Dimming of supernovae by photon-pseudoscalar conversion and the intergalactic plasma

Cédric Deffayet,[1,*] Diego Harari,[2,†] Jean-Philippe Uzan,[3,4,‡] and Matias Zaldarriaga[1,5,§]
[1] *NYU Department of Physics, 4 Washington Place, New York, New York 10003*
[2] *Departamento de Física, FCEyN–Universidad de Buenos Aires, Ciudad Universitaria–Pab. 1, 1428 Buenos Aires, Argentina*
[3] *Laboratoire de Physique Théorique, UMR 8627 du CNRS, Université Paris XI, Bâtiment 210, F-91405 Orsay Cedex, France*
[4] *Institut d'Astrophysique de Paris, 98bis Bd. Arago, 75014 Paris, France*
[5] *Institute for Advanced Study, Einstein Drive, Princeton, New Jersey 08540*
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It has been suggested recently that the observed dimming of distant type Ia supernovae may be a consequence of mixing of the photons with very light axions. We point out that the effect of the plasma, in which the photons are propagating, must be taken into account. This effect changes the oscillation probability and renders the dimming frequency dependent, contrary to observations. One may hope to accommodate the data by averaging the oscillations over many different coherence domains. We estimate the effect of coherence loss, either due to the inhomogeneities of the magnetic field or of the intergalactic plasma. These estimates indicate that the achromaticity problem can be resolved only with very specific, and probably unrealistic, properties of the intergalactic medium.

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I. INTRODUCTION

The mixing between photons and axions, in an external magnetic field, is a well studied mechanism [1–4] as is its analogue with gravitons [5,6]. It is experimentally used, since the pioneering work by Sikivie [2], to constrain the axion parameters [7–11] (see also e.g. Ref. [12] for an up to date review on such experiments). The astrophysical and cosmological implications of this mechanism have also been studied [12], and it was recently advocated to be a possible explanation for the observed dimming of distant type Ia supernovae [13,14] by Csáki *et al.* [15] (see also [16] for the analysis of supernovae data with the model of [15]). The underlying idea is simply that the luminosity of distant supernovae can be diminished due to the decay of photons into very light pseudoscalar particles induced by intergalactic magnetic fields over cosmological distances, and thus that the luminosity distance-redshift relationship can mimic the one of a universe with a nonzero cosmological constant without need for a cosmological constant. The pseudoscalar particles must have an electromagnetic coupling similar to axions, and a specific and very small mass ($m\sim10^{-16}$ eV) to avoid affecting the cosmic microwave background anisotropy beyond its observed value. Another aspect of photon-pseudoscalar conversion in intergalactic magnetic fields is the change of the polarization properties of distant sources [17] such as supernovae. Conversion in the magnetic field of our own galaxy can be used to place constraints on the photon-pseudoscalar coupling from the lack of observed linear polarization of the cosmic microwave background [17,18], the possibility to see stars through otherwise opaque clouds [18], and from the flux of gamma rays that would have arisen from pseudoscalars emitted by SN 1987A [19,20].

The implications on the cosmic microwave background of the similar effect involving photon-graviton oscillation has been considered in empty space [21,22] and it was shown that it becomes negligible for standard cosmological magnetic fields [23] once the contribution of the intergalactic plasma is properly taken into account (the case of axions is also considered in [24]). The effect of the inhomogeneities of the electron density upon the coherence of the oscillations was also considered by Carlson *et al.* [18].

The purpose of this work is to discuss the effect of this plasma on the proposal of [15].

The paper is organized as follows. We first recall standard results on photon-pseudoscalar oscillations in order to introduce our conventions (see [1], or [24] where contributions from Kaluza-Klein modes were also included). We then discuss specifically the effect of the plasma for the parameters relevant for type Ia supernovae.

II. PHOTON-PSEUDOSCALAR MIXING

We consider the generic action for the pseudoscalar-photon system

$$S_4=\int d^4x\left[-\frac{1}{2}\{\partial^\mu a\,\partial_\mu a+m_a^2a^2\}+\frac{a}{M_a}F_{\mu\nu}\tilde{F}^{\mu\nu}-\frac{1}{4}F_{\mu\nu}F^{\mu\nu}\right], \quad (1)$$

where a is the pseudoscalar field and m_a its mass. $\tilde{F}_{\mu\nu}\equiv\epsilon_{\mu\nu\rho\sigma}F^{\rho\sigma}/2$ is the dual of the electromagnetic tensor, $\epsilon_{\mu\nu\rho\sigma}$ being the completely antisymmetric tensor such that $\epsilon_{0123}=+1$. The pseudoscalar couples to the photon with the coupling $1/M_a$.

We now consider an electromagnetic plane wave in the presence of a magnetic field $\vec{B}_0$ which is assumed constant

*Email address: cjd2@physics.nyu.edu
†Email address: harari@df.uba.ar
‡Email address: uzan@th.u-psud.fr
§Email address: mz31@nyu.edu

on a characteristic scale Λ_c in the sense that its variations in space and time are negligible on scales comparable to the photon wavelength and period. From the three independent vectors

$$\vec{e}\equiv\frac{\vec{k}}{k},\quad \vec{e}_{\parallel}\equiv\frac{\vec{B}_{0,n}}{B_{0,n}},\quad \vec{e}_{\perp}, \tag{2}$$

we define a direct orthonormal basis of the three dimensional space $(\vec{e}_{\parallel},\vec{e}_{\perp},\vec{e})$. $\vec{B}_{0,n}$ is the component of $\vec{B}_0$ perpendicular to the direction of propagation $\vec{k}$.

The electromagnetic wave derives from a potential vector that can be chosen to be of the form $\vec{A}=i(A_{\parallel}(s),A_{\perp}(s),0)e^{-i\omega t}$ where s is the coordinate along the direction of propagation.[1] With this decomposition, the coupled Klein-Gordon and Maxwell equations derived from the action (1) read

$$(\Box-m_a^2)a=\frac{4B_{0,n}}{M_a}\omega A_{\parallel}$$

$$\Box A_{\lambda}=\frac{4B_{0,n}}{M_a}\omega\delta_{\lambda\parallel}a, \tag{3}$$

where $\lambda=\parallel,\perp$ denotes the polarization.

Assuming that the magnetic field varies in space on scales much larger than the photon wavelength, we can perform the expansion $\omega^2+\partial_s^2=(\omega+i\partial_s)(\omega-i\partial_s)=(\omega+k)(\omega-i\partial_s)$ for a field propagating in the $+s$ direction. If we assume a general dispersion relation of the form $\omega=nk$ and that the refractive index n satisfies $|n-1|\ll 1$, we may approximate $\omega+k=2\omega$ and $k/\omega=1$. This approximation can be understood as a WKB limit where we set $A(s)=|A(s)|e^{iks}$ and assumes that the amplitude $|A|$ varies slowly, i.e. that $\partial_s|A|\ll k|A|$. In that limit, the above system (3) reduces to the linearized system

$$(\omega-i\partial_s+\mathcal{M})\begin{bmatrix}A_{\perp}\\ A_{\parallel}\\ a\end{bmatrix}=0. \tag{4}$$

The mixing matrix $\mathcal{M}$ is defined by

$$\mathcal{M}\equiv\begin{pmatrix}\Delta_{\perp} & 0\\ 0 & \mathcal{M}_{\parallel}\end{pmatrix}\quad\text{with}\quad \mathcal{M}_{\parallel}\equiv\begin{pmatrix}\Delta_{\parallel} & \Delta_M\\ \Delta_M & \Delta_m\end{pmatrix}. \tag{5}$$

The coefficients Δ_M and Δ_m are given by

$$\Delta_M=\frac{2B_{0,n}}{M_a},\quad \Delta_m=-\frac{m_a^2}{2\omega}. \tag{6}$$

The terms Δ_{λ} can be decomposed as $\Delta_{\lambda}=\Delta_{\mathrm{QED}}+\Delta_{\mathrm{CM}}+\Delta_{\mathrm{plasma}}$. The first term contains the effect of vacuum polarization giving a refractive index to the photon (see e.g. Ref. [25]) and can be computed by adding the Euler-Heisenberg effective Lagrangian which is the lowest order term of the non-linearity of the Maxwell equations in vacuum (see e.g. [26,27]) to the action (1).[2] The second term describes the Cotton-Mouton effect, i.e. the birefringence of gases and liquids in the presence of a magnetic field and the third term the effect of the plasma (since, in general, the photon does not propagate in vacuum). Their explicit expressions are given by

$$\Delta_{\mathrm{QED}}^{\parallel}=\frac{7}{2}\omega\xi,\quad \Delta_{\mathrm{QED}}^{\perp}=2\omega\xi,$$

$$\Delta_{\mathrm{plasma}}=-\frac{\omega_{\mathrm{plasma}}^2}{2\omega},$$

$$\Delta_{\mathrm{CM}}^{\parallel}-\Delta_{\mathrm{CM}}^{\perp}=2\pi CB_0^2 \tag{7}$$

with $\xi\equiv(\alpha/45\pi)(B_{0,n}/B_c)^2$, $B_c\equiv m_e^2/e=4.41\times10^{13}$ G, m_e the electron mass, e the electron charge and α the fine structure constant. C is the Cotton-Mouton constant [28]; its effect is to give only the difference of the refractive indices and the exact value of C is hard to determine [29]. The plasma frequency ω_{plasma} is defined by

$$\omega_{\mathrm{plasma}}^2\equiv4\pi\alpha\frac{n_e}{m_e}, \tag{8}$$

n_e being the electron density. Note that Δ_m is always negative whereas Δ_{λ} is positive if the contribution of the vacuum dominates and negative when the plasma term dominates.

As seen from Eq. (4), only the component $\parallel$, i.e. parallel to the magnetic field, couples to the pseudoscalars, a first consequence of which is that the polarization plane of a light beam traveling in a magnetic field will rotate.

The solution to the equation of motion (4) is obtained by diagonalizing $\mathcal{M}_{\parallel}$ through a rotation

$$\begin{bmatrix}A'_{\parallel}\\ a'\end{bmatrix}=\begin{pmatrix}\cos\vartheta & \sin\vartheta\\ -\sin\vartheta & \cos\vartheta\end{pmatrix}\begin{bmatrix}A_{\parallel}\\ a\end{bmatrix} \tag{9}$$

with the *mixing angle* ϑ given by

$$\tan2\vartheta\equiv2\frac{\Delta_M}{\Delta_{\parallel}-\Delta_m}. \tag{10}$$

By solving Eq. (4) in this new basis, one can easily compute the probability of oscillation of a photon after a distance of flight s starting from the initial state $[A_{\parallel}(0)=1,a(0)=0]$. It is explicitly given by

$$P(\gamma\to a)\equiv|\langle A_{\parallel}(0)|a(s)\rangle|^2$$

[1]With this convention the electric field of the wave is simply $\vec{E}\equiv-\partial_t\vec{A}=(\omega A_{\parallel}(s),\omega A_{\perp}(s),0)e^{-i\omega t}$.

[2]The equation of motion derived from Eq. (1) is Eq. (4) with $\Delta_{\lambda}=0$. We intentionally omit the Euler-Heisenberg contribution in the presentation for the sake of clarity. Its Lagrangian is explicitly given by $\mathcal{L}_{EH}=(\alpha^2/90m_e^4)[(F^{\mu\nu}F_{\mu\nu})^2+\frac{7}{4}(F^{\mu\nu}\tilde{F}_{\mu\nu})^2]$.

$$=\sin^2(2\vartheta)\sin^2\left(\frac{\Delta_{\rm osc}}{2}s\right), \tag{11}$$

$$=(\Delta_M s)^2\frac{\sin^2(\Delta_{\rm osc}s/2)}{(\Delta_{\rm osc}s/2)^2} \tag{12}$$

with the (reduced) oscillation wave number $\Delta_{\rm osc}$ given by

$$\Delta_{\rm osc}=\frac{\Delta_\parallel-\Delta_m}{\cos 2\vartheta}=\frac{2\Delta_M}{\sin 2\vartheta}. \tag{13}$$

The oscillation length is thus given by $l_{\rm osc}\equiv 2\pi/\Delta_{\rm osc}$. We see that a complete transition between a photon and a pseudoscalar is only possible when the mixing is maximal (*strong mixing regime*) i.e. when $\vartheta\simeq\pi/4$.

III. APPLICATION TO SUPERNOVAE

The quantities required for our discussion are Δ_M, Δ_m, $\Delta_{\rm plasma}$ and $\Delta_{\rm QED}$ respectively given by Eqs. (6) and (7). It is useful to rewrite them as

$$\frac{\Delta_M}{1\ \ {\rm cm}^{-1}}=2\times10^{-26}\left(\frac{B_0}{10^{-9}\ \ {\rm G}}\right)\left(\frac{M_a}{10^{11}\ \ {\rm GeV}}\right)^{-1},$$

$$\frac{\Delta_m}{1\ \ {\rm cm}^{-1}}=-2.5\times10^{-28}\left(\frac{m_a}{10^{-16}\ \ {\rm eV}}\right)^2\left(\frac{\omega}{1\ \ {\rm eV}}\right)^{-1},$$

$$\frac{\Delta_{\rm plasma}}{1\ \ {\rm cm}^{-1}}=-3.6\times10^{-24}\left(\frac{\omega}{1\ \ {\rm eV}}\right)^{-1}\left(\frac{n_e}{10^{-7}\ \ {\rm cm}^{-3}}\right),$$

$$\frac{\Delta_{\rm QED}}{1\ \ {\rm cm}^{-1}}=1.33\times10^{-45}\left(\frac{\omega}{1\ \ {\rm eV}}\right)\left(\frac{B_0}{10^{-9}\ \ {\rm G}}\right)^2, \tag{14}$$

where we have used the facts that $1\ \ {\rm eV}\simeq5\times10^4\ \ {\rm cm}^{-1}$, and $1\ \ {\rm G}\simeq1.95\times10^{-2}\ \ {\rm eV}^2$ in the natural Lorentz-Heaviside units where $\alpha=e^2/4\pi=1/137$.

The parameters chosen in [15] are $M_a\sim4\times10^{11}$ GeV, $m_a\sim10^{-16}$ eV and $B_0\sim10^{-9}$ G. The intergalactic medium (IGM) today is fully ionized, as indicated by the lack of Gunn-Petterson effect [30]. Thus the mean electronic density can be estimated to be (see e.g. [31])

$$n_e\simeq10^{-7}\ \ {\rm cm}^{-3}. \tag{15}$$

One immediately sees that, with this choice of parameters, $\Delta_{\rm QED}$ is always negligible, whereas one has $|\Delta_{\rm plasma}|\gg|\Delta_m|$ so that the plasma effects are always dominant over the pure mass term of the pseudoscalar.

In order to be more specific, let us compare the mixing angles ϑ, oscillation wave numbers $\Delta_{\rm osc}$ and oscillation probabilities $P(\gamma\to a)$ with and without including the effect of the intergalactic plasma, and for similar choice of parameters (we will use a subscript 0 for the values computed without the plasma effect).

When one sets Δ_λ to zero in Eq. (5), the mixing angle reduces to

$$\tan 2\vartheta_0\equiv-2\frac{\Delta_M}{\Delta_m}\sim40\left(\frac{\omega}{1\ \ {\rm eV}}\right), \tag{16}$$

so that for optical photons (with $\omega\sim1.5$ eV to 3 eV) $\tan 2\vartheta_0\gg1$ and $\vartheta_0\sim\pi/4$. This corresponds to a regime in which $\omega\gg m^2/\Delta_M$ and the oscillation probability does not depend on ω so that the oscillation is *achromatic*. The oscillation wave number $\Delta_{\rm osc,0}$ is given then by

$$\Delta_{\rm osc,0}\sim2\Delta_M, \tag{17}$$

which is also independent of ω. With the choice of parameters used in [15], $\Delta_{\rm osc,0}\simeq10^{-26}\ \ {\rm cm}^{-1}$, so that the oscillation length is larger than the size s of the domain of coherence of the magnetic field considered which is of the order of a Mpc ($\sim3\times10^{24}$ cm). The probability of oscillation over a domain of size s is then well approximated by

$$P_0(\gamma\to a)\sim(\Delta_M s)^2, \tag{18}$$

which is of order 10^{-4}. The number of such domains in our Hubble radius H_0^{-1}, and on a given line of sight, is given by H_0^{-1}/s. If one considers that the universe is made by patching together such domains with uncorrelated $\vec{B}_0$, the coherence is lost from domain to domain and one can simply sum up the probability of conversion over each domain to obtain the probability of conversion of a photon on cosmological distances given by (see also Sec. IV)

$$P_{0,tot}(\gamma\to a)\sim\Delta_M^2 sH_0^{-1}. \tag{19}$$

This number is of order 1, and one can thus expect a significant reduction of the luminous flux over cosmological distances. This is the bottom line of the mechanism proposed by Csáki *et al.* [15].

Let us now include plasma effects. Since $|\Delta_{\rm plasma}|\gg|\Delta_m|$, the mixing angle ϑ is now much smaller than ϑ_0 and

$$\vartheta\sim\frac{\Delta_M}{\Delta_{\rm plasma}} \tag{20}$$

of order $10^{-3}-10^{-2}$. With such a low mixing angle (*weak mixing regime*), the probability of oscillation $P(\gamma\to a)$ over a domain of size s can be approximated by

$$P(\gamma\to a)\sim(\Delta_M s)^2\frac{\sin^2(\Delta_{\rm plasma}s/2)}{(\Delta_{\rm plasma}s/2)^2}$$
$$\sim P_0(\gamma\to a)\frac{\sin^2(\Delta_{\rm plasma}s/2)}{(\Delta_{\rm plasma}s/2)^2}. \tag{21}$$

Notice then that the oscillation wave number $\Delta_{\rm osc}$ is given by

$$\Delta_{\rm osc}\sim\Delta_{\rm plasma}, \tag{22}$$

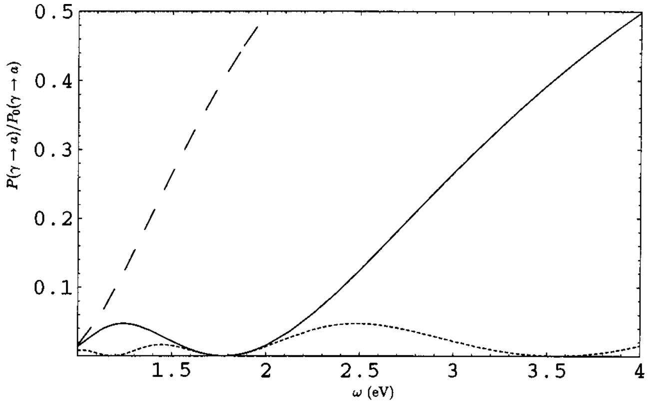

FIG. 1. Ratio between the probability of oscillation of a photon into a pseudoscalar including the effect of the intergalactic plasma and the probability of oscillation when this effect is not considered. The curves are drawn as a function of the photon energy, for $M_a \sim 4\times 10^{11}$ GeV, $m_a \sim 10^{-16}$ eV and $B_0 \sim 10^{-9}$ G, and for a distance of flight 0.5 Mpc (dashed line), 1 Mpc (solid line) and 2 Mpc (dotted line).

so that the oscillation length is smaller than previously and is of the same order as the size of the domain of coherence considered.

It follows that the probability of oscillation is lower than in the previous case (with no plasma effects taken into account) and that it is no longer achromatic (see Fig. 1). Supernovae observations not only argue for a dimming of distant supernovae but also argue for an effect that is achromatic, we will discuss this in more detail in the next section.

It is important to realize that when plasma effects are considered there are two sources for the loss of coherence, spatial fluctuations in the magnetic field and variations in the number density of electrons (i.e. changes in the plasma frequency). For example, in the case of photon-pseudoscalar conversion in the interstellar medium of our galaxy [18], the coherence length is most likely set by the fluctuations in the plasma frequency. When this is the source of coherence loss one expects generically that s in Eq. (21) also depends on frequency.

IV. CHROMATICITY CONSTRAINTS AND COHERENCE LENGTH

As we mentioned in the previous section, once plasma effects are considered the conversion probability can depend on photon frequency. In this section we consider such constraints.

Supernovae observations put a constraint on what is called the color excess between the B and V wavelength bands ($E[B-V]$). The color excess is defined as

$$E[B-V]\equiv -2.5\log_{10}\left[\frac{F^o(B)}{F^e(B)}\,\frac{F^e(V)}{F^o(V)}\right], \qquad (23)$$

where F^o (F^e) is the observed (emitted) flux and the B (V) band corresponds to 0.44 μm (0.55 μm). Observations constrain $E[B-V]$ to be lower than 0.03 [13]. This can be translated to

$$P(\gamma\to a)_V\left[\frac{P(\gamma\to a)_B}{P(\gamma\to a)_V}-1\right]<0.03, \qquad (24)$$

or equivalently to the statement that the conversion probability has to scale with photon wavelength weaker than $\lambda^{0.6}$ near the visual band.

We now consider different limits of Eq. (21) to investigate the chromatic behavior. It is useful to introduce the plasma length scale $\ell_p = \Delta_{\rm plasma}^{-1}$. We can rewrite Eq. (21) as

$$P(\gamma\to a)\sim(2\Delta_M\ell_p)^2\sin^2(s/2\ell_p). \qquad (25)$$

In any astrophysically realistic situation, the effective coherence length is either set by the magnetic field or by the plasma frequency, so that it depends on properties of the IGM which are expected to have a significant dispersion,[3] and will then similarly exhibit such a dispersion.

We start by considering the case in which the coherence length is set by the magnetic field domains. For simplicity we will take s to be distributed as a Gaussian variable with mean s^* and dispersion βs^*. The average of Eq. (25) over s yields

$$\langle P(\gamma\to a)\rangle \sim 2(\Delta_M\ell_p)^2[1-\cos(s^*/\ell_p)e^{-(\beta s^*/\ell_p)^2/2}]. \qquad (26)$$

In the limit $s^* \gg \ell_p$, it reduces to

$$\langle P(\gamma\to a)\rangle \sim 2(\Delta_M\ell_p)^2. \qquad (27)$$

The plasma length ℓ_p scales with the photon frequency so that the conversion probability is proportional to λ^{-2} which is ruled out by the constraint in Eq. (24). Notice also that a conversion probability that grows with frequency would produce a very significant reddening of sources located at distances such that the probability reaches its saturation value for the shorter wavelengths. The absence of such reddening may provide even more stringent constraints on the required achromaticity than that of Eq. (24).

In the opposite limit, $s^* \ll \ell_p$, we get

$$\langle P(\gamma\to a)\rangle \sim (\Delta_M s^*)^2, \qquad (28)$$

which is achromatic if s^* does not depend on frequency. The total probability of conversion will still be given by Eq. (19) with s^* replacing s, however since one must have $s^* \ll \ell_p$ for Eq. (28) to hold, and since $\ell_p \sim 0.1$ Mpc, this means that the total probability of conversion will be lower than the one computed in [15] by at least a factor 10 to 100. One can try to overcome this by changing the coupling of the pseudoscalar, it is however difficult because the coupling is already near the astrophysical bound. Alternatively, the conversion

[3]For example, in the model of Ref. [32], the magnetic field in the intergalactic medium was generated in quasars and expelled in their outflows. The distribution of bubble sizes at redshifts around zero is predicted to be extremely broad with two main components, sizes ranging from 0.5 Mpc to about 5 Mpc, and a mean size of order 1 Mpc. The biggest bubbles have the largest magnetic field.

probability may remain achromatic as well as sufficiently large if the intergalactic magnetic field were stronger in domains of size $s^* \ll \ell_p$. Faraday rotation measurements of distant quasars impose a conservative bound $B_0 \sqrt{s_c} \leqslant 10^{-9}$ G $\sqrt{\mathrm{Mpc}}$ on the strength of an intergalactic magnetic field coherent over scales s_c [33] (it may be somewhat stronger depending on its spatial structure [34]). The conversion probability could be achromatic in the visible band and become of order unity at cosmological distances if the intergalactic magnetic field had some very definite spectral properties, for instance if B_0 were of order 10^{-8} G over domains of average size of the order of 10 kpc, and sufficiently weaker on larger domains. Additional constraints on the spatial distribution of the intergalactic magnetic field compatible with the proposed mechanism arise from preventing excessive dispersion in the observed peak luminosity of distant supernovae. Clusters, for instance, may have magnetic fields significantly strong and extended to make the conversion probability of order unity as photons get across them.

Another possibility is to consider the case $s^* \sim \ell_p$ and to require that the average in Eq. (26) be achromatic enough. This usually does not happen, although it can be accommodated by correctly choosing β and s^*, or by a tight correlation between the strength of the magnetic field and the sizes of the domains, and requires very special, and likely unrealistic, statistical properties of the IGM.

Let us now turn to the case where the coherence length is determined by the spatial variations of the electron density. In this case the required achromaticity would even be more of a fine tuning as the coherence length $\overline{s}$ will also depend on frequency. Just as in the case of the interstellar medium studied in [18], the frequency dependence will be set by the clustering properties of the electron density.

In the case at hand we expect the intergalactic medium to be very clumpy and complicated. Current numerical simulations indicate that the baryons today can be found in several phases. About 30% by mass is in a warm phase ($T \sim 5000$ K) that fills most of the volume of the universe, about another 30% is in a warm-hot phase that resides in non-virialized objects such as filaments and the rest resides in virialized objects such as clusters and in condensed forms such as stars and cool galactic gas (see e.g. [35] and references therein).

The most relevant phase for our study is the warm one because it fills most of the volume. It could be clumpy on scales smaller than ℓ_p, so the clumpiness of the IGM may be the most likely source of coherence loss. However a detailed study of the loss of coherence should probably involve studying lines of sight across these type of simulations. It is important to realize that even in the limit where the clumping scale, L_c, of the warm phase of the IGM is smaller than ℓ_p we still expect that $\overline{s}$ will depend on frequency. For coherence to be lost, the random component of the accumulated phase of the oscillation, ϕ, has to be of order one. The phase on each segment of length L_c is L_c/ℓ_p and accumulates over different segments as a random walk, $\phi \sim L_c/\ell_p \sqrt{\overline{s}/L_c} \sim 1$. In this limit, we estimate $\overline{s} \sim \ell_p^2/L_c$ which will again induce a chromaticity that is ruled out by observations.

It seems that the only natural ways to avoid the chromaticity constraint is (i) to assume that the magnetic field is responsible for setting the coherence length of the oscillation and assume that $s^* \ll \ell_p$, in which case either the coupling of the pseudoscalar needed to accommodate the dimming of the supernovae becomes uncomfortably large or the magnetic field must have very definite strength and spectral features, or (ii) to have very constrained properties of the IGM.

Finally, we note that the situation is not improved by giving a higher mass to the pseudoscalar in order to have $\Delta_m > \Delta_{\rm plasma}$; this choice leads as well to chromaticity (since Δ_m and $\Delta_{\rm plasma}$ have the same spectral dependence) and lowers the probability of oscillation (and the mixing angle). A last logical possibility is that the mass of the pseudoscalar is such that Δ_m and $\Delta_{\rm plasma}$ are of the same order. In this case, a strong mixing regime is possible whenever the electron density is such that Δ_m and $\Delta_{\rm plasma}$ coincide with each other with accuracy Δ_M. This can happen in the IGM from the statistical fluctuations in the electronic density. When this is the case, a strong photon-pseudoscalar conversion can in principle take place (this is very analogous to the resonant MSW effect of neutrino physics). However, for the transition to be significant one has to maintain the resonant condition ($\Delta_m \sim \Delta_{\rm plasma}$ with accuracy Δ_M) over a distance of the order of the oscillation length, π/Δ_M, which given the numerical values (14) is also very unlikely (not to mention the fact that a dimming of SNIa induced by such a resonant conversion would lead to a large dispersion in the observed SNIa magnitude).

V. CONCLUSIONS

We have shown that one cannot ignore the effect of the intergalactic plasma to derive how the luminosity of distant sources, such as supernovae, is affected by a mixing with a hypothetical pseudoscalar particle. In most of the parameter space, this effect either renders the oscillation frequency-dependent or lowers too much the oscillation probability. There is a slight hope to accommodate the mechanism of [15] if the IGM has very specific statistical properties. In [36], the authors of [15] concluded that plasma effects do not strongly influence the oscillations of optical photons if the electron density is smaller than 2.5×10^{-8} cm^{-3}, assumed uniform over magnetic domains of size 1 Mpc. This conclusion is compatible with our results above. Indeed as one can clearly see from Fig. 1, and Eq. (21), one can render the oscillation achromatic by demanding that the product $\Delta_{plasma}s$ be low enough in order to be in a regime analogous to the one discussed after Eq. (28). This is achieved by lowering the electron density below 2.5×10^{-8} cm^{-3}. However, increasing the size of the coherence domain by the same factor one decreases the electron density would lead back to a chromatic regime (as also appears on Fig. 1). This confirms that the mechanism of Ref. [15] is very sensitive to the precise properties of the IGM, and demonstrates the need for much more realistic and detailed studies of the effect of coherence loss due to the electron density and magnetic field spatial variations before any significant conclusion can be drawn.

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[1] G.G. Raffelt and L. Stodolsky, Phys. Rev. D **37**, 1237 (1988).
[2] P. Sikivie, Phys. Rev. Lett. **51**, 1415 (1983).
[3] P. Sikivie, Phys. Rev. Lett. **52**, 695 (1984).
[4] D.E. Morris, Phys. Rev. D **34**, 843 (1986).
[5] M.E. Gertsenshtein, Sov. Phys. JETP **14**, 84 (1962).
[6] Ya.B. Zel'dovich and I.D. Novikov, *The Structure and Evolution of the Universe*, Relativistic Astrophysics Vol. 2 (Chicago University Press, Chicago, 1983).
[7] K. van Bibber *et al.*, Phys. Rev. D **39**, 2089 (1989).
[8] C. Hagmann *et al.*, Phys. Rev. Lett. **80**, 2043 (1998).
[9] J.E. Kim, Phys. Rev. D **58**, 055006 (1998).
[10] P. Sikivie, Phys. Lett. B **432**, 139 (1998).
[11] H. Georgi, P. Ginsparg, and S.L. Glashow, Nature (London) **306**, 765 (1983).
[12] G.G. Raffelt, in Proceedings of Neutrino '98, Japan, edited by Y. Suzuki and Y. Totsuka [Nucl. Phys. B (Proc. Suppl.) **77**, 456 (1999)]
[13] S. Perlmutter *et al.*, Astrophys. J. **517**, 565 (1999).
[14] A.G. Riess *et al.*, Astrophys. J. **116**, 1009 (1998); A.G. Riess *et al.*, *ibid.* **560**, 49 (2001).
[15] C. Csaki, N. Kaloper and J. Terning, Phys. Rev. Lett. **88**, 161302 (2002).
[16] J. Erlich and C. Grojean, Phys. Rev. D **65**, 123510 (2002).
[17] D. Harari and P. Sikivie, Phys. Lett. B **289**, 67 (1992).
[18] E. Carlson and W. Garretson, Phys. Rev. Lett. **336**, 431 (1994).
[19] J.A. Grifols, E. Masso, and R. Toldra, Phys. Rev. Lett. **77**, 2372 (1996).
[20] J.W. Brockway, E.D. Carlson, and G.G. Raffelt, Phys. Lett. B **383**, 439 (1996).
[21] J.C.R. Magueijo, Phys. Rev. D **49**, 671 (1994).
[22] P. Chen, Phys. Rev. Lett. **74**, 634 (1995).
[23] A.N. Cillis and D. Harari, Phys. Rev. D **54**, 4757 (1996).
[24] C. Deffayet and J.-P. Uzan, Phys. Rev. D **62**, 063507 (2000).
[25] S.L. Adler, Ann. Phys. (N.Y.) **67**, 599 (1971).
[26] W. Heisenberg and H. Euler, Z. Phys. **88**, 714 (1936).
[27] C. Itzykson and J.-B. Zuber, *Quantum Field Theory* (McGraw-Hill, New York, 1980).
[28] L.D. Landau and E.M. Lifschitz, *Electrodynamics of Continuous Media* (Mir, Moscow, 1960).
[29] *American Institute of Physics Handbook*, 3rd ed., edited by D.E. Gray (McGraw-Hill, New York, 1972).
[30] J.E. Gunn and B.A. Peterson, Astrophys. J. **142**, 1633 (1965).
[31] P.J.E. Peebles, *Principles of Physical Cosmology* (Princeton University Press, Princeton, NJ, 1993).
[32] S.R. Furlanetto and A. Loeb, Astrophys. J. **556**, 619 (2001).
[33] P. Kronberg, Rep. Prog. Phys. **57**, 325 (1994).
[34] G.R. Farrar and T. Piran, Phys. Rev. Lett. **84**, 3527 (2000).
[35] R. Dave *et al.*, Astrophys. J. **552**, 473 (2001).
[36] C. Csaki, N. Kaloper, and J. Terning, Phys. Lett. B **535**, 33 (2002).