Quantum Observables and a Model of Noncommutativity

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This paper considers a generalization of the notion of quantum observables in ontological models of quantum mechanics. Within this framework it is possible to construct physical models where quantum noncommutativity can arise dynamically. Unlike quantum systems, the basic entities in this model have definite properties. Relations with no-go theorems and other hidden variable theories are also discussed.

I. INTRODUCTION

Quantum mechanics, while being immensely successful, is also very hard to understand. The hidden variable theories (HVT) are attempts to render it more understandable, by holding fast to classical doctrines like determinism and locality. The price to be paid is that, according to these theories, quantum theory is no longer a complete description of the world. However, several no-go theorems seem to show that at least some of these classical concepts have to be abandoned.

First, there is Kochen-Specker theorem [1], [2]. The theorem implies that quantum observables (for a system with Hilbert space dimension greater than two) cannot all possess definite premeasurement values that are faithfully revealed by measurements. One way for an ontological model to avoid the theorem, while maintaining determinism, is to relax the condition that the measured values are preassigned. An example is the contextual hidden variable theory, where the measured values is determined by the complete (hidden) state of the system together with a specific context of measurement. Another way to avoid the theorem is to allow the measurement results to be probabilistic even when complete state of system is known, i.e. a stochastic hidden variable theory. However it is difficult for such theory to account for the correlations shown in entanglement, and Bell’s theorem [3] has shown that they cannot do so without some kind of nonlocality, albeit in the form of probabilistic conditions.

The Kochen-Specker theorem is a consequence of trying to assign measurement values in a certain way to the noncommuting set of observables. The noncommutativity of physical observables is a nonclassical feature of quantum theories, which is arguably the source of the major, if not all, quantum weirdness. From it arises the possibility of superposition of quantum states, and in turn the entanglement [4]. It seems reasonable then to ask where it comes from, and why it is inevitable in the microscopic world. However all the hidden variable theories either takes this for granted, or view it as not demanding any explanation or as an unexplainable brute fact. This paper views this as an unsatisfying state of affair.

Apart from a demand for deeper understanding of quantum theory, the author also views the indefinite properties of quantum systems, which is a consequence of noncommutativity, as particularly worrying. If the quantum theory is taken seriously (i.e. assuming its completeness), quantum system as a physical system is not something that is ontologically well defined: there are no observables that takes definite values at all times. This actually contradicts a long standing fundamental assumption about the world: if something exists, it possesses a set of definite attributes, independently from and prior to any observation, and these attributes completely specifies it’s state. This is such a deep assumption about nature that it virtually remains unchallenged throughout the whole history of science (until quantum theory). This is, of course, not to say that it is a priori true, but that it does seem highly unlikely that it will be challenged in the regime of atoms and electrons, and in a theory that was invented more or less for the purpose of calculating measurement probabilities.

This paper is a proposal for some possible reasons for the apparent inevitability of noncommutativity, in particular, we will look for the reason in dynamics, somewhat analogous to how equilibrium arises in statistical mechanics [5]. Meanwhile in doing so ontological definiteness and determinism are both recovered. To achieve this, we need to revise our view of quantum observables as fundamental and irreducible. This and a discussion of its thermodynamic analogy is discussed in the next section.

Section III discusses the issue of noncommutativity in QM, and how it becomes explainable within the view of quantum observables taken in this paper.

The relations of the proposed model with the no-go theorems described just now are given in section IV, followed by a comparison of this model to other hidden variable theories in the final section.

II. QUANTUM THEORY AS THERMODYNAMICS AND HIDDEN VARIABLE THEORY

Systems in pure quantum states exhibits a kind of objectiveness and ‘rigidity’, which are among the reasons that give rise to appearance that QM is complete and irreducibly indeterministic. The objectiveness stems from the fact that any pure states can always be precisely pro-
duced, via some objectively identified apparatus and procedure. The preparation procedure determines the state of the object. Also, any given pure state can be transformed into any other pure states via some objective and unambiguous procedures. Most importantly, for any pure state, there exist yes/no measurements that will certainly give a “yes” result, and these measurements are solely determined by the facts about the preparation procedure. ‘Rigidity’ means that the information contained in a pure state description is fixed (and is maximal). When the system is in a known pure state, it is not possible that, for example, more knowledge that is not already contained in the pure state is somehow obtained, even by accident. It is also not possible to prepare a quantum object in a more refined state than a pure state.

All these aspects are actually similar to classical macroscopic systems in thermodynamical equilibrium states, when one is confined to the macroscopic thermodynamic variables (e.g. U, V) \[\mathbb{R}\]. In general, when one can manipulate certain macroscopic variables \{A_i\}, there corresponds a unique state (probability distribution over the phase space), which minimally contains the information he can ever obtain by operating at the level of \{A_i\}. This is the equilibrium state that corresponds to these variables. Manipulations over the variables \{A_i\}, for example by changing the values of or preparing certain values for these variables, are objective processes. The ‘rigidity’ of such states arises from the following fact: if the control is entirely within a set of variables, the object cannot be systematically prepared in a state with more refined distributions (even if it did for sometime, one is unable to know it).

This similarity suggests one to view pure quantum states as nothing but equilibrium states corresponding to certain variables that we can control, and which are macroscopic relative to some underlying subquantum variables. In this view, a complete set of quantum observables corresponds to the set of thermodynamic variables that uniquely characterizes the probability distribution. These thermodynamic variables can either be the average values of the extensive variables or, equivalently, the corresponding intensive parameters.

This conception of quantum observables is quite different from that in the usual hidden variable theories. In the latter the observables are taken to be random variables over the space of system’s complete (hidden) states, \(\Lambda\). For deterministic hidden variable theory, the result of measuring observable \(A\) is \(v(A) = A(\lambda)\), entirely determined by the complete state \(\lambda \in \Lambda\); while for stochastic hidden variable theory one can have at most the probability of obtaining the value \(v(A)\) in state \(\lambda\), \(P(v(A)|A, \lambda)\), which is not necessarily 0 or 1.

However by taking the observables to be intrinsic variables, as in this paper, means to take them to correspond instead to certain probability distributions over system’s state space, i.e. \(v(A) = A[p(\lambda)]\). The domain of \(A\) here is now (subset of) the space of distributions over \(\Lambda\), not \(\Lambda\) itself.

When the distribution \(p(\lambda)\) is highly peaked (small standard deviation), measurements of \(A\) made on systems that prepared according to this distribution (which corresponds to the same preparation procedure, i.e. by preparing values of \(A\) to be \(v(A)\)) will almost always give the same result. We can say, as in thermodynamics, that each of the system that lies in the finite region where the distribution is peaked possesses a definite value for the variable \(A\). However the same cannot be said for states that lies outside this region. Besides, when the distribution is not peaked then identical preparation procedures (as mentioned above) will have non-negligible chance of giving rise to different \(A\)-measurement values. In this situation we cannot say that the system possess definite value for the observable, whichever the (complete) state the system is in.

Therefore, in contrast to hidden variable theories, here it is illegitimate to say that for some state \(\lambda\) of the system observable \(A\) takes the value \(v(A)\), or the probability of being so is such and such.

Moreover, there is also the crucial difference regarding the role of explanation. The usual hidden variable theories only models the statistical results of quantum theory, but does not explain why it is so, e.g. it does not give physical reasons as to why classical probability theory is not applicable, why the use of incompatible observables seems inevitable etc.

### III. NONCOMMUTATIVITY

Viewing quantum observables as intensive variables allows for the possibility of constructing physical models where quantum noncommutativity might emerge dynamically. However, any attempt to explain noncommutativity must first supply it with an interpretation. This paper views noncommuting variables as variables that cannot be simultaneously well-defined for all states of the system, where a physical variable is said to be well-defined in case it possesses a definite value for one single system.

In classical thermodynamics there are two situations where the usual thermodynamic intensive variables might become ill-defined. This is when the system is in a nonequilibrium state or when it is small (which correspond respectively to the two cases mentioned in the last section). We will use these to propose two kinds of physical models, as shown below (however the emphasis of this paper will be put on the former situation). In these models noncommutativity is not fundamental, but arises from some mechanisms of underlying physical entities that possess only well-defined attributes.

However since the idea of such physical entities are quite speculative, the method taken in this paper is to
apply the concept of thermodynamic (non)equilibrium on their collective behavior. In doing so we assume the universality of these concepts (which guarantees their applicability in this regime) and then hypothesize on the behavior of underlying entities, under the condition that such model should recover quantum properties. It is therefore important to note that at this stage we are aiming to exhibit the logical possibility of reducing noncommutativity to physical models that satisfy certain intelligible (classical) requirements. Real physical possibility of such models will be left to future papers.

A. Model of Noncommutativity: Small Systems

In this model, all quantum system is assumed to be somehow composed of a (presumably large) fixed number $N$ of subquantum objects or elements (which will be called SQE in this paper). There are many different (possibly continually many) kinds of SQE, and different kinds of SQE are allowed to change to one another. Each kind of SQE would correspond to a complete sets of quantum observables, which is actually a set of intensive parameters describing the (macro)state of SQE.

Consider the simplistic case where there’s only two noncommutative quantum observables $A$ and $B$ for this system, hence it is composed of two kinds of SQE, denoted as SQE$_a$ and SQE$_b$. Corresponding to each of them are intensive parameters $a$ and $b$, respectively. These parameters describe the states of the two different kinds of SQE. Let the number of each of them be $N_a$ and $N_b$, respectively, then $N_a + N_b = N$ and $N$ is fixed at all times. Now as in classical thermodynamics, these parameters are well defined (in the sense described at the beginning of this section III) only when the number of entity is large (see end of section II), and their well-definability is usually quantified as the relative standard deviations of their corresponding extensive parameters, $\frac{\Delta A}{A}$ and $\frac{\Delta B}{B}$, which are proportional to $\frac{1}{\sqrt{N_a}}$ and $\frac{1}{\sqrt{N_b}}$ respectively.

Thus since total $N$ is fixed, if the number of one kind of SQE increases, the other will become less. This therefore captures the intuition that when one variable is more well-defined, the other one becomes less so. All this can be made more precise by casting the relation in a Heisenberg-like inequality, as shown below.

First it is a consequence of Schwartz inequality and $N_a + N_b = N$ that

$$\frac{\Delta A \Delta B}{A B} = k_a k_b \frac{ab}{\sqrt{N_a N_b}} \geq \frac{2}{N} k_a k_b ab$$

(1)

where $k_a, k_b$ are constants for each kind of the SQE’s. We then assume that although as one kind of SQE increases the other kind becomes less, both $N_a$ and $N_b$ are still large enough such that the distribution for both is still concentrated in a small region, and thus that it be taken as nearly uniform in the region, and that within this region both $A$ and $a$ ($B$ and $b$) are approximately linear. Then we will have $\Delta A \propto A$ with proportionality coefficient $(\frac{\partial A}{\partial x})|_{x_0}/(\frac{\partial x}{\partial x})|_{x_0}$, where $x$ is a point in the (complete) state space $A$ of the system, and $x_0$ is the center of the small region.

By letting $l_a = (\frac{\Delta A}{A})|_{x_0}/(\frac{\partial A}{\partial x})|_{x_0}$ (similarly for $l_b$), we then have

$$\frac{\Delta a \Delta b}{a b} = \frac{l_a l_b}{a b} \Delta A \Delta B \geq \frac{2}{N} j_a j_b AB$$

(2)

where $j_a$ is defined to be $l_a k_a$ (and similarly for $j_b$).

Thus if for some such system with the value $\frac{1}{N} ab j_a j_b AB$ is of the same order of magnitude as $\hbar$ we obtain the inequality

$$\Delta a \Delta b \geq \frac{\hbar}{2}$$

(3)

as an approximation (the inequality is, however, a strict consequence for quantum observables $\hat{A}, \hat{B}$ that satisfy $[\hat{A}, \hat{B}] = i \hbar$). Note that it is not required that the value of $\frac{1}{N} ab j_a j_b AB$ to be equal to $\hbar$, but only that it approximately so. This seems possible because $N$ is huge, and thus (2) will tend towards (3) when $N$ is suitably large. In this sense we have Heisenberg-like inequality as an approximation for the relations between $\Delta a$ and $\Delta b$.

From this simple consideration it is found that a Heisenberg-like inequality can be obtained for the variables that we are able to control, under suitable approximations. One possible implication of such derivation is that the number $N$ is now found to be related to Planck’s constant $\hbar$, so it can be estimated if the range of the values of $A$, $B$ is known.

B. Model of Noncommutativity: Nonequilibrium States

1. Ontology

In this model, as in the previous model, the quantum system is also composed of a large number ($N$) of localized entities SQE, however here there is only one kind of SQE for any quantum system, i.e. all the SQE are the same. The key feature of this model is that here each SQE is assumed to take a possibly continuous set of different and independent extensive variables as their properties. For an SQE with index $i$ ($i$ ranges from 1 to $N$), we denote its extensive variables as $A_i(\alpha)$, where $\alpha$ is the parameter ($\alpha$ space is in general a manifold.). Therefore, definite values of $A_i(\alpha)$ for all $\alpha$ fully specifies an SQE.
The SQEs interact in such a way that coupling can exist only among extensive variables of the same kind (i.e., same \( \alpha \)), and the coupling within a set \( \{ A_i(\alpha) \} \), \( g(\alpha) \), tends to bring it towards equilibrium, i.e., a state that can completely be described by a (fixed) ensemble average of the total sum \( \sum_{i=1}^{N} A_i(\alpha) \) at the level of such observables. The couplings is assumed to be local, i.e., they describe contact actions.

Now, the connection to quantum system is made by the assumption that the intensive parameters that are uniquely determined by such equilibrium (the large number \( N \) justifies the use of such intensive parameter) actually corresponds to a quantum observable, and this equilibrium state corresponds to an eigenstate of the quantum observable. As an example, for the spin observable \( \hat{S}_i \), the corresponding set of extensive variables is parameterized by spatial direction \( \vec{n} \), and the equilibrium obtained by such extensive variables corresponds to an eigenstate of the spin observable.

But to reproduce noncommutativity (interpreted as above), some constraints needed to be imposed upon the magnitude of these couplings,

\[
C[g(\alpha)] = 0
\]

with the effect of ensuring that not all sets of extensive variable can achieve equilibrium within some suitable interval of time. Thus although the extensive variables \( A_i(\alpha) \) are independent, their time evolutions are not, and are dependent on each other via the condition. While this condition in general gives a noncommutative theory, to obtain specifically the Hilbert space structure the functional constraint \( C \) should furthermore possess certain properties (we here denote \( g_{\alpha_0}(\alpha) \) as value of coupling for observables \( \{ A_i(\alpha) \} \) when \( \{ A_i(\alpha_0) \} \) is in equilibrium):

\[
\exists \alpha\text{-independent functional } F, \text{ such that }
F[g_{\alpha_0}(\alpha_0), g_{\alpha_0}(\alpha)] \propto P(\alpha, m'|\alpha_0, m)
\]

\[
\quad = |\langle \alpha, m'|\alpha_0, m \rangle|^2
\]

for any \( m \), where \( \{|m\}\} \) and \( \{|m'\}\} \) are the basis states for the two observables corresponding to the parameters \( \alpha_0 \) and \( \alpha \) respectively, and \( |m| \) is the one in the latter basis states that is nearest to \( |m'\rangle \). Equation (5) above is Born’s rule, thus if such functional \( F \) can be found then Born’s rule can be seen as just a codification of the constraint in the Hilbert space framework. Also note that since a function \( g(\alpha) \) with a maximum at \( \alpha = \alpha_0 \) will give rise to an equilibrium state of the variables \( \{ A_i(\alpha_0) \}_{i=1}^{N} \), all such functions that satisfy the above constraints corresponds to the same eigenstate of the equilibrium observable.

In this model, in order to allow for the possibility of quantum correlation, we will need to make an important assumption about space (or vacuum): The ‘empty’ space consists of a vast amount of discrete entities that can interact with the the quantum system’s SQE. Such entities interact by contact action (possibly similar to classical particle interactions) and they are moving in a Newtonian space-time (i.e., no upper limit to their velocities) [13].

2. Dynamics

(i) ‘Measurement’ Process

We will discuss only the case of ideal measurement of observable \( \hat{A}(\alpha_0) \) that gives a definite result \( m(\alpha_0) \) and leaves the system in a pure state \( |m(\alpha_0)\rangle \), where \( m(\alpha_0) \) is an eigenvalue of \( \hat{A}(\alpha_0) \). In this model, quantum measurement process is essentially a (deterministic) local interaction between the measurement apparatus, quantum object and space. Here a measurement apparatus is one which effectively changes the couplings of the object under measurement (with the same constraints \( F \) being satisfied). We say it is a measurement of observable \( \hat{A}(\alpha_0) \) if the new couplings has a maximum at \( g(\alpha_0) \). We also requires the state of the apparatus output reader to be perfectly correlated to the final equilibrium state.

The measurement process is as follows: the interaction between object, apparatus and space in general causes the object to evolve out of its equilibrium state, but the new coupling then allows the system to relax to an equilibrium of the corresponding observable, the value of which is completely determined by the initial states of the three parties. The output reader of the apparatus will then show a reading that is correlated to the final equilibrium state.

Denoting respectively the complete states of apparatus, system and space at time \( \lambda_{\text{M}}(t) \), \( \lambda_{\text{sys}}(t) \) and \( \lambda_{\text{sp}}(t) \), the process is in general of the form:

\[
\lambda_{\text{sys}}(\Delta t) = f(\lambda_{\text{M}}(0), \lambda_{\text{sys}}(0), \lambda_{\text{sp}}(0); \Delta t)
\]

where \( \Delta t \) is the amount of time taken by the measurement interaction that started from \( t = 0 \) and function \( f \) is deterministic. At the end of such interaction, the measurement apparatus is assumed to measure the time average of \( \sum_{i=1}^{N} A_i(\alpha) \), which is the same as the ensemble average \( \langle \sum_{i=1}^{N} A_i(\alpha) \rangle \) because the quantum object is at equilibrium now.

Therefore in our model measurement is not a process that reveals the value of any properties of the quantum object, it is instead a process that forces the object to conform to some properties of the measurement apparatus. Here the randomness of measurement results (which is the reason why we say that quantum mechanics is indeterministic) is a result of our ignorance of or our inability to control the fundamental entities of space, system and apparatus.

Also note that this mechanism allows the possibility of obtaining quantum discreteness, and thus the finiteness of Hilbert space dimension of certain systems, from underlying continuous variables: what is needed is to find
dynamics such that the average of the extensive values belongs to a discrete set of values.

(ii) Unitary Dynamics

Contrary to the usual view where measurement and unitary time evolution are incompatible processes, and that the latter is somehow more ‘fundamental’ than the former, in this model the evolution of pure states is actually a continual series of measurement-like processes. That is, at each moment of such evolution is actually a measurement process whose end results are equilibrium states. Each such states differ only infinitesimally from the previous one [16], in order to ensure that there will be no discontinuous state jump during time evolution.

However the time interval between any two consecutive of these measurement-like processes must be much larger than the relaxation time of the equilibrium states: \( \tau_{\text{relax}} \ll \delta t \), so that at each instant the system can be legitimately considered as being in a pure state.

The reason for adopting such an unconventional view of time evolution is because that a pure state is a stable state (that will not change unless there are interactions with other objects), and that in our model all the interactions are mediated by contact/local interactions via space’s entities. Since we have already seen that ‘measurement’ is a process that couples the system to space and resulting in a state change, for simplicity sake [17], this paper assumes that all possible changes in pure state is due to interactions with space’s entities.

From this discussion it follows that, in the nonequilibrium model, unitary time evolution in quantum theory is not an exact law. It is applicable only in the regime where the relevant physical processes involve time scales much larger than \( \tau_{\text{relax}} \), and where the change in \( \alpha \) is very small for any instant of time with such scale. If this is not satisfied then the system cannot even be described by a pure state, let alone a unitary evolution. As will be seen in section IV.B below, such the system is in a situation similar to an improper mixture [18].

IV. NO-GO THEOREMS

A. On Kochen-Specker Theorem

The Kochen-Specker theorem is avoided in the model, since here a single pure quantum state is generally in a nonequilibrium state for variables that are not the one that corresponds to the pure state eigenvalue. If the nonequilibrium state of the model can be taken as one with local equilibrium, then the system can be visualized as a spatially extended region consisting of many different smaller parts in local equilibrium, each part possesses a definite value of intensive variables (i.e. the eigenvalues). It is natural then to take the fractional volumes to correspond to the quantum probabilities.

Now, in this picture it is then clear that a single (pure state) quantum system possesses all the values of all quantum observables, other than the one of which the system’s state is an eigenstate. It is in this sense that noncommutativity is realized, for all other observables are not well-defined for this system.

Therefore this model is not contextual, the measurement results are stochastic, even when the exact state of the quantum system is known.

B. On Entanglement, Nonlocality and Bell’s Theorem

The model is nonlocal, but the nonlocality is due to the arbitrarily fast propagation of space’s constituent entities, which interacts with one another by local interactions. There is no spontaneous action at a distance.

Let’s see how this model accounts for entanglement. We take that the preparation of an entangled pair to be an ideal measurement followed by a filtering process, and we consider the example of singlet state, which is an eigenstate of \( S_z^{1} \otimes S_z^{2} \) with eigenvalue -1. As in the discussion of section III.B.1, a pure state is one where certain observables reaches equilibrium. It is then possible to have situations like this: for any SQE with index \( i \), there is exactly another one with index \( i' \), such that for any \( \alpha \), the observable \( A_i(\alpha) + A_{i'}(\alpha) \) is a constant, for all \( i \). Then imagine that the quantum object is divided into two in such a way that every SQE and its counterpart is not in the same half. Now any of the halves generally will not be in an equilibrium of any of the observables \( A_i(\alpha) \) (i.e. the average value \( \langle A_i(\alpha) \rangle \) is insufficient to represent the state). Thus this is a situation where the quantum object do not possess any properties that is in equilibrium.

However the observable \( \{ A_i(\alpha) + A_{i'}(\alpha) | i = 1,...,N/2 \} \) is in an equilibrium because it has the same value for all pairs of \( (i, i') \), thus the expectation value \( \langle A_i(\alpha) + A_{i'}(\alpha) \rangle \) (and the corresponding intensive variable) can completely describe the state at this level. Thus in this way the entanglement is explained within the model.

Now, if a measurement is performed on one particle and obtained a certain result (say, spin up in \( S_z \)), this will result in a change of surrounding space’s state. The new local equilibrium then spreads with arbitrarily fast speed and will reach the other particle almost instantaneously. This new equilibrium state of space will interact with the second particle, and the interaction is such that the system will relax into an equilibrium state (pure state). This is so because the whole process is just a preparation (or ideal measurement) process at a distance.

This discussion shows that the model violates the con-
dition of outcome independence \[19\], therefore the Bell’s inequality is not derivable for this model.

V. COMPARISON WITH OTHER NONLOCAL HVT

The nonlocal property of this model is quite distinct from that of Bohm’s theory in some important respects. Simply put, this model satisfies two requirements:

(i) No nonlocality that is nonmaterial/nonphysical;

(ii) Probabilistic behavior of quantum systems is due to our ignorance/uncontrollability of the ontic state, therefore wavefunction is not physical/ontic.

However, Bohm’s theory and any possible variations of it cannot satisfy both \[20\]. This is because that in this theory, the wavefunction is a function of the particle system’s configuration space and yet it entirely determines the probability of their positions. Moreover, the role of wavefunction in the theory also seems to make it inherently more nonlocal than the Newtonian physics. Although the latter contains action at a distance (e.g. gravitational force), the general Newtonian framework need not, because its ontology are local (particles and fields).

This seems to be an advantage of the nonequilibrium model over these hidden variable theories, because this model assumes that all the SQE to be local, and the space entities can be assumed to interact only locally. Besides, the main strength of this model is that it provides an account a possible origin of noncommutativity (given an interpretation of it), which is simply assumed from the outset in all of the hidden variable theories.

VI. SUMMARY AND PROSPECT

This paper briefly describes an outline of a model that explains quantum noncommutativity as a phenomenon that emerges from more fundamental (non-quantum) processes, where the most basic entities involved have well defined properties. Besides this, the interesting features of this model are: (a) randomness (of measurement results) is explained as ignorance about the underlying entities and their interactions; (b) space itself is composed of interacting discrete entities, which plays an important role in explaining entanglement; (c) it is a nonlocal model, but its nonlocality is not due to action at a distance, instead it is transmitted through contact action of the space’s entities; (d) measurement in general do not reveal preexisting values, it is a physical process that alters the state of quantum system (and of space); (e) there is no measurement problem: measurement(-like) process does not clash with the unitary dynamics because it actually gives rise to the latter, this implies that (f) unitary time evolution is not an exact law and a pure state can evolve into an improper mixed state (see also note \[18\]); (g) quantum discreteness is compatible with the underlying variables that are continuous, and might in fact emerges from the underlying dynamics too.

Needless to say, this work is just a start. The framework has not yet been cast in a mathematical form and the paper provides no descriptions for many details, also many topics are left out, for example, the properties of space’s entities, their interactions with SQE etc. However the purpose of this paper is try to show that there are logically possible and at the same time more intelligible explanations for many aspects of quantum theory that many deemed as inevitable. Besides, the model described in this paper is just one kind of all possible nonequilibrium models. Therefore even if some details of the proposed model turns out to be untenable, this do not imply the central idea that noncommutativity originates from nonequilibrium is untenable.

If this model can be developed into a consistent mathematical framework then it has the chance of opening up a new way of looking at the issues of quantum and spacetime. In this view the deeper understanding comes not from combining quantumness and gravity (or spacetime) in any way but from finding out what the underlying entities are, because it is from their interactions that both quantumness and spacetime emerges.

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[6] For an argument against the view that pure states is wholly determined by the objective facts of preparation apparatus, and relevant discussions, see C. M. Caves, C. A. Fuchs, R. Schack, “Subjective probability and quantum certainty”, quant-ph/0608190 (2006).

[7] This paper takes the information theoretic view of thermodynamics. For introduction see: R. Balian, “Information in statistical physics”, cond-mat/0501322 (2005). This view was pioneered by E. T. Jaynes in the papers: “Information Theory and Statistical Mechanics” Phys. Rev., 106, 620 (1957) and “Information Theory and Statistical Mechanics II” Phys. Rev., 108, 171 (1957). In this view, the generality of the form of equilibrium macrostate (a probability distribution, usually the canonical distribution) comes from the fact that it depends only on our knowledge of the system in question, and not on what the components are, eg. whether it is a gas of atoms or a group of stars.

[8] This fact is not in contradiction with the view that these variables represent incomplete information about the microstate of the system. When one changes the variables one ‘changes’ the probability distribution, in the sense that his posterior probability about the system is updated by the new values of the macro-variables.

[9] Systematic preparation here means repeated preparations of the whole ensemble of objects described by a certain distribution.

[10] It should be noted that any theory can be formulated in the operator framework, even for classical theories like Newtonian mechanics. The crucial difference with quantum theory is that in such formulations we can always represent all the physical observables by mutually commuting operators.

[11] For a brief description of the possible interpretations, see Section 2.3 of the entry on the uncertainty principle in Stanford Encyclopedia of Philosophy (SEP), “The Uncertainty Principle”, Jos Uffink, http://plato.stanford.edu/entries/qt-uncertainty.

[12] This is a natural consequence of our informational theoretic viewpoint on thermodynamics, for reference see note [4].

[13] Thus we are actually not touching the situations where one of the \( N_a \) or \( N_b \) is so small that the corresponding intensive parameter is meaningless. We are only using the condition \( N_a + N_b = N \) and that both the numbers are still large, and derived the inequality under such special circumstance.

[14] This system could be a quantum field.

[15] This assumption can be seen to be suggested by recent works on analogue gravity, where it is shown that it is possible to derive (with some qualifications) curved spacetime metric from underlying Newtonian particle dynamics, see eg. Section 2.3 in C. Barcelo, S. Liberati and M. Visser, “Analogue Gravity”, Living Rev. Relativity, 8, (2005), 12. Online Article: http://www.livingreviews.org/lrr-2005-12.

[16] The difference is defined as distance \( d_\alpha \) on the \( \alpha \)-manifold.

[17] Ockam’s razor, or that we do not want to introduce additional entities to explain what can already be explained within our current framework.

[18] This seems to suggest an interesting possibility for a way out from black hole information paradox. Roughly: there’s no information loss in such an evolution at the level of subquantum, only that the system is no longer describable by any pure state. Information contained in the quantum level just ‘flows’ into subquantum level during the interactions with space’s entities, and is irretrievable on the quantum level. Severe distortion of time and space structure such as in black holes may provide a situation where requirements mentioned here for a unitary dynamics is not applicable.

[19] A description of this condition can be found in section 2 of the entry on Bell’s theorem in SEP: “Bell’s Theorem”, Abner Shimony, http://plato.stanford.edu/entries/bell-theorem.

[20] For an introduction to Bohm’s theory, see entry in SEP: “Bohmian Mechanics”, Sheldon Goldstein, http://plato.stanford.edu/entries/qm-bohm.