New Cosmological Constraints on Axions and Axion-Like Particles

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Abstract. We present constraints on the axion-photon coupling scale $M$ for light axion-like particles by combining recent Supernova Type Ia data with the latest measurements of the Hubble expansion at redshifts between 0 and 2. Allowing for a coupling between axions and photons leads to a modification of the inferred luminosity distances for supernovae due to the conversion of photons to axions in the presence of intergalactic magnetic fields. We constrain such couplings by considering deviations from the luminosity-angular diameter distance relation $d_L = d_A(1+z)^2$. We find that for intergalactic magnetic fields of order 1 nG, current supernova and Hubble expansion data rule out a region in the coupling scale $M$ between $10^{10}$ and $10^{11}$ GeV.

1. Introduction
Axions are light pseudo-scalar particles, first introduced by Peccei and Quinn (PQ) [1] as a solution to the strong-CP problem in Quantum ChromoDynamics (QCD). The essence of this problem is that the dimensionless parameter $\theta$, controlling the CP-violating term allowed by QCD:

$$L_{CP} = \frac{\alpha_s}{4\pi} \theta \text{tr} G_{\mu\nu} \tilde{G}^{\mu\nu} \equiv \frac{\alpha_s}{4\pi} \theta \epsilon^{\mu\nu\rho\lambda} \text{tr} G_{\mu\nu} \tilde{G}^{\rho\lambda},$$

(1)

where $G_{\mu\nu}$ is the gluonic field strength ($\alpha_s$ being the strong coupling constant), is constrained experimentally [2] to be unnaturally small, $|\theta| \lesssim 10^{-10}$. Since $\theta$ can a priori take values anywhere in the range between 0 and 2$\pi$, this is a fine-tuning problem with no explanation within the Standard Model. The PQ solution is then to promote $\theta$ to a dynamical field, the axion $a$, with a global shift symmetry $a \rightarrow a + \text{const}$ that can only be broken by anomalous CP-violating terms. The axion thus acquires a potential with minimum at $a = 0$ and so its vacuum expectation value is dynamically driven to small values, thus solving the strong CP-problem. Due to the explicit breaking of the shift symmetry down to a discrete subgroup, the axion becomes a Pseudo-Nambu-Goldstone boson [3] with a parametrically small mass:

$$m_a \approx 0.6 \text{ meV} \times \left( \frac{10^{10}\text{GeV}}{f_a} \right),$$

(2)

where, $f_a$ is the axion decay constant determining the strength of interaction between the axion and the other Standard Model fields. In particular, the axion couples to the photon through a term:

$$L_{a\gamma\gamma} = \alpha \frac{f_a}{8\pi} A_\mu F_{\mu\nu} \tilde{F}^{\mu\nu},$$

(3)
where $F_{\mu\nu}$ is now the electromagnetic field strength and $\alpha$ the fine structure constant.

Even though axions were historically introduced to solve a concrete particle physics problem arising from a puzzling experimental result, as summarised above, it was soon realised that they could play an important role in cosmology and, in particular, provide a significant contribution to Cold Dark Matter (CDM). Indeed, an initial misalignment of the angle $\theta$ from zero, leads to oscillations of the zero-mode of the axion field as the universe cools to $T \ll f_a$. This corresponds to a zero-momentum condensate, behaving as CDM, with a potentially significant contribution to the matter density at the present epoch (for classic reviews see [4, 5]). Further, it is now understood that similar light (pseudo-)scalars, coupling to electromagnetism through two-photon terms like in equation (3), are ubiquitous in fundamental physics (see for example [6, 7, 8]).

Notably, such Axion-Like Particles (ALPs) arise in both Heterotic and Type II string theories as zero-modes of certain antisymmetric form fields.

Of particular cosmological interest is the fact that ALPs occur generically in the very constructions used in Type IIB string theory to build models of cosmic inflation [9, 10, 11]. Indeed, the worldvolume action for a $Dp$-brane, on which these models are based, reads:

$$S_p = -T_p \left( \int d^{p+1}x \sqrt{-g} e^{-\phi} \sqrt{\det \left( g + B + 2\pi \alpha' F \right)} + i \sum_q C_q \wedge e^{B + 2\pi \alpha' F} \right),$$

where $T_p$ is the brane tension, $\phi$ the dilaton, $g$ the induced metric on the brane worldvolume, $B$ the antisymmetric 2-form, $F$ the field strength of $U(1)$ gauge fields living on the worldvolume, and $C_q$ the Ramond-Ramond (RR) $q$-forms.

The inflationary sector is contained in the first term, known as the Dirac-Born-Infeld (DBI) piece, which generates kinetic terms for worldvolume scalars that play the role of inflaton fields. In the most well-studied class of brane inflation models [12, 13, 14, 15], the inflaton is identified with a certain combination of the embedding fields $Y^M$, worldvolume scalars describing the position of the brane in the bulk. This is captured by the $\sqrt{\det g}$ term, where the pullback $g_{\mu\nu} = G_{MN} \partial_\mu Y^M \partial_\nu Y^N$ of the bulk metric $G_{MN}$ on the worldvolume generates kinetic terms for the scalar fields $Y^M$. The gauge field term $F_{\mu\nu}$ generates, to lowest order, Maxwell electrodynamics $F_{\mu\nu} F^{\mu\nu}$. For $p > 3$, the $Dp$-brane can wrap one-cycles, in which case $F_{\mu\nu} F^{\mu\nu}$ contains kinetic terms for scalar fields that correspond to Wilson Lines along a combination of these one-cycles. This term can also be used for inflation in the context of Wilson line models [19, 20]. Finally, the term involving the (pullback) 2-form $B_{\mu\nu}$ has also been successfully used to generate (axion monodromy) inflation [21].

The second term in the action (4) is known as the Wess-Zumino (WZ) piece, generating all possible combinations of wedge products of the RR $q$-forms $C_q$ and powers of $B$ and $F$ that integrate on the $(p+1)$-dimensional worldvolume. In most brane inflation models, its main role is to cancel the constant (energy) term in the DBI expansion, but it is of central interest here because it contains axions. In the case of a $D3$-brane, for example, expanding the exponential to 2nd order in $F$ yields a term:

$$S_{\text{axion}} \propto \int C_0 F \wedge F = \int C_0 F_{\mu\nu} F_{\kappa\lambda} \, dx^\mu \wedge dx^\nu \wedge dx^\kappa \wedge dx^\lambda =$$

1. The potential $V(Y)$ comes from a Coulomb-like interaction between different branes and can be computed through a cylinder string amplitude [16]. It receives, however, important corrections from moduli stabilisation effects [14, 17, 18].

2. The dilaton, $\phi$, even though was historically the first field to be considered as an inflaton candidate in string theory, generically has too steep (exponential) potential and is unsuitable for inflation [22].

3. Another important role of the WZ piece is in understanding the relation between cosmic strings produced at the end of brane inflation and $D1$-branes. Expanding the exponential to linear order yields a term $C_2 \wedge F$, which can be used to argue that topological defects produced by Higgsing the gauge field $F$ couple to the RR 2-form $C_2$ and so can be identified with $D1$-branes [23].
\[ = \int C_0 F_{\mu\nu} F_{\kappa\lambda} \epsilon^{\mu\nu\kappa\lambda} d^4x = \int C_0 F_{\mu\nu} \tilde{F}^{\mu\nu} d^4x. \] (5)

This is the same term as in (3) with the axion identified with the RR scalar \( C_0 \), while the constant of proportionality (coupling strength) depends on a combination of fundamental parameters like the string scale and string coupling. Thus, even though axions are not necessarily part of the inflationary sector, they are generically present in the underlying theory and could be used to provide additional constraints on fundamental hi-energy physics parameters through their coupling (5) to photons. Indeed, such couplings can lead to the conversion of ordinary photons into axions (or ALPs) in the presence of a magnetic field (cf the Primakoff effect [24]).

Here, we present constraints on axion-photon couplings by allowing for the conversion of photons originating from Type Ia Supernovae (SN) into axions in the presence of intergalactic magnetic fields. This gives rise to an opacity effect at cosmological redshifts, which can be constrained by confronting luminosity and angular diameter distance measurements with the Etherington relation [25]

\[ d_L(z) = (1 + z)^2 d_A(z), \] (6)

where \( d_L(z) \) and \( d_A(z) \) are respectively the luminosity and angular diameter distances at redshift \( z \). This “distance-duality” relation is valid in any cosmological background where photons travel on null geodesics, provided that local Lorentz invariance is respected and, crucially, photon number is conserved. Since, for the optical frequencies we are interested in (SN observations), deviations from Lorentz invariance are strongly constrained [26], one can interpret any such opacity effect in terms of photon-axion conversions and use the data to constrain the fundamental coupling scale quantifying terms like (3) and (5). We find that, for intergalactic magnetic fields of the order of 1 nG, current SN and cosmic expansion data rule out a region in the axion-photon coupling between \( 10^{10} \) and \( 10^{11} \) GeV.

2. Constraining cosmic opacity

Let us first describe the general method and place constraints on cosmic opacity \( \tau(z) \) (defined below) without specifying the details of the underlying theory [27]. In section 3 we will interpret these constraints as a photon-axion mixing effect, which actually imposes a particular parametrisation of \( \tau \) [28]. We use Type Ia SN brightness data (namely the SCP Union 2008 compilation [29]) in combination with measurements of cosmic expansion \( H(z) \) from differential ageing of luminous red galaxies (LRGs) [30, 31, 32] to constrain opacity at redshifts up to \( z \approx 2 \). The basic idea is to study possible deviations from the Etherington relation (6), the distance duality between luminosity and angular diameter distance. As discussed in section 1, systematic violations of this relation can be interpreted as an opacity effect in the observed luminosity distance, which we quantify in terms of an opacity function \( \tau(z) \):

\[ d_{L,obs}^2 = d_{L,true}^2 e^{\tau(z)}. \] (7)

Therefore, the inferred distance moduli for the observed SNe pick a term linear in \( \tau(z) \):

\[ DM_{obs}(z) = DM_{true}(z) + 2.5[\log e]\tau(z). \] (8)

On the other hand, we can also use other observations independent of \( \tau \), to determine the right-hand-side (\( d_A \)) in equation (6). This approach was initiated in [33] (see also [34, 35, 36, 37]) where the authors used measurements [38] of the baryon acoustic oscillation (BAO) scale at two redshifts, \( z = 0.20 \) and \( z = 0.35 \), obtaining an upper-bound for the difference in opacity between these two redshifts, \( \Delta \tau < 0.13 \) at 95% confidence. In reference [27] we improved this constraint by one order of magnitude and extended it over a wider redshift range by using
instead measurements of cosmic expansion $H(z)$ from differential ageing of LRGs at redshifts $z \lesssim 2$. Introducing a parameter $\epsilon$, we studied deviations from the Etherington relation of the form:

$$d_L(z) = d_A(z)(1 + z)^{2+\epsilon},$$  \hspace{1cm} (9)

and constrained this parameter to be $\epsilon = -0.01^{+0.08}_{-0.09}$ at 95% confidence. Using the latest $H(z)$ data [32] combined with the most recent determination of $H_0$ (Riess et al [39]) improves this constraint to $\epsilon = -0.04^{+0.08}_{-0.07}$, again at 95% confidence [28].

Fig. 1 shows our joint constraints combining the $H(z)$ and SN datasets described above. On the left, the dark blue contours correspond to the (two-parameter) 68% and 95% confidence levels obtained from SN data alone, while lighter blue contours are the corresponding confidence levels for $H(z)$ data. Solid-line transparent contours are for joint SN+$H(z)$ data. On the right we show one-parameter (marginalized over all other parameters) constraints on $\epsilon$. Note on the left panel that even the SN contours alone close, excluding a flat universe without a cosmological constant by more than 2-$\sigma$. This is a known feature of high-redshift SN data. The statistical significance is dramatically improved by including the $H(z)$ data.

**Figure 1.** Left: Two-parameter constraints on the $\epsilon - \Omega_m$ plane. Dark blue contours correspond to 68% and 95% confidence levels obtained from SN data alone, light blue contours are for $H(z)$ data, and solid line transparent contours are for joint SN+$H(z)$. Right: One-parameter joint constraints on $\epsilon$. The dashed line shows the 95% confidence level, $\Delta \chi^2 = 2$. The dashed (left) and dotted (right) lines show the corresponding constraints using the HST Key project determination of $H_0$ [40] instead of [39].

We now move on to study the implications of these opacity constraints on models of (pseudo-)scalar ALPs coupling to photons.

### 3. Opacity from Axions and Axion-Like-Particles

Scalar or pseudo-scalar particles from physics beyond the standard model, here denoted as $\phi$, may couple to photons through

$$\mathcal{L}_{\text{scalar}} = \frac{1}{4M} F_{\mu\nu} F^{\mu\nu} \phi$$  \hspace{1cm} (10)

and

$$\mathcal{L}_{\text{pseudo-scalar}} = \frac{1}{8M} \epsilon_{\mu\nu\lambda\rho} F^{\mu\nu} F^{\lambda\rho} \phi$$  \hspace{1cm} (11)
where $M$ is the energy scale of the coupling, $F_{\mu\nu}$ the electromagnetic field strength and $\epsilon_{\mu\nu\lambda\rho}$ the Levi-Civita symbol in four dimensions. As described in the introduction, a coupling of the form (11) arises for the axion, introduced by Peccei and Quinn (PQ) to solve the strong CP problem [1], but, more general (pseudo-)scalar ALPs are ubiquitous in fundamental theory and particular string theory [7]. Interestingly, they arise in the constructions used to build models of brane inflation, thus providing a new potential observational window to the physics of the early universe.

In the context of PQ theory, the global shift symmetry is explicitly broken and the axion mass is set by the coupling (11), see equation (2). For generic ALPs, however, the mass is in principle independent of the coupling strength, and can be zero if the related shift symmetry remains intact. Here, we will focus on very light ALPs, for which opacity constraints from $H(z)+$SN data are strongest.

As briefly discussed in the introduction, the mixing terms (10) and (11) imply that ALPs can affect photon propagation. Indeed, if photons traverse a magnetic field there is a finite probability that they will oscillate into ALPs [41]. Notice however that only photons polarized perpendicular (parallel) to the magnetic field mix with scalar (pseudo-scalar) particles. Therefore, the interactions between photons and ALPs in the presence of a magnetic field not only imply that photon number is not conserved, but can also alter the polarization of the light beam. Both effects have been exploited in searches for ALPs both in the laboratory and astronomical observations, see [8] for a recent review.

3.1. Effects of ALPs on SN brightness

ALPs can have an impact on SN observations if photons from SNe are observed through intergalactic magnetic fields. This would give rise to changes in the observed SN luminosities, in a redshift-dependent way.

Indeed, the probability of a suitably polarized photon converting into an ALP after traveling a distance $L$ in a constant, coherent magnetic field is given by [41]:

$$P(L) = \sin^2(2\theta) \sin^2 \left( \frac{\Delta(L)}{\cos 2\theta} \right), \quad (12)$$

where

$$\Delta(L) = \frac{m_{\text{eff}}^2 L}{4\omega}, \quad (13)$$

$$\tan 2\theta = \frac{2B\omega}{Mm_{\text{eff}}^2}. \quad (14)$$

Here, $\omega = 2\pi\nu$ is the photon energy, $B$ is the strength of the magnetic field and $m_{\text{eff}}^2 = |m_\phi^2 - \omega_P^2|$, with $m_\phi$ the mass of the ALP and $\omega_P^2 = 4\pi^2\alpha n_e/m_e$ the plasma frequency of the medium acting as an effective mass for the photons ($\alpha$ is the fine structure constant, $n_e$ the local number density of electrons, and $m_e$ the mass of the electron). In what follows we restrict our attention to very light fields, $m_\phi^2 < \omega_P^2$, for which constraints on cosmic opacity have the most power for constraining the strength of the coupling between ALPs and photons.

The probability (12) depends strongly on the intergalactic magnetic field and its coherence length. Unfortunately, the astrophysical uncertainties for both of these quantities are large. Magnetic fields coherent over the whole Hubble volume are limited, by observations of the CMB and Faraday rotation, to $B \lesssim 10^{-9}$ G [42, 43], but fields with shorter coherence lengths are also constrained. In particular, fields coherent on scales $\sim 50$ Mpc must satisfy $B \lesssim 6 \times 10^{-9}$ G, while fields coherent on scales $\sim$ Mpc are constrained to $B \lesssim 10^{-8}$ G [43]. These values are relevant
because, to explain the origin of observed galactic magnetic fields, intergalactic magnetic fields with coherence lengths $\sim$ Mpc are needed [44].

When considering axion effects on SN brightness, one is interested in the integrated effect of the probability (12) along the line of sight. The intergalactic magnetic field is not expected to be coherent all the way from the earth to supernovae sources and the standard approach for modelling its fluctuations over cosmological scales is to use the cell magnetic field model\(^4\). In this model, one assumes that the magnetic field is split up into a large number of equally sized domains of length $L$ each. Within each domain the magnetic field is taken to be constant, its strength being the same in all domains, while its orientation is chosen randomly. The photon-axion conversion probability in each domain is then given by equation (12) and one must sum over all domains along the line of sight to each given supernova.

In addition, one also needs to know the plasma frequency of the intergalactic medium along the line of sight, a complicated issue given that no measurements of the electron density are available in the large voids of the interstellar medium. A large-scale average can be easily inferred from the total amount of electrons determined by the CMB estimation of the baryon to photon ratio, giving $\omega_p \simeq 1.8 \times 10^{-14}$ eV in the present era, see e.g. [47]. Note, however, that average values up to a factor of 15 smaller may be plausible [48]. The accepted approach is to assume that $\omega_p^2$ is homogeneous given by a fixed value in this range. To check the dependence of our results on this assumption, we will allow a range of a couple of orders of magnitude around the average.

### 3.2. The integrated effect

As we are interested in opacity effects out to redshifts $z \gtrsim O(1)$ we must also take into account the redshift evolution of the environment that causes mixing between photons and (pseudo-)scalars. Assuming the magnetic fields are frozen into the plasma, their strength scales as $B(z) = B_0(1+z)^2$ [44] while $\omega(z) = \omega_0(1+z)$ and $\omega_p^2(z) = \omega_p^2(1+z)^3$ (since it is proportional to the electron density). Here, the subscript 0 indicates values in the present epoch. The physical length of a magnetic domain scales as $L(z) = L_0(1+z)^{-1}$ as long as it is smaller than the Hubble radius. Then, the two quantities appearing in the probability of conversion (12), $\theta$ and $\Delta$, redshift as $\Delta(z) = \Delta_0(1+z)$ and $\tan 2\theta(z) = \tan 2\theta_0$.

Ignoring this redshift evolution, the photon survival probability for a source at redshift $z$ was first calculated in [49]. In [28] we computed this probability including the above redshift evolution. The overall effect can be modelled by a product of $N$ redshift-dependent matrices, which can be diagonalised by a redshift-independent transformation. If the number of domains $N$ is large, this product can be approximated by an integral, yielding for the survival probability $P(z)$ [28]:

$$P(z) = A + (1 - A) \exp \left( -\frac{P}{H_0 L} \frac{H(z) - H_0}{\Omega_m H_0} \right),$$

where we have allowed for an initial flux of axions $I_{\phi}(z_I)$ and defined:

$$P(z) = \frac{I_{\phi}(0)}{I_{\gamma}(z_I)} ; \quad A = \frac{2}{3} \left( 1 + \frac{I_{\phi}(z_I)}{I_{\gamma}(z_I)} \right).$$

Note that with this “averaging” procedure one loses sensitivity on whether the effect comes from scalar or pseudo-scalar particles, i.e. form terms like (10) or (11) respectively.

\(^4\) A more accurate choice for a model of the magnetic field would be to assume a power spectrum of its fluctuations. However, at high frequencies, $\Delta \ll \pi/2$, the cell and power spectrum models give the same results, and at lower frequencies the cell model captures all the qualitative features of ALP-photon interactions, but underestimates the probability of conversion [45, 46]. Therefore, using the cell magnetic field model will give rise to conservative constraints.
Equation (15) is valid for small $P$; in the case where $P$ is of order unity, one should replace $3P/2 \rightarrow -\ln(1 - 3P/2)$.

Also note that the redshift dependence in (15) is absent when $A = 1$. When this is the case, the initial flux of photons and scalars is thermalised, $I_{\gamma}(z_I) = 2I_{\phi}(z_I)$, so that, on average, the effect of photons converting into scalars is compensated by that of scalars converting into photons, and thus there is no propagation effect.

As explained above, the finite probability of conversion gives rise to an apparent change in luminosity distance. Thus, if photons are converted to ALPs along the line of sight, then the inferred and true luminosity distance squared (cf equation (7)) will differ by the (redshift-dependent) factor $P(z)$, so:

$$d_L \rightarrow d_L/\sqrt{P(z)}.$$  \hspace{1cm} (17)

This allows us to search for systematic deviations from the Etherington relation (6) using the methodology of section 2, and interpret them in terms of ALPs, placing constraints on the axion-photon coupling strength $M$ in equations (10) and (11). Therefore, interpreting opacity as coming from ALPs imposes a special parametrization for $\tau(z)$:

$$\tau(z) = -\ln P(z),$$  \hspace{1cm} (18)

where $P(z)$ is given by equation (15) above.

We now proceed to use this result and the methodology of section 2 to examine constraints on the ALP parameters $A$ and $P$, that can in turn be translated into a constraint for the scalar-photon coupling strength $M$.

4. Constraints on Axion-Like-Particles

In Fig. 2 we show our results marginalized over $\Omega_m$ and $H_0$. The dark and light contours represent 68% and 95% joint confidence levels respectively, using the SN data only (left) and joint SN+$H(z)$ data (right). For very small conversion probability $P \lesssim 10^{-5}$ (weak mixing) our constraints become weak because of the lack of photon-ALP mixing. On the other extreme, when the probability is very strong ($P \gtrsim \text{few} \times 10^{-2}$) the photons and axions mix until thermalisation and the redshift dependence of the opacity is lost, so, again, our constraints become weak. Finally we can also observe that a band around $A = 1$ (yellow line) again cannot be constrained. $A = 1$ corresponds to an equilibrated photon-ALP flux at the SNe such that photon$\rightarrow$ALP conversions are compensated by the reverse process, making the photon number redshift-independent.

The first notable feature is that our 2-$\sigma$ bounds (white regions) are significantly improved when including the $H(z)$ data (right panels) compared to using SN data alone (left panels). This improvement is particularly visible in the weak mixing regime.

An interesting feature appears in the upper left plot, corresponding to a redshift-dependent background using SN data only. Here, the 68% C.L. contour shows not one, but three regions where photon-ALP mixing improves the fitting of the data compared to standard $\Lambda$CDM cosmology. The lower-left region corresponds to the parameters invoked by Csaki et al. [50] to explain SN dimming without cosmic acceleration. As we show below, the SN data alone, exclude a flat $\Lambda$CDM model with zero cosmological constant at $> 3$-$\sigma$. Most importantly, the joint SN+$H(z)$ analysis (upper right panel) completely rules out this region.

The preferred region at large $A$ and small $P/L$ also deserves some comments. A value of $A$ greater than unity produces an increase of the SN luminosity with redshift because the SNe would shine more ALPs than photons. In Fig. 1 we have also seen a slight preference for this scenario for the SN data alone, as the 1-$\sigma$ region lies in negative values of $\epsilon$. In the joint SN+$H(z)$ analysis, this trend is softened but, still, small negative values (slight SN brightening) are slightly preferred, even though this is not statistically significant. In Fig. 2 this possibility cannot be excluded, but the statistical preference for this region decreases compared to the rest of the allowed parameter space.
Figure 2. Two parameter constraints in the $A$-$P/L$ plane for general ALPs. Contours represent the 68% (dark) and 95% (light) confidence levels. For the left panels we have used SN data only, while in the right panels we show the joint SN+$H(z)$ analysis. The contours are marginalised over cosmologies and $H_0$.

4.1. Cosmological Constraints

Let us now interpret our results by keeping dependence of our likelihood contours on cosmological parameters, in particular $\Omega_m$. We will finally express our results in terms of a constraint on the axion-photon coupling strength $M$.

ALPs having no other interactions than the two-photon coupling were shown to contribute very little to the SN luminosity [49], corresponding to the case $A \simeq 2/3$. It is evident from Fig. 2 that only a range of conversion probabilities around $P/L \sim O(10^{-3})$ can be excluded. Fig. 3 shows 1- and 2-$\sigma$, two-parameter likelihood contours on the $P/L-\Omega_m$ plane after marginalisation over $H_0$. Note that even the SN constraints alone (dark blue contours) rule out this model as an alternative to a cosmological constant, at greater than 3-$\sigma$ significance. However, there is still a notable degeneracy in the weak mixing regime, and a value for $\Omega_m$ greater than 0.8 is still allowed at the 2-$\sigma$ level, this constraint being slightly weaker when the background redshift dependence is taken into account (left). Including the $H(z)$ data (light blue contours) breaks this degeneracy, yielding strong joint constraints in $P/L-\Omega_m$ (solid line contours).

We can now translate our bounds on $P/L$ into constraints on the strength of the ALP coupling to photons. Since the coupling always appears multiplied by the magnetic field (which is also unknown) we quote bounds on the combination $B/M$. Let us also define appropriately normalised values of the magnetic field strength $B$ and the energy-scale of the the axion-photon coupling $M$, as

$$B_{\text{nG}} = \frac{B}{1\text{nG}} ; \quad M_{10} = \frac{M}{10^{10}\text{GeV}}.$$ (19)

In Fig. 4 we show our constraints for the case $L = 1$ Mpc as a function of the uncertain value of the average electron density, or, equivalently, the plasma frequency$^5$. The exclusion limit is a horizontal band which bends upwards around $n_e \simeq 0.2 \times 10^{-7}$ cm$^{-3}$. In the same figure, we

$^5$ To get rid of the oscillations of equation (12), which not only will be averaged out by energy binning but also by small fluctuations in the sizes of the domains and the values of the plasma frequencies, we used the substitution $\sin^2 x \rightarrow (1 - \exp(-2x^2))/2$. 

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Figure 3. Confidence levels (68% and 95%) on the $P/L - \Omega_m$ plane for the simplest axion-like-particle model ($A = 2/3$). The small and large $P$ regions correspond to the weak and strong mixing regimes respectively. Dark blue contours show constraints from SN data only, light blue from $H(z)$ data, and solid line contours from joint SN+$H(z)$. In the left panel, the redshift dependence of the background is taken into account, while in the right panel these effects are ignored.

have also reproduced the constraints of Mirizzi et al [51] from distortions of the CMB (region above the dashed line) and those of [52] (blue region) from QSO spectra. For $n_e \gtrsim 10^{-9}$cm$^{-3}$ our bounds are stronger than the CMB ones while still competitive with the QSO bounds. Our approach provides a complementary, independent way to obtain these constraints. Each of the three approaches reported in the figure is affected by different, unrelated, systematics: their agreement promotes one’s confidence in these results. Significant improvement in these bounds is expected with future SN and BAO data, see reference [28] for forecasts.

5. Conclusions
If new particles from physics beyond the standard model couple to photons then the propagation of light may be altered. Here we have focused on Axion-Like-Particles coupling to electromagnetism through 2-photon interactions. Measurements of cosmic opacity provide a strong tool to constraining such scenarios, as interactions between photons and (pseudo-)scalar ALPs in the magnetic fields of the intergalactic medium leads to a new source of cosmic opacity.

Indeed, deviations from cosmic transparency can be constrained through their effects on distance duality, by considering possible deviations from the Etherington relation: the luminosity distance is $(1 + z)^2$ times the angular diameter distance if photon number is conserved. Both luminosity distance and angular diameter distance depend on the Hubble parameter $H(z)$. We have thus combined direct measurements of cosmic expansion (from the latest determinations of the Hubble parameter) at redshifts $0 < z < 2$ and recent SN data yielding the luminosity distance. SN-inferred luminosity distances are affected by violation of photon conservation, but the $H(z)$ measurements we use are not. Assuming an underlying flat $\Lambda$CDM model, we have placed tight limits on possible deviations from photon-conservation. Photon-conservation can be violated by simple astrophysical effects which give uniform attenuation such as grey dust and we have reported constraints on this effect.
We have considered the implications of these constraints on more exotic sources of photon number violation involving a coupling of photons to ALPs. Remarkably, this allows us to use optical photons at energies as low as $\sim 1$ eV to constraint the axion-photon coupling strength, a fundamental parameter of the Lagrangian describing physics at much higher energies $\sim 10^{10}$ GeV. Indeed, assuming intergalactic magnetic fields of strength $1$ nG and coherent over lengthscales of order $1$ Mpc, our constraints rule out a region of the axion-photon coupling $M$ between $10^{10}$ and $10^{14}$ GeV (Fig. 4).

It is also possible to search for the effects of new physics on the propagation of photons in the laboratory, for a review see [8]. In regions of parameter space of these models where the transparency of the universe imposes constraints, the bounds from the analysis given here are at least two orders of magnitude better than what can currently be achieved with laboratory experiments. Future measurements of BAOs, and an increase in the number of observations of high redshift supernovae will lead to further improvements of constraints on ALPs.

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Figure 4. Bounds for the product of ALP coupling times magnetic field $B_{nG}/M_{10}$ as a function of the average electron density or corresponding plasma frequency. The grey regions are excluded at 95% C.L. for $L = 1$ Mpc in the redshift-independent (dark) and dependent (light) cases. Also shown are the regions constrained from CMB [51] (above the dashed line), which dominate at low $n_e$, and QSO [52] (blue) spectral distortions.
