immune to the challenge represented by the Kochen-Specker theorem, since this theorem imposes no restriction on possibilities. Moreover, it seems reasonable to expect that this conception of individual supplies the basis for solving the problem of the indistinguishability of “identical particles”, [72] introduced in the formalism as an *ad hoc* restriction on the set of states. At present, we are working on this problem: if the traditional assumption of substantial objects, which preserve their individuality when considered in collections, is the main obstacle to explain quantum statistics, the conception of the quantum system as a bundle of possible properties seems to offer a promising starting point in the search for a solution of the problem.

Summing up, from our interpretational perspective, the talk of individual entities as electrons or photons and their interactions can be retained only in a metaphorical sense. In fact, in the quantum framework even the number of particles is represented by an observable N , which is subject to the same theoretical constraints as any other observable of the system; this leads, specially in quantum field theory, to the possibility of states that are superpositions of different particle numbers. [73] Therefore, the number of particles N has an actual definite value only in some cases, but it is indefinite in others. This fact, puzzling from an ontology populated by substantial objects, is deprived of mystery when viewed from our ontological perspective. The quantum system is not a substantial individual, but a bundle of possible properties. The particle picture, with a definite number of particles, is only a contextual picture valid exclusively when the observable N is picked out by the preferred context. In this case, we could metaphorically retain the idea of a composite system composed of individual particles that interact to each other. But in the remaining cases, this idea proves to be completely inadequate, even in a metaphorical sense.

V. CONCLUDING REMARKS

We hope that this journey through the main contributions of Castagnino in the field of the foundations of QM supplies an idea of the active work that he and his research group are developing. Nevertheless, we do not want to finish this review without recalling the rest of the areas of the philosophy of physics where he has fruitfully produced: time’s arrow, [74]–[81] time-asymmetric QM, [82], [83] quantum chaos, [84], [85] and even philosophy

of chemistry.[87] Not bad for a baby philosopher. However, this is not surprising when coming from an even-eager spirit as Mario Castagnino.

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