hep-ph/0302030
CITA-2003-03

Super-GZK Photons from Photon-Axion Mixing

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ABSTRACT

We show that photons with energies above the GZK cutoff can reach us from very distant sources if they mix with light axions in extragalactic magnetic fields. The effect which enables this is the conversion of photons into axions, which are sufficiently weakly coupled to travel large distances unimpeded. These axions then convert back into high energy photons close to the Earth. We show that photon-axion mixing facilitates the survival of super-GZK photons most efficiently with a photon-axion coupling scale $M \gtrsim 10^{11}$ GeV, which is in the same range as the scale for the photon-axion mixing explanation for the dimming of supernovae without cosmic acceleration. We discuss possible observational consequences of this effect.

1 Introduction

Observations suggest that Ultra High Energy Cosmic Rays (UHECRs) with energies $\sim$ few$\times$ 10^{20} eV may exist in Nature (for reviews, see e.g. [1, 2, 3]). Usually it is assumed that such particles are of astrophysical origin.* While different candidate sources of UHECRs have been proposed, there is a consensus that they should be at distances farther than $\mathcal{O}(100)$ Mpc. Such distances are difficult to reconcile with the observed energies of UHECRs in simple scenarios. Indeed, if the UHECRs are protons, any flux above $\sim 5 \times 10^{19}$ eV – the well-known Greisen-Zatsepin-Kuzmin (GZK) cutoff [7] – is strongly attenuated by the photo-pion production on the Cosmic Microwave Background (CMB) photons. If the UHECRs are photons, their flux is in turn strongly attenuated by the scattering against the extragalactic radio background, seemingly ruling them out as the candidates for the highest energy UHECRs [8].

A way around the GZK cutoff might be a mechanism where the cosmic rays mostly consist of some very weakly interacting particles, which could travel long distances without being significantly absorbed. The only experimentally observed particles which could do this are neutrinos. A few mechanisms with neutrinos as the mediators, traveling most of the distance between a source and us have been suggested. Such proposals, however, suffer from the "converse" problem: how does a sufficient fraction of the neutrinos convert into nucleons or photons near the Earth. It has been argued [9, 10, 11, 12, 13] that, for neutrino masses of order eV, the UHE neutrinos can interact significantly with the cosmological background neutrinos through the Z boson resonance. This mechanism however may require very high incoming fluxes and a very strong local clustering of the background neutrinos to conform with the observations [12].

It is therefore natural to ask if other particles (from beyond the standard model (SM) of particle physics) could be such mediators. Recently, pseudoscalars arising from the spontaneous breaking of an axial symmetry, or axions for short (similar to the QCD-axion [14]), have been considered in an alternative explanation for the apparent dimming of distant type Ia supernovae (SNe) [15]. The idea was that a fraction of the photons emitted by a supernova could convert into axions in the presence of an intergalactic magnetic field en route to the Earth, thereby dimming the supernova and making it look farther away than it is [15]. Since these axions can travel very far without significant interactions, and they can convert to photons (and vice-versa) in background magnetic fields (see, e.g. [16]), they are a natural candidate for a mediator which could help evade the GZK cutoff.

Ref. [17] investigated whether the QCD-axion, which could arise from a solution of the strong CP problem [14], could act as a mediator of UHECRs. Because the photon-axion

*Alternatively, it has been suggested that they may originate from the nearby decay of very massive particles [4, 5] or topological defects [6].

coupling is of the form

$$\frac{a}{4M}\tilde{F}\,F\ , \tag{1}$$

where a is the axion field, and F the electromagnetic field strength, the evolution in the presence of a background magnetic field B induces "flavor" (photon-axion) conversions. The trilinear interaction (1) gives rise to a bilinear term which is not diagonal in "flavor" space. The probability for the photon-axion oscillation in a homogeneous magnetic field is

$$P_{a\to\gamma} \;=\; \frac{1}{1+\frac{\Delta m^4 M^2}{4\,\mathcal{E}^2\,B^2}}\sin^2\left[\frac{y}{2}\sqrt{\frac{B^2}{M^2}+\frac{\Delta m^4}{4\,\mathcal{E}^2}}\right]\ , \tag{2}$$

$$\Delta m^2 \;\equiv\; m_\gamma^2 - m_a^2 \tag{3}$$

where m_γ is the effective photon mass in the intergalactic plasma, m_a is the axion mass, $\mathcal{E}$ the energy of the particles, and y the path length. For a QCD axion, $10^{-5}\,\mathrm{eV} \lesssim m_{a,\mathrm{QCD}} \lesssim 10^{-1}$ eV, so the conversion $a \to \gamma$ is completely negligible [17]. One way around this problem pursued in [17] was to assume that the axion does not convert into photons, but that it interacts with the Earth's atmosphere due to its relatively large coupling to gluons. Here we will follow a different direction. In the photon-axion model for the dimming of the SNe a much lighter axion is invoked [15] (see also [18]). Then the axion mass will not suppress the mixing even for the optical photons contributing to the SNe luminosity, let alone for the UHECRs with energies larger by nearly 20 orders of magnitude. This observation is the starting point of this paper, where we will show that the coupled photon-axion system could indeed give rise to the highest energy cosmic rays. We will also compare the preferred region of the parameter space where the transmission of ultra-high energy photons is maximal to the parameters needed to explain the dimming of the SNe, and show that the UHECR flux due to the photon-axion mixing is especially enhanced for the choice of the parameters needed to account for the SNe dimming in [15]. Since the mean free path l_{GZK} of the photons is very sensitive to their energy, and since the energy of the UHECRs is nearly 20 orders of magnitude higher than the ones of the SNe, this coincidence is very intriguing.

Our model for the transmission of UHECRs works as follows. Initially, a faraway astrophysical source releases a flux of very high energy unpolarized photons. A small fraction of these photons converts into axions before being depleted by the interactions with the background (radio and CMB) photons[†]. At distances $d \gg l_{\mathrm{GZK}}$, defined as the mean free path of the photons in absence of the mixing with axions, the beam is mostly comprised of axions[‡]. These axions will gradually convert back into photons, but most of these photons will be depleted by the interaction with background photons. However, the conversion rate per distance traveled is low, allowing most of the axions to travel unimpeded, and keep

[†]The photons which do interact with the background photons give rise to secondary photons which can also propagate along the beam with a lower energy than the primary ones. To simplify the present discussion we will neglect those and concentrate on particles which have never interacted with the background photons.

[‡]Alternatively, one can start with only axions from the very beginning, as in ref. [17], assuming that they are produced directly at the source.

replenishing the photon beam. This will ensure that a fraction of the photons will survive over distances much larger than one would expect from the GZK length of UHE photons.

The paper is organized as follows: in Section 2 we describe the evolution of the photon-axion system in the presence of the GZK cutoff, making the (unrealistic) assumption of a homogeneous magnetic field of order $\sim 10^{-9}$ Gauss. In this case the evolution of the system can be calculated exactly, and we show that there indeed exists a long-lived component in the photon beam. In Section 3 we examine a more realistic situation, where the magnetic field consists of domains of size $\sim$ Mpc, with a random orientation of the field inside each domain, again with a strength $\sim 10^{-9}$ Gauss [20]. We calculate the surviving intensity of the photon beam as a function of the distance traveled and the parameters of the photon-axion system, and show that the highest survival probabilities are expected for the same parameters that are preferred for the SNe dimming mechanism of [15]. We also discuss possible peculiar signatures of this UHECR mediation mechanism arising from the fact that most of the observed photons would be generated from a predominantly axion beam within the last few magnetic domains around the Earth. We estimate the initial photon flux needed for this mechanism to work, and finally conclude in Section 4.

2 Evolution of the photon-axion system

The evolution of the photon-axion system propagating along the y direction is governed by

$$\left\{\frac{d^2}{dy^2} + \mathcal{E}^2 + \begin{pmatrix} i\Gamma(\mathcal{E}) & -i\mathcal{E}\frac{B}{M} \\ i\mathcal{E}\frac{B}{M} & 0 \end{pmatrix}\right\} \begin{pmatrix} |\gamma\rangle \\ |a\rangle \end{pmatrix} = 0 \tag{4}$$

where we have Fourier-transformed the fields to the energy ($\mathcal{E}$) picture. Here $\vec{B}$ is the extra-galactic magnetic field and $B = \vec{e}\cdot\vec{B} \sim |\vec{B}|$ is its projection on the photon polarization $\vec{e}$. In a constant magnetic field $|\gamma\rangle$ denotes the photon polarization component along $\vec{B}$, $|a\rangle$ the axion, while the photon polarization $|\gamma_\perp\rangle$ orthogonal to $\vec{B}$ decouples from $|a\rangle$ and evolves independently. In the next section we will discuss the evolution of the photon-axion beam through a sequence of domains with a different orientation of the magnetic field in each domain, where both photon polarizations must be retained. In both cases we take $|\vec{B}| \sim \text{few}\cdot 10^{-9}$ G. Throughout our analysis we can ignore the axion mass and the effective photon mass, since they are much smaller* than the beam energy $\mathcal{E} \sim 10^{20} eV$, and $\mathcal{E}\,B/M$. Instead we include the (energy dependent) decay rate Γ, which parameterizes the decrease of the photon intensity due to their interaction with the background photons. If we rewrite $\Gamma \sim \mathcal{E}/l_{\rm GZK}$, in the absence of the coupling to the axions $l_{\rm GZK}$ gives the mean free path of the photons. For $\mathcal{E} = 3\cdot 10^{20} eV$ (which we will consider in our numerical examples later on), $l_{\rm GZK} \simeq 6.4$ Mpc [8].

*We assume that $m_a^2 \ll \mathcal{E}\,B/M$; this ensures that the photon-axion oscillations are unsuppressed, but also excludes the QCD axion.

For a constant B, the solution of Eqs. (4) is

$$\begin{aligned}\begin{pmatrix}|\gamma\rangle\\|a\rangle\end{pmatrix} &= \begin{pmatrix}\frac{1+\sqrt{1-4\delta^2}}{2\sqrt{1-4\,\delta^2}} & -\frac{\delta}{\sqrt{1-4\,\delta^2}}\\ \frac{\delta}{\sqrt{1-4\,\delta^2}} & -\frac{1-\sqrt{1-4\delta^2}}{2\sqrt{1-4\,\delta^2}}\end{pmatrix} \mathrm{e}^{i\,\mathcal{E}\,y+\lambda_1\,y}\begin{pmatrix}|\gamma\rangle\\|a\rangle\end{pmatrix}_0 + \\ &+\begin{pmatrix}-\frac{1-\sqrt{1-4\delta^2}}{2\sqrt{1-4\,\delta^2}} & \frac{\delta}{\sqrt{1-4\,\delta^2}}\\ -\frac{\delta}{\sqrt{1-4\,\delta^2}} & \frac{1+\sqrt{1-4\delta^2}}{2\sqrt{1-4\,\delta^2}}\end{pmatrix} \mathrm{e}^{i\,\mathcal{E}\,y+\lambda_2\,y}\begin{pmatrix}|\gamma\rangle\\|a\rangle\end{pmatrix}_0 ,\end{aligned} \tag{5}$$

where the subscript zero denotes the initial amplitudes at the source at $y=0$, and where we use

$$\begin{aligned}\lambda_1 &\equiv -\frac{1}{4\,l_{\mathrm{GZK}}}\left[1+\sqrt{1-4\,\delta^2}\right] ,\\ \lambda_2 &\equiv -\frac{1}{4\,l_{\mathrm{GZK}}}\left[1-\sqrt{1-4\,\delta^2}\right] ,\\ \delta &\equiv \frac{B\,l_{\mathrm{GZK}}}{M} \simeq 0.11\left(\frac{B}{1\,\mathrm{nG}}\right)\left(\frac{10^{11}\,\mathrm{GeV}}{M}\right)\left(\frac{l_{\mathrm{GZK}}}{\mathrm{Mpc}}\right) .\end{aligned} \tag{6}$$

Notice that (5) is regular for all values of δ. One can see this by rewriting (5) as

$$\begin{aligned}&\begin{pmatrix}|\gamma\rangle\\|a\rangle\end{pmatrix} = \mathrm{e}^{i\mathcal{E}y}\,\mathrm{e}^{-\frac{y}{4l_{\mathrm{GZK}}}}\begin{pmatrix}\mathcal{C}-\mathcal{S} & 2\,\delta\mathcal{S}\\ -2\,\delta\mathcal{S} & \mathcal{C}+\mathcal{S}\end{pmatrix}\begin{pmatrix}|\gamma\rangle\\|a\rangle\end{pmatrix}_0\\ &\mathcal{C}\equiv\cosh\left[\sqrt{1-4\delta^2}\frac{y}{4l_{\mathrm{GZK}}}\right] \;,\; \mathcal{S}\equiv\frac{\sinh\left[\sqrt{1-4\delta^2}\frac{y}{4l_{\mathrm{GZK}}}\right]}{\sqrt{1-4\delta^2}} ,\end{aligned} \tag{7}$$

with $\mathcal{C}$ and $\mathcal{S}$ real numbers for any choice of δ.

We are interested in the mean free path of the photons in the presence of the coupling to the axions. The initial photon beam is taken to be composed of unpolarized photons, that is it contains an equal mixture of $|\gamma\rangle$ and $|\gamma_\perp\rangle$, with the total intensity normalized to unity. Then the intensities of the surviving photons and axions after the distance y away from the source, using eq. (5) I_γ and I_a are

$$\begin{aligned}I_\gamma &= \frac{1}{2}\left\{\mathrm{e}^{-\frac{y}{l_{\mathrm{GZK}}}}+\frac{1}{4\,(1-4\delta^2)}\left[\left(1+\sqrt{1-4\delta^2}\right)\,\mathrm{e}^{-\frac{1+\sqrt{1-4\delta^2}}{4\,l_{\mathrm{GZK}}}y}\right.\right.\\ &\left.\left.-\left(1-\sqrt{1-4\delta^2}\right)\,\mathrm{e}^{-\frac{1-\sqrt{1-4\delta^2}}{4\,l_{\mathrm{GZK}}}y}\right]^2\right\} ,\\ I_a &= \frac{\delta^2}{2\,(1-4\delta^2)}\left[\mathrm{e}^{-\frac{1+\sqrt{1-4\delta^2}}{4\,l_{\mathrm{GZK}}}y}-\mathrm{e}^{-\frac{1-\sqrt{1-4\delta^2}}{4\,l_{\mathrm{GZK}}}y}\right]^2 .\end{aligned} \tag{8}$$

The first term in (8) gives the intensity of the photons orthogonal to $\vec{B}$, which do not mix with the axion and thus simply decay away as is expected from the GZK mechanism, while the remaining terms give the intensities of the coupled axion-photon system. They have an easy interpretation both in the limit of large and small δ. For a large mixing ($\delta \gg 1$),

$$I_\gamma \simeq \frac{1}{2}\,\mathrm{e}^{-\frac{y}{l_{\rm GZK}}} + \frac{1}{2}\,\mathrm{e}^{-\frac{y}{2\,l_{\rm GZK}}}\cos^2\left(\frac{\delta y}{2\,l_{\rm GZK}}\right)$$

$$I_a \simeq \frac{1}{2}\,\mathrm{e}^{-\frac{y}{2\,l_{\rm GZK}}}\sin^2\left(\frac{\delta y}{2\,l_{\rm GZK}}\right) . \qquad (9)$$

This displays the presence of an eigenstate which rapidly oscillates between the original $|\gamma\rangle$ and $|a\rangle$ states. In terms of the original parameters, the oscillation length is $L_{\rm osc} \sim M/B$ and, not surprisingly, it is independent of $l_{\rm GZK}$. Thus, $\delta \sim l_{\rm GZK}/L_{\rm osc}$ gives the number of the photon-axion oscillations within the mean free path $l_{\rm GZK}$. Thus we can view the coupled system as particles which are photon-like half of the time, and axion-like half of the time, so that their probability to interact with the background photons is one-half of the probability of the decoupled photons $|\gamma_\perp\rangle$. This explains why the mean free path of this state is $2l_{\rm GZK}$.[†]

In the case of small mixing, $\delta \ll 1$, eq. (8) reduces to

$$I_\gamma \simeq \frac{1}{2}\,\mathrm{e}^{-\frac{y}{l_{\rm GZK}}} + \frac{1}{2\,(1-4\delta^2)}\left[\left(1-\delta^2\right)\mathrm{e}^{-\frac{1-\delta^2}{2l_{\rm GZK}}y} - \delta^2\mathrm{e}^{-\frac{\delta^2 y}{2\,l_{\rm GZK}}}\right]^2$$

$$I_a \simeq \frac{\delta^2}{2}\left[\mathrm{e}^{-\frac{y}{2\,l_{\rm GZK}}} - \mathrm{e}^{-\frac{\delta^2 y}{2l_{\rm GZK}}}\right]^2 . \qquad (10)$$

This shows the presence of a long-lived mode, characterized by the mean free path $l_{\rm GZK}/\delta^2 \gg l_{\rm GZK}$. To see how it arises, we consider a simple example depicted in Fig. 1, where we plot $I_{\gamma,a}$ as functions of y, taking the parameters: $l_{\rm GZK} = 6.4$Mpc, $B = 1$nG, $M = 4 \cdot 10^{11}$ GeV (this value of the photon-axion coupling is close to its experimental bound [19]) corresponding to $\delta \simeq 0.18$. These are also the values of B and M that have been used in [15]. The photon intensity I_γ initially decreases[‡] as $\exp\left(-y/l_{\rm GZK}\right)$. Meanwhile, some photons convert into axions, and the intensity I_a (initially taken to be zero[§]) increases. For a small mixing, the quantity $\delta \sim l_{\rm GZK}/L_{\rm osc}$ gives the amplitude for the process $|\gamma\rangle \to |a\rangle$ to occur within the mean free path of the photons. Thus, δ^2 is approximately the fraction of the original photons which converts into axions before being depleted by the interaction with the background photons. This is confirmed by Eq. (10) and by the numerical evolution shown, which indicate that the axion intensity peaks at $\sim \delta^2$. From this point on, $I_a > I_\gamma$ and the subsequent photon-axion oscillations lead to a decrease of I_a. The number of axions decreases by a fraction δ^2 per

[†]In the case of a magnetic field randomly oriented along different domains and of a strong mixing δ, we find numerically a mean free path of $3l_{\rm GZK}/2$, showing that the quanta are equally "shared" between the axion and the two photon polarizations states.

[‡]We neglect the secondary photons formed by the scattering of the primaries against the background photons.

[§]This is a conservative assumption, since if initial axions are also produced at the source, then our mechanism will still work, the only difference being that the required sources can have lower intensities.

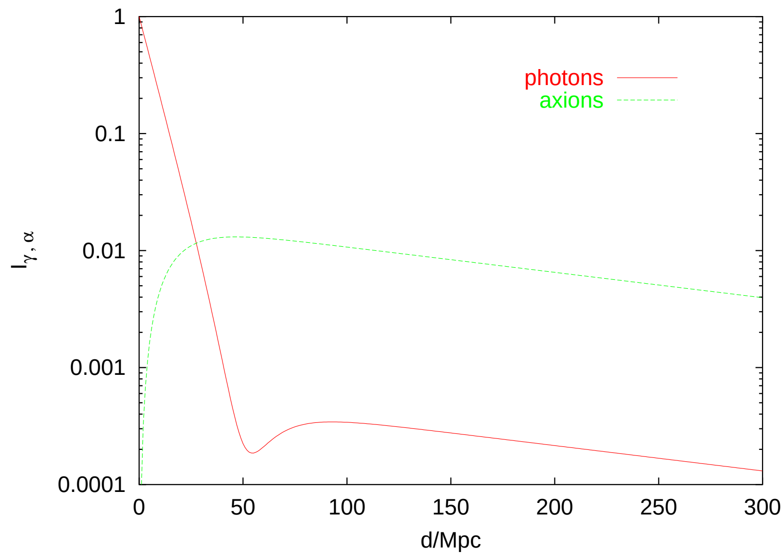

Figure 1: Intensity in photons (initial normalization $I_\gamma = 1$) and axions traveling along a constant magnetic field over a distance of 300 Mpc.

each distance $\Delta y = l_{\rm GZK}$ traveled. To leading order in δ^2, this effect is not compensated by those photons which convert back into axions, since in the case of small mixing the photons in the beam are more likely to be dissipated by the background photons than oscillate back into axions. Thus, $I_a \propto I_{a,max}\,(1-\delta^2)^{y/l_{\rm GZK}} \sim \delta^2 \exp\left(-\delta^2 y/l_{\rm GZK}\right)$.

This argument explains the presence of an eigenstate with the mean free path $l_{\rm GZK}/\delta^2 \gg l_{\rm GZK}$ in Eq. (10). The axion intensity in this mode relative to the total initial photon intensity is small, of order δ^2, due to the small probability for $|\gamma\rangle \to |a\rangle$ oscillations. The photon-like component of this mode is suppressed by a further power of δ^2, because these photons arise from the double conversion $|\gamma\rangle \to |a\rangle \to |\gamma\rangle$. Most of these photons disappear quickly, within a distance $y \sim l_{\rm GZK}$ because of the interaction with background photons. However, the incessant conversion of the surviving axions in the beam back into photons provides for a long-lasting source of photons, yielding a slowly decreasing photon intensity $I_\gamma \sim \delta^4 \exp\left(-\delta^2 y/l_{\rm GZK}\right)$. The rise of this "second population" of photons is clearly visible in Fig. 1. The distance y_* at which these photons reach the level of the surviving intensity of the primary photons in the beam which never underwent any interactions is roughly given by $e^{-y_*/l_{GZK}} \simeq \delta^4$. In the present example, $y_* \simeq 45$ Mpc, in good agreement with the value shown in the figure.

3 Consequences for the UHECR

Until now we have considered the evolution of the photon-axion system in a constant magnetic field. However, the intergalactic magnetic field is unlikely to be completely uniform. Instead, it is more reasonable to take the magnetic fields with a strength close to the observed upper bound of few $\times 10^{-9}$ G to be homogeneous within domains of a typical size $\sim$ Mpc [20], but with their orientation randomly varying from domain to domain (the same assumptions were made in [15]). We will now look at the evolution of the photon-axion beam in such a universe. Photons with polarization perpendicular to $\vec{B}$ do not couple to axions, so the photon-axion mixing involves photons whose helicity changes from domain to domain [16]. If it were not for the interaction with the background photons, the quanta in the beam would be roughly equipartitioned between the axion component and the two photon polarizations after traversing several domains. This is what is expected to happen to the visible photons released in SNe explosions, as explained in [15]. For the range of energies we are interested in, the interaction with the background photons cannot be neglected, and so the depletion effect described in the previous section must be taken into account.

It is easy to extend the formulae from the previous section to the general situation. If the beam is still taken to travel along the y direction, it can be described by a vector in the basis $(|\gamma_x\rangle, |\gamma_z\rangle, |a\rangle)$, where $|\gamma_{\{x,z\}}\rangle$ denote photons polarized along the $\{x, z\}$ axis. Inside each domain, the transfer matrix is such that it can be rotated to the 2×2 matrix (5) for the coupled axion-photon components, plus a third term $\exp(-y/2\, l_{\rm GZK})$ describing the evolution of the decoupled photon polarization orthogonal to $\vec{B}$.

The resulting transfer equation is:

$$\begin{pmatrix} \gamma_x \\ \gamma_z \\ a \end{pmatrix} = \mathrm{e}^{i\,E\,y}\left[T_0\,\mathrm{e}^{\lambda_0\,y} + T_1\,\mathrm{e}^{\lambda_1\,y} + T_2\,\mathrm{e}^{\lambda_2\,y}\right]\begin{pmatrix} \gamma_x \\ \gamma_z \\ a \end{pmatrix}_0 \tag{11}$$

where $\lambda_{1,2}$ have been given above,

$$\lambda_0 \equiv -\frac{1}{2\,l_{\rm GZK}}\ , \tag{12}$$

and the three transfer matrices are

$$T_0 \equiv \begin{pmatrix} \sin^2\theta & -\cos\theta\sin\theta & 0 \\ -\cos\theta\sin\theta & \cos^2\theta & 0 \\ 0 & 0 & 0 \end{pmatrix}\ , \tag{13}$$

$$T_1 \equiv \begin{pmatrix} \frac{1+\sqrt{1-4\,\delta^2}}{2\,\sqrt{1-4\,\delta^2}}\cos^2\theta & \frac{1+\sqrt{1-4\,\delta^2}}{2\,\sqrt{1-4\,\delta^2}}\cos\theta\sin\theta & -\frac{\delta}{\sqrt{1-4\,\delta^2}}\cos\theta \\ \frac{1+\sqrt{1-4\,\delta^2}}{2\,\sqrt{1-4\,\delta^2}}\cos\theta\sin\theta & \frac{1+\sqrt{1-4\,\delta^2}}{2\,\sqrt{1-4\,\delta^2}}\sin^2\theta & -\frac{\delta}{\sqrt{1-4\,\delta^2}}\sin\theta \\ \frac{\delta}{\sqrt{1-4\,\delta^2}}\cos\theta & \frac{\delta}{\sqrt{1-4\,\delta^2}}\sin\theta & -\frac{1-\sqrt{1-4\,\delta^2}}{2\,\sqrt{1-4\,\delta^2}} \end{pmatrix} \tag{14}$$

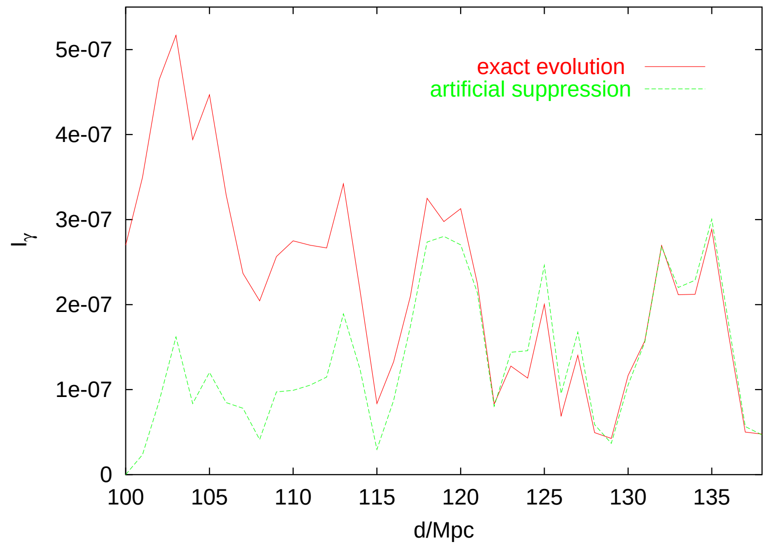

Figure 2: Evolution of the photon intensity across several domains with coherent magnetic field of length 1 Mpc each, for $l_{\rm GZK} = 6.4$ Mpc. The horizontal axis indicates the distance from the source. The exact evolution is compared to the case in which the intensity of photons is (artificially) set to zero at $d = 100$ Mpc. The two evolutions converge within a few $l_{\rm GZK}$ distances.

$$T_2 \quad \equiv \quad \begin{pmatrix} -\frac{1-\sqrt{1-4\,\delta^2}}{2\,\sqrt{1-4\,\delta^2}}\cos^2\theta & -\frac{1-\sqrt{1-4\,\delta^2}}{2\,\sqrt{1-4\,\delta^2}}\cos\theta\sin\theta & \frac{\delta}{\sqrt{1-4\,\delta^2}}\cos\theta \\ -\frac{1-\sqrt{1-4\,\delta^2}}{2\,\sqrt{1-4\,\delta^2}}\cos\theta\sin\theta & -\frac{1-\sqrt{1-4\,\delta^2}}{2\,\sqrt{1-4\,\delta^2}}\sin^2\theta & \frac{\delta}{\sqrt{1-4\,\delta^2}}\sin\theta \\ -\frac{\delta}{\sqrt{1-4\,\delta^2}}\cos\theta & -\frac{\delta}{\sqrt{1-4\,\delta^2}}\sin\theta & \frac{1+\sqrt{1-4\,\delta^2}}{2\,\sqrt{1-4\,\delta^2}} \end{pmatrix}\,, \tag{15}$$

with θ being the angle between the magnetic field and the x axis. The overall transfer matrix is then simply the product of the transfer matrices in each domain.

The evolution of the system is qualitatively similar to the one for a constant magnetic field illustrated in Fig. 1. However, the interference effects modify the quantitative results whenever the domain size is comparable to, or smaller than, the absorption length $l_{\rm GZK}$. For example, consider the evolution of the photon intensity illustrated in Fig. 2. The curve denoted "exact evolution" describes I_γ as a function of the distance from the source. All parameters are as in Fig. 1 except that $\vec{B}$ is randomly oriented in each domain. The random orientation of $\vec{B}$ is responsible for the stochastic character of I_γ seen in the figure. At these distances from the source the intensity in axions I_a is much greater than the one in photons, decreasing smoothly as in Fig. 1. The overall intensity $I_\gamma + I_a$ is thus also smoothly decreasing. By comparing the value of I_γ of the two figures (recalling that in both cases the photon intensity is normalized to unity at the source), one notices that I_γ is more suppressed by a randomly pointing $\vec{B}$. This effect might be interpreted as a decrease of the effective mixing parameter δ. As a consequence, fewer photons will convert into axions before being depleted by the interaction with the background photons. However, for the very same reason, the decrease of I_γ at distances $\gg l_{\rm GZK}$ appears to be much milder in the numerical simulations than the one computed for a constant magnetic field.

An interesting consequence for the transmission of UHECRs is that for sufficiently distant sources the only photons in the beam are the ones which have been generated by the axions within the previous $\sim 4-5\, l_{\rm GZK}$ distances. Otherwise they would have been depleted by the interactions with the background photons.* Since within each domain only the polarization along the magnetic field is actually generated, the UHECRs reaching the Earth would be very sensitive to the orientation of $\vec{B}$ in domains within the sphere of radius $\sim 5\, l_{\rm GZK}$ from us. Consequently the incoming photons should have slightly different polarizations along different lines of sight. The magnitude of this effect depends on the relative size of the domains and on the value of $l_{\rm GZK}$. If the beam crosses several domains in the last few $l_{\rm GZK}$ distances, the differences in the polarization may average out.

We also note that the presence of domains of $\vec{B}$ with different orientation along different lines of sight can lead to differences in the intensities of UHECRs from different directions, even if the distribution of sources were homogeneous and isotropic. This is because only the magnetic field perpendicular to the line of travel contributes to the photon-axion mixing in each domain. Thus, we expect lower fluxes from regions in which the magnetic field is mostly oriented in the direction pointing toward the Earth. At present there is not enough experimental data about UHECRs and extragalactic magnetic fields to enable us to test these predictions. However, this is a very distinctive signature that may eventually be tested experimentally. We illustrate these effects in Fig. 2, where we have artificially removed the photons in the beam at a distance $d = 100$ Mpc from the source, and compared the subsequent evolution to the exact one, without this artificial suppression. One can easily see that the two curves approach each other within a few $l_{\rm GZK}$ distances. Since in this case the evolution is stochastic, the curves themselves behave very differently according to different random choices for the orientations of $\vec{B}$. However, in all the different simulations the two curves approach each other very quickly, irrespectively of the particulars of the magnetic configuration.

In Fig. 2 the mild decrease of the photon intensity at large distances is hidden by its stochastic behavior. To see the decrease, one has to average the evolution over several configurations of $\vec{B}$. This averaging procedure is appropriate if we discuss UHECRs independently of their arrival directions, which have crossed different, uncorrelated configurations of the magnetic field. The remaining figures that we show have been obtained using this averaging procedure.

In Fig. 3 we display the evolution of the photon intensity as a function of the distance from the source (vertical axis) and of the mixing parameter δ (horizontal axis). As before $l_{\rm GZK} = 6.4$ Mpc. The photon intensity drops rapidly both at low and high values of δ. We

*We remark that this concerns the *primary* photons, i.e. the ones which have not interacted with the background photons.

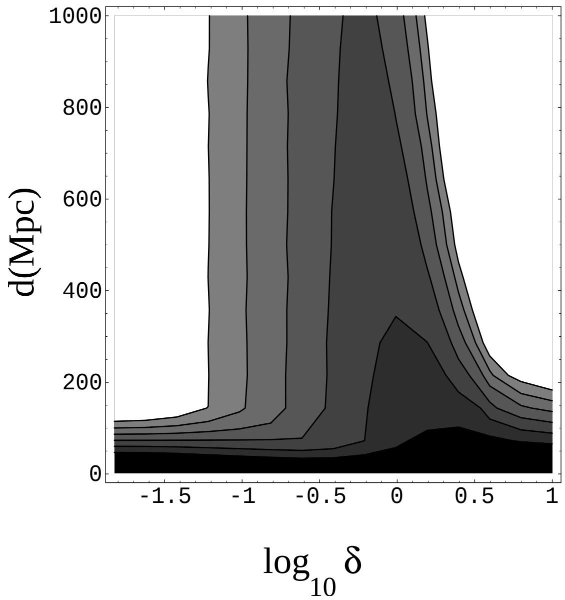

Figure 3: Intensity of photons I_γ (logarithmic units) from a single source at a distance d (given in Mpc on the vertical axis) for an arbitrary value of δ ($\mathrm{Log}_{10}\delta$ shown on the horizontal axis). The intensity at the source is normalized to unity and different contour lines are at $I_\gamma = 10^{-3}\,, 10^{-4}\,, \ldots\,, 10^{-8}\,.$ Darker regions correspond to greater intensity.

note that the surviving photon intensity is maximized for δ between ~ 0.1 and order unity. This gives the region of preferred values of the photon-axion coupling parameter M and the magnetic field B where our mechanism is efficient (see Eq. (6)). Interestingly, the choice of parameters adopted in [15] to explain the SNe dimming in terms of the photon-axion conversion is precisely within this region.

Ultimately we would like to find how the photon-axion mixing affects the spectrum of UHECRs observed on Earth. A precise computation, which requires the inclusion of the secondary particles produced when a photon in the beam interacts with a background photon, is beyond the scope of this paper. However, a rough but illustrative estimate can be obtained relatively simply. We begin by considering an explicit example of the evolution of a primary photon spectrum from a single source at $d = 500$ Mpc from the Earth, shown in Fig 4. The intensity at the source is assumed to scale as E^{-3} (the overall normalization is irrelevant for the present discussion). Even without the mixing with axions the spectrum of primary photons reaching the Earth is strongly modified due to their interaction with the background photons. The reason is that the mean free path l_{GZK} is strongly dependent on their energy. For the energy range shown in Fig. 4, it ranges from ~ 0.3 Mpc for $E = 10^{18}$ eV up to ~ 100 Mpc for $E = 10^{22}$ eV (we have used the results of [8] summarized in [3], see these references for details). The strong increase of l_{GZK} in this interval explains why the second

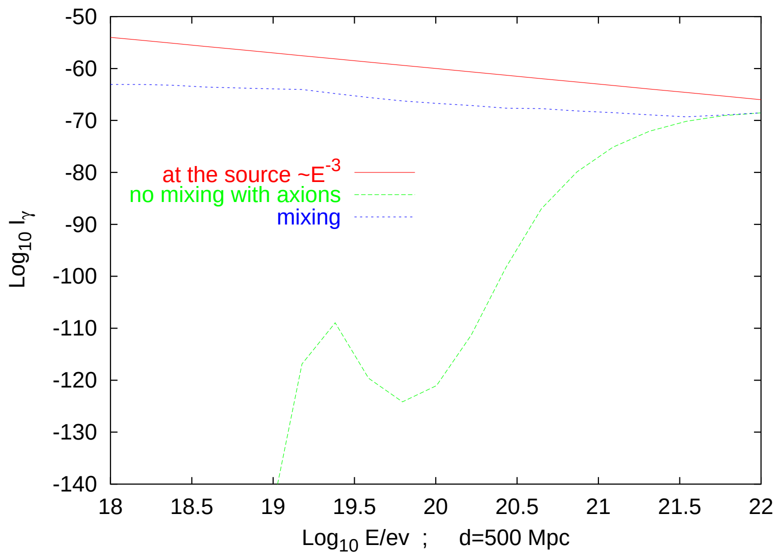

Figure 4: Photons from a source at $d = 500$ Mpc from the Earth. The initial spectrum is assumed to have an energy dependence $I_\gamma \propto E^{-3}$. The spectrum of primary photons reaching the Earth is shown both with and without the photon-axion mixing. See the main text for details.

spectrum in the figure is peaked at high energies, although the original spectrum[†] scaled as E^{-3} .The curves shown in Fig. 4 confirm a much lower depletion of primary photons when the mixing with axions is taken into account (as in all the previous cases, the axion-photon mass parameter is taken to be $M = 4 \cdot 10^{11}$ GeV).

We finally relate this decreased depletion of I_γ to an apparent absence of a GZK cutoff in the UHECR spectrum. In Fig. 5 we combined a spectrum of protons of extragalactic origin with the spectrum of primary photons. Both spectra have been computed for an isotropic and homogeneous distribution of sources at distances $100\,\mathrm{Mpc} < d < 1000\,\mathrm{Mpc}$ from the Earth. For both populations, the intensity at the source is again taken to have a $I \propto E^{-3}$ dependence. Notice that the proton spectrum I_p exhibits the GZK cut-off,[‡] while – due to their coupling with the axions – this is not the case for the photon spectrum I_γ . By an appropriate relative choice of I_p and I_γ at the source, it is thus possible to match the two spectra such that the photon flux "catches up" precisely where the proton flux exhibits the GZK cutoff. The overall spectrum can thus be extended at energies well above E_{GZK} .

Before concluding, we stress the limitations of the computations summarized in the last two figures. The main one is related to the fact that I_γ does not include the secondary

[†]The local maximum at $E \sim 2 \cdot 10^{19}$ eV corresponds to a local maximum of l_{GZK}, due to the fact that at smaller energies the incoming photons are mostly affected by CMB photons, while at higher energies they are most affected by radio photons. See [8, 3] for details.

[‡]For the proton evolution, we have followed the results of [21] summarized in [3].

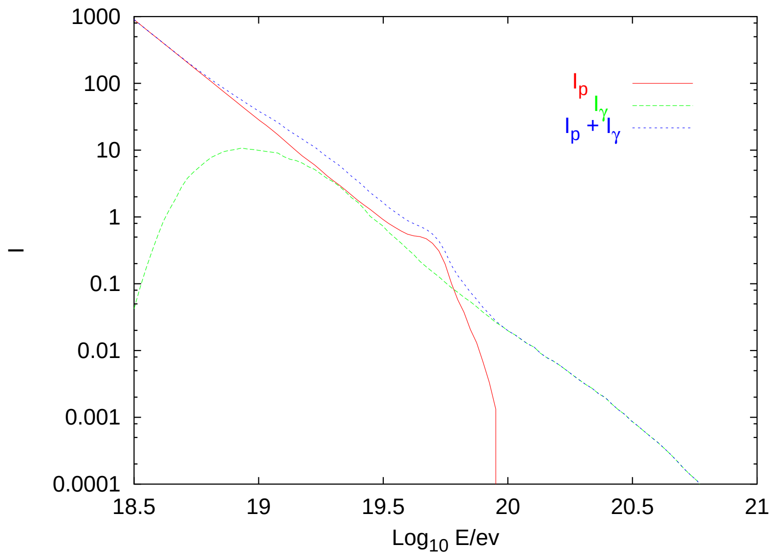

Figure 5: Combined spectra of high energy protons and primary photons, for a homogeneous and isotropic distribution of sources at distances between 100 and 1000 Mpc from the Earth. An energy dependence $I \propto E^{-3}$ is assumed at the source. See the main text for details.

photons arising when the photons in the beam scatter against the background photons. This will lead to a significant modification of I_γ. Fig. 5 has been obtained for a ratio $I_\gamma/I_p = 25,000$ at the source. This would amount to a flux of about 10^{-23} $(m^2\, sr\, s\, GeV)^{-1}$ on Earth, if the photons were travelling unimpeded. We expect that a smaller ratio will be required once secondary photons are taken into account. Secondly, in all of the above computations we have neglected the decrease of the energy (with a consequent change of l_{GZK}) of the quanta in the beam due to the cosmological redshift (for sources at $d = 1000$ Mpc, this correspond to about a 30% decrease of the energy). Thus, the computation summarized in Fig. 5 should be understood only as an order of magnitude estimate and as an illustration of the effect which can be expected. The significance of this estimate will be highly dependent on which sources for the photons one is assuming. The problem of accelerating particles at such energies in different astrophysical environments is completely independent from the problem of their propagation across intergalactic media. Thus it is sensible to separate these problems from each other. Therefore in the present paper we have focused only on the latter, the propagation problem.

We close this section with a note on the nature of the UHECRs. Once UHECRs hit the atmosphere of the Earth, they give rise to extensive air showers. Different incident particles generate distinct longitudinal profiles, that is the amount of particles as a function of the amount of atmosphere penetrated by the cascade. Several analyses, for example the one reported by the Fly's Eye Collaboration in [22], suggest a change from an iron dominated composition of the incoming particles at $\sim 10^{17}$ eV to a proton dominated composition near 10^{19} eV. At higher energies, the current statistics are too poor for a conclusive claim, and one still cannot exclude photons as a relevant fraction of UHECRs above the GZK cut-off [23,24,3]. The most accurate method to determine the nature of UHECR of such energies is the study of the muon content of inclined showers (that is, the ones with incident zenith angle $\theta > 60^0$). The strongest bounds quoted in the literature [3] have been obtained in [23], where inclined air showers recorded by the Haverah Park [25] detector have been analyzed. The results of [23] are affected by uncertainties in the parameterization of the flux of cosmic rays. Assuming the parameterization given in [2], the authors of [23] obtain that less than 48% of the observed events above 10^{19} eV can be photons with a 95% confidence level. For energies higher than 4×10^{19} eV, they find instead that the content of photons has to be less than 50% of the overall flux. Different results are obtained based on the parameterization given in the more recent work [26]. The two upper bounds change to 25% for energies above 10^{19} eV, while only to 70% for energies greater than 4×10^{19} eV [23].

Hence, we see that a mixed composition of proton and photons in the UHECR above the GZK cutoff is compatible with the present experimental limits. Thus the mechanism we propose here is in agreement with the data at the moment: it relies precisely on the assumption that both protons and photons give a sizeable contribution to the UHECR at energies around 10^{19} eV. The present experimental results do not yield clear-cut constraints at higher energies, as manifested by the discrepancy between the analysis based on the different parameterizations of [2] and [26]. On the contrary, the higher statistics and better energy determination expected in the forthcoming Auger Observatory [27] have the power of greatly improving the present bounds. Therefore it may either discover photon UHECRs or yield significantly stronger bounds on the dynamics of photon-axion mixing we employ here. It will thus be of interest to subject the mechanism we propose to more scrutiny.

One more concern about photons being UHECRs could be that most of the photons that do interact with the background (the secondary photons) would cascade down to GeV energies and contribute to the diffuse gamma ray background. If the survival probability of the UHE photons were too low, then the model would have predicted too many secondary photons and thus a gamma ray background that is too large. The average survival probability from Fig. 3 for the optimal values of δ are of order 10^{-4}. The ratio of total energies contained in the 1 GeV gamma ray background to the energies contained in UHECRs at 10^{20} eV can be obtained from comparing the measured EGRET flux [28] to the various measurements of UHECR fluxes. This ratio turns out to be of order $10^{-3} - 10^{-4}$, depending on whether a high or low number for the UHECR flux is adopted. The lower value seems to fit in our model very well, while for the higher value there could seemingly be a marginal conflict (of at most an order of magnitude) with the gamma ray background bounds. However, the ratio of primary to secondary photons can be higher than the survival probabilities in Fig. 3 for various different reasons. First, there could be an initial axion component in the beam which would dramatically change the prediction for the ratio of primary to secondary photons. Note that if the initial beam contained an order unity fraction of the energy in axions, the primary photon to secondary photon ratio would also be of order unity. Thus it is clear that even a relatively small admixture of axions in the initial beam can drastically improve the primary to secondary photon ratio. Second, the conversion rate within our galaxy is much larger since the galactic magnetic field is $\sim 10^3$ times bigger than the assumed value of the intergalactic magnetic field. If a somewhat larger value of magnetic field persists for a few kiloparsecs, then the survival probabilities will be increased by a factor of few, and this on its own could be sufficient to bring down the number of secondary photons. Lastly, as we have stressed above, only a fraction of UHECRs may be photons. Then the total number of secondary photons would fall below the diffuse gamma ray bounds even if the primary to secondary photon ratio stays the same. In this way the signatures of our mechanism would still persist without any violations of the diffuse gamma ray background bounds.

Finally, a photon component of UHECRs may be related to the modification of their spectrum by the interaction with the background photons, once the secondary photons are also taken into account in the propagation codes. After scattering with the background photons, an incoming photon of energy $\mathcal{E}_0$ produces secondary photons which accumulate at energies $< \mathcal{E}_0$. Starting with some flux at the source and normalizing the final spectrum to the observed amount of events above the GZK cutoff, one would typically find an excess of events at lower energies [8]. However here the mixing with the axions can significantly modify the final spectrum precisely in the direction of "balancing out" this excess at lower energies. Indeed, the energy of the photons continuously supplied by the axions is given by the initial energy at the source decreased only by the expansion of the Universe, and it is greater than that of the secondary photons present in the beam. A detailed computation of the final spectrum, including both primary and secondary photons will be an important test of the viability of the mechanism proposed here.

4 Conclusions

An interesting property of axions, which arise in many extensions of the SM, is that they mix with photons in an external magnetic field. Photon-axion conversions in extragalactic magnetic fields can have relevant consequences for astrophysics and cosmology. We have shown that the same parameters (axion mass, coupling scale, magnetic field configurations) which provide an alternative explanation for the dimming of distant SNe [15] could also imply an enhancement of the flux of super-GZK photons arriving to the Earth from faraway sources. Whether this enhancement is observable clearly depends on the distribution of the sources and on the intensity of the high-energy photon spectrum they emit. We have seen here that, for realistic assumptions, super-GZK photons can be detectable, and may provide an explanation for the origin of super-GZK events in the UHECR spectrum, if the required sources of very high energy photons exist. The viability of this mechanism can be tested by improved simulations of the photon propagation, and ultimately by the increase of UHECR data expected over the next few years.

Acknowledgments

We thank J. Arons, J. Feng, H.P. Nilles, and G. Raffelt for useful discussions, and to Dimitry Semikoz for useful comments on the manuscript. C.C. is supported in part by the DOE OJI grant DE-FG02-01ER41206 and in part by the NSF grant PHY-0139738. N.K. is supported in part by a Research Innovation Award from the Research Corporation. J.T. is supported by the U.S. Department of Energy under contract W-7405-ENG-36.

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