Quantum decoherence of photons in the presence of hidden U(1)s

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Many extensions of the standard model predict the existence of hidden sectors that may contain unbroken Abelian gauge groups. We argue that in the presence of quantum decoherence photons may convert into hidden photons on sufficiently long time scales and show that this effect is strongly constrained by CMB and supernova data. In particular, Planck-scale suppressed decoherence scales $D \propto \omega^2/M_{\rm Pl}$ (characteristic for noncritical string theories) are incompatible with the presence of even a single hidden U(1). The absence of photon decoherence in this simple standard model extension complements other strong bounds derived from solar, reactor, and atmospheric neutrinos.

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I. INTRODUCTION

Physics beyond the standard model (SM) typically predicts the existence of hidden sectors containing extra matter with new gauge interactions. There is no compelling reason why these extra sectors should be very massive if their interaction with standard model matter is sufficiently weak. This can be accomplished, e.g., in extra dimensional extensions with a geometric separation of sectors or in models where tree-level interactions are absent and higher order corrections are suppressed by the scale of messenger masses.

The presence of these light hidden sectors can have various observable effects. For example, in the case of kinetic and mass mixing [1,2] between a hidden U(1) and the electromagnetic U(1), hidden sector matter can receive small fractional electromagnetic charges. These "minicharged" particles can have a strong influence on early Universe physics and astrophysical environments [3]. If the hidden sector U(1) is only slightly broken by a Higgs or Stückelberg mechanism we can have a situation analogous to neutrino systems with characteristic oscillation patterns between photons and hidden photons over sufficiently long baselines [4].

Here we concentrate on another effect induced by the presence of these light hidden sectors: the sensitivity of photon propagation to sources of quantum decoherence, e.g. quantum gravity effects [5]. A heuristic picture describes space-time at the Planck scale as a foamy structure [6], where virtual black holes pop in and out of existence on a time scale allowed by Heisenberg's uncertainty principle [7]. This can lead to a loss of quantum information across their event horizons, providing an "environment" that might induce quantum decoherence of apparently isolated matter systems [5,8].

It is an open matter of debate whether quantum decoherence induced by a quantum theory of gravity would simultaneously preserve Poincaré invariance and locality [5,7,9–11]. A violation of energy and momentum conservation by particle reactions with a space-time foam could be reflected by an (energy dependent) effective refractive index in vacuum [12]. This could be tested, e.g., by the measurement of the arrival time of gamma rays or high-energy neutrinos at different energies or by the propagation of ultrahigh-energy cosmic rays [13].

If, on the other hand, Poincaré invariance is preserved, the presence of a nontrivial space-time vacuum can still be signaled by decoherence effects in systems of stable elementary particles [9]. A particularly interesting and well-studied case are neutrino systems, where the interplay between mixing, mass oscillation, and decoherence can influence atmospheric, solar, and reactor neutrino data [14–17], as well as flavor composition of astrophysical high-energy neutrino fluxes [18,19].

If hidden sectors contain unbroken Abelian gauge groups it is also feasible that the system of the electromagnetic photon (γ_0) and hidden photons ($\{\gamma_i\}$ with $i \geq 1$) experience energy and momentum conserving decoherence effects: transitions between different photon "species," $\gamma_\mu \to \gamma_\nu$, are allowed by gauge and Poincaré invariance. This is the only alternative system to neutrinos involving a stable and neutral elementary particle of the standard model, which can experience these decoherence effects on extremely long (cosmological) time scales.

The outline of this paper is as follows. We will start in Sec. II with an outline of the Lindblad formalism of quantum decoherence. We will then discuss in Sec. III the effect of photon decoherence on the Planck spectrum of the cosmic microwave background (CMB) and the luminosity distance of type Ia supernovae, respectively. This enables us to derive strong limits on various decoherence models in this scenario. We comment in Sec. IV on the interplay of decoherence and photon interactions and outline a possible mechanism to extend the survival probability of extra-galactic TeV gamma rays. We finally summarize in Sec. V.

II. QUANTUM DECOHERENCE

The Lindblad formalism is a general approach to quantum decoherence that does not require any detailed knowledge of the environment [20]. In the presence of decoherence the modified Liouville equation can be written in the form

$$\frac{\partial \rho}{\partial t} = -i[H, \rho] + \mathcal{D}[\rho]. \tag{1}$$

The Hamiltonian H can include possible background contributions, e.g. plasma effects for the photon. However, this does not affect the evolution of the density matrix in the absence of mixing between the gauge bosons. The decoherence term $\mathcal{D}$ in the modified Liouville equation (1) can be written as

$$\mathcal{D}[\rho] = \frac{1}{2}\sum_j ([b_j, \rho b_j^\dagger] + [b_j \rho, b_j^\dagger]), \tag{2}$$

where $\{b_j\}$ is a sequence of bounded operators acting on the Hilbert space of the open quantum system, $\mathcal{H}$, and satisfying $\sum_j b_j^\dagger b_j \in \mathcal{B}(\mathcal{H})$, where $\mathcal{B}(\mathcal{H})$ indicates the space of bounded operators acting on $\mathcal{H}$. The dynamical effects of space-time on a microscopic system can then be interpreted as the existence of an arrow of time which in turn makes possible the connection with thermodynamics via an entropy. The monotonic increase of the von Neumann entropy, $S(\rho) = -\mathrm{Tr}(\rho \ln \rho)$, implies the Hermiticity of the Lindblad operators, $b_j = b_j^\dagger$ [21]. In addition, the conservation of energy and momentum can be enforced by taking $[P_\mu, b_j] = 0$.

We will assume in the following that there is a total of N Abelian gauge bosons γ_μ with $\mu = 0, \ldots, N-1$ including the photon γ_0. The solution to Eq. (1) is outlined in the Appendix. For simplicity, we assume degeneracy of the decoherence parameters $D_i = D$ which simplifies the photon survival probability after a distance $x = t = L$ significantly (see the Appendix),

$$P_{\gamma\to\gamma} = \frac{1}{N} + \frac{N-1}{N} e^{-DL}. \tag{3}$$

The energy behavior of D depends on the dimensionality of the operators b_j. We can estimate the energy dependence from gauge invariance. Possible combinations of the field strength tensors $F_{\mu\nu}$ and $G_{\mu\nu}$ of the two U(1)s are [5] $b_j \propto (F_{\mu\nu}G^{\mu\nu})^j \propto \omega^j$. This restriction of the energy behavior to non-negative powers of ω may possibly be relaxed when the dissipative term is directly calculated in the most general space-time foam background [19].

An interesting example is the case where the dissipative term is dominated by the dimension-4 operator b_1 yielding the energy dependence $D \propto \omega^2/M_{\rm Pl}$. This is characteristic of noncritical string theories where the space-time defects of the quantum gravitational environment are taken as recoiling D-branes, which generate a cellular structure in the space-time manifold [22].

III. OBSERVATIONAL CONSTRAINTS

In the following we will investigate the limits on quantum decoherence in the presence of hidden *massless* U(1)s from cosmological and astrophysical observations. Unless otherwise stated, we will make the conservative assumption that there exists only a single hidden U(1) in addition to the standard model ($N = 2$). We will parametrize the decoherence rate of photons as

$$D(z, \omega) = (1+z)^{p+1} D_* \frac{\omega^{p+1}}{M_{\rm QD}^p} \equiv (1+z)^{p+1}\gamma_0\left(\frac{\omega}{\rm GeV}\right)^{p+1}, \tag{4}$$

where $M_{\rm QD}$ is the scale of quantum decoherence, not necessarily the Planck scale, and following the notation of Ref. [16] we have introduced the effective parameter

$$\gamma_0 = D_*\left(\frac{\rm GeV}{M_{\rm QD}}\right)^p \ \mathrm{GeV}. \tag{5}$$

This parametrization approximates the energy dependence of decoherence as a power law and we will consider a wide range of indices $p = 0, \pm1, \pm2$ in the following.

The differential flux of photons from a source at redshift z is reduced by the exponential factor

$$P_{\gamma\to\gamma}(z, \omega) = \frac{1}{2} + \frac{1}{2}\exp\left[-\int_0^z d\ell' D(z', \omega)\right], \tag{6}$$

where the propagation distance ℓ is given by $dz = H(z) \times (1+z) d\ell$ with Hubble parameter H. The Hubble parameter at redshift z is given by $H^2(z) = H_0^2[(1 - \Omega_m - \Omega_r) + \Omega_m(1+z)^3 + \Omega_r(1+z)^4]$ where $\Omega_m h^2 \simeq 0.128$ and $\Omega_r h^2 \simeq 2.47 \times 10^{-5}$. The present Hubble expansion is $H_0 = h 100\ \mathrm{km\,s^{-1}\,Mpc^{-1}}$ with $h \simeq 0.73$ [23].

A. CMB distortions

In the standard big bang cosmology the CMB forms at a redshift of about $z_{\rm CMB} \simeq 1100$ after recombination of electrons and (mostly) protons in the expanding and cooling universe. The CMB is well described by a Planck spectrum with a temperature of $T = 2.725 \pm 0.001$ K [24] and spectral intensity ($\mathrm{W\ m^{-2}\,Hz^{-1}\,sr^{-1}}$)

$$\frac{d^2 I}{d\nu d\Omega} = \frac{1}{2\pi^2}\frac{\omega^3}{\exp(\frac{\omega}{T}) - 1}. \tag{7}$$

The high degree of accuracy (better than 1 in 10^4 around 1 meV) between the CMB measurement and cosmological predictions is an ideal probe for exotic physics that could have affected the CMB photons in the redshift range $0 <$

$z < z_{\rm CMB}$ with energies from meV ($z = 0$) to eV ($z = z_{\rm CMB}$).

The influence of light particles coupling to the CMB has been studied previously for the case of axionlike particles [25], minicharged particles [26], and massive hidden photons with kinetic mixing [27]. In the case of photon absorption or decoherence the observed spectrum is modified as

$$\frac{d^2 I^{\rm obs}}{d\nu d\Omega} = P_{\gamma\to\gamma}(z_{\rm CMB}, \omega)\frac{d^2 I}{d\nu d\Omega}. \tag{8}$$

As an illustration of the effect of decoherence, Fig. 1 shows the distortion of the CMB spectrum for the case $p = 1$ and $M_{\rm QD} = M_{\rm Pl} \simeq 1.2 \times 10^{19}$ GeV for $D_* = 0.01,\ 0.1,\ 1$ (upper panel) and $D_* = 10^{-4}$ (lower panel). The 3σ limits from a χ^2 fit of various models to the COBE/FIRAS data [24] are shown in Table I.

B. SN dimming

Also the luminosity distance measurements of cosmological standard candles like type Ia supernovae (SNe) [28,29] are able to test feeble photon absorption and decoherence effects. The luminosity distance d_L is defined as

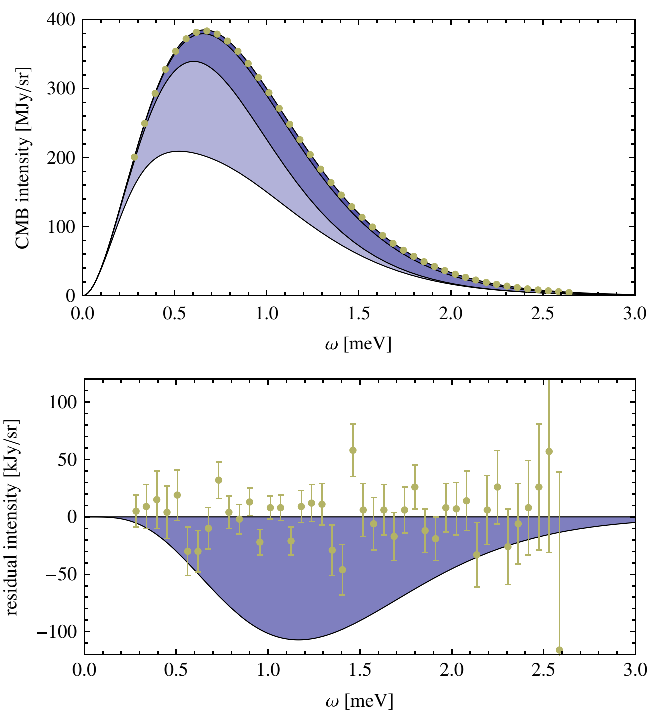

FIG. 1 (color online). Modification of the CMB spectrum by decoherence effects. The dots show the residual CMB spectrum measured by COBE/FIRAS [24] (1 Jy(Jansky) $\equiv 10^{-26}$ W m^{-2} Hz^{-1}). The upper panel shows the effect of $D_* = 0.01, 0.1, 1$ and the lower panel $D_* = 10^{-4}$. The curves show the accumulated deficit in photons compared to the "no decoherence" case. (A different shadowing is included to guide the eye.)

TABLE I. The 3σ limits on the parameter combination $\gamma_0 = D_*({\rm GeV}/M_{\rm QD})^p$ (GeV) derived from the COBE/FIRAS data [24] and SN data [28] assuming one extra U(1) ($N = 2$). Note that the limits become stronger by up to a factor 2 for $N > 2$. For illustration, the last two columns show the limits obtained from reactor and solar [16] as well as atmospheric neutrinos [14,17].

Model	CMB	SNe	Reactor & solar ν	Atmospheric ν
$p = -2$	1.4×10^{-58}	5.6×10^{-52}	$7.8 \times 10^{-27\,\rm a}$	$1.5 \times 10^{-21\,\rm b}$
$p = -1$	1.2×10^{-46}	4.1×10^{-43}	$6.7 \times 10^{-25\,\rm a}$	...
$p = 0$	3.9×10^{-35}	2.9×10^{-34}	$5.8 \times 10^{-23\,\rm a}$	$1.2 \times 10^{-27\,\rm c}$
$p = 1$	1.1×10^{-24}	1.9×10^{-25}	$4.7 \times 10^{-21\,\rm a}$	$1.3 \times 10^{-31\,\rm c}$
$p = 2$	2.4×10^{-15}	1.3×10^{-16}	...	$5.3 \times 10^{-36\,\rm c}$

[a]Bounds are at the 95% C.L. Their notation corresponds to $n = p + 1$ [16].
[b]The authors consider the case $p = -2$ with $\mu^2/({\rm GeV})^2 = 2\gamma_0/{\rm GeV}$ [14].
[c]Bounds are at the 90% C.L. Their notation corresponds to $n = p + 1$ and $D^*({\rm GeV})^p = \gamma_0/{\rm GeV}$ [17].

$$d_L(z) \equiv \sqrt{\frac{\mathcal{L}}{4\pi F}}, \tag{9}$$

where $\mathcal{L}$ is the luminosity of the standard candle (assumed to be sufficiently well known) and F the measured flux. In a homogeneous and isotropic universe this is predicted to be

$$d_L(z) = (1 + z)a_0\Phi\left(\int_0^z \frac{dz'}{a_0 H(z')}\right), \tag{10}$$

with $a_0^{-1} = H_0\sqrt{|1 - \Omega_{\rm tot}|}$ and $\Phi_k(\xi) = (\sinh\xi, \xi, \sin\xi)$ for spatial curvature $k = -1, 0, 1$, respectively. If the photon flux of a source, located at distance z and observed in a (small) frequency band centered at $\omega_\star$, is attenuated by photon interactions or quantum decoherence the observed luminosity distance *increases* as

$$d_L^{\rm obs}(z) = \frac{d_L(z)}{\sqrt{P_{\gamma\to\gamma}(z, \omega_\star)}}. \tag{11}$$

The apparent extension of the luminosity distance by photon interactions and oscillations has been investigated in the context of axionlike particles [25,30], hidden photons [31], chameleons [32], and minicharged particles [33]. One of the main attractions of these models is the possibility that the conclusions about the energy content of our Universe drawn from the Hubble diagram can be dramatically altered. We will briefly comment on this possibility at the end of this section.

As an example, the upper panel of Fig. 2 shows the effect on quantum decoherence in the ΛCDM model assuming a single hidden U(1) ($N = 2$). The luminosity distance of the SNe is shown as the difference between their measured apparent magnitude m and their known absolute magnitude M,

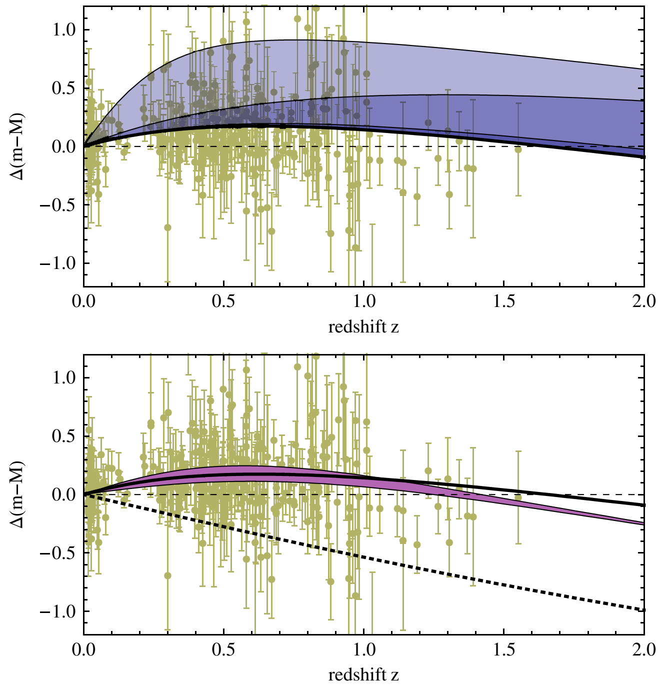

FIG. 2 (color online). Upper panel: Enhanced SNe dimming by decoherence effects assuming $N = 2$, $p = 1$, $M_{QD} = M_{Pl}$ and $D_* = 10^{-6}$, 10^{-5}, 10^{-4} assuming sources observed in a frequency interval centered at $\omega_\star = 1$ eV. The dots show the SNe Ia "union" compilation from Ref. [28]. The luminosity distance d_L is shown by the difference $\Delta(m - M)$ to an empty ($\Omega_{tot} = 0$) flat universe. The thin lines show the accumulated deficit in photons compared to the "no decoherence" case (standard ΛCDM model) indicated by the thick line. Lower panel: As the upper panel, but now showing also a flat CDM model with $\Omega_m = 1$ and $\Omega_\Lambda = 0$ (dotted line). For illustration, we consider a hidden photon model with $p = 1$, $N = 2$, and $D_* = 6 \times 10^{-6}$ and show the dimming for the B ($\lambda_\star \simeq 440$ nm, upper line) and V ($\lambda_\star \simeq 550$ nm, lower line) bands. Whereas the overall dimming effect is practically indistinguishable from the ΛCDM model, the strong *reddening* of the starlight from the energy dependence of $D \propto \omega^2/M_{Pl}$ is challenged by the data [34].

$$m - M = 5\log_{10} d_{L,\mathrm{Mpc}} + 25. \tag{12}$$

As in the previous case we can derive 3σ limits for various decoherence models that are shown in Table I.

However, we would like to emphasize that these limits depend on the cosmological model and the normalization of the SN data. If the SN dimming by photon decoherence is strong this can have an effect on the evaluation of cosmological data. We give an example in the lower panel of Fig. 2 where the decoherence effect could even be able to reproduce the observed SN luminosities from a CDM model. Note that this model is not excluded by the corresponding CMB limits, though it is not compatible with the analogous bound from atmospheric neutrinos. In addition, one has to keep in mind that the photon frequency dependence of the parameter D causes a color excess (unless $p = -1$) which is in conflict with observation [28]. For the example shown in the right panel of Fig. 2 the color excess between the B and V band of the form, $E[\mathrm{B} - \mathrm{V}] \equiv \Delta(m - M)_\mathrm{B} - \Delta(m - M)_\mathrm{V}$, is larger than 0.8 for redshifts $0.2 \lesssim z \lesssim 0.8$, more than twice the median color excess observed from SNe at this distance [34]. Moreover, photon absorption as a SN dimming mechanism would violate the cosmic distance duality, i.e. the luminosity and angular diameter distance relation $d_L/d_A = (1 + z)^2$ [35]. Hence, it is unlikely that decoherence can fully account for SN dimming. Nevertheless, it could have an effect on the evaluation of cosmological data.

The 3σ limits on the parameter γ_0 [cf. Eq. (5)] shown in Table I translate into bounds on D_* as a function of the decoherence scale M_{QD} [cf. Eq. (4)]. The dimensionless parameter D_* is naturally expected to be of order 1. Hence, limits on photon decoherence in the presence of hidden U(1)s require $D_* \ll 1$ for $p \leq 1$. In reverse, a quantum theory of gravity predicting Planck-scale suppressed decoherence is not compatible with the existence of unbroken hidden U(1)s unless, *unnaturally*, $D_* \ll 1$.

The two rightmost columns in Table I show decoherence limits derived from solar and reactor neutrino data [16] as well as atmospheric neutrino results [14,17]. One must bear in mind that the bounds from neutrino data and the ones derived in this work are not directly comparable without assumptions. First, our bounds rely on the presence of massless hidden U(1)s while the neutrino bounds hold in the absence of those. Second, it is still feasible that decoherence effects are not universally present in both, neutrino and photon/hidden photon systems (though if quantum decoherence has a common origin, e.g. quantum gravity, it seems reasonable to assume a universal decoherence scale M_{QD}). But most importantly the simple power law parametrization (4) does not necessarily approximate the decoherence effects at the very different experimental energy scales probed by the different data from meV (CMB) up to TeV (atmospheric ν).

Generically, high (low) energy data—either photons or neutrinos—are more sensitive to decoherence effects that grow (decrease) with energy. This tendency can be observed in Table I. Consequently, higher power quantum decoherence with $p \geq 2$ are only weakly constrained by photon/hidden photon systems in the meV to eV energy range. We will speculate in the following section, how high-energy gamma ray sources could possibly explore this unconstrained parameter space.

IV. PHOTON PROPAGATION

Decoherence effects can also have interesting effects on the propagation of photons in the presence of photon absorption. We can account for photon absorption effects in the Liouville equation (1) by a contribution

$$\mathcal{D}_\Gamma[\rho] = -\frac{\Gamma}{2}\{\Pi_0, \rho\}, \tag{13}$$

where $\Gamma = 1/\lambda$ is the photon absorption rate,[1] e.g. in the intergalactic photon background (BG) via $\gamma + \gamma_{BG} \rightarrow e^+ + e^-$, and Π_0 is the photon projection operator.

The photon survival probability in the presence of a single hidden U(1) can be readily solved from the modified Liouville equation and gives (see the Appendix for details)

$$P_{\gamma\rightarrow\gamma} = e^{-(L/2)(D+\Gamma)}\bigg[\cosh\bigg(\frac{L}{2}\sqrt{D^2+\Gamma^2}\bigg) - \frac{\Gamma}{\sqrt{D^2+\Gamma^2}}\sinh\bigg(\frac{L}{2}\sqrt{D^2+\Gamma^2}\bigg)\bigg]. \tag{14}$$

We can estimate the sensitivity of extra-galactic TeV gamma ray sources to decoherence effects as

$$D_* \sim 10^{-38+16p}\bigg(\frac{M_{QD}}{M_{Pl}}\bigg)^p\bigg(\frac{\text{TeV}}{\omega}\bigg)^{p+1}\bigg(\frac{\text{kpc}}{L}\bigg). \tag{15}$$

In particular for $M_{QD} \sim M_{Pl}$, $D_* \sim 10^{-22}$ for $p = 1$ or $D_* \sim 10^{-6}$ for $p = 2$. Hence, this probe has the potential to be more sensitive to Planck-scale suppressed decoherence than existing neutrino data [14,17] by many orders of magnitude. We will discuss in the following possible signals of decoherence in the spectra of TeV gamma ray sources.

Figure 3 shows the survival probability for three different values of D. The functional behavior can be easily understood as follows: The hidden U(1) serves as an invisible "storage" of photons during propagation. If $D\lambda \gg 1$ decoherence quickly equalizes the number of photons and hidden photons. This results into a *decrease* of the photon survival probability to $1/2$ (for $N = 2$) for $L \lesssim \lambda$. On the other hand, inelastic scattering only affects photons. Therefore at larger L photons can be "replenished" by the decoherence of the unabsorbed hidden photons into photons. Hence, for propagation distances $L > \lambda$ the presence of hidden photons increases the photon survival probability.

For the general case of N U(1)s and strong decoherence $D \gg \Gamma$ we can express the photon survival probability at a distance much larger than the decoherence scale $L \gg 1/D$ as (see the Appendix for details)

$$P_{\gamma\rightarrow\gamma} \simeq \frac{1}{N}\exp\bigg(-\frac{L}{N\lambda}\bigg). \tag{16}$$

This can have an important effect on the spectra of TeV gamma ray sources.[2] First, if the onset of decoherence appears at energies covered by the spectra one could observe a steplike drop of the flux by a factor $1/N$. And second, as long as the absolute source emissivity of photons is unknown [such that the prefactor $1/N$ in Eq. (16) gets renormalized], the expected spectral cutoffs of the sources could be shifted according to an extended photon interaction length $\lambda_{eff} = N\lambda$. These effects can be clearly seen in Fig. 3 for the case $N = 2$. For a TeV gamma ray source at 100 Mpc and Planck-scale suppressed decoherence this requires $D_* \gg 10^{-27}$ ($D_* \gg 10^{-11}$) for $p = 1$ ($p = 2$). In comparing this with Table I and Eq. (5) we see that there is ample room for models that could have such an effect on the spectra.

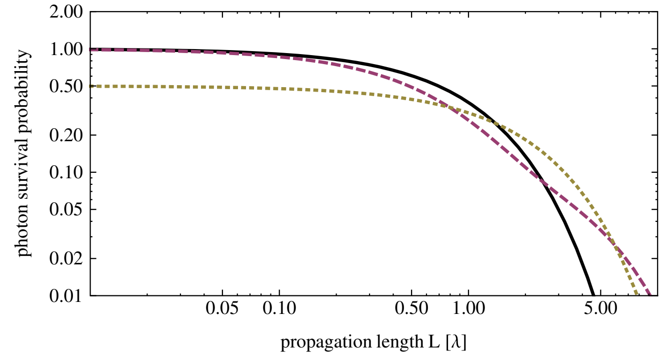

FIG. 3 (color online). Photon survival probability in the presence of decoherence with a single hidden U(1) ($N = 2$). The black line shows the photon attenuation, e.g. in the intergalactic photon background via $\gamma + \gamma_{BG} \rightarrow e^+ + e^-$, in the absence of decoherence ($D = 0$). The dashed curve shows decoherence with $D\lambda = 1$ and the dotted curve $D\lambda \gg 1$.

V. CONCLUSIONS

We have discussed the effects of quantum decoherence of photons in the presence of hidden sector U(1)s. Quantum decoherence in the system of photon and hidden photons could be induced by a foamlike structure of spacetime in a quantum theory of gravity, where virtual black holes pop in and out of existence at scales allowed by Heisenberg's uncertainty principle.

We have shown that these decoherence effects are strongly constrained by the absence of photon disappearance in the cosmic microwave background. Furthermore

[1]Since we consider point-source fluxes in the following we will treat the reaction $\gamma + \gamma_{BG} \rightarrow e^+ + e^-$ as an absorption process of the photon. Subsequent electromagnetic interactions of secondary $e^\pm$ will contribute to the *diffuse* GeV-TeV background.

[2]Note that in the absence of hidden sectors, as long as $\omega \lesssim v$ (v being the scale of electroweak symmetry breaking) the photon survival probability is 1 independently on whether the Higgs potential around virtual black holes has its minimum at the trivial vacuum $v = 246$ GeV or at the unbroken vacuum $v = 0$. For photon energies beyond the electroweak breaking scale, one may theorize over a possible decoherence effect between the neutral SM gauge bosons. In this case the conservation of energy and momentum in decoherence effects, which forbids transitions of the form $\gamma \rightarrow Z$, is a crucial assumption.

quantum decoherence as an additional source of starlight dimming can also be constrained by the luminosity distance measurements of type Ia supernovae. Consequently based on the standard ΛCDM model we can derive constraints on the decoherence in the presence of hidden U(1)s. In principle, the decoherence effect could be strong enough to influence the evaluation of cosmological data. However, color dependencies in the dimming via photon decoherence are not favored by the data and could be used to derive further constraints.

Our main results are summarized in Table I. We observe that Planck-scale suppressed decoherence scales $D \propto \omega^2/M_{\rm Pl}$ are incompatible with the presence of even a single hidden U(1). These results complement other strong decoherence limits derived from solar, reactor, and atmospheric neutrinos.

We have also discussed the interplay between photon absorption and decoherence. This effect can become important for distant TeV gamma ray point sources if the decoherence length is much smaller than the photon interaction length. Assuming $N-1$ additional U(1)s this could leave characteristic features in gamma ray spectra in the form of steplike drops by factors $1/N$ or by an effective increase of the absorption length λ to $\lambda_{\rm eff} = N\lambda$.

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APPENDIX: LINDBLAD FORMALISM

We outline the solution to the Liouville equation (1) in the presence of quantum decoherence. The density matrix ρ and Lindblad operators b_j can be expanded in a basis of Hermitian matrices F_μ that satisfy the orthonormality condition $\mathrm{Tr}(F^\dagger_\mu F_\nu) = \delta_{\mu\nu}/2$. Without loss of generality we consider a basis with $(F_0)_{ij} = \delta_{ij}/\sqrt{2N}$. Explicitly, we have

$$\rho = \sum_\mu \rho_\mu F_\mu, \qquad b_j = \sum_\mu b^{(j)}_\mu F_\mu. \tag{A1}$$

The free propagation of photons and hidden photons ($H = -i\partial/\partial x$) can be readily solved in terms of "light-cone" coordinates $\hat{x} = (x-t)/2$ and $\hat{t} = (x+t)/2$. In these new coordinates Eq. (1) can be written $\partial\rho/\partial\hat{t} = \mathcal{D}[\rho]$.

Hence, the coefficients of the free equations of motion satisfy the differential equation

$$\frac{\partial}{\partial t}\rho_\mu = -\sum_\nu D_{\mu\nu}\rho_\nu, \tag{A2}$$

with $D_{\mu 0} = D_{0\mu} = 0$ and

$$D_{ij} = \frac{1}{2}\sum_{k,l,m,n} b^{(n)}_m f_{iml} b^{(n)}_k f_{jkl}, \tag{A3}$$

where f_{ijk} are structure constants defined by $[F_i, F_j] = i\sum_k f_{ijk}F_k$.

The solution of $\partial_t\rho_0 = 0$ is trivial and requires $\rho_0(t) = \rho_0 = 2\,\mathrm{Tr}(\Pi_\alpha F_0) = \sqrt{2/N}$ for all species α. If D_{ij} is diagonalizable by a matrix M, $(M^{-1}DM)_{ij} = D_i\delta_{ij}$, we can write the final solution as

$$P_{\gamma\to\gamma} = \frac{1}{N} + \frac{1}{2}\sum_{i,k,j} e^{-D_k t}\rho_i(0)M_{ik}M^{-1}_{kj}\rho_j(0). \tag{A4}$$

This reduces to Eq. (3) in the case $D_i = D$ using $2 = \sum_\mu \rho_\mu^2$ following from $\rho^2(0) = \rho(0)$ and $\mathrm{Tr}[\rho] = 1$.

We can extend the Liouville equation (1) by a photon decay term of the form (13). The modified equation can be readily solved in the case of $N = 2$, taking $F_i = \frac{1}{2}\sigma_i$ with Pauli matrices σ_i. In this basis the photon projection operator in the term (13) has the form $\Pi_0 = F_0 + F_3$. And the general solution of the photon survival probability is given in Eq. (14).

In the case of strong decoherence, i.e. at propagation distances $L \gg 1/D$ and $D \gg \Gamma$, we can derive the asymptotic solution (16) of the general survival probability with N U(1)s in the following way. In the presence of strong decoherence we can assume that at any step during the evolution $\mathrm{Tr}[\Pi_\mu\rho] \simeq \mathrm{Tr}[\Pi_\nu\rho]$ and, in particular, $\mathrm{Tr}[\rho] \simeq N\,\mathrm{Tr}[\Pi_0\rho]$. Since the trace of the decoherence term vanishes we can derive the asymptotic differential equation

$$\mathrm{Tr}[\dot\rho] \simeq N\,\mathrm{Tr}[\Pi_0\dot\rho] \simeq -\Gamma\,\mathrm{Tr}[\Pi_0\rho]. \tag{A5}$$

This has the solution (16), noting that $P_{\gamma\to\gamma} = \mathrm{Tr}[\Pi_0\rho]$ and initially $\mathrm{Tr}[\Pi_0\rho(0)] = 1/N$.

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