Dimming of Supernovae by Photon-Pseudoscalar Conversion and the Intergalactic Plasma

Cédric Deffayet\textsuperscript{a,} 1, Diego Harari\textsuperscript{b,} 2, Jean-Philippe Uzan\textsuperscript{c,d,} 3 and Matias Zaldarriaga\textsuperscript{a,} 4

\textsuperscript{a}NYU Dept of Physics, 4 Washington Place, New York, NY10003 (USA).
\textsuperscript{b}Departamento de Física, FCEyN - Universidad de Buenos Aires, Ciudad Universitaria - Pab. 1, 1428 Buenos Aires (Argentina).
\textsuperscript{c}Laboratoire de Physique Théorique, UMR 8627 du CNRS, Université Paris XI, Bât. 210 F–91405 Orsay Cedex (France).
\textsuperscript{d}Institut d’Astrophysique de Paris, 98bis Bd. Arago, 75014 Paris (France).

Abstract

Recently Csáki, Kaloper and Terning (hep-ph/0111311) suggested 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.

1 Introduction

The mixing between photons and axions, in an external magnetic field, is a well studied mechanism [1, 2, 3, 4] as is its analog with gravitons [5, 6]. It is experimentally used, since the pioneer work by Sikivie [4], to constrain the axion parameters [7, 8] (see also e.g. Ref. [2] for an up to date review on such

\begin{footnotesize}
\begin{itemize}
\item[1]cjd2@physics.nyu.edu
\item[2]harari@df.uba.ar
\item[3]uzan@th.u-psud.fr
\item[4]mz31@nyu.edu
\end{itemize}
\end{footnotesize}
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 non zero 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 ($m \sim 10^{-16}$ 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.

The implications on the cosmic microwave background of the similar effect involving photon-graviton oscillation has been considered in empty space [19, 20] and it was shown that it becomes negligible for standard cosmological magnetic fields [21] once the contribution of the intergalactic plasma is properly taken into account (the case of axions is also considered in [22]). The effect of the inhomogeneities of the electron density upon the coherence of the oscillations was also considered by Carlson et al. [18].

The effect of this plasma was not addressed in [15], and it is our purpose to discuss it in this work.

The paper is organized as follows: we first remind standard results on photon-pseudoscalar oscillations in order to introduce our conventions (see [1], or [22] 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.

## 2 Photon–pseudoscalar mixing

We consider the generic action for the pseudoscalar-photon system

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

(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 $1/M_a$.
wavelength and period. From the three independent vectors
\[ \vec{e} \equiv \frac{\vec{k}}{|k|}, \quad \vec{e}_\parallel \equiv \frac{\vec{B}_{0,n}}{B_{0,n}}, \quad \vec{e}_\perp, \] (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. With this decomposition, the coupled Klein–Gordon and Maxwell equations derived from the action (1) read
\[ \left( \Box - m_a^2 \right) a - \frac{4B_{0,n}}{M_a} A_\parallel \]
\[ \Box A_\lambda = \frac{4B_{0,n}}{M_a} \omega \delta_{\lambda\parallel} a, \] (3)
where \( \lambda = ||, \perp \) denotes the polarization.

Assuming that the magnetic field varies in space on scales much larger than the
photonic 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
\[ (\omega - i\partial_s + \mathcal{M}) \begin{bmatrix} A_\perp \\ A_\parallel \\ a \end{bmatrix} = 0. \] (4)

The mixing matrix \( \mathcal{M} \) is defined by
\[ \mathcal{M} \equiv \begin{pmatrix} \Delta_\perp & 0 \\ 0 & \mathcal{M}_|| \end{pmatrix} \quad \text{with} \quad \mathcal{M}_|| \equiv \begin{pmatrix} \Delta_\parallel & \Delta_M \\ \Delta_M & \Delta_m \end{pmatrix}. \] (5)

The coefficients \( \Delta_M \) and \( \Delta_m \) are given by
\[ \Delta_M = \frac{2B_{0,n}}{M_a}, \quad \Delta_m = \frac{-m_a^2}{2\omega}. \] (6)

The terms \( \Delta_\lambda \) can be decomposed as \( \Delta_\lambda = \Delta_{\text{QED}} + \Delta_{\text{CM}} + \Delta_{\text{plasma}} \). The first term contains the effect of vacuum polarization giving a refractive index to the photon

\begin{footnote}{with this convention the electric field of the wave is simply \( \vec{E} \equiv -\partial_t \vec{A} = (\omega A_\parallel(s), \omega A_\perp(s), 0)e^{-i\omega t}. \) }\end{footnote}
(see e.g. Ref. [23]) 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. [24, 25]) to the action (1). The second term describes the Cotton–Mouton effect, i.e. the birefringence of gases and liquids in 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

\[ \Delta^\parallel_{\text{QED}} = \frac{7}{2} \omega \xi, \quad \Delta^\perp_{\text{QED}} = 2 \omega \xi, \]

\[ \Delta_{\text{plasma}} = -\frac{\omega^2_{\text{plasma}}}{2\omega}, \]

\[ \Delta^\parallel_{\text{CM}} - \Delta^\perp_{\text{CM}} = \frac{2}{\pi} \frac{C B_0^2}{B_c} \]

with \( \xi \equiv \frac{\alpha}{45\pi} (B_0 n_c/B_c)^2 \), \( B_c \equiv m_e^2/e = 4.41 \times 10^{13} \text{G} \), \( m_e \) the electron mass, \( e \) the electron charge and \( \alpha \) the fine structure constant. \( C \) is the Cotton–Mouton constant [26]; its effect is to give only the difference of the refractive indices and the exact value of \( C \) is hard to determine [27]. The plasma frequency \( \omega_{\text{plasma}} \) is defined by

\[ \omega^2_{\text{plasma}} \equiv 4\pi \alpha \frac{n_e}{m_e}, \]

\( n_e \) being the electron density. Note that \( \Delta_m \) is always negative whereas \( \Delta_\lambda \) is positive if the contribution of the vacuum dominates and negative when the plasma term dominates.

As seen from Eq. (4), only the component ||, 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 \( \mathcal{M}_{||} \) through a rotation

\[ \begin{pmatrix} A'_{||} \\ a' \end{pmatrix} = \begin{pmatrix} \cos \vartheta & \sin \vartheta \\ -\sin \vartheta & \cos \vartheta \end{pmatrix} \begin{pmatrix} A_{||} \\ a \end{pmatrix} \]

with the mixing angle \( \vartheta \) given by

\[ \tan 2\vartheta \equiv \frac{2 \Delta M}{\Delta|| - \Delta_m}. \]

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 \( (A_{||}(0) = 1, a(0) = 0) \). It is explicitly given by

\[ P(\gamma \rightarrow a) \equiv | \langle A_{||}(0) | a(s) \rangle |^2 = \sin^2(2\vartheta) \sin^2 \left( \frac{\Delta_{\text{osc}} s}{2} \right), \]

\[ = (\Delta M s)^2 \frac{\sin^2(\Delta_{\text{osc}} s/2)}{(\Delta_{\text{osc}} s/2)^2} \]

\[ \]
with the (reduced) oscillation wavenumber $\Delta_{osc}$ given by
\[
\Delta_{osc} = \frac{\Delta_\| - \Delta_m}{\cos 2\vartheta} = \frac{2\Delta_M}{\sin 2\vartheta},
\]
(13)
The oscillation length is thus given by $\ell_{osc} \equiv 2\pi/\Delta_{osc}$. 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 $\vartheta \simeq \pi/4$.

3 Application to Supernovae

The quantities required for our discussion are $\Delta_M$, $\Delta_m$, $\Delta_{\text{plasma}}$ and $\Delta_{\text{QED}}$ respectively given by equations (6) and (7). It is useful to rewrite them as
\[
\begin{align*}
\frac{\Delta_M}{1\text{ cm}^{-1}} &= 2 \times 10^{-26} \left( \frac{B_0}{10^{-9}\text{G}} \right) \left( \frac{M_a}{10^{11}\text{ GeV}} \right)^{-1}, \\
\frac{\Delta_m}{1\text{ cm}^{-1}} &= -2.5 \times 10^{-28} \left( \frac{m_a}{10^{-16}\text{eV}} \right)^2 \left( \frac{\omega}{1\text{ eV}} \right)^{-1}, \\
\frac{\Delta_{\text{plasma}}}{1\text{ cm}^{-1}} &= -3.6 \times 10^{-24} \left( \frac{\omega}{1\text{ eV}} \right)^{-1} \left( \frac{n_e}{10^{-7}\text{ cm}^{-3}} \right), \\
\frac{\Delta_{\text{QED}}}{1\text{ cm}^{-1}} &= 1.33 \times 10^{-45} \left( \frac{\omega}{1\text{ eV}} \right) \left( \frac{B_0}{10^{-9}\text{G}} \right)^2,
\end{align*}
\]
(14)
where we have used the facts that $1\text{ eV} \simeq 5 \times 10^4\text{ cm}^{-1}$, $1\text{ G} \simeq 1.95 \times 10^{-2}\text{ eV}^2$ in the natural Lorentz-Heaviside units where $\alpha = e^2/4\pi = 1/137$.

The parameters chosen in [15] are $M_a \sim 4 \times 10^{11}\text{GeV}$, $m_a \sim 10^{-16}\text{eV}$ and $B_0 \sim 10^{-9}\text{G}$. The intergalactic medium (IGM) today is fully ionized, as indicated by the lack of Gunn-Petterson effect [28]. Thus the mean electronic density can be estimated to be (see e.g. [29])
\[
n_e \simeq 10^{-7}\text{ cm}^{-3}.
\]
(15)

One immediately sees that, with this choice of parameters, $\Delta_{\text{QED}}$ is always negligible, whereas one has $|\Delta_{\text{plasma}}| \gg |\Delta_m|$ 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 $\vartheta$, oscillation wavenumbers $\Delta_{osc}$ 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 $\Delta_\lambda$ to zero in Eq. (3), as implicitly done in [14], the mixing angle reduces to
\[
\tan 2\vartheta_0 \equiv -2\frac{\Delta_M}{\Delta_m} \sim 40 \left( \frac{\omega}{1\text{ eV}} \right),
\]
(16)
so that for optical photons (with $\omega \sim 1.5$ eV to 3 eV) $\tan 2\vartheta_0 \gg 1$ 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 $\omega$ so that the oscillation is *achromatic*. The oscillation wavenumber $\Delta_{\text{osc},0}$ is given then by

$$\Delta_{\text{osc},0} \sim 2\Delta_M,$$  

(17)

which is also independent of $\omega$. With the choice of parameters used in [15], $\Delta_{\text{osc},0} \simeq 10^{-26}\text{cm}^{-1}$, so that the oscillation length is larger than the size $s$ of the domain of coherence of the magnetic field considered which is of order of a Mpc ($\sim 3 \times 10^{24}\text{cm}$). The probability of oscillation over a domain of size $s$ is then well approximated by

$$P_0(\gamma \rightarrow a) \sim (\Delta_M s)^2,$$  

(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 section 4)

$$P_{0,\text{tot}}(\gamma \rightarrow a) \sim \Delta^2_M sH_0^{-1}.$$  

(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 $|\Delta_{\text{plasma}}| \gg |\Delta_m|$, the mixing angle $\vartheta$ is now much smaller than $\vartheta_0$ and

$$\vartheta \sim \frac{\Delta_M}{\Delta_{\text{plasma}}},$$  

(20)

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

$$P(\gamma \rightarrow a) \sim (\Delta_M s)^2 \frac{\sin^2 (\Delta_{\text{plasma}} s/2)}{(\Delta_{\text{plasma}} s/2)^2} \sim P_0(\gamma \rightarrow a) \frac{\sin^2 (\Delta_{\text{plasma}} s/2)}{(\Delta_{\text{plasma}} s/2)^2}.$$  

(21)

Notice then that the oscillation wavenumber $\Delta_{\text{osc}}$ is given by

$$\Delta_{\text{osc}} \sim \Delta_{\text{plasma}},$$  

(22)

so that the oscillation length is smaller than previously and is of the same order as the size of the domain of coherence considered.
Figure 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 (as done in [15]). 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).

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 details 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 equation (21) also depends on frequency.
4 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

$$E[B - V] \equiv -2.5 \log_{10} \left[ \frac{F^o(B)}{F^e(B)} \frac{F^e(V)}{F^o(V)} \right],$$  \hspace{1cm} (23)

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

$$P(\gamma \to a)_V \left[ \frac{P(\gamma \to a)_B}{P(\gamma \to a)_V} - 1 \right] < 0.03,$$  \hspace{1cm} (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 equation (21) to investigate the chromatic behavior. It is useful to introduce the plasma length scale $\ell_p = \Delta^{-1}_{\text{plasma}}$. We can rewrite (21) as

$$P(\gamma \to a) \sim (2\Delta_M \ell_p)^2 \sin^2(s/2\ell_p).$$  \hspace{1cm} (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\footnote{For example, in the model of Ref. 30, 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.} 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 $\beta s^*$. The average of equation (25) over $s$ yields,

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

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

$$\langle P(\gamma \to a) \rangle \sim 2(\Delta_M \ell_p)^2.$$  \hspace{1cm} (27)
The plasma length $\ell_p$ scales with the photon frequency so that the conversion probability is proportional to $\lambda^{-2}$ which is ruled out by the constraint in equation (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

$$\langle P(\gamma \rightarrow a) \rangle \sim (\Delta M s^*)^2,$$

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 (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 Ref. [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 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 $B_0 \sqrt{s_c} \leq 10^{-9}$ G $\sqrt{\text{Mpc}}$ on the strength of an intergalactic magnetic field coherent over scales $s_c$ [31] (it may be somewhat stronger depending on its spatial structure [32]). 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 equation (29) be achromatic enough. This usually does not happen, although it can be accommodated by correctly choosing $\beta$ 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 $\bar{s}$ will also depend on frequency. Just as in the case of the interstellar medium studied in Ref. [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 \((T \sim 5000K)\) 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. [33] 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 \(\ell_p\), 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 \(\ell_p\) we still expect that \(\bar{s}\) will depend on frequency. For coherence to be lost, the random component of the accumulated phase of the oscillation, \(\phi\), has to be of order one. The phase on each segment of length \(L_c\) is \(L_c/\ell_p\) and accumulates over different segments as a random walk, \(\phi \sim L_c/\ell_p \sqrt{\bar{s}/L_c} \sim 1\). In this limit, we estimate \(\bar{s} \sim \ell_p^2/L_c\) 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 \(\Delta_m > \Delta_{\text{plasma}}\); this choice leads as well to chromaticity (since \(\Delta_m\) and \(\Delta_{\text{plasma}}\) 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 \(\Delta_m\) and \(\Delta_{\text{plasma}}\) are of the same order. In this case, a strong mixing regime is possible whenever the electron density is such that \(\Delta_m\) and \(\Delta_{\text{plasma}}\) coincide with each other with accuracy \(\Delta_M\). 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 \((\Delta_m \sim \Delta_{\text{plasma}}\) with accuracy \(\Delta_M\)) over a distance of the order of the oscillation length, \(\pi/\Delta_M\), 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).
5 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.

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