Novel type of CPT violation for correlated EPR states

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We discuss modifications to the concept of an "antiparticle", induced by a breakdown of the CPT symmetry at a fundamental level, realized within an extended class of quantum gravity models. The resulting loss of particle-antiparticle identity in the neutral-meson system induces a breaking of the EPR correlation imposed by Bose statistics. The latter is parametrized by a complex parameter controlling the amount of contamination by the "wrong" symmetry state. The physical consequences are studied, and novel observables of CPT-violation in ϕ factories are proposed.

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The CPT theorem is one of the most profound results of quantum field theory [1]. It is a consequence of Lorentz invariance, locality, as well as quantum mechanics (specifically unitary evolution of a system). One implication of CPT invariance is the equality of the masses between particles and antiparticles. In this respect, the best experimental tests of the CPT symmetry so far have been in the neutral Kaon system, where the equality of particle – antiparticle masses has been confirmed to better than one part in 10^{17} [2]. However, this is not the end of the story, given that CPT violation may manifest itself in many subtle ways, thus motivating further experimental searches in various directions.

The possibility of a violation of CPT invariance has been considered in a number of theoretical contexts that go beyond conventional local quantum field theory. In several models of quantum gravity (QG), for example, the axioms of quantum field theory, as well as conventional quantum mechanical behaviour, may not be maintained [3] in the presence of special field configurations, such as wormholes, microscopic (Planck size) black holes, and other topologically non-trivial solitonic objects, such as *geons* [4]. Such configurations are collectively referred to as *space time foam*, a terminology coined by J.A. Wheeler [5], who first conceived the idea that the structure of *quantum* space-time at Planckian scales (10^{-35} m) may actually be fuzzy, characterised by a "foamy" nature. Given that such "objects" cannot be accessible to low-energy observers, it has been argued that a mixed state description must be employed (*QG-induced decoherence*) [3, 6], "tracing" over them in the context of an effective field theory. In the case of microscopic black holes, for example, the decoherence arises due to the loss of information across microscopic event horizons, leading to complications in defining proper asymptotic state-vectors and thus a Heisenberg scattering matrix. As a corollary of this, it has been argued [7] that, in general, CPT invariance in its strong form must be abandoned in quantum gravity. Since in such models the breakdown of the CPT symmetry happens at a fundamental level, it would imply that a proper CPT operator is *ill defined*. This in turn would lead to possible deviations from standard quantum mechanical evolution of states [8], which may not be necessarily associated with the mass difference between particle and antiparticle.

In addition to such effects on the quantum mechanical evolution of a state, however, the violation of CPT invariance leads to a modified concept of what one calls an *antiparticle state*. This is an aspect that has not been discussed previously, and, as we shall argue in the present article, leads to novel observables that could parametrize the CPT violation. Usually the antiparticle is defined as the state with quantum numbers such that, upon interaction with the corresponding particle it produces a state with the quantum numbers of the *vacuum* (annihilation). If the CPT operator is well defined, such a state is obtained by the action of this operator on the corresponding particle state. If, however, the operator is ill defined, the particle and antiparticle spaces should be thought of as *independent* subspaces of matter states. In this case, the usual assumption for *identical states*, when supplemented by particle-antiparticle conjugation, in the case of the electrically neutral mesons K^0 and $\overline{K}^0$ (or B^0 and $\overline{B}^0$), which requires their symmetry under the exchange operator $\mathcal{P}$ as a natural consequence of Bose statistics, *is relaxed*. This, in turn, modifies the description of (neutral) meson entangled states, and may bring about significant deviations to their EPR correlations. The purpose of this paper is to explore these issues, and propose novel CPT-violating observables for the $\phi-$ and $B-$ factories.

In conventional formulations of *entangled* meson states [9, 10, 11] one imposes the requirement of *Bose statistics* for the state $K^0\overline{K}^0$ (or $B^0\overline{B}^0$), which implies that the physical neutral meson-antimeson state must be *symmetric* under the combined operation $C\mathcal{P}$, with C the charge conjugation and $\mathcal{P}$ the operator that permutes the spatial coordinates. Specifically, assuming *conservation* of angular momentum, and a proper existence of the *antiparticle state* (denoted by a bar), one observes that, for $K^0\overline{K}^0$ states which are C-conjugates with $C = (-1)^\ell$ (with ℓ the angular momentum quantum number), the system has to be an eigenstate of $\mathcal{P}$ with eigenvalue $(-1)^\ell$. Hence, for $\ell = 1$, we have that $C = -$, implying $\mathcal{P} = -$. As a consequence of Bose statistics this ensures that for $\ell = 1$ the state of two identical bosons is forbidden [9]. As a result, the initial entangled state $K^0\overline{K}^0$ produced in a ϕ factory can be written as:

$$|i> = \frac{1}{\sqrt{2}}\left(|K^0(\vec{k}), \overline{K}^0(-\vec{k})> - |\overline{K}^0(\vec{k}), K^0(-\vec{k})>\right) \quad (1)$$

This is the starting point of all formalisms known to date, either in the K-system [9, 10] or in the B-system [11], including those [12] where the evolution of the entangled state is described by non-quantum mechanical terms, in the formalism of [6]. In fact, in all these works it has been claimed that the expression in Eq.(1) is actually independent of any assumption about CP, T or CPT symmetries.

However, as has been alluded above, the assumptions leading to Eq.(1) may not be valid if CPT symmetry is violated. In such a case $\overline{K}^0$ cannot be considered as identical to K^0, and thus the requirement of $C\mathcal{P} = +$, imposed by Bose-statistics, is relaxed. As a result, the initial entangled state (1) can be parametrised in general as:

$$\begin{aligned} |i> \ &= \ \frac{1}{\sqrt{2}}\left(|K^0(\vec{k}), \overline{K}^0(-\vec{k})> - |\overline{K}^0(\vec{k}), K^0(-\vec{k})>\right) \\ &+ \ \frac{\omega}{\sqrt{2}}\left(|K^0(\vec{k}), \overline{K}^0(-\vec{k})> + |\overline{K}^0(\vec{k}), K^0(-\vec{k})>\right) \quad (2) \end{aligned}$$

where $\omega = |\omega|e^{i\Omega}$ is a *complex* CPT violating (CPTV) parameter, associated with the non-identical particle nature of the neutral meson and antimeson states. This parameter describes a *novel* phenomenon, not included in previous analyses.

Notice that an equation such as the one given in (2) could also be produced as a result of deviations from the laws of quantum mechanics, during the initial decay of the ϕ or Υ states. Thus, Eq.(2) could receive contributions from two different effects, and can be thought off as simultaneously parametrizing both of them. In the present article we will assume that Eq.(2) arises solely due to deviations from the identical-particle nature of the neutral Kaon and Antikaon states, while the Hamiltonian evolution of the entangled state is governed entirely by the laws of quantum mechanics. Of course, in an actual QG decoherening situation one may have to invoke non-quantum-mechanical, open-system evolution aĺa [6, 8, 12]; however, this lies beyond the scope of the present work, and will be addressed elsewhere.

We now proceed to study the possible consequences of Eq.(2). To this end, we should first express the initial state in terms of CP eigenstates, and also in terms of mass eigenstates, which will be useful when we discuss decays. In terms of CP eigenstates $K_\pm = \frac{1}{\sqrt{2}}\left(|K^0> \pm |\overline{K}^0>\right)$, the initial entangled state (2) reads (for definiteness we concentrate from now on on the ϕ/Kaons case, although our formalism is generic and applies equally to B^0-mesons, *etc*):

$$\begin{aligned} |i> \ &= \ \frac{1}{\sqrt{2}}\left(|K_-(\vec{k}), K_+(-\vec{k})> - |K_+(\vec{k}), K_-(-\vec{k})>\right) \\ &+ \ \frac{\omega}{\sqrt{2}}\left(|K_+(\vec{k}), K_+(-\vec{k})> - |K_-(\vec{k}), K_-(-\vec{k})>\right) (3) \end{aligned}$$

Notice the appearance of K_+K_+ or K_-K_- combinations, as a result of the CPTV parameter ω, which would not exist if the conventional expression (1) had been used instead of (2).

Let us express now (2) in terms of physical (mass) eigenstates, defined as $K_S = \frac{1}{\sqrt{1+|\epsilon_1^2|}}\left(|K_+> + \epsilon_1|K_->\right)$, $K_L = \frac{1}{\sqrt{1+|\epsilon_2^2|}}\left(|K_-> + \epsilon_2|K_+>\right)$, where ϵ_1, ϵ_2 are complex parameters, and such that, if CPT invariance of the Hamiltonian is assumed (within a quantum mechanical framework), $\epsilon_1 = \epsilon_2$, otherwise the quantity $\delta \equiv \epsilon_1 - \epsilon_2$ parametrizes the CPT violation within quantum mechanics. It is convenient to use the CP-violating parameters δ and $\epsilon \equiv |\epsilon|e^{i\phi_\epsilon} = \frac{\epsilon_1+\epsilon_2}{2}$ to parametrize CPT and T violation in a quantum mechanical framework.

In terms of such physical eigenstates, the state (2) is written as (we keep linear terms in the small parameters ω, δ, i.e. in the following we ignore higher-order terms $\omega\delta$, δ^2 *etc.*)

$$\begin{aligned} |i> \ &= \ C\Bigg[\left(|K_S(\vec{k}), K_L(-\vec{k})> - |K_L(\vec{k}), K_S(-\vec{k})>\right) \\ &+ \ \omega\left(|K_S(\vec{k}), K_S(-\vec{k})> - |K_L(\vec{k}), K_L(-\vec{k})>\right)\Bigg] \quad (4) \end{aligned}$$

with $C = \frac{\sqrt{(1+|\epsilon_1|^2)(1+|\epsilon_2|^2)}}{\sqrt{2}(1-\epsilon_1\epsilon_2)} \simeq \frac{1+|\epsilon^2|}{\sqrt{2}(1-\epsilon^2)}$. Notice again the presence of combinations K_SK_S and K_LK_L states, proportional to the novel CPTV parameter ω. As we will see, such terms become important when one considers decay channels.

Specifically, consider the decay amplitude corresponding to the appearance of a final state X at time t_1 and Y at time t_2, as illustrated in fig. 1. One assumes (4) for the initial two-Kaon system, after the ϕ decay. The time is set $t = 0$ at the moment of the decay. Denoting the corresponding amplitude by $A(X, Y)$ we have [9, 10, 11]:

$$A(X, Y) = \langle X|K_S\rangle\langle Y|K_S\rangle\, \mathcal{C}\, (\mathcal{A}_\infty + \mathcal{A}_\in) \quad (5)$$

with

$$\begin{aligned} A_1 \ &= \ e^{-i(\lambda_L+\lambda_S)t/2}[\eta_X e^{-i\Delta\lambda\Delta t/2} - \eta_Y e^{i\Delta\lambda\Delta t/2}] \\ A_2 \ &= \ \omega[e^{-i\lambda_S t} - \eta_X\eta_Y e^{-i\lambda_L t}] \quad (6) \end{aligned}$$

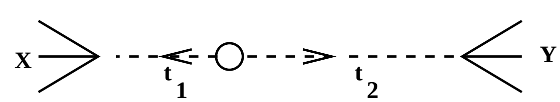

FIG. 1: A typical amplitude corresponding to the decay of, say, a ϕ state into final states X, Y; $t_i, i = 1, 2$ denote the corresponding time scales for the appearance of the final products of the decay.

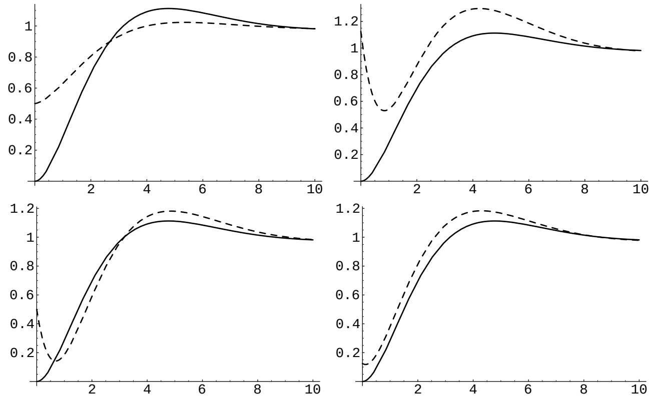

FIG. 2: Characteristic cases of the intensity $I(\Delta t)$, with $|\omega| = 0$ (solid line) vs $I(\Delta t)$ (dashed line) with (from top left to right): (i) $|\omega| = |\eta_{+-}|$, $\Omega = \phi_{+-} - 0.16\pi$, (ii) $|\omega| = |\eta_{+-}|$, $\Omega = \phi_{+-} + 0.95\pi$, (iii) $|\omega| = 0.5|\eta_{+-}|$, $\Omega = \phi_{+-} + 0.16\pi$, (iv) $|\omega| = 1.5|\eta_{+-}|$, $\Omega = \phi_{+-}$. Δt is measured in units of τ_S (the mean life-time of K_S) and $I(\Delta t)$ in units of $|C|^2|\eta_{+-}|^2|\langle\pi^+\pi^-|K_S\rangle|^4\tau_S$.

the CPT-allowed and CPT-violating parameters respectively, and $\eta_X = \langle X|K_S\rangle/\langle X|K_L\rangle$ and $\eta_Y = \langle Y|K_S\rangle/\langle Y|K_L\rangle$. Next, one integrates the square of the amplitude over all accessible times $t = t_1 + t_2$, keeping the difference $\Delta t = t_2 - t_1$ as constant. This defines the "intensity" $I(\Delta t)$:

$$I(\Delta t) \equiv \frac{1}{2}\int_{|\Delta t|}^{\infty} dt\, |A(X,Y)|^2 \tag{7}$$

In what follows we concentrate on identical final states $X = Y = \pi^+\pi^-$, because as we shall argue they are the most sensitive channels to probe the novel effects associated with the CPTV parameter ω. Indeed [2] the amplitudes of the CP violating decays $K_L \to \pi^+\pi^-$ are suppressed by factors of order $\mathcal{O}(10^{-3})$, as compared to the principal decay mode of $K_S \to \pi^+\pi^-$. In the absence of CPTV ω, (1), due to the $K_S K_L$ mixing, such decay rates would be suppressed. This would not be the case, however, when the CPTV ω (2) parameter is non zero, due to the existence of a separate $K_S K_S$ term in that case ((4)). This implies that the relevant parameter for CPT violation in the intensity is ω/η_X, which enhances the potentially observed effect.

Substituting in Eq.(7) $|A(\pi^+\pi^-,\pi^+\pi^-)|^2 = |\langle\pi^+\pi^-|K_S\rangle|^4(|A_1|^2 + |A_2|^2 + 2\Re e\{A_1 A_2^*\})$ and integrating over t we obtain

$$I(\Delta t) = |\langle\pi^+\pi^-|K_S\rangle|^4|C|^2|\eta_{+-}|^2\Big[I_1 + I_2 + I_{12}\Big] \tag{8}$$

with

$$I_1(\Delta t) = \frac{e^{-\Gamma_S\Delta t} + e^{-\Gamma_L\Delta t} - 2e^{-(\Gamma_S+\Gamma_L)\Delta t/2}\cos(\Delta M\Delta t)}{\Gamma_L + \Gamma_S}$$

$$I_2(\Delta t) = \frac{|\omega|^2}{|\eta_{+-}|^2}\frac{e^{-\Gamma_S\Delta t}}{2\Gamma_S}$$

$$\begin{aligned} I_{12}(\Delta t) = & -\frac{4}{4(\Delta M)^2 + (3\Gamma_S + \Gamma_L)^2}\frac{|\omega|}{|\eta_{+-}|}\times \\ & \Big[2\Delta M\Big(e^{-\Gamma_S\Delta t}\sin(\phi_{+-} - \Omega) - \\ & e^{-(\Gamma_S+\Gamma_L)\Delta t/2}\sin(\phi_{+-} - \Omega + \Delta M\Delta t)\Big) \\ & -(3\Gamma_S + \Gamma_L)\Big(e^{-\Gamma_S\Delta t}\cos(\phi_{+-} - \Omega) - \\ & e^{-(\Gamma_S+\Gamma_L)\Delta t/2}\cos(\phi_{+-} - \Omega + \Delta M\Delta t)\Big)\Big] \end{aligned}$$

where we have set $\Delta M = M_S - M_L$ and $\eta_{+-} = |\eta_{+-}|e^{i\phi_{+-}}$.

The effects of the CPTV ω on such intensities $I(\Delta t)$ are indicated in figure 2. The order of the CPTV effect is highly model dependent, and hard to evaluate. In line with other generic approaches to QG-decoherening evolution [6, 8, 12], which is also associated with an ill definition of the concept of a CPT operator, and thus of the antiparticle, in view of the lack of a well-defined scattering matrix [7], one might expect situations in which ω is of similar order as, say, the QG-decoherening (dimensionless) parameters [6, 8] $\widehat{\alpha}, \widehat{\beta}, \widehat{\gamma}$, where the $\widehat{\dots}$ implies division of the corresponding parameter (with dimensions of energy) with $\Delta\Gamma = \Gamma_S - \Gamma_L$. In optimistic scenaria of QG-induced decoherening situations, the relevant effects are of order E^2/M_{QG}, where E a typical average energy of the Kaon system (or rest-mass, in the Lorentz-invariant case), and M_{QG} the QG scale (which could be taken to be the Planck scale $M_P \sim 10^{19}$ GeV). For Kaons, such effects imply that the dimensionless (hatted) quantities are expected to be of order $10^{-3} - 10^{-4}$, thereby being well within the sensitivity of ϕ factories [13]. Indeed, with $|\omega| \sim 10^{-3} - 10^{-4}$ the new CPTV effects become comparable to those associated with $|\eta_{+-}| \sim 10^{-3}$; therefore, a precision of 10^{-3} in $I(\Delta t)$, which is needed in order to observe ϵ' effects, would probe sensitivities up to $|\omega| \sim 10^{-6}$. It is understood that a similar analysis can be done for the $X = Y = \pi^0\pi^0$ case.

We continue with a brief discussion concerning the distinguishability of the ω effect (2),(4) from non-quantum mechanical effects associated with the evolution, as in [6]. The ω effect can be distinguished from those of the QG-decoherening evolution parameters α, β, γ, when the formalism is applied to the entangled states ϕ [12, 14]. A non-quantum mechanical evolution of the entangled Kaon state with $\omega = 0$ has been considered in [12]. In such a case the resulting density-matrix ϕ state $\tilde{\rho}_\phi = \mathrm{Tr}|\phi><\phi|$ can be written as

$$\begin{aligned} \tilde{\rho}_\phi = & \rho_S\otimes\rho_L + \rho_L\otimes\rho_S - \rho_I\otimes\rho_{\overline{I}} - \rho_{\overline{I}}\otimes\rho_I \\ & - \frac{2\beta}{d}(\rho_I\otimes\rho_S + \rho_S\otimes\rho_I) - \frac{2\beta}{d^*}(\rho_{\overline{I}}\otimes\rho_S + \rho_S\otimes\rho_{\overline{I}}) \\ & + \frac{2\beta}{d}(\rho_{\overline{I}}\otimes\rho_L + \rho_L\otimes\rho_{\overline{I}}) + \frac{2\beta}{d^*}(\rho_I\otimes\rho_L + \rho_L\otimes\rho_I) \end{aligned}$$

$$- \frac{i\alpha}{\Delta M}(\rho_I \otimes \rho_I - \rho_{\overline{I}} \otimes \rho_{\overline{I}}) - \frac{2\gamma}{\Delta\Gamma}(\rho_S \otimes \rho_S - \rho_L \otimes \rho_L)$$

where the standard notation $\rho_S = |S >< S|$, $\rho_L = |L >< L|$, $\rho_I = |S >< L|$, $\rho_{\overline{I}} = |L >< S|$ has been employed, $d = -\Delta M + i\Delta\Gamma/2$, and an overall multiplicative factor of $\frac{1}{2}\frac{(1+2|\epsilon|^2)}{1-2|\epsilon|^2\cos(2\phi_\epsilon)}$ has been suppressed. On the other hand, the corresponding density matrix description of the ϕ state (4) in our case reads:

$$\begin{aligned}\rho_\phi &= \rho_S \otimes \rho_L + \rho_L \otimes \rho_S - \rho_I \otimes \rho_{\overline{I}} - \rho_{\overline{I}} \otimes \rho_I \\ &- \omega(\rho_I \otimes \rho_S - \rho_S \otimes \rho_I) - \omega^*(\rho_{\overline{I}} \otimes \rho_S - \rho_S \otimes \rho_{\overline{I}}) \\ &- \omega(\rho_{\overline{I}} \otimes \rho_L - \rho_L \otimes \rho_{\overline{I}}) - \omega^*(\rho_I \otimes \rho_L - \rho_L \otimes \rho_I) \\ &- |\omega|^2(\rho_I \otimes \rho_I + \rho_{\overline{I}} \otimes \rho_{\overline{I}}) + |\omega|^2(\rho_S \otimes \rho_S + \rho_L \otimes \rho_L)\end{aligned}$$

with the same multiplicative factor suppressed. It is understood that the evolution of both $\tilde{\rho}_\phi$ and ρ_ϕ is governed by the rules given in [6, 8, 12]. As we can see by comparing the two equations, the terms linear in ω in our case are *antisymmetric* under the exchange of particle states 1 and 2, *in contrast* to the *symmetry* of the corresponding terms linear in β in the case of [12]. Similar differences characterize the terms proportional to $|\omega|^2$, and those proportional to α and γ, which involve $\rho_I \otimes \rho_I$, $\rho_{\overline{I}} \otimes \rho_{\overline{I}}$, $\rho_S \otimes \rho_S$, $\rho_L \otimes \rho_L$. Such differences are therefore important in disentangling the ω CPTV effects proposed here from non-quantum mechanical evolution effects [6, 8, 12, 14].

We next comment on the distinguishability of the ω effect from conventional background effects. Specifically, the mixing of the initial state due to the non-identity of the antiparticle to the corresponding particle state has similar form to that induced by a non-resonant background with $C = +$ [9]. This latter effect is known to have a small size; estimates based on unitarity bounds give a size of many orders of magnitude smaller than the $C = -$ effect in the ϕ decays [9, 13]. Terms of the type $K_S K_S$ (which dominate over $K_L K_L$) coming from the ϕ-resonance as a result of CPTV can be distinguished from those coming from the $C = +$ background because they interfere differently with the regular $C = -$ resonant contribution (i.e. Eq.(4) with $\omega = 0$). Indeed, in the CPTV case, the $K_L K_S$ and $\omega K_S K_S$ terms have the same dependence on the center-of-mass energy s of the colliding particles producing the resonance, because both terms originate from the ϕ-particle. Their interference, therefore, being proportional to the real part of the product of the corresponding amplitudes, still displays a peak at the resonance. On the other hand, the amplitude of the $K_S K_S$ coming from the $C = +$ background has no appreciable dependence on s and has practically vanishing imaginary part. Therefore, given that the real part of a Breit-Wigner amplitude vanishes at the top of the resonance, this implies that the interference of the $C = +$ background with the regular $C = -$ resonant contribution vanishes at the top of the resonance, with opposite signs on both sides of the latter. This clearly distinguishes experimentally the two cases.

Finally we close with a comment on the application of this formalism to the B factories. Although, formally, the situation is identical to the one discussed above, however the sensitivity of the CPTV ω effect for the B system is much smaller. This is due to the fact that in B factories there is no particularly "good" channel X (with $X = Y$) for which the corresponding η_X is small. The analysis in that case may therefore be performed in the equal sign dilepton channel, where the branching fraction is more important, and a high statistics is expected. Results will appear in future work.

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