Implementing rules for the measurement of Markovian and Bayesian time operations, establishes the measurement of charged and uncharged boundary conditions – without the need for observers.

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Abstract:

In this model, we establish rules for the measurement of Lorentz invariant potential using the separation of Markovian and Bayesian boundary conditions. By establishing a symmetry between Markovian and Bayesian measurements, we create relative states limited by quantum potential or quantum states limited by relative potential. We use the collapse of the wave function, and our rules for measurement, to define the boundary between Markovian and Bayesian least-time operators and demonstrate that quantum least-time operators can act as an uncharged binary boundaries for charged, Lorentz invariant potential. Because we implement rules for the measurement of these potentials, we can measure them without the need for an observer, closing any loopholes in violations of Bell's inequalities based on observed measurement. This results in a model that is capable of measuring both wave-based and entangled states, while still maintaining adherence to Lorentz invariance.

Keywords: *Aharanov-Bohm effect, Josephson Potential, Statistical Mechanics, Nambu Dynamics, Particle Physics, Quantum Time, Gravitational Potential, Quantum Gravity, Quantum Optics, Quantum Geometry, Optical Statistics, Statistical Geometry,*

1.00 - Introduction

In this investigation into boundary conditions for Lorentz invariant coherencies, we employ rules for measurement which create a hard line between quantum and relative measurements. In effect, we create a symmetry between quantum and relative boundaries by separating Markovian and Bayesian time operators. Because rules for measurement hold for all measurements, they do not require an observer. As we will attempt to demonstrate, the separation of Markovian from Bayesian potential allows us to describe both local and non-local boundary conditions for measurement; allowing us to satisfy violations of Bell's inequality[3,4] and to model actions beyond the collapse of the wave function – like entanglement. It has become increasingly clear that any accurate modeling of our universe must include any violations of Bell's inequalities, as well as incorporating Lorentz invariant descriptions of space time[5,6,7] and light speed boundaries. We must also be able to describe the small world of quantum mechanics[8,9,11,12,13] using the same principles. While these seem to be incompatible and often competing ideas, there is one language that they all share – that is the language of statistics. Statistical measurements are not limited by the speed of light and can be used to model both local and non-local actions[3,4]. Quite often, objects, (or particles) can be so far apart that the probability for their interaction is almost non-existent but, in a purely statistical model, their *potential* for interaction can never be zero. Therefore, all particles in our universe can be said to be connected through their *unmeasured* statistical potential. Even if that potential is infinitely small, it can never be zero in a statistical model. As we will attempt to demonstrate; *in a system determined by measurement there are no zero states or infinities to renormalize because both zero and infinity are not measurable states*. Therefore, in this investigation into the boundaries for statistical potential, we will always establish quantum minimums for the measurement of thermodynamic, electrodynamic and classic potentials. It is the parameters of the measurement that will determine the quantum scalar for all unmeasured potential. We apply the principle of minimum time and action[14,15,18] to help us define rules for the measurement, and interaction, of quantum and relative time operations. The separation of Markovian and Bayesian statistical boundaries will allow us to describe quantum states with *relative potential* or relative states with *quantum potential*. The basic difference between the two potentials is the boundary condition created by distinct differences in the way we measure quantum or relative time operations. Bayesian time is how we experience time. Bayesian time is the time of Einstein and is held to all the boundaries created by general relativity[3,4]. It flows based on the concept of cause and effect. All wave functions operate using this time operator and are bounded by the speed of light. This includes all the Bob and Alice scenarios used to describe the relative measurements taken by different observers. In this investigation we treat Lorentz invariance as a set of coherencies that we translate between boundaries created between Bayesian and Markovian requirements. Markovian time operators have very different boundaries and are not limited by the speed of light. Markovian time operators are quantum in nature and carry specific rules for time transition operators[16]. Entanglement is best described using Markovian time operators. These are the basic rules for the measurement and interaction of time operators that we will employ for this discussion:

- *Two or more simultaneous measurements constitute a single relative measurement*[1]
- *Quantum measurements require a quantum time operator*
- *Quantum measurements cannot be simultaneous*
- *Simultaneous measurements that combine two or more different quantum time operators must use a shared relative time operator*.

[1](Einstein places limits on simultaneous measurement based on the limits of relative observation. In this model, all simultaneous measurements must establish a common relative time operator. This is essentially a re-statement of the same principle. We are bounded by the same limits established by Lorentz invariance and Minkowski space. This is discussed further in Sections-1.01 and 1.02)*

These simple rules will help us to define boundary conditions for unmeasured vector potential which establish the boundaries for measured potential. *In this model, based purely on measured statistical states, there are no zero states or infinities to renormalize, because both zero and infinity are unmeasurable states*. Discreet Markovian fundamental states are then defined by how we measure their time operations. This mirrors the basic requirement of any Markovian transition matrix[11,16] or wave-packet[12,13,18]. Combining statistical states to describe both wave and quantum potential is, of course, the essence of Quantum Statistical Mechanics (QSM) and the use of Feynman path integrals[10]:

$$|\psi(t_2)\rangle \ = \ \breve{\mathrm{U}}(t_2, t_1)|\psi(t_1)\rangle \tag{01}$$

The time evolution operator $\breve{\mathrm{U}}(\mathrm{t}_2, \mathrm{t}_1)$ satisfies the equation of motion:

$$\hbar i \frac{\partial}{\partial t_2} \breve{\mathrm{U}}(t_2, t_1) \ = \breve{\mathcal{H}} \breve{\mathrm{U}}(t_2, t_1) \tag{02}$$

In our investigation, using the rules for measurement, we establish a hard line between quantum Markovian discreet states and relative, and wave-based, Bayesian potentials. This requires us to fulfill the discrete requirements of any Markov chain and balance them against the continuous cause and effect time represented by Bayesian wave functions. This will allow us to model any violations of Bell's inequalities as Lorentz invariant potentials that exist both in and outside the collapse of the wave function.

1.01 - Defining Markovian and Bayesian time operators

General relativity tells us that there are no simultaneous measurements and the relative velocity of observers limits the context of the measurement. Our first rule for measurement seems to conflict directly with this basic limit to any simultaneous measurement imposed by general relativity. In fact, our first rule states that any simultaneous measurements must take into account the relative time operators of each measurement. This is actually in agreement with Einstein's principles regarding simultaneous measurement. In this discussion, we preserve individual measurements of for both Bob and Alice, but any interaction between them requires the establishment of a shared relative time operator. We constrain Lorentz invariant measurements to the same boundaries established by general relativity and constrain Lorentz invariant quantities to the same boundary conditions that are established by Minkowski space or de Sitter space[5,6,7] . However, in this model, fundamental quantum states of potential are determined by the measuring of unmeasured potential. Each measurement of "*unmeasured*" potential determines the minimum quantum scalar for each type of potential energy. In effect, we create quanta that respond to measurement to determine quantum states. Quantum measurements simply require a quantum time operator. This single requirement serves to separate quantum and relative time operations and implies the need for a discrete Markov chain in balance with a Bayesian wave function. We then draw a hard line between these two time functions and the center of least time through the use of a Nambu-Poisson symmetry and phased Hamiltonians representing the center of least time for measured potential. Generalized Nambu-Hamilton equations of motion involve two Hamiltonians and an evolutionary time operator[8,9]. It will establish how the system evolves using the principle of least time represented by the Hamiltonians within the equation:

$$\frac{df}{dt} = \{\mathcal{H}_1, \mathcal{H}_2, f(t_2 \rightarrow t_1)\} \qquad (03)$$

Where, $f(t_2 \rightarrow t_1)$ represents the least-time evolution operator of the two symmetric Hamiltonians.

$$-\{\mathcal{H}_1, \mathcal{H}_2, f(t_1)\ \} = \{\mathcal{H}_1, \mathcal{H}_2, f(t_2)\} \qquad (04)$$

We establish Markovian boundaries to the quantum (left) side of the Nambu-Poisson symmetry by equating an identity matrix to a time operator. Markovian transition matrices allow us to create boundaries for least-time potential based on the restrictions of the matrix. One of the primary advantages of working with *unmeasured* potential is that it allows us to maintain positive boundaries for *measured* potentials in Markovian transition matrices. Unmeasured potential can also placed in the complex plane as an asymptotic boundary condition for measured potential. We can represent the boundaries for unmeasured potential using the set of complex numbers and Hermitian matrices[11]. We start with a transition matrix using real complex numbers representing the opposite potentials needed to reflect conservation boundaries:

$$\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix} \qquad (05)$$

This transition matrix establishes boundary conditions for all measured binary potential. We can associate this basic binary requirement to uncharged Markovian transition probability matrices:

$$f\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix} = f\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix}^2 = f\begin{bmatrix} +1 & -1 \\ -1 & +1 \end{bmatrix} = f(t_n) \qquad (06)$$

Where the value $f\begin{bmatrix} +1 & -1 \\ -1 & +1 \end{bmatrix}$ represents the least-time boundaries for unmeasured Markovian potential and $f(t_n)$ is the quantum evolutionary time operator. In our separation of quantum and relative operations, required by our rules for measurement, any Markovian binary boundary condition requires this basic identity as a time operator:

$$-\{\mathcal{H}_1, \mathcal{H}_2, f(t_n)\} = -\left\{\mathcal{H}_1, \mathcal{H}_2, f\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix}\right\} \qquad (07)$$

Using this matrix as the Markovian transition time operator allows us to create asymptotic boundaries for unmeasured potential that will set the positive boundaries for measured potential. We use an integer-based Markovian operator as our quantum time operator and a Bayesian wave packet as the relative time operator:

$$-\{\mathcal{H}_1, \mathcal{H}_2, f(t_n)\ \} = \{\mathcal{H}_1, \mathcal{H}_2, f(t_\psi)\} \qquad (08)$$

Where the value (t_n) represents a discrete Markovian time operator and (t_ψ) represents a wave-based Bayesian time operator. In our model, quantum time operators, act as uncharged, and asymptotic, binary boundaries for Bayesian charge and wave-based potentials. Charged potentials combine separate quantum measurements into a single measurement and, therefore, based on our rules, we must use a Bayesian wave function as the time operator for any charged potential. This means that all charges, positive, negative and neutral must be contained on the Bayesian side of the NP symmetry. We can show this best by returning to Markovian transition matrices representing charged and wave-based potential. If we include all three charges of the standard model on the Bayesian side of our symmetry we create a transition identity for the wave packet. We can now use the set of complex numbers to represent neutral charge potential:

$$t_\psi = \begin{bmatrix} 1^- & i & 1^+ \\ i & i & i \\ 1^- & i & 1^+ \end{bmatrix} = \begin{bmatrix} 1^+ & i & 1^- \\ i & i & i \\ 1^+ & i & 1^- \end{bmatrix} \quad (09)$$

In our model we use $(1^+, 1^-, 1)$ to represent positive, negative and neutral currents. We use a quantum fundamental state as the carrier for neutral currents, because it carries electrodynamic unmeasured potential without the boundaries set by Coulomb rules and charges. Vector potential works as a quantum operator. By assigning the appropriate quantum or wave-based time operations we can tie these equations to our Nambu-Poisson symmetry. This can be related back to the measured quantum vector potential for the Aharanov-Bohm effect[1, 2] :

$$\hat{A}(x, t_n) = \{\mathcal{H}_{\theta_{sin}}, \mathcal{H}_{\theta_{cos}}, f(t_n)\ \} = \{\mathcal{H}_{\theta_{sin}}, \mathcal{H}_{\theta_{cos}}, f(t_\psi)\} = \psi(r, t_\psi) \quad (10)$$

If we examine the diagonal representing this time operator as a position vector (r) then the wave function represents the boundary conditions for the general Schrodinger equation[13,18]:

$$\hbar i \frac{\partial}{\partial t} |\psi(r, t) \ = \ \mathcal{H} |\psi(r, t) \quad (11)$$

$$\hbar i \frac{\partial}{\partial t} |\psi(r, t) \ = \ \mathcal{H} |t_\psi \quad (12)$$

$$t_\psi = \begin{bmatrix} 1^- & 1 & 1^+ \\ 1 & 1 & 1 \\ 1^- & 1 & 1^+ \end{bmatrix} = \psi(r, t) \quad (13)$$

$$\mathcal{H} |t_\psi = \{\mathcal{H}_1, \mathcal{H}_2, f(t_\psi)\} = -\{\mathcal{H}_1, \mathcal{H}_2, f(t_n)\ \} \quad (14)$$

$$f(t_n) = f \begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix} \quad (15)$$

$$f(t_\psi) = f \begin{bmatrix} 1^- & i & 1^+ \\ i & i & i \\ 1^- & i & 1^+ \end{bmatrix} \quad (16)$$

$$-\left\{\mathcal{H}_1, \mathcal{H}_2, f \begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix}\ \right\} = \left\{\mathcal{H}_1, \mathcal{H}_2, f \begin{bmatrix} 1^- & i & 1^+ \\ i & i & i \\ 1^- & i & 1^+ \end{bmatrix}\right\} = i \frac{\partial}{\partial t} |\psi(r, t) = \mathcal{H} |\psi(r, t) \quad (17)$$

In the next sections we will explore how measurement of unmeasured potential, (electrodynamic, thermodynamic and classical), determines the boundary conditions for the interaction of quantum and relative measured potential . We will examine how the measurement of quantum potential defines relative boundary conditions and the measurement of Bayesian relative potential is bounded by least-time Markovian limits. We will then quickly explore how a quantum center of least time establishes boundaries in relative space for thermodynamic, electrodynamic and classical potentials.

1.02 - Measuring statistical potential

In the standard model, photons are the carriers of electromagnetic force, but can only be described using statistics due to Heisenberg's uncertainty principle. Electromagnetic potentials are also limited by the Planck constant, Berry circuits, and Josephson potentials. In contrast, thermodynamic boundaries are defined purely within the context of Boltzmann limits. In this model, we use measurement to define the context for all statistical potential. However, each type of potential requires a different type of measurement. We can demonstrate how each different variation of potential is determined by measurement using Hermitian identity matrices[11,16] to represent the unmeasured vector potential for each type of statistical measurement. Returning to the NP symmetry and the Hermitian identity matrices for both uncharged and charged boundaries we get:

$$-\left\{\mathcal{H}_1, \mathcal{H}_2, f\begin{bmatrix}1 & i\\ i & 1\end{bmatrix}\right\} = \left\{\mathcal{H}_1, \mathcal{H}_2, f\begin{bmatrix}1^- & i & 1^+\\ i & i & i\\ 1^- & i & 1^+\end{bmatrix}\right\} \qquad (18)$$

By associating the set of complex numbers and arguments in the complex plane with unmeasured potential we can satisfy the requirements of any Markovian probability transition matrices as well as supplying the structure for measurement of that potential. We introduce three variables to the matrix on the right of our equation:

$$\begin{matrix}p_e\\ p_k\\ p_c\end{matrix}\begin{bmatrix}1^- & i & 1^+\\ i & i & i\\ 1^- & i & 1^+\end{bmatrix} \qquad (19)$$

Where $\{p_e, p_k, p_c\}$ represent electrodynamic, thermodynamic and classic potentials and the boundaries for that *measured* potential required by least-time conservation laws. This allows us to represent imaginary squared values in the complex plane as vector potential measurement. Take for example the measurement of electro dynamic potential:

$$\begin{bmatrix}1^+ & 1 & 1^-\\ 1 & 1 & 1\\ 1^+ & 1 & 1^-\end{bmatrix} = \begin{bmatrix}1^- & p_e & 1^+\\ p_e & p_e & p_e\\ 1^- & p_e & 1^+\end{bmatrix} \qquad (20)$$

When measuring electrodynamic potential, the top row represents all three charges as well as Coulomb rules:

$$\langle -1|\mathrm{p}_e| + 1\rangle \qquad (24)$$

The diagonal of each matrix represents the time operation for that measurement of potential as well as incorporating it into a wave packet. For example; when measuring electrodynamic potential we need to be able to represent all three charges in the standard model, (positive, negative, neutral). In our matrix we represent statistical potential using the neutral charge:

$$\begin{matrix} p_e \\ p_k \\ p_c \end{matrix} \begin{bmatrix} 1^- & & \\ & p_e & \\ & & 1^+ \end{bmatrix} or \begin{bmatrix} 1^+ & & \\ & p_e & \\ & & 1^- \end{bmatrix} \qquad (21)$$

Returning to the NP symmetry, we can represent both phases of the wave function with the Hamiltonians using sin and cos values:

$$-\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(t_n)\ \} = \{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(t_\psi)\} \qquad (22)$$

All charged electrodynamic potential can be divided by $\hbar$, therefore electromagnetic measurement adds another boundary condition for both sides of the equation:

$$-\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(\hbar)(t_n)\} = \{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(\hbar)(t_\psi)\} \qquad (23)$$

$$-\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(\hbar)\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix}\right\} = \left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(\hbar)\begin{bmatrix} 1^- & p_e & 1^+ \\ p_e & p_e & p_e \\ 1^- & p_e & 1^+ \end{bmatrix}\right\} \qquad (24)$$

Returning us again to the general Schrodinger equation[18] :

$$-\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(\hbar)\begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix}\right\} = -\hbar \frac{\partial}{\partial t}|\psi(r,t) = -\mathcal{H}|\psi(r,t) \qquad (25)$$

In the next section, we will discuss how measurement creates Markovian and Bayesian boundaries for quantum and wave-based electrodynamic potentials.

1.03 - Electrodynamic potential

In the standard model, photons are described as force carriers for electrodynamic potential. Photons exhibit particle and wave duality and can operate in quantum time or relative time depending on measurement. *By equating the rules of measurement to the interaction of Markovian and Bayesian time operators; we can demonstrate how measurement determines the observation of particle or wave phenomenon*. The original two-slit experiment[12] demonstrates, quite clearly, how probability waves can describe the interference patterns generated by either photons or electrons:

$$S = {\psi_S}^2 \rightarrow {\psi_a}^2 + {\psi_b}^2 + \ 2\,\psi_a|\,\psi_b cos\delta = \ interference \qquad (26)$$

Where ${\psi_S}^2$ describes the source of electrons or photons, and the quantity $({\psi_a}^2)$ and $({\psi_b}^2)$ describe the probabilities generated at slits (a,b). δ is the phase difference between $\psi_a\ and\ \psi_b$. We get an interference pattern, because the single measurement at detector screen has combined two quantum time operators into a relative measurement. Assigning separate time operators to each component of the experiment allows us to move between quantum and relative measurements. When we remove the single measurement at the detector screen and replace it with two separate, and sequential, measurements at the slits; we are adding the quantum operators without interference. These measurements never occur simultaneously, because the two probabilities operate in sequential and additive time. When detectors at each slit count single photons or electrons as they pass through, they effectively create Markovian quantum measurements which result in a "particle". The result of the time measurements being additive, and not relative, is that it does not generate an interference pattern at the detector screen.

$$\psi_S{}^2(t_S) \rightarrow \psi_A{}^2(t_A) = \ no\ interference \quad (27)$$

$$\psi_S{}^2(t_S) \rightarrow \psi_B{}^2(t_B) = no\ interference \quad (28)$$

Boundary conditions for measured potential are really no different than the path integrals associated with QED or the vector potential of the (AB) effect. In fact, we can relate this idea of statistical potential directly to the electrodynamic vector potential represented by the Aharonov-Bohm effect or the quantum actions represented by Berry circuits[17] and Josephson energy[24]. Electrodynamic potential in our model is no different than the standard model. The advantage we have is the ability of unmeasured potential to exceed light speed limits and to model violations of Bell's inequalities – like entanglement. We describe the boundary between quantum and wave based phenomenon much like Dr. Julian Schwinger describes the use of boundary conditions for electromagnetic waves[21]:

"We have, therefore, the boundary condition that the Green's function, in its dependence upon the latest of all times, contains only positive frequencies, and in its dependence upon the earliest of all times, contains only negative frequencies. In effect we have a description in terms of waves which can be considered as moving in the space-time region in such a way that if we have a number of such points in space-time, the waves are moving always out of the region in question. When we are on the boundary of the region in the sense of considering the time coordinate that is later than all the others, the frequencies are positive and the waves move out; if it is the earliest of all times, the frequencies are negative, and the waves move out again. In short, we are dealing with a generalization of the Green's function originally introduced by Feynman which corresponds precisely to the boundary condition of outgoing waves. The waves are in a time sense, running out of the region in question." Julian Schwinger, Nobel Lecture [21]

At the line between Dr. Schwinger's latest and earliest times, he uses a wave operator that is always running out of a region of negative potential. We see this as a representation of our separation of quantum and relative states at the collapse of the wave function. In our model, quantum time operations exist outside of the collapse of the wave function and represent boundaries established by the center for least time for each measurement. The Green's functions will allow us to represent the center of least time as a geometric boundary condition representing the spherically symmetric nature of electro dynamic potential:

$$-\nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(t_n)\} \longleftarrow \nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(t_\psi)\} \quad (29)$$

$$-\nabla^2 G(t_2) f_n \begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix} \leftarrow \nabla^2 G(t_1) f_\psi \begin{bmatrix} 1^- & p_e & 1^+ \\ p_e & p_e & p_e \\ 1^- & p_e & 1^+ \end{bmatrix} \qquad (30)$$

Our geometric boundaries for electrodynamic potential also reflect the statistical probabilities associated with the two-state system described by Dr. Feynman[12], as well as, the two-state EM field operator described by Roy Glauber in his work regarding electromagnetic field coherence and quantification[23]. Quantum electrodynamic potential, representing the concept of least time, can also be defined by Josephson energy and Josephson currents[24]. Tunneling works because the Josephson current passes a barrier using a quantum phase jump. This is a quantum effect based on the *potential* of Josephson energy to make the jump as current with zero voltage and with the least definable time. The relationship between Josephson phase and Josephson current can be stated as:

$$J_s = J_c \, sin(\theta_1 - \theta_2) \qquad (31)$$

The potential energy at a Josephson junction can be stated as:

$$J_p = \frac{\theta_0 \, I_c}{2\pi} (1 - cos\varphi) \qquad (32)$$

J_s = Josephson phase
J_p = Josephson potential
θ_0 = quantum flux
I_c = critical current

The ability of a Josephson current to jump the tunneling barrier can directly related to its quantum potential. We can use it to define a quantum time limit for electrodynamic wave-based potential.

$$-\nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(t_n)\} \longleftarrow \nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(t_\psi)\} \qquad (33)$$

$$-\nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n(J_p)\} \longleftarrow \nabla^2 G \left\{ \mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi \begin{pmatrix} 1^- & p_e & 1^+ \\ p_e & p_e & p_e \\ 1^- & p_e & 1^+ \end{pmatrix} \right\} \qquad (34)$$

Therefore, the Josephson effect and the Aharanov-Bohm Effect are the minimum *measured* potentials for *unmeasured* electrodynamic vector potential. Just as in our model, these effects do not require a field to represent electrodynamic potential. The difference is our ability to create a hard line between Markovian and Bayesian boundaries and to describe actions beyond the collapse of the wave function – like entanglement. These vector potentials are two of the clearest examples of a quantum effect that can also be described using a wave function. In the next section we shall discuss Markovian and Bayesian thermodynamic potential defined by quantum thermal microstates which are also based on the principle of least time and action.

1.04 - Thermodynamic potential

Thermodynamic quantum states, defined by a center of least time, can be described by thermal microstates and Eigen state thermalization[25]. Thermodynamic microstates represent the minimum time for any thermodynamic state in equilibrium. Thermodynamic potential and Maxwell-Boltzmann statistics can therefore, operate using quantum or wave-based boundaries. Each microstate is a quantum snapshot of thermal equilibrium and, therefore, it acts as an uncharged boundary based on conservation limits. This fits well with the laws of thermodynamic equilibrium and the description of thermal microstates given by Goldstein, Lebowitz, and Lieb[26,27]:

$$\langle\psi|\mathrm{P}_{eq}|\psi\rangle \approx 1 \tag{35}$$

"...Let $\mu(d\psi)$ be the uniform measure on the unit sphere in $\mathcal{H}$ (13,14). It follows from (7*) that most ψ relative to μ are in thermal equilibrium. Indeed,*

$$\int\langle\psi|\mathrm{P}_{eq}|\psi\rangle\mu(d\psi) = \frac{1}{D}Tr\mathrm{P}_{eq} = \frac{D_{eq}}{D} \approx 1 \tag{36}$$

Since the quantity $\langle\psi|\mathrm{P}_{eq}|\psi\rangle$ is bounded from above by 1, most ψ must satisfy $(8)^2$."

(7)T. Kato: A short introduction to perturbation theory for linear operators. New York:Springer-Verlag, 1982.[13*] M. Rigol, V. Dunjko, M. Olshanii: Thermalization and its mechanism for generic isolated quantum systems. Nature 452, 854–858, 2008.*
[14] E. Schrodinger: Statistical Thermodynamics. Second Edition, Cambridge University Press, 1952.*

Thermodynamic equilibrium acts dynamically to restrain thermodynamic expansion. Thermodynamic microstates represent the center of least time for any thermodynamic potential. Thermodynamic potential moves from quantum to relative states using Maxwell-Boltzmann statistics. As we approach higher and higher temperatures we reduce the contributions from electrodynamic momentum and classical momentum and Maxwell-Boltzmann statistics become the primary measurement of statistical potential. The Boltzmann constant now drives all time operations:

$$\langle\psi^2|\mathrm{P}_{eq}|\psi^2\rangle = \left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f(k)\begin{bmatrix}1 & i\\ i & 1\end{bmatrix}\right\} \tag{37}$$

Where (k) is the Boltzman constant. To represent the spherically symmetric boundary for thermodynamic equilibrium we can, again, use the Green's functions and the equilibrium microstate as the quantum operator:

$$\langle\psi^2|\mathrm{P}_{eq}|\psi^2\rangle = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi(k)\begin{bmatrix}1 & \mathrm{P}_{eq}\\ \mathrm{P}_{eq} & 1\end{bmatrix}\right\} \tag{38}$$

In the final section we will discuss how the rules for measurement, the separation of Markovian and Bayesian requirements, and the concept of least time can dictate the asymptotic boundaries for gravitational potential.

1.05 - Gravitational potential

In this model, gravity is positioned as a quantum asymptotic boundary condition for classic relative potential. We establish the same asymptotic boundaries established by Lorentz invariance using the stress tensor and the gravitational constant as least-time boundaries. This seems like a departure from general relativity, but it actually is just a rephrasing of the boundaries without the need for fields. This is a direct reflection of the boundaries we associate with the vector potential of the AB effect. Gravimetric *measured* potential is defined using the same boundaries established by Einstein and Minkowski and de Sitter[5,6,7]. In fact, if we think of classic potential in its most recognizable form we can turn to Einstein's famous equivalency of energy to mass:

$$E = mc^2 \tag{39}$$

This equivalency of energy to mass is basically a statement that equates energy to potential. Our use of NP symmetries allows us to translate Lorentz invariant potential between quantum and relative states. In effect, we treat Lorentz invariance as a quantity that we can translate between quantum and relative time operations using the hard line set by our rules for measurement. This allows us to fulfill the unique requirements of special relativity as well as working outside the collapse of the wave function and with quantum boundaries. Let's return to our original equivalency of charged potential and uncharged binary boundary conditions. In this instance we will measure the asymptotic boundary for classical potential momentum based on a quantum center of least time:

$$-\nabla^2 G f_n \begin{bmatrix} 1 & i \\ i & 1 \end{bmatrix} \leftarrow \nabla^2 G f_\psi \begin{bmatrix} 1^- & i & 1^+ \\ i & i & i \\ 1^- & i & 1^+ \end{bmatrix} \tag{40}$$

While relative gravitational potential, and the relative time operator, are always limited by the wave function and the speed of light, quantum gravitational potential is not. We define this quantum minimum for gravitational potential as an uncharged asymptotic boundary for classic potential. Lorentz invariant quantities define the minimum measurement for any classic potential. By attaching vectors to Lorentz invariant classic potential, we can portray the center for least time as an uncharged quantum binary boundary condition for relative potential based on conservation of momentum:

$$-\nabla^2 G f_n \begin{bmatrix} 1 & mc^2 \\ mc^2 & 1 \end{bmatrix} \leftarrow \nabla^2 G f_\psi \begin{bmatrix} 1^- & mc^2 & 1^+ \\ mc^2 & mc^2 & mc^2 \\ 1^- & mc^2 & 1^+ \end{bmatrix} \tag{41}$$

Beginning with our description of classic momentum, we can use our NP symmetry, and the gravitational constant, to define a center for least time for gravitational statistical potential[5,14,15]. We can use the Einstein's momentum and energy stress tensor to construct a minimum for quantum measurement of these momentum vectors and boundaries. Einstein uses the energy stress tensor ($T_{\mu\nu}$) to represent the asymptotic source of gravitational potential. Because it represents the center of least time for any Lorentz-invariant momentum, we can use the energy-stress tensor as our quantum time operator for vectored momentum. We can also use it on the charged side of our symmetry, but we change to the opposite vector constrained by the binary uncharged boundary condition:

$$-\nabla^2 G f_n \begin{bmatrix} 1 & \overleftarrow{T_{\mu\nu}} \\ \overleftarrow{T_{\mu\nu}} & 1 \end{bmatrix} = \nabla^2 G f_\psi \begin{bmatrix} 1^- & \overrightarrow{T_{\mu\nu}} & 1^+ \\ \overrightarrow{T_{\mu\nu}} & \overrightarrow{T_{\mu\nu}} & \overrightarrow{T_{\mu\nu}} \\ 1^- & \overrightarrow{T_{\mu\nu}} & 1^+ \end{bmatrix} \tag{42}$$

Returning to our original NP symmetry:

$$-\nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n(\overleftarrow{T_{\mu\nu}})\} = \nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi(\overrightarrow{T_{\mu\nu}})\} \tag{43}$$

In our model we represent gravity as an uncharged boundary condition for vectored classic momentum. To set a hard line between quantum and relative time operators, let us return to the use of Einstein's gravitational constant[5]:

$$-\nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n\left(\frac{8\pi G_g}{c^4}\left(\overleftarrow{T_{\mu\nu}}\right)\right)\right\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{8\pi G_g}{c^4}\left(\overrightarrow{T_{\mu\nu}}\right)\right)\right\} \quad (44)$$

We can use the gravitational coupling constant (α_G) to help us move all charge to one side of our equation:

$$G_g = \frac{\alpha_G hc}{2\pi m_e{}^2} \quad (45)$$

Where (G_g) is the gravitational constant and α_G is the gravitational coupling constant:

$$-\nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n\left(\frac{8\pi G_g}{c^4}\, T_{\mu\nu}\right)\right\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{8\pi \alpha_G hc}{2\pi m_e{}^2 c^4}\, T_{\mu\nu}\right)\right\} \quad (46)$$

$$-\nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n\left(\frac{8\pi G_g}{c^4}\, T_{\mu\nu}\right)\right\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{4\alpha_G h}{m_e{}^2 c^3}\, T_{\mu\nu}\right)\right\} \quad (47)$$

Moving all light speed and charge boundaries to the Bayesian side of the equation:

$$-\nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n(8\pi G_g\, T_{\mu\nu})\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{4\alpha_G hc}{m_e{}^2}\, T_{\mu\nu}\right)\right\} \quad (48)$$

$$-\nabla^2 G\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n(2\pi G_g\, T_{\mu\nu})\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{\alpha_G hc}{m_e{}^2}\, T_{\mu\nu}\right)\right\} \quad (49)$$

$$-\nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_n\left(\frac{\alpha_G}{2\pi}\right)\right\} = \nabla^2 G\left\{\mathcal{H}_{sin}, \mathcal{H}_{cos}, f_\psi\left(\frac{G_g m_e{}^2}{hc}\right)\right\} \quad (50)$$

Adding back the stress tensor allows us to establish a geometric center of least time and gravitational potential that also separates quantum Markovian and wave-based Bayesian boundaries and uncharged and charged potentials:

$$f_n\left(\frac{\alpha_G}{2\pi}\right) T_{\mu\nu} \leftarrow f_\psi\left(\frac{G_g m_e{}^2}{hc}\right) T_{\mu\nu} \quad (51)$$

Unfortunately, we must end this discussion of the quantum limits to gravity in order to draw some "boundary conditions" of our own for the length of this paper. We hope we have demonstrated that this model is not in conflict with special relativity, but has the advantage of working beyond light speed limits and beyond the collapse of the wave function.

1.6-Discussion

Working with only our the rules for measurement, and the interaction of time operators, we have reduced many of the barriers between quantum theory and relativity. We have demonstrated that this model can be used to measure the potential associated with the Standard Model with the additional advantage of being able to model actions beyond the collapse of the wave function –like entanglement. We have tied this model directly to the vector potentials associated with the Aharanov-Bohm effect and, therefore, original Schrodinger equations and demonstrated that least time principles can then be extended to thermodynamic microstates and the modeling of quantum gravitational potential. We are currently completing an investigation of charge and time symmetries at the event horizon of a black hole which will greatly expand on these basic concepts as well as providing recent astrophysical evidence to support this model.

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(Authors note to editors: Often, when we see a list of references from the foundational papers of science, we tend to dismiss the list as assumed. However, the current standard model of physics was built using these same bricks. This model is a direct reflection a decision to return to the original sources and see where reinterpretation could deliver new insights. I have added some more recent experimental verifications of these basic principles at the bottom of the reference list (28-39). They are not included in the body of the paper, but they can incorporated or expanded, if required.)

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