Semiclassical world is one of infinite many cloneworlds in common spacetime

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We consider N clones of the quantized world, interacting with each other via quantum gravity, coupled by the downscaled Newton constant G/N. In the limit $N\to\infty$, we obtain the semiclassical Einstein equation for every single cloneworld. In the nonrelativistic limit, De Filippo had already obtained the semiclassical Schrödinger-Newton equation; we present an alternative elementary proof. In the general-relativistic case, we complete the semifinished derivation of Hartle and Horowitz. We compare our simple correlated cloneworlds with Stamp's more complicated proposal of correlated worldlines and show why the two constructions differ despite the conceptual similarity.

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I. INTRODUCTION

Although quantum theory was initially conceived for the atomic world, it has come to be thought of as the universal theory of the whole Universe. We often use the semiclassical theory where gravity remains classical (unquantized) and it interacts with the quantized matter. The simplest semiclassical theory goes back to the 1960s [1,2].

Consider a given foliation of the spacetime in spacelike hypersurfaces Σ. The state vector of the quantized matter evolves with the Tomonaga-Schwinger equation (1) where the Hamiltonian density depends on the classical metric g_{ab}, which is the solution of the semiclassical Einstein equation (2) with the Einstein tensor on the left and the expectation value of the energy-momentum operator on the right:

$$\frac{\delta|\Psi_\Sigma\rangle}{\delta\Sigma(x)}=-i\hat{\mathcal{H}}(x)|\Psi_\Sigma\rangle, \tag{1}$$

$$G_{ab}(x)=8\pi G\langle\Psi_\Sigma|\hat{T}_{ab}(x)|\Psi_\Sigma\rangle, \qquad (x\in\Sigma). \tag{2}$$

In the Newtonian limit, the semiclassical theory becomes much simpler. The state vector of the quantized nonrelativistic matter evolves with the Schrödinger equation (3), where $\hat{H}$ is the self-Hamiltonian and the nonrelativistic mass distribution operator $\hat{\mu}$ couples to the classical Newton potential Φ, which is the solution of the Poisson-Newton equation (4):

$$\frac{d|\Psi_t\rangle}{dt}=-i\left(\hat{H}+\int\Phi(\mathbf{r},t)\hat{\mu}(\mathbf{r})d\mathbf{r}\right)|\Psi_t\rangle, \tag{3}$$

$$\Delta\Phi(\mathbf{r},t)=-4\pi G\langle\Psi_t|\hat{\mu}(\mathbf{r},t)|\Psi_t\rangle. \tag{4}$$

This equation, unlike its general-relativistic form (2), is easy to solve. Let us insert the solution into Eq. (3):

$$\frac{d|\Psi\rangle}{dt}=-i\hat{H}|\Psi\rangle+iG\int\int\hat{\mu}(\mathbf{r})\langle\Psi|\hat{\mu}(\mathbf{s})|\Psi\rangle\frac{d\mathbf{r}d\mathbf{s}}{|\mathbf{r}-\mathbf{s}|}|\Psi\rangle. \tag{5}$$

Although this equation had already been used for quantized stellar masses [3], its possible relevance in foundations and its features in the quantized motion of nanomasses were revealed by the present author in 1984 and by Penrose, who called it the Schrödinger-Newton equation [4,5].

In 1981, Hartle and Horowitz considered N identical bosonic fields coupled by quantized gravity at downscaled coupling G/N [6]. The leading order in $1/N$ yielded an approximation closely related to semiclassical gravity:

$$G_{ab}(x)=8\pi G\frac{\langle 0_+|\hat{T}^H_{ab}(x)|0_-\rangle}{\langle 0_+|0_-\rangle}, \tag{6}$$

where $|0_\pm\rangle$ are the asymptotic initial/final bosonic vacuum states, respectively, and $\hat{T}^H_{ab}(x)$ is in the Heisenberg picture. This differs from the correct semiclassical equation $G_{ab}(x)=8\pi G\langle 0_-|\hat{T}^H_{ab}(x)|0_-\rangle$, as noticed by the authors.

Twenty years later, and being apparently unaware of the construction [6], De Filippo discussed its nonrelativistic special case [7]. Let our system of interest exist in N identical copies. Let them interact via the Newton pair potential with the downscaled Newton constant G/N. Consider the Schrödinger equation and take an uncorrelated initial state $|\Psi\rangle^{\otimes N}$:

$$\frac{d|\Psi\rangle^{\otimes N}}{dt}=-i\left(\sum_n\hat{H}_n+\frac{1}{N}\sum_{n<m}\hat{V}^{\mathrm{G}}_{nm}\right)|\Psi\rangle^{\otimes N}, \tag{7}$$

$$\hat{V}^{\mathrm{G}}_{nm} = -G \int\int \frac{\hat{\mu}_n(\mathbf{r})\hat{\mu}_m(\mathbf{s})}{|\mathbf{r}-\mathbf{s}|} d\mathbf{r}d\mathbf{s}. \quad (8)$$

$\hat{H}_n$ stands for the same Hamiltonian $\hat{H}$ acting on the nth subsystem and, similarly, $\hat{\mu}_n(\mathbf{r})$ is the mass distribution operator of the nth subsystem. The *reduced dynamics* of any one of the N components is the same; let us take the first one:

$$\frac{d}{dt}(|\Psi\rangle\langle\Psi|) = \left(\prod_{n\neq 1}\mathrm{tr}_n\right)\frac{d}{dt}(|\Psi\rangle\langle\Psi|)^{\otimes N}. \quad (9)$$

In a lengthy path-integral proof, De Filippo showed that in the $N \to \infty$ limit this reduced state remains a pure state and its evolution is governed by the Schrödinger-Newton equation (5).

Here we present a much simpler proof. It starts with the reduced dynamics of a fixed number $k < N$ of copies:

$$\frac{d}{dt}(|\Psi\rangle\langle\Psi|)^{\otimes k} = \left(\prod_{n>k}\mathrm{tr}_n\right)\frac{d}{dt}(|\Psi\rangle\langle\Psi|)^{\otimes N}. \quad (10)$$

Using Eqs. (7) and (8) together with (4), the r.h.s reads

$$-i\sum_{n=1}^{k}\left[\hat{H}_n + \frac{N-k}{N}\int \Phi(\mathbf{r},t)\hat{\mu}_n(\mathbf{r})d\mathbf{r}, (|\Psi\rangle\langle\Psi|)^{\otimes k}\right] - \frac{i}{N}\sum_{n,m=1}^{k}[\hat{V}^{\mathrm{G}}_{nm}, (|\Psi\rangle\langle\Psi|)^{\otimes k}]. \quad (11)$$

In the limit $N \to \infty$, Eq. (10) reduces to

$$\frac{d}{dt}(|\Psi\rangle\langle\Psi|)^{\otimes k} = -i\sum_{n=1}^{k}\left[\hat{H}_n + \int \Phi(\mathbf{r},t)\hat{\mu}_n(\mathbf{r})d\mathbf{r}, (|\Psi\rangle\langle\Psi|)^{\otimes k}\right]. \quad (12)$$

We have thus proved that a constant number of copies will evolve separately by the Schrödinger-Newton equation (3) each. Note that the whole composite of N copies becomes entangled by the Schrödinger equation (7); only the constant-size subsystems remain disentangled in the limit $N \to \infty$. This explains the general asymptotic mechanism of semiclassicality's emergence from unitary dynamics.

Section II contains our main result. We complete the semifinished proof of Hartle and Horowitz [6] and suggest the narrative of correlated cloneworlds (CCWs) following the perspective of Refs. [6,7]. More recently, Stamp also proposed infinitely many clones of physical fields coupled by Einstein gravity [8]. Our work enjoys strong motivations by his correlated worldlines (CWLs) theory. Section III compares it briefly with our CCW theory.

II. CORRELATED CLONEWORLDS

The following narrative can be imagined behind the model. Suppose that in the *same* quantized spacetime there exist infinitely many *identical* quantized worlds (*cloneworlds*) and we live in one of them. Which one does not matter; they are all identical. Of course, we must replace Newton's coupling G between matter and spacetime curvature by G/N, while the number N of cloneworlds is going to infinity.

Exact methods are hopeless because the quantization of gravity is not yet solved. We restrict ourselves to the naive form of Feynman's path integrals and disregard the unsolved problems like nonrenormalizability, and we disregard even solved ones like the diffeomorphism ambiguity of the metric g.

For simplicity, we consider bosonic matter fields $\phi(x)$ only. We begin with N cloneworlds. Let a spacelike foliation of the spacetime be given and let the (bosonic) matter in each cloneworld have the same initial wave function(al) $\psi_{\Sigma_0}[\phi_{\Sigma_0}]$ on a hypersurface Σ_0. Then, the joint initial state of the N cloneworlds and their common spacetime reads

$$\mathbf{\Psi}_{\Sigma_0}[\phi^1_{\Sigma_0},\ldots,\phi^N_{\Sigma_0},g_{\Sigma_0}] = \left(\prod_{n=1}^{N}\Psi_{\Sigma_0}[\phi^n_{\Sigma_0}]\right)\Psi^G_{\Sigma_0}[g_{\Sigma_0}], \quad (13)$$

where $\Psi^G_{\Sigma_0}[g_{\Sigma_0}]$ is the initial wave function of the spacetime metric g. The following naive Feynman integral expresses the state on a later hypersurface Σ:

$$\Psi_\Sigma[\phi^1_\Sigma,\ldots,\phi^N_\Sigma,g_\Sigma] = \int \exp\left(iNS_G[g] + i\sum_n S_M[\phi^n,g]\right) \times\left(\prod_n \Psi_{\Sigma_0}[\phi^n_{\Sigma_0}]\mathcal{D}\phi^n\right)\Psi^G_{\Sigma_0}[g_{\Sigma_0}]\mathcal{D}g. \quad (14)$$

Here S_G is the Einstein-Hilbert action and S_M is the action of the matter in each cloneworld. Consider the following standard Feynman integral:

$$\Psi_\Sigma[\phi_\Sigma; g] = \int \exp\left(iS_M[\phi,g]\right)\Psi_{\Sigma_0}[\phi_{\Sigma_0}]\mathcal{D}\phi. \quad (15)$$

This expresses the unitary evolution of the matter's wave function in one world in the background metric g. It is known that Ψ_Σ satisfies the Tomonaga-Schwinger equation (1), where $\hat{\mathcal{H}}$ depends on g. We are going to show in the limit $N \to \infty$ that the state of quantized matter in each cloneworld keeps to be in the pure state $\Psi_\Sigma[\phi_\Sigma; g]$, where the metric g depends on these pure states via the semiclassical Einstein equation (2). Just to be clear: the evolution (15) is unitary as long as g is an independently fixed geometry. The evolution is not unitary and not even linear in the semiclassical model.

Recognizing the structures (15) in the expression (14) of the total state Ψ_Σ, we can rewrite the r.h.s of (14):

$$\mathbf{\Psi}_\Sigma[\phi_\Sigma^1, \ldots, \Phi_\Sigma^N, g_\Sigma] = \int \exp\left(iNS_G[g]\right)\left(\prod_n \Psi_\Sigma[\phi_\Sigma^n; g]\right)\Psi_{\Sigma_0}^G[g_{\Sigma_0}]\mathcal{D}g. \tag{16}$$

At this very stage, we modify the method of effective action used by Hartle and Horowitz [6]. We do it in such a way that we can easily calculate the reduced dynamics of a single cloneworld; let it be the first one. Its reduced density matrix is defined by

$$\rho_\Sigma[\phi_\Sigma^1, \phi'^1_\Sigma] = \int \Psi_\Sigma[\phi_\Sigma^1, \ldots, \phi_\Sigma^N, g_\Sigma]\bar{\Psi}_\Sigma[\phi'^1_\Sigma, \ldots, \phi_\Sigma^N, g_\Sigma] \times \left(\prod_{n\neq 1}\mathcal{D}\phi_\Sigma^n\right)\mathcal{D}g_\Sigma. \tag{17}$$

We insert Ψ_Σ and $\bar{\Psi}_\Sigma$ from Eq. (16):

$$\rho_\Sigma[\phi_\Sigma, \phi'_\Sigma] = \int \exp\left(iNS_G[g] - iNS_G[g']\right) \times \Psi_\Sigma[\phi_\Sigma; g]\bar{\Psi}_\Sigma[\phi'_\Sigma; g']\langle\Psi_\Sigma; g'|\Psi_\Sigma; g\rangle^{N-1} \times \Psi_{\Sigma_0}^G[g_{\Sigma_0}]\bar{\Psi}_{\Sigma_0}^G[g'_{\Sigma_0}]\mathcal{D}g\mathcal{D}g', \tag{18}$$

where $|\Psi_\Sigma; g\rangle$ stands for the state vector of the wave functional $\Psi_\Sigma[\phi_\Sigma; g]$ [Eq. (15)]. Because of the trace over the gravity subspace, it is understood that this time g and g' have the same final boundary values $g_\Sigma = g'_\Sigma$, and the path integrals over g, g' will extend for the final boundary Σ.

Now we turn on the limit $N \to \infty$. The factor $\langle\Psi_\Sigma; g'|\Psi_\Sigma; g\rangle^{N-1}$ vanishes if $g \not\equiv g'$. To avoid such degeneracy, we assume that $g' - g = \delta g$ is a finite small function, and then we take the limits $N \to \infty$ and $\delta g \to 0$ in this order. Using the leading-order Taylor expansion $S_M[\Phi, g + \delta g] = -\frac{1}{2}\int T^{ab}\delta g_{ab}\sqrt{|g|}dx$, we can derive the following relationship:

$$\langle\Psi_\Sigma; g + dg|\Psi_\Sigma; g\rangle = 1 + \frac{i}{2}\int_{\Sigma_0}^{\Sigma}\langle\Psi_{\Sigma_0}|\hat{T}_H^{ab}(x)|\Psi_{\Sigma_0}\rangle\delta g_{ab}(x)\sqrt{|g|}dx. \tag{19}$$

The $(N-1)^{th}$ power in Eq. (18) yields a phase factor diverging with N:

$$\langle\Psi_\Sigma; g + \delta g|\Psi_\Sigma; g\rangle^{N-1} = \exp\left(i\frac{N-1}{2}\int\langle\Psi_{\Sigma_0}|\hat{T}_H^{ab}(x)|\Psi_{\Sigma_0}\rangle\delta g_{ab}(x)\sqrt{|g|}dx\right). \tag{20}$$

Fortunately, there is another diverging phase on the r.h.s. of Eq. (18):

$$\exp\left(iNS_G[g] - iNS_G[g + \delta g]\right) = \exp\left(-i\frac{N}{16\pi G}\int G^{ab}(x)\delta g_{ab}(x)\sqrt{|g|}dx\right). \tag{21}$$

For the two divergent phases to cancel each other out, we must require that

$$G_{ab}(x) = 8\pi G\langle\Psi_{\Sigma_0}|\hat{T}_{ab}^H(x)|\Psi_{\Sigma_0}\rangle, \tag{22}$$

which is the semiclassical Einstein equation (2) in the Heisenberg picture.

Now we set $\delta g \equiv 0$. The double path integral (18) reduces to a single one, $\int\mathcal{D}g$. It reduces further to the integral $\int\mathcal{D}g_\Sigma$ over the initial conditions since the rest of g is determined by the semiclassical Einstein equation (2). Equation (18) becomes as simple as

$$\rho_\Sigma[\phi_\Sigma, \phi'_\Sigma] = \int \Psi_\Sigma[\phi_\Sigma; g]\bar{\Psi}_\Sigma[\phi'_\Sigma; g]|\Psi_{\Sigma_0}^G[g_{\Sigma_0}]|^2\mathcal{D}g_{\Sigma_0},$$
$$\hat{\rho}_\Sigma = \int |\Psi_\Sigma; g\rangle\langle\Psi_\Sigma; g||\Psi_{\Sigma_0}^G[g_{\Sigma_0}]|^2\mathcal{D}g_{\Sigma_0}, \tag{23}$$

where the lower equation rewrites the upper one into the Dirac formalism.

According to this, the quantum state of the matter on hypersurface Σ is the statistical mixture of pure states weighted by the probability distribution of the initial values g_{Σ_0} of the spacetime structure. If a single configuration g_{Σ_0} is chosen, then the quantum state remains the pure state $\Psi_\Sigma[\phi_\Sigma; g]$ [Eq. (15)], which, as said there, satisfies the Tomonaga-Schwinger equation (1). We already showed that the metric is determined by the semiclassical Einstein equation (2). This completes the proof that in the limit $N \to \infty$ the emergent dynamics of any single cloneworld is semiclassical.

III. CORRELATED WORLDLINES

The original realization [8] of Stamp's concept that infinitely many (clone)fields are correlated by gravity has changed over the years [9]; the updated theory was reviewed in Ref. [10]. There, the generator functional for N clones in the same quantized spacetime was defined by the following ring path integral:

$$Z_N[J] = \oint e^{iNS_G[g]}(Z_1[g, J])^N\mathcal{D}g, \tag{24}$$

where

$$Z_1[g, J] = \oint e^{iS_M[\phi,g]+i\int J\phi dx}\mathcal{D}\phi \tag{25}$$

is the standard generator functional of a single field in fixed metric g. The functional $Z_N[J]$ does not generate all correlations of the N clones $\phi^1, \phi^2, \ldots, \phi^N$ on the same quantized metric g, but rather the correlations of the collective variables $\sum_{n=1}^N \phi^n$. The functional $Z_N[J]$ generates what we call the reduced dynamics of the summed field. [We think that a plausible choice might be $Z_N[J/N]$, yielding the reduced dynamics of the *average* $(1/N)\sum_{n=1}^N \phi^n$ of the N fields, avoiding divergences in the limit $N \to \infty$.] But the CWL theory keeps on building. It constructs the above reduced dynamics of N clones for $N = 1, 2, \ldots, \infty$ and considers the uncorrelated composition of all of them:

$$Z[J] = \prod_{N=1}^{\infty} Z_N[J]. \tag{26}$$

This generator functional means a further reduction: it only generates the subdynamics of the "particular collective variables," i.e., the sum of all fields, as observed in Ref. [9].

If we cut the product at finite N, it contains $\nu_N = N(N-1)/2$ fields. (We think again that a plausible choice might be a straightforward rescaling of the current in $Z_N[J]$ in the above product yielding the reduced dynamics of the *averages* of the ν_N fields to avoid divergencies in the limit $N \to \infty$.) CWL proposes the following rescaling of the generator functional itself:

$$Z_{\mathrm{CWL}}[J] = \prod_{N=1}^{\infty} (Z_N[J])^{1/\nu_N}. \tag{27}$$

In the ultimate form of CWL theory, this scaled generator constitutes the dynamics of the physical fields.

However, the above rescaling is not standard in field theory. The unscaled generator (26) described (the limit $N \to \infty$ of) the standard reduced dynamics of the "particular collective variables," legitimate in field theory at least formally. The rescaled generator is problematic. It corresponds no more to the subdynamics of collective observables and it is unknown what the new observables could be. CWL theory *postulates* that the rescaled generator generates the correlations of the physical fields. Apparently, this interpretation is used in applications as well [11].

Just for comparison, let the generator functional representation of the CCW theory (Sec. II) stand here:

$$Z_{\mathrm{CCW}}[J] = \lim_{N\to\infty} \oint e^{iNS_G[g]} Z_1[g, J](Z_1[g, 0])^{N-1} \mathcal{D}g. \tag{28}$$

This is equivalent to the reduced dynamics (18) up to the limit $N \to \infty$ (the irrelevant -1 after N is kept for full conformity). This generator functional formalism offers an alternative way to see and prove how the infinite power of $Z_1[g, 0] \equiv \langle \Psi_\Sigma; g' | \Psi_\Sigma; g \rangle$ under a ring integral will make the metric g classical.

The CCW theory is the $N \to \infty$ limit of field theory of N cloneworlds (N-CCW) of quantized matter in the same quantized spacetime. The physical world is a standard reduction to one of these replica worlds. The CWL theory is the field theory of the *uncorrelated composition* of all N-CCW from $N = 1$ to $N = \infty$. The physical world is a postulated dynamics that does not follow from standard field theory. CCW yields exact semiclassical gravity and it contains the remarkable self-attraction, best illustrated by the solitons of the nonrelativistic Schrödinger-Newton equation [4]. Self-attraction is a known mechanism of what is considered a key feature of "path bunching" in CWL. Path bunching and semiclassical self-attraction are not exactly the same, but they are very similar. This is not too surprising since CCW and CWL are based on related concepts and operate on similar mathematical structures.

IV. SUMMARY

We showed that the semiclassical gravity can be derived within standard quantum field theory of infinitely many copies of the quantized matter in common quantized spacetime, which we call (gravitationally) CCWs. In the reduced dynamics of a single copy, the quantumness of spacetime vanishes exactly and we obtain the equations of semiclassical gravity. Before his work with an infinite number of copies [7], De Filippo discussed the Newton interaction with a single mirror of the quantized system, leading to entanglement between the physical world and its mirror [12]; see also Refs. [13,14]. The case of an infinite number of clones is remarkable in that the entanglement between the clones disappears asymptotically: we get the semiclassical Einstein (or Schrödinger-Newton) equation for the single physical world. The model requires cloneworlds, which is a rather unnatural technical assumption. Yet, the merit of CCW stands: the semiclassical equations are not mere approximate equations, but rather exact consequences of standard quantum theory. This raises questions immediately since semiclassical gravity has long been known to be inconsistent [15]. A closer look shows that it is inconsistent with quantum measurements and the statistical interpretation of the wave function [16]. And, indeed, the CCW theory cannot accommodate measurements. The random measurement outcomes make the cloneworlds different, and hence the postmeasurement derivation of the semiclassical equations breaks down. Nevertheless, it is thought provoking whether nonselective measurements, or instead Everett's branchings [17], could have some status in CCW, and thus in semiclassical gravity.

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DATA AVAILABILITY

No data were created or analyzed in this study.

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