Primordial space-time foam as an origin of cosmological matter-antimatter asymmetry*

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Abstract

The possibility is raised that the observed cosmological matter-antimatter asymmetry may reside in asymmetric space-time fluctuations and their interplay with the Stöckelberg-Feynman interpretation of antimatter. The presented thesis also suggests that the effect of space-time fluctuations is to diminish the fine structure constant in the past. Recent studies of the QSO absorption lines provide a 4.1 standard deviation support for this prediction. Our considerations suggest that in the presence of space-time fluctuations, the principle of local gauge invariance, and the related notion of parallel transport, must undergo fundamental changes.

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1 Introduction

The idea put forward by Stückelberg and Feynman that antimatter is nothing more than matter propagating backward in time acquires a special significance in cosmology. In the absence of gravity, space-time is a kinematic arena. On the one side, it constrains the form of the laws of nature by preserving their form under Poincaré transformations. On the other side, it determines the equations of motion and introduces the notion of matter and antimatter in terms of particular finite dimensional representations of the Lorentz group.

As soon as one begins to take effects of gravitation into account, and allows for an interplay of the quantum and gravitational realms, space-time becomes a dynamical entity. In particular, in the early universe this entity suffered violent quantum-induced fluctuations as was first written down at length by Wheeler [1]. Once the temporal component of this primordial space-time foam, as the Planck-era space-time with quantum-induced fluctuations has come to be known, is allowed to be endowed with fluctuations in its direction, the question immediately arises as to what implications does Stückelberg’s and Feynman’s framework hold for the early universe.

The evidence in favor of the cosmic matter-antimatter asymmetry has been strengthened by a series of recent observations. For instance, the antihelium-helium ratio, as reported by the Alpha Magnetic Spectrometer collaboration [2], is

\[
\frac{\Phi_{\text{He}}}{\Phi_{\text{He}}} < 1.7 \times 10^{-6}. \tag{1}
\]

The AMS detected no antihelium, or any \( Z \geq 2 \) anti-nuclei (which could have been created in an antimatter star). When this information is coupled with the observed isotropy of the cosmic microwave background, on the one hand, and the information on the cosmic diffuse \( \gamma \)-ray background, on the other hand, one comes to the conclusion that we live in a universe which contains no significant amount of antimatter [3]. This asymmetry is known as the cosmic antimatter problem, or problem of the cosmological CP violation.

The latter interpretation implicitly assumes CPT symmetry to hold in the early universe. However, as long as the early universe contains an inherent element of non-locality, CPT might have suffered a substantial violation [4]. For this reason, it is perhaps not quite appropriate to refer to the cosmological
matter-antimatter asymmetry as the cosmological CP violation. Stated in the textbook language, the cosmological matter–antimatter problem reads [5]:

At \( kT \gtrsim m_p c^2 \), there exist in equilibrium roughly equal numbers of photons, protons and antiprotons. Today, \( N_p/N_\gamma \sim 10^{-9} \), but \( N_p \simeq 0 \). Conservation of baryon number would imply that \( N_p/N_\gamma = 1 + O(10^{-9}) \) at early times. Where did this initial asymmetry come from?

Our answer, in this essay, is: *This initial asymmetry is a signature of the violent space-time fluctuations in the early universe.*

While referring to space-time fluctuations one needs to clearly identify the gravitational environment one is embedded in. For the present epoch, there are at least four important studies which explore the observability of the space-time foam effects. In Ref. [8] observability of the spatial component of the space-time foam is studied. For certain models, the space-time foam has been found to carry detectable effects in gravity-wave interferometers. In Ref. [9], Amelino-Camelia et al. also explore the possibility of studying quantum-gravity effects via analysing the fine-scale structure and hard spectra of gamma-ray bursters. The work of Klapdor-Kleingrothaus et al. [10] investigates the effects of space-time foam on the neutrino-less double \( \beta \) decay and neutrino oscillations within the astrophysical context. Finally, a series of papers [11, 12, 13] have attempted to understand the existing data on solar and atmospheric neutrinos in terms of quantum-gravity decoherence induced by space-time fluctuations.

LoSecco has analysed the SN1987a “neutrino” events observed in the Kamiokande and Irvine-Michigan-Brookhaven (IMB) detectors. Assuming that these events contain \( \nu_e e \) scattering events (i.e., \( \nu_e e \rightarrow \nu_e e \)), as well as \( \bar{\nu}_e p \) capture events (i.e., \( \bar{\nu}_e p \rightarrow e^+ n \)), he verified CP invariance in general relativity to a few parts in \( 10^6 \) [see, Ref. [14] for exact meaning of this number]. However, it is to be noted that the LoSecco limit on CP violation in general relativity applies to the present epoch only, i.e., when the space-time fluctuations are “tiny.” It does not hold for the early universe where one expects violent fluctuations of space-time. It is that very epoch of the universe that is of prime interest to us here.

To establish the thesis that the observed cosmological matter-antimatter asymmetry arises, and carries a significant contribution from asymmetric
space-time fluctuations and their interplay with the Stückelberg-Feynman interpretation of antimatter, it is essential to first undertake an *ab initio* look at the fermionic (baryonic/leptonic) representation space from the point of view of the Lorentz symmetry without space-time fluctuations. This is done in Section 2. Therein we shall discover a new quantum-mechanical phase in Sec. 2.1. In Section 3 we shall connect this phase with the matter-antimatter symmetry. The stated answer is then established in Section 4 by invoking the Stückelberg-Feynman interpretation of antimatter \[6, 7\]. Section 4.2 briefly studies the effect of the space-time fluctuations on the fine structure constant and concludes that it must have been smaller in the past. Encouragingly, recent studies of the QSO absorption lines provide a 4.1 standard deviation support for this prediction. The essay closes by a few concluding remarks in Section 5.

2 Spin-1/2 matter-antimatter without space-time fluctuations

In order to study implications of space-time fluctuations on the notion of matter and antimatter we here undertake a critical and detailed look at the \((1/2, 0) \oplus (0, 1/2)\) representation space.

2.1 Matter-antimatter phase factor

The spin-1/2 fermions of interest are described by the \((1/2, 0) \oplus (0, 1/2)\) representation space of the Lorentz group (in the usual notation of Ref. \[15\]). This assumption shall be taken to provide sufficiently accurate description of baryons and leptons even in the early universe — its validity may, at most, be confined to sufficiently small local domains. In general, however, an important distinction must be made whether we consider the \((1/2, 0) \oplus (0, 1/2)\) representation space of the Dirac type, or that of Majorana type. Fortunately, the arguments that we present are equally valid for the Dirac case as for the Majorana case. Because some of the readers may be more familiar with the usual Dirac construct, and because certain additional subtleties enter the Majorana construct \[16\], we shall confine our attention to the \((1/2, 0) \oplus (0, 1/2)\) Dirac construct.
To describe charged particles in the Dirac sense the appropriate \((1/2, 0) \oplus (0, 1/2)\) spinors have the form:

\[
\psi(\vec{p}) \equiv \begin{pmatrix} \phi_R(\vec{p}) \\ \phi_L(\vec{p}) \end{pmatrix}.
\]  

(2)

The \(\phi_R(\vec{p})\) transforms as a \((1/2, 0)\) spinor, and boosts as

\[
\phi_R(\vec{p}) = \exp \left( \frac{\vec{\sigma} \cdot \vec{\varphi}}{2} \right) \phi_R(\vec{0}).
\]  

(3)

In the above equation, the \(\vec{\sigma}\) stand for the standard Pauli matrices, and \(\vec{\varphi}\) is the boost parameter. The \(\vec{0}\) corresponds to the momentum of a particle at rest. The definition of \(\vec{\varphi}\) is motivated by the fact that \(E^2 - \vec{p}^2 = m^2\), and that \(\cosh^2 \alpha - \sinh^2 \alpha = 1\) (as an identity). Thus, one can parameterize Lorentz boosts via \(\vec{\varphi}\),

\[
\cosh \varphi = \frac{E}{m}, \quad \sinh \varphi = \frac{|\vec{p}|}{m}, \quad \vec{\varphi} = \frac{\vec{p}}{|\vec{p}|}.
\]  

(4)

We shall assume that \(m \neq 0\). On the other hand, \(\phi_L(\vec{p})\) transforms as a \((0, 1/2)\) spinor, and boosts with the opposite sign in the exponent:

\[
\phi_L(\vec{p}) = \exp \left( -\frac{\vec{\sigma} \cdot \vec{\varphi}}{2} \right) \phi_L(\vec{0}).
\]  

(5)

As we shall see in Sec. 2.2, the wave equation satisfied by the \((1/2, 0) \oplus (0, 1/2)\) spinor defined in Eq. (2) follows from: (a) The transformation properties of \(\phi_R(\vec{p})\) and \(\phi_L(\vec{p})\) as contained in Eqs. (3) and (5), and very importantly, (b) The relative phase between the \(\phi_R(\vec{0})\) and \(\phi_L(\vec{0})\). The latter observation is of special importance to us here and it shall be found to lie at the heart of our arguments.

Due to the isotropy of the zero-momentum vector, \(\vec{0}\), one may argue that \(\phi_R(\vec{0}) = \phi_L(\vec{0})\). In fact, that is precisely what is done in the standard

\footnote{Note that identical transformation properties under Lorentz group do not necessarily imply that other transformations properties will be identical as well. The latter, e.g., may refer to transformations under C, P, and T. Thus, the Dirac- and Majorana-\((1/2, 0) \oplus (0, 1/2)\) constructs carry different physical characteristics under operations of C, P, and T, while carrying identical transformations under the Lorentz group.}
textbooks \[15\] \[17\]. Ryder’s classic book on the theory of quantum fields, in fact, argues:

Now when a particle is at rest, one cannot define its spin as either left- or right-handed, so

\[
\phi_R(\vec{0}) = \phi_L(\vec{0}). \tag{6}
\]

Same assumption has been adopted by Hladik \[17\], and we have encountered it in literature elsewhere as well (for which we have not kept a complete record). The assertion (6) is justified if one was to confine to a purely classical framework. However, if one is to invoke a quantum framework for the interpretation of these spinors then this equality can only be claimed up to a phase:

\[
\phi_R(\vec{0}) = \zeta \phi_L(\vec{0}). \tag{7}
\]

We find that a convenient choice for \(\zeta\) is:

\[
\zeta = \pm \exp(\pm i\phi). \tag{8}
\]

The + sign is to be taken for “particle” spinors, while the − sign is to be used for the “antiparticle” spinors (see, Ref. \[18\]). Otherwise, i.e., if one ignores this phase, one misses anti-particles\footnote{In order to place this observation within a proper context, we wish to stress that this phase becomes manifest only within the group-theoretical derivations of the Dirac equation. Further, in the relevant limit it is in agreement with the remarks of Gaioli and Garcia Alvarez \[19\] on Ryder’s derivation of the Dirac equation. Moreover, if C, P, and T covariances are invoked, it is implicitly contained in the more traditional treatments of the Dirac equation. However, to the best of our knowledge, nowhere in literature has the full physical content of this phase been understood; nor has it been explored in the context of space-time fluctuations. The first hint for the possible existence of this phase appears in a footnote of a 1993 paper on the Bargmann-Wightman-Wigner type quantum field theory \[20\].}.

The most general form of \(\phi\) is a \(2 \times 2\) space-time dependent matrix. That is,

\[
\exp(i\phi) = \exp(i\sigma^a \phi_a(t, \vec{x})). \tag{9}
\]

where \(\sigma^a\) forms the set \((I_2, \vec{\sigma})\) with \(I_2\) as a \(2 \times 2\) identity matrix, and the \(\phi_a(t, \vec{x})\) are a set of four space-time dependent, i.e., local, parameters. The
demand that the resulting spinor of Eq. (2) still transform as an \((1/2, 0) \oplus (0, 1/2)\) object yields, \(\phi_a = 0\), for \(a = 1, 2, 3\). Consequently, \(\phi\) is reduced to an identity matrix multiplied by a space-time dependent real parameter \(\phi_0\):

\[
\phi = I_2 \times \phi_0(t, \vec{x})
\]  

(10)

In the absence of space-time fluctuations, a glimpse into the physical and mathematical content of \(\phi\) can already be obtained by treating \(\phi_0\) as a real c-number with no space-time dependence. By doing so, the effect of space-time fluctuations — a subject of our immediate interest — shall become easier to understand.

### 2.2 Derivation of a Dirac-type wave equation

Equations (3), (5), and (7) contain essentially the entire kinematic structure of the Dirac-type spin-1/2 particles. To see this we follow the footsteps of Lewis Ryder [15], but we now carefully incorporate the important ingredient embedded in a non-vanishing \(\phi\).

1. On the right-hand side of Eq. (3), substitute for \(\phi_R(\vec{0})\) from (7). This gives

\[
\phi_R(\vec{p}) = \zeta \exp \left( \frac{\vec{\sigma} \cdot \vec{\varphi}}{2} \right) \phi_L(\vec{0}).
\]  

(11)

2. From Eq. (3) obtain,

\[
\phi_L(\vec{0}) = \exp \left( \frac{\vec{\sigma} \cdot \vec{\varphi}}{2} \right) \phi_L(\vec{p}),
\]  

(12)

and insert it into the right-hand side of Eq. (11). This yields:

\[
\phi_R(\vec{p}) = \zeta \exp (\vec{\sigma} \cdot \vec{\varphi}) \phi_L(\vec{p}).
\]  

(13)

3. Similarly, starting from Eqs. (3) and (7) we obtain:

\[
\phi_L(\vec{p}) = \zeta^{-1} \exp (-\vec{\sigma} \cdot \vec{\varphi}) \phi_R(\vec{p}).
\]  

(14)
4. Now, because $(\vec{\sigma} \cdot \vec{p})^2 = 2 \times 2$ Identity matrix, $I_2$

\[
(\vec{\sigma} \cdot \vec{p})^n = \begin{cases} I_2 & \text{for } n \text{ even} \\ \vec{\sigma} \cdot \vec{p} & \text{for } n \text{ odd} \end{cases}
\]  

(15)

This leads to the identities:

\[
\exp(\pm \vec{\sigma} \cdot \vec{\varphi}) = \frac{EI_2 \pm \vec{\sigma} \cdot \vec{p}}{m}
\]  

(16)

5. Next, substitute these identities in Eqs. (13) and (14), and re-arrange to obtain:

\[
\begin{pmatrix}
-m\zeta^{-1} & EI_2 + \vec{\sigma} \cdot \vec{p} \\
EI_2 - \vec{\sigma} \cdot \vec{p} & -m\zeta
\end{pmatrix}
\begin{pmatrix}
\phi_R(\vec{p}) \\
\phi_L(\vec{p})
\end{pmatrix} = 0.
\]  

(17)

6. Finally, with $p_\mu = (p^0, -\vec{p})$, $E = p^0$, read off the Weyl-representation gamma matrices:

\[
\gamma^0 = \begin{pmatrix} 0 & I \\ I & 0 \end{pmatrix}, \quad \gamma^i = \begin{pmatrix} 0 & -\sigma^i \\ \sigma^i & 0 \end{pmatrix},
\]  

(18)

and introduce

\[
\Phi(\phi) = \begin{pmatrix} \zeta^{-1}(\phi) & 0 \\ 0 & \zeta(\phi) \end{pmatrix}.
\]  

(19)

This yields,

\[
(\gamma^\mu p_\mu - m\Phi(\zeta))\psi(\vec{p}) = 0.
\]  

(20)

The obtained equation is Poincaré covariant and indeed carries the solutions with the correct dispersion relations $E = \pm \sqrt{\vec{p}^2 + m^2}$, because

\[
\det [\gamma^\mu p_\mu - m\Phi(\zeta)] = (\vec{p}^2 + m^2 - E^2)^2.
\]  

(21)

Now, if $\Phi(\zeta)$ is allowed to carry space-time dependence, the same method proceeds (with $\phi_0$ becoming a local parameter), and Eq. (20) becomes the CP-violating Dirac equation postulated in Ref. [21].

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\(^3\)Note, it is precisely the property expressed by Eq. (14) that is responsible for the linearity of the Dirac equation as will become obvious below.

\(^4\)We now abbreviate $I_2$ by $I$. The zeros below stand for $2 \times 2$ null matrices, and $\sigma^1 = \sigma_x$, etc.
3  New phase and matter-antimatter symmetry

The reader may have already noted that \( \text{Det} [\gamma^\mu p_\mu - m\Phi(\phi)] \) is independent of \( \phi \). Consequently, Poincaré covariance alone cannot constrain \( \phi \). To constrain this phase angle one needs to invoke additional symmetries. If we demand Eq. (20) to be covariant under the operations of \( C \), \( P \), and \( T \), then we find that

\[
\phi_0 = n \times 2\pi, \quad \text{with} \quad n = 0, 1, 2, \ldots \tag{22}
\]

That is, for a kinematic structure that is covariant under \( C \), \( P \), and \( T \) symmetries, Eq. (7) reduces to

\[
\phi_R(\vec{0}) = \pm \phi_L(\vec{0}). \tag{23}
\]

The \( \psi(\vec{p}) \) of Eq. (2) constructed with a plus sign in the above equation results in the “particle \( u \)-spinors”, while the minus sign in Eq. (23) results in the “antiparticle \( v \)-spinors.” Explicitly, the particle spinors are

\[
\begin{align*}
&u_{+1/2}(\vec{p}) = \kappa(\vec{p}) \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}, \quad u_{-1/2}(\vec{p}) = \kappa(\vec{p}) \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix}, \\
&v_{+1/2}(\vec{p}) = \kappa(\vec{p}) \begin{pmatrix} 1 \\ 0 \\ -1 \end{pmatrix}, \quad v_{-1/2}(\vec{p}) = \kappa(\vec{p}) \begin{pmatrix} 0 \\ 1 \\ -1 \end{pmatrix}.
\end{align*}
\tag{24}
\]

and the antiparticle spinors read

\[
\begin{align*}
&u_{+1/2}(\vec{p}) = \kappa(\vec{p}) \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}, \quad u_{-1/2}(\vec{p}) = \kappa(\vec{p}) \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix}, \\
&v_{+1/2}(\vec{p}) = \kappa(\vec{p}) \begin{pmatrix} 1 \\ 0 \\ -1 \end{pmatrix}, \quad v_{-1/2}(\vec{p}) = \kappa(\vec{p}) \begin{pmatrix} 0 \\ 1 \\ -1 \end{pmatrix}.
\end{align*}
\tag{25}
\]

Here the \((1/2, 0) \oplus (0, 1/2)\) boost, \( \kappa(\vec{p}) \), as contained in Eqs. (3) and (4), is given by

\[
\kappa(\vec{p}) = \begin{pmatrix} \exp(\vec{\sigma} \cdot \vec{\varphi}/2) & 0 \\ 0 & \exp(-\vec{\sigma} \cdot \vec{\varphi}/2) \end{pmatrix}.
\tag{26}
\]

The particle-antiparticle spinors found in the standard textbooks, see, e.g., Bjorken and Drell’s (BD) well-known classic [22], are related to those obtained here by:
\[ \psi^{BD}(\vec{p}) = \frac{1}{\sqrt{2}} \begin{pmatrix} I & I \\ I & -I \end{pmatrix} \psi(\vec{p}) \] (27)

Parenthetically, one may be interested to note that different fermions in Nature, i.e., quarks and leptons, need not carry same \( \phi_0 \) dependence for the underlying spinorial structure. In the event, say, the \( s \)-quark carried a different \( \phi_0 \) than given by Eq. (22), the resulting kinematic structure might not be intrinsically CP respecting. We envisage the above considerations, when incorporated in the kinematic structure of the standard model particles, to hold serious potential to provide a source for CP violation in particle physics. But, here in this essay our interest is mainly in the cosmological matter-antimatter asymmetry — a violation which we now argue arises from an interplay of the \( \phi_0 \) and space-time fluctuations.

At this stage in our essay we have arrived at the standard \((1/2, 0) \oplus (0, 1/2)\) Dirac construct. The merit of this apparently simple exercise lies in having exposed a CPT-related hidden phase in the \((1/2, 0) \oplus (0, 1/2)\) representation space and having deciphered that the Poincaré symmetry alone does not entirely constrain the \((1/2, 0) \oplus (0, 1/2)\) representation space. One needs additional requirements of P, C, and T symmetries. Further, invoking the standard conventions we reach the key understanding. It reads:

Spin-1/2 matter of the Dirac type corresponds to the plus sign in Eq. (23), while the antimatter is characterized by the minus sign.

Now, in order to arrive at our thesis, we make the needed observations explicitly. In the standard textbook treatments, the Dirac’s hole theory and Stückelberg-Feynman’s interpretation of antiparticles as particles going backward in time are considered equivalent. However, as pointed out by Hatfield [23], and confirmed explicitly by the work on \((1, 0) \oplus (0, 1)\) representation space [24], the former framework applies only to the fermions while the latter is applicable to fermions as well as to bosons. From a parenthetic remark of Feynman contained in his paper entitled “The theory of positrons” it is also apparent that he was aware of this advantage of his theory (see, p. 750, in [7]). For this reason, we shall here adopt the Stückelberg-Feynman framework for the matter-antimatter solutions of the \((1/2, 0) \oplus (0, 1/2)\) representation space.
It is then seen that the $+ \exp(+i\phi) = +1$ corresponds to a particle propagating in the forward direction of time while the $- \exp(-i\phi) = -1$ corresponds to the particle propagating backward in time (which is then interpreted as antiparticle).

4 Matter-antimatter asymmetry and Space-time fluctuations

In the presence of space-time fluctuations, Stöckelberg-Feynman framework suggests a natural extension. We propose this to be:

A particle moving forward in time on encountering a region of space-time, with a backward temporal fluctuation, becomes an antiparticle.

That is, temporal fluctuations interchange $(1/2, 0) \oplus (0, 1/2)$ sectors associated with the phases $+ \exp(+i\phi)$ and $- \exp(-i\phi)$, and thereby transforms a particle into an antiparticle:

**Temporal fluctuations:**

$(1/2, 0) \oplus (0, 1/2)$ Sector with $+ \exp(+i\phi)$

$\leftrightarrow (1/2, 0) \oplus (0, 1/2)$ Sector with $- \exp(-i\phi)$

(28)

In this process, the conservation of charges connected by the $C$ operator (such as, electric charge and baryon number) are violated. With CP symmetry and baryon number conservation violated in this manner these processes must proceed in an environment with a slight lack of thermal equilibrium in order to account for the cosmological matter-antimatter asymmetry. We underline “slight” in order to accommodate for the high degree of temperature isotropy of the cosmic microwave background.

In the cosmological realm the advantage of the Stöckelberg-Feynman over the Dirac’s hole theory is more than formal. In fact, the the former framework appears crucial to understanding the matter-antimatter asymmetry in the early universe. To argue this assertion we note:

1. In the standard big bang scenario, at $t = 0$, it is not possible to define propagation backward in time.
2. The interplay of the quantum and gravitational realms implies that space-time itself must carry fluctuations. The simplest argument leading to this observation is that if gravitational effects associated with a quantum measurement process are incorporated, then space-time measurements become non-commutative. This has been discussed in detail, e.g., in Refs. [4, 23, 26, 27, 28, 29, 30, 31].

Now photons are self-conjugate under the operation of charge conjugation, C, and thus are insensitive to the arrow of time. On the other hand, in the Stückelberg-Feynman framework, spin-1/2 matter fields of the Dirac type are endowed with a sensitivity to the arrow of time. To understand the cosmological matter-antimatter asymmetry, the observation in item (1) above suggests to confine one’s attention to a universe with only fermions (and no anti-fermions) at \( t = 0 \). This is precisely what we shall do next. The possibility of a photonic birth shall be taken up in Sec. 4.1.

Given this circumstance, and once the reality of quantum-induced space-time fluctuations is accepted, it is readily seen that from the vantage point of an hypothetical observer riding the temporal component of these fluctuations, the Stückelberg-Feynman interpretation implies that matter-antimatter identity of the primordial spin-1/2 Dirac type particles varies with time. That is, as seen by the indicated hypothetical observer, the phases in Eq. (23) is no longer uniquely plus, or uniquely minus, for a given particle. But, instead, it oscillates, \( + \leftrightarrow - \), with time in an observer-dependent manner.

As determined by the averaged observations of a large number of indicated hypothetical observers, the ratio defining the cosmic antimatter problem

\[
\frac{N_p}{N_T}_{kT \approx m_pc^2} = 1 + O(10^{-9})
\]

is then an indication that the temporal fluctuations were dramatically violent even at \( kT \approx m_pc^2 \). The early universe at \( kT \approx m_pc^2 \) was characterized by a forward-backward asymmetry in time of the order \( O(10^{-9}) \). This asymmetry

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5 An alternate view may be taken that an \( e^- \), say [described by a \( u(\vec{p}) \)], propagating forward in time encounters a space-time region undergoing a backward temporal fluctuation. The effect of this encounter is to transform an \( e^- \) into an \( e^+ \) [described by a \( v(\vec{p}) \)]. It is be noted that in every space-time region there are, with respect to the local arrow of time, both forward and backward in time propagating particles.

6 Whose ride on the space-time fluctuations is assumed incoherent.
was perhaps even much smaller at still higher temperatures. The cosmological arrow of time was still in the birth pangs. However, it is apparent that as soon as the cosmological arrow of time becomes well pronounced, i.e., the backward temporal fluctuations loose their amplitude in relation to the forward fluctuations, the universe appears as composed of matter.

In brief, the considered scenario, the primordial universe begins with a maximal matter-antimatter asymmetry in favor of matter. However, the violent space-time fluctuations soon transform it into a matter-antimatter universe with roughly equal densities of both. At \( t \approx m_p c^2 \), the matter-antimatter asymmetry is roughly one part in a billion. At this stage, matter and antimatter annihilate and in the process an attempt towards a thermal equilibrium of various then-existing radiation and particle components can proceed. To what extent a thermal equilibrium is reached is determined by \( t_u \) (see Eq. (31) below), and various related thermodynamic considerations.

Adding to thermal equilibrium is the fact that before a distinct cosmological arrow of time is born, cosmic time is really not a good measure of the cosmic age. To attend to this unique cosmological circumstance, we suggest that the cosmic age should be defined as the sum of the cosmic time plus a thermodynamic time. Around \( t = 0 \), we define the thermodynamic time, \( t_{th} \), as

\[
t_{th} \approx \bar{A} \times N(t).
\]

Here, \( \bar{A} \) stands for the average amplitude of the temporal fluctuations (around \( t = 0 \)), while \( N(t) \) equals the number of the accumulated fluctuations at cosmic time \( t \). From that, the age of the universe \( t_u \) can be read off as

\[
t_u \approx t + t_{th}.
\]

In order that a cosmological arrow of time comes into a distinct existence we surmise that (and a full theory of quantum gravity should account for) as evolution of the universe proceeds, the amplitude associated with the backward temporal fluctuation begins to diminish compared to the amplitude for forward fluctuations. As this happens, the universe once again begins to appear as matter dominated, but now it has an additional radiation component. \( N(t) \) can, in principle, be determined from the constraints imposed by directional uniformity of the temperature of the cosmic microwave background radiation, and additional sources of thermalization that may emerge.
from “opening up” of the light cone in the theory of relativity that carries an invariant length \[32\].

4.1 Cosmological matter-antimatter asymmetry with a photonic birth

An alternate scenario would be a photonic birth of the universe, followed by a pair-created matter and antimatter. As a result of the asymmetric temporal fluctuations, the matter-antimatter ratio will undergo local changes resulting into excess of one type of matter over the other. This can happen when, e.g., a \(e^+\) encounters a region of space-time where a temporal fluctuation has changed the direction of time. In that process the \(e^+\) becomes an \(e^-\), which can thus, in one of the possibility (say), annihilate an \(e^+\) which has either not suffered any reversal in temporal fluctuation, or, has suffered an even number of them. In considering such a process one may envisage a rigid mountain form placed on two spatial dimensions and one time dimension — the latter representing the height. In such processes space-time fluctuations result in a net creation/destruction of electric charge (and/or, baron number). Note that a similar fate meets if we begin with an \(e^-\) in the above example. A one part in a billion imbalance in such processes could then give rise to the observed matter-antimatter asymmetry. This scenario is conceptually more intricate and calls for a detailed Monte Carlo simulations to explore its viability.

4.2 Fine Structure constant and space-time fluctuations

Above considerations suggest a natural and testable prediction:

As already noted, it is a general expectation of every model of space-time foam that space-time fluctuations were much more violent at the big bang, and that these have slowly become less intense. As a consequence, far from big bang, and still in the distant past (to become more precise below), the magnitude of the effective charge of an electron would appear smaller than its present value. For, while the time fluctuates backward, a particle that was interpreted as an electron would appear as a positron. Larger backward time fluctuations in the past would thus effectively reduce the magnitude of the
observed electronic charge. The same result applies to a proton. Thus, under the assumption of time-constancy for $\hbar$ and $c$ over the period of interest, one would expect the fine structure constant $\alpha = e^2/\hbar c$ to be smaller in the past.

As this essay was written, we learned that this prediction is consistent with the latest results on the variation of the fine structure constant \cite{33}. With the definition, $\Delta \alpha/\alpha = (\alpha_z - \alpha_0)/\alpha_0$, where $\alpha_0$ is the present day value of $\alpha$ and $\alpha_z$ is the value at absorption redshift $z$ [see \cite{33} for details], they find:

$$\frac{\Delta \alpha}{\alpha} = (-0.72 \pm 0.18) \times 10^{-5} \quad (32)$$

over the redshift range $0.5 < z < 3.5$. The result (32) confirms their earlier study \cite{34, 35}, and now represents a statistically significant 4.1 standard deviation effect in favor of the conclusion that $\alpha$ was slightly smaller in the past.

## 5 Concluding remarks

The notion of charge conservation, and in particular that of baryon number conservation, is an empirically derived concept from observations in an cosmic epoch when the cosmological arrow of time is a well established object. However, as we have seen, these conservations law carry little meaning when this very important object was still in its formative stages and suffered violent fluctuations. This circumstance suggests that, in the presence of space-time fluctuations, the principle of local gauge invariance (whether in some internal space, or space-time), and the related notion of parallel transport, must undergo fundamental changes.

Given a specific form of the space-time fluctuations (an input to come from a future full theory of quantum and gravitational realms), one may derive the time-dependent matter-antimatter ratio beginning with, say, matter alone. The considerations presented above strongly indicate that even at $kT \approx m_p c^2$ the space-time fluctuations were so violent that for every observer the matter-antimatter ratio appeared as to be roughly unity. As the backward fluctuations in time became less dominant, and the cosmic arrow of time acquired a physical reality, the same universe began to appear as matter dominated. The violent interplay of the quantum and gravitational realms
did not respect conservation of electric charge, and baryon number. Interestingly, all this is a natural consequence if one adopts the St"uckelberg-Feynman interpretation of matter-antimatter and allows for quantum-induced fluctuations of space-time. If all this provides a fundamental space-time origin of the cosmological matter-antimatter asymmetry, the CP violation in heavy quark systems may be understood if the light quarks carried a $\phi_0$ given by Eq. (22), while the heavy quarks, say the strange one, carried a CP violating kinematic structure.

In the last two years it has become apparent that quantum-gravity induced fluctuations in spatial distances can be studied using gravity-wave interferometers \cite{8, 36, 37, 38}. This essay indicates that a careful analysis of the cosmological evolution of the fine structure constant, along the lines of the pioneering work of Webb \textit{et al.} \cite{33}, could become a powerful probe of the temporal fluctuations in the space-time foam.

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