Tachyonic neutrinos?

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It is shown that tachyons are associated with unitary representations of Poincaré mappings induced from $SO(2)$ little group insted of $SO(2,1)$ one. This allows us to treat more seriously possibility that neutrinos are fermionic tachyons according to the present experimental data.

I. INTRODUCTION

Almost all recent experiments, measuring directly or indirectly the electron and muon neutrino masses, have yielded negative values for the mass square [1]. It suggests that these particles might be fermionic tachyons. This intriguing possibility was written down some years ago by Chodos et al. [2]. Furthermore, possible new experiments are presented in the papers by Kostelecky [3] and Chodos et al [4,5].

On the other hand, in the current opinion, there is no satisfactory theory of superluminal particles. This persuasion creates a psychological barrier to take such possibility seriously. Even if we consider eventuality that neutrinos are tachyons, the next problem arises; namely a modification of the theory of electro-weak interaction will be necessary in such a case. But, as we known, in the standard formulation of special relativity, the unitary representations of the Poincaré group, describing fermionic tachyons, are induced from infinite dimensional unitary representations of the noncompact $SO(2,1)$ little group. Consequently, the neutrino field should be infinite-component one so a construction of an acceptable local interaction is extremally difficult.

In this paper we suggest a solution to the above dilemma. To do this we use the formalism developed in the paper [6] based on the earlier works [7,8], where it was proposed a consistent description of tachyons on both classical and quantum level. The basic idea is to extend the notion of causality without a change of special relativity. This can be done by means of a freedom in the determination of the notion of the one-way light velocity, known as the “conventionality thesis” [9,10]. The main results obtained in [6] can be summarized as follows:

• The relativity principle is formulated in the framework of a nonstandard synchronization scheme (the Chang–Tangherlini (CT) scheme). The absolute causality holds for all kinds of events (time-like, light-like, space-like).
• For bradyons and luxons our scheme is fully equivalent to the standard formulation of special relativity.
• For tachyons it is possible to formulate covariantly proper initial conditions.
• There exists a covariant lower bound of energy for tachyons.
• The paradox of “transcendental” tachyons is solved.
• Tachyonic field can be consistently quantized using the CT synchronization scheme.
• Tachyons distinguish a preferred frame via mechanism of the relativity principle breaking [8,6].

In this paper we make the next step in this direction by classification of all possible unitary Poincaré mappings for space-like momenta. The main and unexpected result of the present work is that unitary mappings for space-like momenta are induced from the $SO(2)$ little group. This holds because we have a bundle of Hilbert spaces rather than a single Hilbert space of states. Therefore unitary operators representing Poincaré group act in irreducible orbits in this bundle. Consequently, elementary states are labelled by helicity, in an analogy with the light-like case. This fact is extremally important because we have no problem with infinite component field.

Now, let us begin with a brief review of the theory proposed in [6,8].

It is rather evident that a consistent description of tachyons lies in a proper extension of the causality principle. Note that interpretation of the space-like world lines as physically admissible tachyonic trajectories favour the constant-time

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initial hyperplanes. It follows from the fact that only such surfaces intersect each world line with locally nonvanishing slope once and only once. Notwithstanding, the instant-time hyperplane is not a Lorentz-covariant notion which is just the source of many troubles with causality. In the standard framework of the special relativity, space-like geodesics do not have their physical counterparts. This is an immediate consequence of the assumed causality principle which admits time-like and light-like trajectories only.

It is important to stress the following two well known facts from special relativity:

- The definition of a coordinate time depends on the synchronization scheme [9,11,12].
- Synchronization scheme is a convention, because no experimental procedure exists which makes it possible to determine the one-way velocity of light without use of superluminal signals [10].

Therefore a choice of a synchronization scheme does not affect the assumptions of special relativity but evidently it can change the notion of causality, depending on the definition of the coordinate time.

Following Einstein, intrasystemic synchronization of clocks in their “setting” (zero) requires a definitional or conventional stipulation (for discussion see Jammer [10] and Sjödin [13]). Indeed, to determine one-way light speed it is necessary to use synchronized clocks at rest in their “setting” (zero). On the other hand to synchronize clocks we should know the one-way light velocity. Thus we have a logical loophole. Therefore no experimental procedure exists (if we exclude superluminal signals) which makes possible to determine unambiguously and without any convention the one-way velocity of light. Consequently, only the average value of the light velocity around closed paths has an operational meaning. This statement is known as the conventionality thesis [10]. However, the requirement of causality, logically independent on the requirement of the Lorentz covariance, can contradict the conventionality thesis and consequently it can prefer a definite synchronization scheme, namely CT scheme.

In the papers by Chang [14–16], it was introduced four-dimensional version of the Tangherlini transformations [17], termed the Generalized Galilean Transformations (GGT). In [7] it was shown that GGT, extended to form a group, are nonlinear form of the Lorentz group transformations with SO(3) as a stability subgroup. The coordinate transformations should be supplemented by transformations of a vector-parameter interpreted as the velocity of a privileged frame. It was also shown [7] that the above family of frames is equivalent to the Einstein–Lorentz one. A difference lies in another synchronization procedure for clocks. As a consequence a constant-time hyperplane is a covariant notion in our formalism. Hereafter we call this procedure of synchronization the Chang–Tangherlini synchronization scheme.

In the papers [8,6] these ideas was developed and applied to the description of tachyons.

II. FORMALISM

Let us start with a simple observation that the description of a family of relativistic inertial frames in the Minkowski space-time is not so natural. Instead, it seems that the geometrical notion of bundle of frames is much more natural. Base space is identified with the space of velocities; each velocity marks out a coordinate frame. Indeed, from the point of view of an observer (in a fixed inertial frame), all inertial frames are labelled by their velocities with respect to him. Therefore, in principle, to define the transformation rules between frames, we can use, except of coordinates, also this vector parameter, possibly related to velocities of frames with respect to a distinguished observer.

According to the paper [6], transformation between two coordinate frames $x^\mu$ and $x'^\mu$ has the following form

\begin{align}
    x' &= D(\Lambda, u)(x + a), \quad (1a) \\
    u' &= D(\Lambda, u)u. \quad (1b)
\end{align}

Here $\Lambda$ belongs to the Lorentz group $L$, whilst $u$ is a fourvelocity of a privileged inertial frame measured in the coordinate frame $x^\mu$. The $a^\mu$ are translations. The transformations (1) have standard form for rotations i.e. $D(R, u) = R$, whereas for boosts the matrix $D$ takes the form

\begin{itemize}
    \item A necessity of a presence of a preferred frame for tachyons was stressed by many authors (see, for example, [4,18]).
\end{itemize}
\[
D(\vec{V}, \vec{\sigma}) = \begin{pmatrix}
\gamma & 0 \\
\frac{-\vec{V}}{c}\gamma^{-1} & I + \frac{\vec{V}\otimes\vec{V}^T}{c^2} \gamma^{-2} - \frac{\vec{V}\otimes\vec{\sigma}^T}{c^2\gamma\gamma_0}
\end{pmatrix}
\]  

(2)

where we have used the following notation

\[
\gamma_0 = \left[\frac{1}{2} \left(1 + \sqrt{1 + \left(\frac{2\vec{\sigma}}{c}\right)^2}\right)\right]^{1/2},
\]  

(3a)

\[
\gamma(\vec{V}) = \left(\frac{1 + \vec{\sigma}\vec{V}^2}{c^2}\gamma_0^{-2} - \left(\frac{\vec{V}}{c}\right)^2\right)^{1/2},
\]  

(3b)

\[
\frac{\vec{\sigma}}{c} = \frac{\vec{u}}{u^0},
\]  

(3c)

Here \(\vec{V}\) is a relative velocity of \(x'\) frame with respect to \(x\) whilst \(\vec{\sigma}\) is the velocity of a preferred frame. The transformations (2) remain unaffected the line element

\[
ds^2 = g_{\mu\nu}(u)dx^\mu dx^\nu
\]

(4)

with

\[
g(u) = \begin{pmatrix}
1 & \frac{\vec{\sigma}^T}{c} \gamma_0^{-2} \\
\frac{\vec{\sigma}}{c} \gamma_0^{-2} & -I + \frac{\vec{\sigma}\otimes\vec{\sigma}^T}{c^2} \gamma_0^{-4}
\end{pmatrix}
\]

(5)

Notice that \(u^2 = g_{\mu\nu}(u)u^\mu u^\nu = c^2\).

From (5) we can calculate the velocity of light propagating in a direction \(\vec{n}\)

\[
\vec{c} = \frac{c\vec{n}}{1 - \frac{\vec{n}\vec{\sigma}}{c} \gamma_0^{-2}}
\]

(6)

It is easy to verify that the average value of \(\vec{c}\) over a closed path is always equal to \(c\).

Now, according to our interpretation of the freedom of the realization of the Lorentz group as freedom of the synchro-nization convention, there should exists a relationship between \(x^\mu\) coordinates and the Einstein-Poincaré (EP) ones denoted by \(x_E^\mu\). Indeed, we observe, that the coordinates

\[
x_E = T^{-1}(u)x,
\]  

(7a)

\[
u_E = T^{-1}(u)u,
\]

(7b)

where the matrix \(T\) is given by

\[
T(\vec{\sigma}) = \begin{pmatrix}
1 & -\frac{\vec{\sigma}^T}{c} \gamma_0^{-2} \\
0 & I
\end{pmatrix}.
\]

transform under the Lorentz group standardly i.e. (1) and (7) imply

\[
x_E' = \Lambda x_E,
\]  

(8a)
\[ u'_E = \Lambda u_E. \]  

It holds because \( D(\Lambda, u) = T(u')\Lambda T^{-1}(u) \). Moreover, \( ds^2 = ds_E^2 \), \( \vec{c} = c\vec{n} \), \( u_E^2 = c^2 \) and \( g_L = \text{diag}(+,-,-,-) \). Therefore the CT synchronization scheme defined by the transformations rules \([1]\) is at first glance equivalent to the EP one. In fact, it is a different choice of the convention of the one-way light propagation (see \([2]\)); in other words it is another formulation of special relativity. Notwithstanding, the above statement is true only if we exclude superluminal signals. Indeed, the causality principle, logically independent of the requirement of Lorentz covariance, is not invariant under change of the synchronization \([1]\). It is evident from the form of the boost matrix \([2]\); the coordinate time \( x^0 \) is rescaled by a positive factor \( \gamma \) only. Therefore \( \varepsilon(dx^0) \) is an invariant of \([1]\) and this factor allow us to introduce an absolute notion of causality, generalizing the EP causality. Consequently, as was shown in \([2]\), all inconsistencies of the standard formalism, related to the superluminal propagation, disappear in this formulation of special relativity.

If we exclude tachyons then, as was mentioned above, physics cannot depend of synchronization. Thus in this case any inertial frame can be chosen as the preferred frame, determining a concrete CT synchronization. This statement is in fact the relativity principle articulated in the CT synchronization language.

What happens, when tachyons do exist? In such a case the relativity principle is obviously broken: If tachyons exist then only one inertial frame must be a true privileged frame. Therefore, in this case, the EP synchronization is inadequate to description of reality; we must choose the synchronization defined by \([1]\) \([1]\). Moreover the relativity principle is evidently broken in this case as well as the conventionality thesis: The one-way velocity of light becomes (a priori) a really measured quantity.

To formalize the above analysis, in \([1]\) it was introduced notion of the synchronization group \( L_S \). It connects different synchronizations of the CT–type and it is isomorphic to the Lorentz group:

\[ x' = T(u'T^{-1}(u))x = D(\Lambda_S, u)T(u)\Lambda_S^{-1}T^{-1}(u)x, \]  

\[ u' = D(\Lambda_S, u)u, \]  

with \( \Lambda_S \in L_S \).

For clarity we write the composition of transformations of the Poincaré group \( L_K T^4 \) and the synchronization group \( L_S \) in the EP coordinates

\[ x'_E = \Lambda(x_E + a_E), \]  

\[ u'_E = \Lambda_S \Lambda u_E. \]  

Therefore, in a natural way, we can select three subgroups:

\[ L = \{(I, \Lambda)\}, \quad L_S = \{(\Lambda_S, I)\}, \quad L_0 = \{(\Lambda_0, \Lambda_0^{-1})\}. \]

By means of \([1]\) it is easy to check that \( L_0 \) and \( L_S \) commute. Therefore the set \( \{(\Lambda_S, I)\} \) is simply the direct product of two Lorentz groups \( L_0 \otimes L_S \). The intersystemic Lorentz group \( L \) is the diagonal subgroup in this direct product. From the composition law \([1]\) it follows that \( L \) acts as an authomorphism group of \( L_S \).

Now, the synchronization group realizes in fact the relativity principle: If we exclude tachyons then transformations of \( L_S \) are canonical ones. On the other hand, if we include tachyons then the synchronization group \( L_S \) is broken to the \( SO(3)_u \) subgroup of \( L_S \); \( SO(3)_u \) is the stability group of \( u^a \). In fact, transformations from the \( L_S/SO(3)_u \) do not leave the absolute notion of causality invariant. On the quantum level \( L_S \) is broken down to \( SO(3)_u \) subgroup i.e. transformations from \( L_S/SO(3)_u \) cannot be realized by unitary operators \([1]\).

### III. QUANTIZATION

The following two facts, true only in CT synchronization, are extremely important for quantization of tachyons:

- Invariance of the sign of the time component of the space-like fourmomentum i.e. \( \varepsilon(k^0) = inv \),

- Existence of a covariant lower energy bound.

This is the reason why an invariant Fock construction can be done in our case \([3]\). In the paper \([3]\) it was constructed a quantum free field theory for scalar tachyons. Here we classify unitary Poincaré mappings in the bundle of Hilbert spaces \( H_u \) for a space-like fourmomentum. Furthermore we find the corresponding canonical commutation relations. As result we obtain that tachyons correspond to unitary mappings which are induced from \( SO(2) \) group rather than \( SO(2,1) \) one. Of course, a classification of unitary representations for time-like and light-like fourmomentum is the same as in EP synchronization; this holds because the relativity principle is working in this case.
A. Tachyonic representations

As usually, we assume that a basis in a Hilbert space $H_u$ (fibre) of one-particle states consists of the eigenvectors $|k, u; \ldots\rangle$ of the fourmomentum operators $P^\mu$ namely

$$P^\mu |k, u; \ldots\rangle = k^\mu |k, u; \ldots\rangle$$

(11)

where

$$(k', u; \ldots |k, u; \ldots\rangle = 2k_0^0 \delta^3(k' - k)$$

(12)
i.e. we adopt a covariant normalization. The $k_0^0 = g^0\mu k^\mu_+$ is positive and $k^\mu_+$ is the corresponding solution of the dispersion relation

$$k^2 \equiv g_{\mu\nu} k^\mu k^\nu = -\kappa^2.$$  

(13)

Namely

$$k_{0+} = -\frac{\sigma}{c} k + \gamma_0^2 \omega_k$$

(14)

with

$$\omega_k = \gamma_0^{-2} \sqrt{\left(\frac{\sigma k}{c}\right)^2 + (|k|^2 - \kappa^2) \gamma_0^2}.$$  

(15)

Notice that $k^0_+ = \omega_k$ and the range of the covariant momentum $k_\perp$ is determined by the following inequality

$$|\vec{k}| \geq \kappa \left(1 + (\gamma_0^2 - 1) \left(\frac{\sigma k}{|\sigma| |k|}\right)^2\right)^{-1/2},$$  

(16)
i.e. values of $\vec{k}$ lie outside the oblate spheroid with halfaxes $a = \kappa$ and $b = \kappa \gamma_0^{-1}$. The covariant normalization is possible because in CT synchronization the sign of $k^0$ is an invariant. Thus we have no problem with an indefinite norm in $H_u$.

Now, $ku \equiv k_\mu u^\mu$ is an additional invariant. Indeed, because the transformations of $L_S$ are restricted to $SO(3)_u$ subgroup by causality requirement, and $SO(3)_u$ does not change $u$ nor $k$, our covariance group reduces to the Poincaré mappings (realized in the CT synchrony). Summarizing, irreducible family of unitary operators $U(\Lambda, a)$ in the bundle of Hilbert spaces $H_u$ acts on an orbit defined by the following covariant conditions

- $k^2 = -\kappa^2$;
- $\varepsilon(k^0) = \text{inv}$; for physical representations $\varepsilon(k^0) = 1$ which guarantee a covariant lower bound of energy [3].
- $q \equiv \frac{uk}{c} = \text{inv}$; it is easy to see that $q$ is an energy of tachyon measured in the privileged frame.

As a consequence there exists an invariant, positive definite measure

$$d\mu(k, \kappa, q) = d^4k \theta(k^0) \delta(k^2 + \kappa^2) \delta(q - \frac{uk}{c})$$  

(17)
in a Hilbert space of wave packets.

Let us return to the problem of classification of irreducible unitary mappings $U(\Lambda, a)$:

$$U(\Lambda, a) |k, u; \ldots\rangle = |k', u'; \ldots\rangle;$$

Notice that we have contravariant as well as covariant fourmomenta related by $g_{\mu\nu}$; the physical energy and momentum are covariant because they are generators of translations.

In fact $SO(3)_u$ acts like an $SO(3)$ intrinsic symmetry!
Here the pair \((k, u)\) is transported along trajectories belonging to an orbit fixed by the above mentioned invariant conditions. To follow the familiar Wigner procedure of induction one should find a stability group of the double \((k, u)\). To do this, let us transform \((k, u)\) to the preferred frame by the Lorentz boost \(L_u^{-1}\). Next, in the privileged frame, we rotate the spatial part of the fourmomentum to the \(z\)-axis by an appropriate rotation \(R^{−1}_\pi\). As a result, we obtain the pair \((k, u)\) transformed to the pair \((\tilde{k}, \tilde{u})\) with

\[
\tilde{k} = \begin{pmatrix}
q \\
0 \\
0 \\
\sqrt{k^2 + q^2}
\end{pmatrix}, \quad \tilde{u} = \begin{pmatrix}
c \\
0 \\
0 \\
0
\end{pmatrix}.
\]

It is easy to see that the stability group of \((k, u)\) is the \(SO(2) = SO(2,1) \cap SO(3)\) group. Thus tachyonic unitary representations should be induced from the \(SO(2)\) instead of \(SO(2,1)\) group! Recall that unitary representations of the \(SO(2,1)\) noncompact group are infinite dimensional (except of the trivial one). As a consequence, local fields was necessarily infinite component ones (except of the scalar one). On the other hand, in the CT synchronization case unitary representations for space-like fourmomenta in our bundle of Hilbert spaces are induced from irreducible, one dimensional representations of \(SO(2)\) in a close analogy with a light-like fourmomentum case. They are labelled by helicity \(\lambda\), by \(\kappa\) and by \(q\) \((\varepsilon(k^0) = \varepsilon(q)\) is determined by \(q\); of course a physical choice is \(\varepsilon(q) = 1\).

Now, by means of the familiar Wigner trick we determine the Lorentz group action on the base vectors; namely

\[
U(\Lambda) \vert k, u, \kappa, \lambda, q \rangle = e^{i\lambda \varphi(\Lambda; k, u)} \vert k', u'; \kappa, \lambda, q \rangle
\]

where

\[
e^{i\lambda \varphi(\Lambda; k, u)} = U \left( R^{-1}_{\Omega \Lambda} \right)
\]

with

\[
\Omega = L_{u'} \Lambda L_u.
\]

Here \(k\) and \(u\) transform according to the law [1]. The rotation \(R_\pi\) connects \(\tilde{k}\) with \(D(L_u^{-1}, u)k\), i.e.

\[
R_\pi \tilde{k} = D(L_u^{-1}, u)k.
\]

It is easy to check that \(R^{-1}_{\Omega \Lambda} \Omega R_\pi\) is a Wigner-like rotation belonging to the stability group \(SO(2)\) of \((k, u)\) and determines the phase \(\varphi\). By means of standard topological arguments \(\lambda\) can take integer or halfinteger values only i.e. \(\lambda = 0, \pm 1/2, \pm 1, \ldots\).

Now, the orthogonality relation [12] reads

\[
\langle k', u'; \kappa', \lambda', q' \vert k, u, \kappa, \lambda, q \rangle = 2\omega_k \delta^3(k' - k) \delta_{\lambda', \lambda}.
\]

### B. Canonical quantization

Following the Fock procedure, we define canonical commutation relations

\[
[a_\lambda(k_+, u), a_\tau(p_+, u)]_{\pm} = [a_\lambda^\dagger(k_+, u), a_\tau^\dagger(p_+, u)]_{\pm} = 0,
\]

\[
[a_\lambda(k_+, u), a_\tau^\dagger(p_+, u)]_{\pm} = 2\omega_k \delta(k - p) \delta_{\lambda, \tau},
\]

where \(-\) or \(\) means the commutator or anticommutator and corresponds to the bosonic (\(\lambda\) integer) or fermionic (\(\lambda\) halfinteger) case respectively. Furthermore, we introduce a Poincaré invariant vacuum \(|0\rangle\) defined by

\[
\langle 0 | 0 \rangle = 1 \quad \text{and} \quad a_\lambda(k_+, u) \langle 0 | 0 \rangle = 0.
\]

Therefore the one particle states

\[
a_\lambda^\dagger(k_+, u) |0\rangle
\]
are the base vectors belonging to an orbit in our bundle of Hilbert spaces iff

\[ U(\Lambda)a\lambda^\dagger(k_+, u)U(\Lambda^{-1}) = e^{i\lambda\varphi(\Lambda, k, u)}a\lambda^\dagger(k_+', u'), \]

(27a)

\[ U(\Lambda)a\lambda(k_+, u)U(\Lambda^{-1}) = e^{-i\lambda\varphi(\Lambda, k, u)}a\lambda(k_+', u'), \]

(27b)

and

\[ [P_\mu, a\lambda^\dagger(k_+, u)]_\mp = k^\dagger_\mu a\lambda^\dagger(k_+, u), \]

(28)

Notice that

\[ P_\mu = \int d^4k \theta(k^0) \delta(k^2 + \kappa^2) k_\mu \left( \sum_\lambda a\lambda^\dagger(k, u)a\lambda(k, u) \right) \]

(29)

is a solution of (28).

Finally we can deduce also the form of the helicity operator:

\[ \hat{\lambda} = -\frac{W^\mu u_\mu}{c\sqrt{(Pu)^2 - P^2}} \]

(30)

where

\[ W^\mu = \frac{1}{2}\varepsilon^{\mu\sigma\lambda\tau} J_\sigma a_\lambda P_\tau \]

is the Pauli-Lubanski fourvector.

The next step is to construct local free fields and the corresponding field equations. An example of such construction was presented in [6] for a free scalar tachyonic field. Here we give a local field equation for Dirac tachyons; it can be possibly associated with neutrinos

\[ -i\partial_\mu \gamma^\mu + \frac{1}{2\xi} \left( \kappa^2 + \left( \frac{u^\mu}{c} \partial_\mu + \xi \right)^2 \right) \frac{u^\mu \gamma^\mu}{c} - \frac{1}{2\xi} \left( \kappa^2 + \left( \frac{u^\mu}{c} \partial_\mu + \xi \right) \left( \frac{iu^\mu}{c} \partial_\mu - \xi \right) \right) I \psi(x) = 0 \]

(31)

with \( \xi \) is a real parameter of inverse of length dimensionality; here \( \psi(x) \) is a Dirac bispinor, whereas

\[ \gamma^\mu = T(u)^\mu_\nu \gamma^\nu, \]

i.e.,

\[ \{\gamma^\mu, \gamma^\nu\} = 2g^{\mu\nu}(u)I \]

(32)

From (31) it follows the Klein-Gordon equation

\[ (g^{\mu\nu}(u)\partial_\mu \partial_\nu - \kappa^2) \psi = 0, \]

(33)

so in the momentum representation we have the tachyonic dispersion relation

\[ k^2 = -\kappa^2. \]

(34)

Recall that in the standard approach it is impossible to introduce a mass term without breaking of hermicity of the Lagrangian describing fermionic tachyon.

Now, the helicity operator (30) takes the form

\[ \hat{\lambda} = \frac{1}{4c\sqrt{(Pu)^2 - P^2}} \gamma^5 [P\gamma, u\gamma], \]

(35)

Because \( \hat{\lambda}^2 = \frac{1}{2}I \) and \( \text{Tr} \hat{\lambda} = 0 \), the spectrum of \( \hat{\lambda} \) equals \( \{-\frac{1}{2}, -\frac{1}{2}, \frac{1}{2}, \frac{1}{2}\} \). Furthermore \( \hat{\lambda} \) commutes with the Dirac operator defined by (31). This allows us to construct solutions of (31) by means of standard procedure, i.e., by using the projection operators constructed from helicity and Dirac operators.
IV. CONCLUSIONS

The main result of this work is that tachyons are classified according to the unitary representations of $SO(2)$ rather than $SO(2,1)$ group. Together with the fact that they can be consistently described under some appropriate choice of synchronization, this shows that there are no serious theoretical obstructions to interpret the experimental data about square of mass of neutrinos as a signal that they can be fermionic tachyons.

A more exhaustive discussion of the Dirac-like equation for fermionic tachyon will be given in the forthcoming article.

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