Feynman’s Decoherence

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Abstract

Gell-Mann’s quarks are coherent particles confined within a hadron at rest, but Feynman’s partons are incoherent particles which constitute a hadron moving with a velocity close to that of light. It is widely believed that the quark model and the parton model are two different manifestations of the same covariant entity. If this is the case, the question arises whether the Lorentz boost destroys coherence. It is pointed out that this is not the case, and it is possible to resolve this puzzle without inventing new physics. It is shown that this decoherence is due to the measurement processes which are less than complete.

I. INTRODUCTION

The superposition principle is the backbone of quantum mechanics. We are thus led to think the transition from quantum to classical mechanics is accompanied by a loss of coherence. The transition from classical to quantum mechanics thus requires that particles somehow gain coherence. Indeed, these have been outstanding problems ever since the present form of quantum mechanics was formulated in the late 1920s.

We are quite familiar with coherent and incoherent light waves. If the phases become random, we say that the light waves are incoherent. Thus, we are naturally led to say that the process of randomizing phases is a form of decoherence. During this process, the entropy of the system increases. We know also that the transition from quantum to classical mechanics has a deeper content than randomizing phases. Indeed, we still have to define the precise meaning of the word “decoherence.” In the meantime, we have to review critically the existing processes where quantum coherence disappears.

In this report, we examine a concrete case of Feynman’s parton picture. The question is very simple. According to Gell-Man [1], hadrons like protons and mesons, are bound states

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of quarks capable of quantum excitations like the hydrogen atom. This picture is valid if
the hadrons are in their rest frame or move slowly. If the hadron moves very fast, it appears
as a collection of free independent particles. This observation was made first by Feynman,
who named those free particles as “partons”.

It is a widely accepted view that the quark model and the parton model are two different
manifestations of one covariant entity, with the quark model valid when the hadron is at
rest and the parton model observed for the hadron moving with a speed close to that of
light. Quarks are coherent and partons are incoherent. Then does the Lorentz boost wash
out the coherence as in randomizing phase differences for light waves? No! Because we
are not adding any physical processes. We are simply looking at the same thing from
different Lorentz frames. Therefore, we are led to conclude that there are limitations on our
observation processes.

The parton picture was invented by Feynman based on phenomenology in high-energy
physics. In order to appreciate fully its implications, it is important to study his papers
written on other subjects. We also should take into account the fact that Feynman could
come up with creative physical ideas without mathematical rigor or sometimes with wrong
mathematics. Indeed, in order to understand the parton picture, we need a broader physical
and mathematical base.

The issue is that Feynman’s original parton picture is valid only for fast-moving hadrons.
Is it possible to construct a covariant model which produces the parton model as a limit?
We can ask the same question about Gell-man’s quark model. It is based on non-relativistic
quantum mechanics of bound states. Is there a covariant picture of the quantum bound
state? Indeed, Feynman raised this question and published a paper on this subject with
Kislinger and Ravndal.

Feynman was also interested in the concept of entropy coming from measurement pro-
cesses which are less than complete. This subject was of course originated by von Neu-
mann in his classic book on mathematical foundations of quantum mechanics. There
are many different forms of measurement theory, and they are currently debated in the
literature in connection with the EPR problems and Bell’s inequality. The letter ”E” in
EPR stands for Einstein. Then we should take into account another contribution made by
Einstein, namely Lorentz covariance. It is therefore essential to study the entropy effects
when we deal with the covariance, because our the measurement theory for non-relativistic
quantum mechanics is necessarily incomplete if we introduce another variable, namely the
time variable.

We shall consider all of the above items in this report. We shall start with Feynman’s
research style in Sec. II. Feynman was quite fond of using harmonic oscillators to find a
clue to new physics. In Sec. III, we use coupled harmonic oscillators to illustrate Feynman’s
physical ideas. The oscillators are useful in illustrating Feynman’s rest of the universe and
Feynman’s relativistic oscillators.

In Sec. IV, it is noted that the time-separation variable does not exist in the present
form of quantum mechanics, while it clearly exists physically wheresoever the space-like
separation, like the Bohr radius, exists if we believe in Einstein’s special relativity. If the
system is Lorentz-boosted, the time-separation variable becomes prominent because it is
linearly mixed with the space-separation variable. If we do not measure this variable, there
will be an increase in entropy.
In Sec. V, we consider a hadron consisting of two quarks bound together by a harmonic oscillator potential. It is shown that the covariant harmonic oscillator formalism gives non-relativistic quantum mechanics in the Lorentz system where the hadron is at rest. It is then shown that the same formalism produces Feynman’s parton model in the frame where the hadron moves with a speed close to that of light. Finally, in Sec. VII we study carefully why the partons appear as incoherent particles. It is shown that it is due to our measurement process which is less than complete in both the rest frame and in the frame in which the hadron moves with a speed close to that of light.

II. FEYNMAN’S WORLD

Feynman was quite fond of using harmonic oscillators to probe new territories of physics. In this section, we examine which oscillator formalism is most suitable to interpret some of Feynman’s papers during the period 1969 – 1972. This formalism should accommodate special relativity and quantum mechanics of extended objects.

Let us start with a simple physical system. Two coupled harmonic oscillators play many important roles in physics. In group theory, it generates symmetry group as rich as \(O(3,3)\) \[12\]. It has many interesting subgroups useful in all branches of physics. The group \(O(3,1)\) is of course essential for studying covariance in special relativity. It is applicable to three space-like variables and one time-like variable. In the harmonic oscillator regime, those three space-like coordinates are separable. Thus, it is possible to separate longitudinal and transverse coordinates. If we leave out the transverse coordinates which do not participate in Lorentz boosts, the only relevant variables are longitudinal and time-like variables. The symmetry group for this case is \(O(1,1)\) easily derivable from the Hamiltonian for the two-oscillator system.

It is widely known that this simple mathematical device is the basic language for two-photon coherent states known as squeezed states of light \[13,14\]. However, this \(O(1,1)\) device plays a much more powerful role in physics. According to Feynman, the adventure of our science of physics is a perpetual attempt to recognize that the different aspects of nature are really different aspects of the same thing \[15\]. Feynman wrote many papers on different subjects of physics, but they are coming from one paper according to him. We are not able to combine all of his papers, but we can consider three of his papers published during the period 1969-72.

In this paper, we would like to consider Feynman’s 1969 report on partons \[2\], the 1971 paper he published with his students on the quark model based on harmonic oscillators \[3\], and the chapter on density matrix in his 1972 book on statistical mechanics \[4\]. In these three different papers, Feynman deals with three distinct aspects of nature. We shall see whether Feynman was saying the same thing in these papers. For this purpose, we shall use the \(O(1,1)\) symmetry derivable directly from the Hamiltonian for two coupled oscillators \[16\]. The standard procedure for this two-oscillator system is to separate the Hamiltonian using normal coordinates. The transformation to the normal coordinate system becomes very simple if the two oscillators are identical. We shall use this simple mathematics to find a common ground for the above-mentioned articles written by Feynman.

First, let us look at Feynman’s book on statistical mechanics \[4\]. He makes the following
statement about the density matrix. When we solve a quantum-mechanical problem, what we really do is divide the universe into two parts - the system in which we are interested and the rest of the universe. We then usually act as if the system in which we are interested comprised the entire universe. To motivate the use of density matrices, let us see what happens when we include the part of the universe outside the system.

In order to see clearly what Feynman had in mind, we use the above-mentioned coupled oscillators. One of the oscillators is the world in which we are interested and the other oscillator as the rest of the universe. There will be no effects on the first oscillator if the system is decoupled. Once coupled, we need a normal coordinate system in order separate the Hamiltonian. Then it is straightforward to write down the wave function of the system. Then the mathematics of this oscillator system is directly applicable to Lorentz-boosted harmonic oscillator wave functions, where one variable is the longitudinal coordinate and the other is the time variable. The system is uncoupled if the oscillator wave function is at rest, but the coupling becomes stronger as the oscillator is boosted to a high-speed Lorentz frame.

We shall then note that for two-body system, such as the hydrogen atom, there is a time-separation variable which is to be linearly mixed with the longitudinal space-separation variable. This space-separation variable is known as the Bohr radius, but we never talk about the time-separation variable in the present form of quantum mechanics, because this time-separation variable belongs to Feynman’s rest of the universe.

If we pretend not to know this time-separation variable, the entropy of the system will increase when the oscillator is boosted to a high-speed system. Does this increase in entropy correspond to decoherence? Not necessarily. However, in 1969, Feynman observed the parton effect in which a rapidly moving hadron appears as a collection of incoherent partons. This is the decoherence mechanism of current interest.

### III. TWO COUPLED OSCILLATORS

Two coupled harmonic oscillators serve many different purposes in physics. It is well known that this oscillator problem can be formulated into a problem of a quadratic equation in two variables. To make a long story short, let us consider a system of two identical oscillators coupled together by a spring. The Hamiltonian is

$$H = \frac{1}{2m} \{p_1^2 + p_2^2\} + \frac{1}{2} \left\{ K \left( x_1^2 + x_2^2 \right) + 2Cx_1x_2 \right\}. \tag{1}$$

We are now ready to decouple this Hamiltonian by making the coordinate rotation:

$$y_1 = \frac{1}{\sqrt{2}} (x_1 - x_2), \quad y_2 = \frac{1}{\sqrt{2}} (x_1 + x_2). \tag{2}$$

In terms of this new set of variables, the Hamiltonian can be written as

$$H = \frac{1}{2m} \{p_1^2 + p_2^2\} + \frac{K}{2} \left\{ e^{2\eta}y_1^2 + e^{-2\eta}y_2^2 \right\}, \tag{3}$$

with
The world in which we do physics

FIG. 1. Feynman’s rest of the universe. The vertical axis is not measurable, and thus belongs to the rest of the universe. The horizontal axis corresponds to the world in which we do physics.

\[ \exp (\eta) = \sqrt{\frac{K + C}{K - C}}. \] (4)

Thus \( \eta \) measures the strength of the coupling. If \( y_1 \) and \( y_2 \) are measured in units of \((mK)^{1/4}\), the ground-state wave function of this oscillator system is

\[ \psi_\eta(x_1, x_2) = \left( \frac{1}{\pi} \right)^{1/2} \exp \left\{ -\frac{1}{2} (e^{\eta} y_1^2 + e^{-\eta} y_2^2) \right\}. \] (5)

The wave function is separable in the \( y_1 \) and \( y_2 \) variables. However, for the variables \( x_1 \) and \( x_2 \), the story is quite different.

The key question is how quantum mechanical calculations in the world of the observed variable are affected when we average over the other variable. The \( x_2 \) space in this case corresponds to Feynman’s rest of the universe, if we only consider quantum mechanics in the \( x_1 \) space. As was discussed in the literature for several different purposes \[14,8\], the wave function of Eq. (5) can be expanded as

\[ \psi_\eta(x_1, x_2) = \frac{1}{\cosh \eta} \sum_k \left( \frac{\tanh \eta}{2} \right)^k \phi_k(x_1) \phi_k(x_2). \] (6)

This expansion corresponds to the two-photon coherent states in Yuen’s paper \[13\], and the wave function of Eq. (5) is a squeezed wave function \[14\].

The question then is what lessons we can learn from the situation in which we average over the \( x_2 \) variable. In order to study this problem, we use the density matrix. From this wave function, we can construct the pure-state density matrix

\[ \rho(x_1, x_2; x_1', x_2') = \psi_\eta(x_1, x_2) \psi_\eta(x_1', x_2'). \] (7)

If we are not able to make observations on the \( x_2 \), we should take the trace of the \( \rho \) matrix with respect to the \( x_2 \) variable. Then the resulting density matrix is
\[ \rho(x, x') = \int \psi_\eta(x, x_2) \{ \psi_\eta(x', x_2) \}^* dx_2. \]  

(8)

We have simplicity replaced \( x_1 \) and \( x'_1 \) by \( x \) and \( x' \) respectively. If we perform the integral of Eq.(8), the result is

\[ \rho(x, x') = \left( \frac{1}{\cosh(\eta/2)} \right)^2 \sum_k \left( \tanh \frac{\eta}{2} \right)^{2k} \phi_k(x) \phi_k^*(x'), \]  

(9)

which leads to \( Tr(\rho) = 1 \). It is also straightforward to compute the integral for \( Tr(\rho^2) \). The calculation leads to

\[ Tr(\rho^2) = \left( \frac{1}{\cosh(\eta/2)} \right)^4 \sum_k \left( \tanh \frac{\eta}{2} \right)^{4k}. \]  

(10)

The sum of this series is \((1/\cosh \eta)\), which is smaller than one if the parameter \( \eta \) does not vanish.

This is of course due to the fact that we are averaging over the \( x_2 \) variable which we do not measure. The standard way to measure this ignorance is to calculate the entropy defined as

\[ S = -Tr(\rho \ln(\rho)), \]  

(11)

where \( S \) is measured in units of Boltzmann’s constant. If we use the density matrix given in Eq.(9), the entropy becomes

\[ S = \cosh^2 \left( \frac{\eta}{2} \right) \ln \left( \cosh \frac{\eta}{2} \right)^2 - \sinh^2 \left( \frac{\eta}{2} \right) \ln \left( \sinh \frac{\eta}{2} \right)^2. \]  

(12)

This expression can be translated into a more familiar form if we use the notation

\[ \tanh \frac{\eta}{2} = \exp \left( -\frac{\hbar \omega}{kT} \right), \]  

(13)

where \( \omega \) is given in Eq.(4) [17].

It is known in the literature that this rise in entropy and temperature causes the Wigner function to spread wide in phase space causing an increase of uncertainty [16]. Certainly, we cannot reach a classical limit by increasing the uncertainty. On the other hand, we are accustomed to think this entropy increase has something to do with decoherence, and we are also accustomed to think the lack of coherence has something to do with a classical limit. Are they compatible? We thus need a new vision in order to define precisely the word “decoherence.”

IV. TIME-SEPARATION VARIABLE IN FEYNMAN’S REST OF THE UNIVERSE

Quantum field theory has been quite successful in terms of perturbation techniques in quantum electrodynamics. However, this formalism is basically based on the S matrix for
scattering problems and useful only for physically processes where a set of free particles becomes another set of free particles after interaction. Quantum field theory does not address the question of localized probability distributions and their covariance under Lorentz transformations. The Schrödinger quantum mechanics of the hydrogen atom deals with localized probability distribution. Indeed, the localization condition leads to the discrete energy spectrum. Here, the uncertainty relation is stated in terms of the spatial separation between the proton and the electron. If we believe in Lorentz covariance, there must also be a time separation between the two constituent particles.

Before 1964 [4], the hydrogen atom was used for illustrating bound states. These days, we use hadrons which are bound states of quarks. Let us use the simplest hadron consisting of two quarks bound together with an attractive force, and consider their space-time positions \( x_a \) and \( x_b \), and use the variables

\[
X = \frac{(x_a + x_b)}{2}, \quad x = \frac{(x_a - x_b)}{2\sqrt{2}}.
\] (14)

The four-vector \( X \) specifies where the hadron is located in space and time, while the variable \( x \) measures the space-time separation between the quarks. According to Einstein, this space-time separation contains a time-like component which actively participates as can be seen from

\[
\begin{pmatrix}
\cosh \eta \\
\sinh \eta
\end{pmatrix}
\begin{pmatrix}
z' \\
t'
\end{pmatrix} =
\begin{pmatrix}
\cosh \eta & \sinh \eta \\
\sinh \eta & \cosh \eta
\end{pmatrix}
\begin{pmatrix}
z \\
t
\end{pmatrix},
\] (15)

when the hadron is boosted along the \( z \) direction. In terms of the light-cone variables defined as [18]

\[
u = (z + t)/\sqrt{2}, \quad v = (z - t)/\sqrt{2}.
\] (16)

The boost transformation of Eq. (15) takes the form

\[
u' = e^\eta \nu, \quad v' = e^{-\eta} v.
\] (17)

The \( u \) variable becomes expanded while the \( v \) variable becomes contracted.

Does this time-separation variable exist when the hadron is at rest? Yes, according to Einstein. In the present form of quantum mechanics, we pretend not to know anything about this variable. Indeed, this variable belongs to Feynman’s rest of the universe. In this report, we shall see the role of this time-separation variable in decoherence mechanism.

Also in the present form of quantum mechanics, there is an uncertainty relation between the time and energy variables. However, there are no known time-like excitations. Unlike Heisenberg’s uncertainty relation applicable to position and momentum, the time and energy separation variables are c-numbers, and we are not allowed to write down the commutation relation between them. Indeed, the time-energy uncertainty relation is a c-number uncertainty relation [19], as is illustrated in Fig. 3.

How does this space-time asymmetry fit into the world of covariance [4]. This question was studied in depth by the present authors. The answer is that Wigner’s \( O(3) \)-like little group is not a Lorentz-invariant symmetry, but is a covariant symmetry [3]. It has been shown that the time-energy uncertainty applicable to the time-separation variable fits perfectly into the \( O(3) \)-like symmetry of massive relativistic particles [8].
FIG. 2. Lorentz boost in the light-cone coordinate system.

FIG. 3. Space-time picture of quantum mechanics. There are quantum excitations along the space-like longitudinal direction, but there are no excitations along the time-like direction. The time-energy relation is a c-number uncertainty relation.
The c-number time-energy uncertainty relation allows us to write down a time distribution function without excitations \[8\]. If we use Gaussian forms for both space and time distributions, we can start with the expression

\[
\left(\frac{1}{\pi}\right)^{1/2} \exp\left\{-\frac{1}{2} \left(z^2 + t^2\right)\right\}
\]

for the ground-state wave function. What do Feynman et al. say about this oscillator wave function?

In his classic 1971 paper \[3\], Feynman et al. start with the following Lorentz-invariant differential equation.

\[
\frac{1}{2} \left\{x^2 - \frac{\partial^2}{\partial x^2}\right\} \psi(x) = \lambda \psi(x).
\]

This partial differential equation has many different solutions depending on the choice of separable variables and boundary conditions. Feynman et al. insist on Lorentz-invariant solutions which are not normalizable. On the other hand, if we insist on normalization, the ground-state wave function takes the form of Eq.(18). It is then possible to construct a representation of the Poincaré group from the solutions of the above differential equation \[8\].

If the system is boosted, the wave function becomes

\[
\psi_\eta(z, t) = \left(\frac{1}{\pi}\right)^{1/2} \exp\left\{-\frac{1}{2} \left(e^{-2\eta u^2} + e^{2\eta v^2}\right)\right\}.
\]

This wave function becomes Eq.(18) if \(\eta\) becomes zero. The transition from Eq.(18) to Eq.(20) is a squeeze transformation. The wave function of Eq.(18) is distributed within a circular region in the \(uv\) plane, and thus in the \(zt\) plane. On the other hand, the wave function of Eq.(20) is distributed in an elliptic region with the light-cone axes as the major and minor axes respectively. If \(\eta\) becomes very large, the wave function becomes concentrated along one of the light-cone axes. Indeed, the form given in Eq.(20) is a Lorentz-squeezed wave function. This squeeze mechanism is illustrated in Fig. 4.

It is interesting to note that the Lorentz-invariant differential equation of Eq.(19) contains the time-separation variable which belongs to Feynman’s rest of the universe. Furthermore, the wave function of Eq.(18) is identical to that of Eq.(3) for the coupled oscillator system, if the variables \(z\) and \(t\) are replaced \(x_1\) and \(x_2\) respectively. Thus the entropy increase due to the unobservable \(x_2\) variable is applicable to the unobserved time-separation variable \(t\).

**V. FEYNMAN’S PARTON PICTURE**

It is a widely accepted view that hadrons are quantum bound states of quarks having localized probability distribution. As in all bound-state cases, this localization condition is responsible for the existence of discrete mass spectra. The most convincing evidence for this bound-state picture is the hadronic mass spectra which are observed in high-energy laboratories \[3,8\]. However, this picture of bound states is applicable only to observers in the Lorentz frame in which the hadron is at rest. How would the hadrons appear to observers in
other Lorentz frames? To answer this question, can we use the picture of Lorentz-squeezed hadrons discussed in Sec. IV.

The radius of the proton is $10^{-5}$ of that of the hydrogen atom. Therefore, it is not unnatural to assume that the proton has a point charge in atomic physics. However, while carrying out experiments on electron scattering from proton targets, Hofstadter in 1955 observed that the proton charge is spread out [20]. In this experiment, an electron emits a virtual photon, which then interacts with the proton. If the proton consists of quarks distributed within a finite space-time region, the virtual photon will interact with quarks which carry fractional charges. The scattering amplitude will depend on the way in which quarks are distributed within the proton. The portion of the scattering amplitude which describes the interaction between the virtual photon and the proton is called the form factor.

Although there have been many attempts to explain this phenomenon within the framework of quantum field theory, it is quite natural to expect that the wave function in the quark model will describe the charge distribution. In high-energy experiments, we are dealing with the situation in which the momentum transfer in the scattering process is large. Indeed, the Lorentz-squeezed wave functions lead to the correct behavior of the hadronic form factor for large values of the momentum transfer [21].

Furthermore, in 1969, Feynman observed that a fast-moving hadron can be regarded as a collection of many “partons” whose properties do not appear to be quite different from those of the quarks [2]. For example, the number of quarks inside a static proton is three, while the number of partons in a rapidly moving proton appears to be infinite. The question then is how the proton looking like a bound state of quarks to one observer can appear different to an observer in a different Lorentz frame? Feynman made the following systematic observations.

a. The picture is valid only for hadrons moving with velocity close to that of light.
b. The interaction time between the quarks becomes dilated, and partons behave as free independent particles.

c. The momentum distribution of partons becomes widespread as the hadron moves fast.

d. The number of partons seems to be infinite or much larger than that of quarks.

Because the hadron is believed to be a bound state of two or three quarks, each of the above phenomena appears as a paradox, particularly b) and c) together.

In order to resolve this paradox, let us write down the momentum-energy wave function corresponding to Eq.\((20)\). If the quarks have the four-momenta \(p_a\) and \(p_b\), we can construct two independent four-momentum variables \[ P = p_a + p_b, \quad q = \sqrt{2}(p_a - p_b), \] (21)
where \(P\) is the total four-momentum and is thus the hadronic four-momentum. \(q\) measures the four-momentum separation between the quarks. Their light-cone variables are
\[ q_u = (q_0 - q_z)/\sqrt{2}, \quad q_v = (q_0 + q_z)/\sqrt{2}. \] (22)

The resulting momentum-energy wave function is
\[ \phi_\eta(q_z, q_0) = \left(\frac{1}{\pi}\right)^{1/2} \exp\left\{-\frac{1}{2} \left(e^{-2\eta q_u^2} + e^{2\eta q_v^2}\right)\right\}. \] (23)

Because we are using here the harmonic oscillator, the mathematical form of the above momentum-energy wave function is identical to that of the space-time wave function. The Lorentz squeeze properties of these wave functions are also the same. This aspect of the squeeze has been exhaustively discussed in the literature \[8–10\].

When the hadron is at rest with \(\eta = 0\), both wave functions behave like those for the static bound state of quarks. As \(\eta\) increases, the wave functions become continuously squeezed until they become concentrated along their respective positive light-cone axes. Let us look at the z-axis projection of the space-time wave function. Indeed, the width of the quark distribution increases as the hadronic speed approaches that of the speed of light. The position of each quark appears widespread to the observer in the laboratory frame, and the quarks appear like free particles.

The momentum-energy wave function is just like the space-time wave function, as is shown in Fig.\(\text{5}\). The longitudinal momentum distribution becomes wide-spread as the hadronic speed approaches the velocity of light. This is in contradiction with our expectation from nonrelativistic quantum mechanics that the width of the momentum distribution is inversely proportional to that of the position wave function. Our expectation is that if the quarks are free, they must have their sharply defined momenta, not a wide-spread distribution.

However, according to our Lorentz-squeezed space-time and momentum-energy wave functions, the space-time width and the momentum-energy width increase in the same direction as the hadron is boosted. This is of course an effect of Lorentz covariance. This indeed is the key to the resolution of the quark-parton paradox \[8–10\].
FIG. 5. Lorentz-squeezed space-time and momentum-energy wave functions. As the hadron’s speed approaches that of light, both wave functions become concentrated along their respective positive light-cone axes. These light-cone concentrations lead to Feynman’s parton picture.

VI. PARTONS AS INCOHERENT PARTICLES

The most puzzling problem in the parton picture is that partons in the hadron appear as incoherent particles, while quarks are coherent when the hadron is at rest. Does this mean that the coherence is destroyed by the Lorentz boost? The answer is NO, and here is the resolution to this puzzle.

When the hadron is boosted, the hadronic matter becomes squeezed and becomes concentrated in the elliptic region along the positive light-cone axis, as is illustrated in Figs. 4 and 5. The length of the major axis becomes expanded by $e^\eta$, and the minor axis is contracted by $e^{-\eta}$.

This means that the interaction time of the quarks among themselves become dilated. Because the wave function becomes wide-spread, the distance between one end of the harmonic oscillator well and the other end increases. This effect, first noted by Feynman \[2\], is universally observed in high-energy hadronic experiments. The period is oscillation increases like $e^\eta$.

On the other hand, the interaction time with the external signal, since it is moving in the direction opposite to the direction of the hadron, travels along the negative light-cone axis, as illustrated in Fig. 6. If the hadron contracts along the negative light-cone axis, the interaction time decreases by $e^{-\eta}$. The ratio of the interaction time to the oscillator period becomes $e^{-2\eta}$. The energy of each proton coming out of the Fermilab accelerator is 900 GeV. This leads the ratio to $10^{-6}$. This is indeed a small number. The external signal is not able
FIG. 6. Quarks interact among themselves and with external signal. The interaction time of the quarks among themselves become dilated, as the major axis of this ellipse indicates. On the other hand, the external signal, since it is moving in the direction opposite to the direction of the hadron, travels along the negative light-cone axis. To the external signal, if it moves with velocity of light, the hadron appears very thin, and the quark's interaction time with the external signal becomes very small.

to sense the interaction of the quarks among themselves inside the hadron.

Indeed, Feynman's parton picture is one concrete physical example where the decoherence effect is observed. As for the entropy, the time-separation variable belongs to the rest of the universe. Because we are not able to observe this variable, the entropy increases as the hadron is boosted to exhibit the parton effect. The decoherence is thus accompanied by an entropy increase.

Let us go back to the coupled-oscillator system. The light-cone variables in Eq. (20) correspond to the normal coordinates in the coupled-oscillator system given in Eq. (2). According to Feynman’s parton picture, the decoherence mechanism is determined by the ratio of widths of the wave function along the two normal coordinates.

VII. MEASUREMENT PROBLEM

In this report, we introduced a covariant model which produced the quark model in the rest frame and the parton model where the hadron moves with a speed close to that of light. Unlike the case of incoherent light arising from random phases, the present case does not involve additional physical processes. The decoherence we discussed here is simply due to a frame change, and looking at the same thing from different places.

Indeed, in the hadronic rest frame, our measurement process ignores the existence of the time-separation variable, and our measurement process is less than complete. When the hadron moves very fast, we are measuring the system through the light-cone coordinate system. Here again, the measurement is made through only one of the two coordinate axes. Again the measurement process is incomplete. Thus we have to conclude that the
decoherence effect observed in Feynman’s parton picture is purely due to our observational process. It has nothing to with dissipation of randomness of phases.
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