AXION DECAY AND ANISOTROPY OF NEAR-IR EXTRAGALACTIC BACKGROUND LIGHT

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ABSTRACT

The extragalactic background light (EBL) is composed of the cumulative radiation from all galaxies and active galactic nuclei over cosmic history. In addition to point sources, the EBL also contains information from diffuse sources of radiation. The angular power spectra of the near-infrared intensities could contain additional signals, and a complete understanding of the nature of the infrared (IR) background is still lacking in the literature. Here we explore the constraints that can be placed on particle decays, especially candidate dark matter (DM) models involving axions that trace DM halos of galaxies. Axions with a mass around a few electronvolts will decay via two photons with wavelengths in the near-IR band and will leave a signature in the IR background power spectrum. Using recent power spectra measurements from the Hubble Space Telescope and the Cosmic Infrared Background Experiment, we find that the 0.6–1.6 $\mu$m power spectra can be explained by axions with masses around 4 eV. The total axion abundance $\Omega_a \simeq 0.05$, and it is comparable to the baryon density of the universe. The suggested mean axion mass and abundance are not ruled out by existing cosmological observations. Interestingly, the axion model with a mass distribution is preferred by the data, which cannot be explained by the standard quantum chromodynamics theory and needs further discussion.

Key words: cosmology: theory – diffuse radiation – large-scale structure of universe

1. INTRODUCTION

The extragalactic background light (EBL) from the ultraviolet (UV) to far-infrared (far-IR) bands contains the cumulative radiation emitted from all galaxies and active galactic nuclei (AGNs) over cosmic history. There are two peaks on the EBL intensity spectrum located at near-infrared (near-IR) and far-IR bands, which are expected to be contributed by direct emission from stars in galaxies and reemission from dust heated by UV and optical photons, respectively (Baldry & Glazebrook 2003; Fukugita & Peebles 2004; Franceschini et al. 2008; Finke et al. 2010; Stecker et al. 2012; Gong & Cooray 2013; Inoue et al. 2013). At NIR wavelengths, there is a large discrepancy between measurements of absolute EBL intensity and the expectations from models of galaxy formation and evolution, in that the absolute measurements are higher than the galaxy counts by at least an order of magnitude at a wavelength around 1 $\mu$m. This difference is likely due to large systematic uncertainties in the absolute measurements involving the foreground contamination from zodiacal light and to a lesser extent from our Galaxy (Hauser & Dwek 2001).

Instead of absolute intensity, the recent measurements focus on the anisotropies of the background. This is because the foregrounds such as zodiacal light have smooth spatial distributions and have a different correlation function than the fluctuations generated by extragalactic signals. The current measurements of the NIR intensity fluctuations suggest an excess in the clustering signal at scales of a few arcminutes, relative to galaxy clustering (Kashlinsky et al. 2005, 2007, 2012; Thompson et al. 2007; Matsumoto et al. 2011; Cooray et al. 2012a, 2012b; Zemcov et al. 2014). This has allowed new ideas involving diffuse sources, such as intrahalo light (IHL) (Cooray et al. 2012b; Zemcov et al. 2014) or direct-collapse black holes (DCBHs) from very high redshifts (Yue et al. 2013). The existing measurements and model development are such that we still have room to explore additional signals. Here we focus on the possibility involving the decay products of axion particles that can also be considered as a viable candidate for dark matter (DM).

The axion is a hypothetical particle that was proposed to solve the strong color parity problem (Weinberg 1978; Wilczek 1978). It is created during the breaking of the Peccei–Quinn (PQ) symmetry at an energy scale that is yet to be determined experimentally, thus labeled as $f_a$ hereafter (Peccei & Quinn 1977). It is also a viable candidate for cold DM. Axions decay into two photons with a lifetime that depends on the mass $m_a$ and the coupling strength $g_{a\gamma\gamma}$ (Ressell 1991). For $m_a \lesssim 10^{-2}$ eV, the axion is nonthermal, and its lifetime is much longer than the age of the universe. In such a scenario, axions are totally invisible. On the other hand, for $m_a > 10^{-2}$ eV, axions can be produced by thermal mechanisms in the early universe (Turner 1987), and their presence is detectable in the multielectronvolt window (Overduin & Wesson 2004).

In this paper we use axion–photon decay to interpret the existing measurements of intensity power spectra at optical to NIR wavelengths. This in turn will constrain the properties of the electronvolt-mass axions, including the allowed mass range and coupling strength to photons. Although thermal axions are not as cold as nonthermal axions, we assume here that they can and will be captured by the gravitational potential well primarily formed by another cold DM candidate, such as weakly interacting massive particles (WIMPs). These halos are the locations of galaxies and galaxy clusters (Kephart & Weiler 1987). The decay of axions in such DM halos will lead to a signal that is similar to IHL (Cooray et al. 2012b) in terms of spatial structure. The differences will be in the redshift evolution of the signal. The fluctuation signal is such that it
may explain the excess clustering that is seen in the NIR anisotropy power spectrum.

To compare our model predictions with the data and to constrain the overall model related to the presence of axions, we make use of the latest intensity power spectra measurements from the Hubble Space Telescope (HST) and the Cosmic Infrared Background Experiment (CIBER). These measurements lead to EBL anisotropies in seven optical and NIR bands (five bands of HST and two bands of CIBER) from 0.6 to 1.65 μm (Zemcov et al. 2014; Mitchell-Wynne et al. 2015). Our overall model includes contributions from axion–photon decay and red-shock (z < 6) and high-redshift (z > 6) galaxies. We assume that the suggested signals such as IHL and DCBHs are zero for the purposes of obtaining the maximally allowed axion abundance. In reality, with such signals present, the actual axion abundance may be lower. Given that we have large uncertainties on the exact level of the IHL and DCBH signals, we do not have a reliable way to constrain the exact abundance here apart from an upper limit. The upper limit, however, as we discuss below, is still lower than the total energy density of DM, suggesting that the axion of 4 eV mass scale alone cannot explain the total cold DM content of the universe. We make use of the Markov chain Monte Carlo (MCMC) method to model-fit the data and to obtain constraints on the axion mass and coupling strength with photons. We assume the flat ΛCDM model with ΩM = 0.27, Ωb = 0.046, and h = 0.71 for the calculations throughout the paper.

This paper is organized as follows. In Section 2, we discuss the details of our calculation to derive the luminosity, mean intensity, and angular power spectra of axion–photon decay; in Section 3, we show power spectrum data from the HST and CIBER observations we use and the details of the model we adopted in the MCMC fitting process; we show the fitting results in Section 4 and give a summary and discussion in Section 5.

2. MEAN INTENSITY AND ANISOTROPIES OF AXION DECAY

In this section, we discuss the method that we use to estimate the mean intensity and anisotropies from the axion–photon decay. First, we need to estimate the luminosity of axion decay in a halo with mass M. The luminosity of a halo contributed by axion–photon decay is given by (Kephart & Weiler 1987)

$$ L_a = \frac{M_a c^2}{\tau_a}, \quad (1) $$

where $M_a$ is the total axion mass in the halo, and we assume the fraction of axion density in a halo is the same as the mean fraction in the universe, which gives

$$ M_a = \Omega_a M. \quad (2) $$

Here M is the total halo mass, $\Omega_a = \rho_a / \rho_{crit,0}$ is the present axion mean density parameter of the universe, and $\rho_a$ and $\rho_{crit,0}$ are the present axion mass density and critical density of the universe. Also, $\Omega_M = \Omega_m + \Omega_\Lambda + \Omega_b$ is the total matter density parameter, and $\Omega_b$ and $\Omega_m$ are the mean density parameters for cold DM (e.g., WIMPs) and baryons, respectively. Since the axion mass range we consider is effectively thermally cold, its energy density evolves as $(1 + z)^3$, which is the same as for cold DM, and Equation (2) is accurate for all redshifts.

Axions with mass greater than $\sim 10^{-2}$ eV are produced thermally and are in thermal equilibrium with the other particle species in the early universe (Turner 1990). These thermal axions decouple when they are still relativistic, and their comoving number density is effectively “frozen in” during the subsequent evolution of the universe (Turner 1990). Thus, the present-day density parameter of thermal axions $\Omega_a$ can be obtained by solving the Boltzmann equation. If the number of relativistic degrees of freedom when axions “froze out” is around 15, $\Omega_a$ is given by (Turner 1990; Overduin & Wesson 2004)

$$ \Omega_a = 5.2 \times 10^{-3} h^{-2} \frac{m_a}{eV}. \quad (3) $$

where $m_a$ is the axion particle mass in eV, and $h = H(100 \text{ km s}^{-1} \text{ Mpc}^{-1})$ where $H_0$ is the Hubble constant today. We can find an upper limit for $m_a$ using Equation (3) and assuming axions contribute all the dark matter density such that $\Omega_a = \Omega_M - \Omega_b$. This gives $m_a \lesssim 22$ eV.

In Equation (1), $\tau_a$ is the axion lifetime for the process of an axion decaying to a photon pair and is given by (Ressell 1991)

$$ \tau_a = (6.8 \times 10^{24}) \left( \frac{m_a}{eV} \right)^5 \zeta^{-2}, \quad (4) $$

where $\zeta = [E/N - 1.95]/0.72$ is the normalized coupling constant. Here E and N are the values of the electromagnetic anomaly and the color anomaly of PQ symmetry, respectively. We have $\zeta = 1$ for the simplest unified models of strong and electroweak interactions (Ressell 1991). The $\zeta$ is proportional to the axion–photon coupling strength $g_{a\gamma\gamma}$, and the relation is

$$ g_{a\gamma\gamma} = 0.72 \frac{\alpha_t}{2\pi f_a}. \quad (5) $$

Here $\alpha_t = 1/137$ is the fine-structure constant, and $f_a = f_{QQ}/N = 6.2 \times 10^6/(m_a/eV)$ GeV where $f_{QQ}$ is the energy scale at which axions are created when U(1) PQ symmetry breaks down (Ressell 1991). In principle, $\zeta$ can be much smaller than one and can even vanish completely in some stable axion models (Kaplan 1985; Ressell 1991; Overduin & Wesson 2004). Here we take $0 < \zeta \leq 1$ as a prior range in our model fits described later. In Equation (4), we find $m_a \approx 27.5$ eV if $\zeta = 1$ and the age of the universe $t_{uni} \approx 4.33 \times 10^{17}$ s as the axion lifetime. Of course, $m_a$ can be much smaller if $\tau_a$ is much larger than $t_{uni}$.

From Equations (1)–(4), the luminosity contributed by the axion–photon decay in a halo with mass M is

$$ L_a(M) = 2.6 \times 10^{-6} \left( \frac{m_a}{eV} \right)^6 \zeta^2 \frac{M}{M_\odot} L_\odot. \quad (6) $$

Hence, approximately, axions with mass around 10 eV with $\zeta \approx 0.1$ can provide a luminosity comparable to the galaxy emission in a halo, leading to $L_a \sim 10^{-2} (M/M_\odot) L_\odot$ or mass-to-light ratios of 100 (Kephart & Weiler 1987). We also note
that, according to Equation (6), $m_a$ and $\zeta$ are not sensitive to the fraction of axions in DM halos. This is because $L_a$ is proportional to the axion fraction in halos from Equations (1) and (2), and we have $L_a \sim m_a^2 \zeta^2$, which indicates that $m_a \sim L_a^{1/6}$ and $\zeta \sim L_a^{1/2}$. With such low dependency, $m_a$ and $\zeta$ are mostly independent of the axion fraction.

Here, instead of a single mass for axions, we also consider a mass distribution. Inspired by string theory (e.g., Svrcek & Witten 2006; Arvanitaki et al. 2010) and experiments of high-energy physics, we assume a general form of axion mass distribution as

$$P(m_a) = \frac{\alpha}{m_0 \Gamma \left(1 + \frac{1}{\alpha}\right)} \left(\frac{m_a}{m_0}\right)^{\alpha} \exp\left[-\left(\frac{m_a}{m_0}\right)^{\alpha}\right],$$

where $m_0$ and $\alpha$ are free parameters, and $\Gamma(x)$ is the Gamma function. Note that this form is phenomenological, and our purpose is to derive the mass distribution from the observational data and minimize theoretical assumptions. We find this form is good enough for the constraints, and more general forms with more free parameters would not change the results. The mean mass of axion particles, under such a distribution, is

$$m_a = \int m_a P(m_a) dm_a.$$  

In this case, we calculate $\Omega_a$ and mean axion lifetime $\tau_a$ by replacing $m_a$ with $m_a$ in Equations (3) and (4).

To estimate the mean intensity of axion–photon decay at different wavelengths, we also need the energy distribution (SED) of the decay photons. Since each photon in the decay photon pair has the same energy $\frac{1}{2}m_a c^2$, the wavelength of each decay photon should be

$$\lambda_a = \frac{h p c}{2 m_a c^2} \approx \frac{2.48}{m_a/\text{eV}} \mu\text{m},$$

where $h_p$ is the Planck constant. This wavelength is around the NIR and optical bands when $1 \lesssim m_a \lesssim 10 \text{ eV}$, which is the mass range of the axions that we will be interested in this paper.

For the case of a single axion mass, following Ressell (1991) and Overduin & Wesson (2004), we use a Gaussian rest-frame SED to consider axions bound in the galaxy cluster halo:

$$F(\lambda) = \frac{1}{\sqrt{2\pi} \sigma_\lambda} \exp\left[-\frac{1}{2} \left(\frac{\lambda - \lambda_a}{\sigma_\lambda}\right)^2\right].$$

where $\sigma_\lambda = 2(v_c/c)\lambda_a$ is the standard deviation, which can be derived from the velocity dispersion of the axions $v_c$ in the halo. Note that $v_c$ is a function of halo mass $M$, and we assume $v_c(M) \sim \sqrt{M}$ according to the general relation of velocity dispersion and halo mass. However, we find the intensity spectrum of axion–photon decay is insensitive to the value of $v_c$ unless $v_c$ is close to $10^4 \text{ km s}^{-1}$. This is because the velocity dispersions for galaxies or galaxy clusters always provide relatively small $\sigma_\lambda$, and this results in narrow SED profiles and makes Equation (10) approach a delta function. For instance, for large galaxy clusters with a mass of $10^{15} \text{ M}_\odot$, we have $v_c = 1300 \text{ km s}^{-1}$, which leads to $\sigma_\lambda \approx 220 \text{ km} / (m_a/\text{eV})$ (Ressell 1991; Overduin & Wesson 2004). In the NIR band

with $\lambda \sim \mathcal{O}(1) \mu\text{m}$, this $\sigma_\lambda$ is about two orders of magnitude smaller than $\lambda$ with $m_a \sim \mathcal{O}(1) \text{ eV}$, and it would be smaller for galaxy clusters with smaller masses. Hence, for simplicity and considering the accuracy and efficiency of our calculation (especially in the MCMC process), we adopt a uniform $v_a$ for all halo masses. Our results with this assumption are identical to that with $v_a(M)$ for the accuracy required in this work.

For the case with an axion mass distribution, we find

$$F(\lambda) = P(\lambda),$$

where $P(\lambda)$ is the normalized probability distribution for the wavelength $\lambda = \lambda_a$, which can be derived from $P(m_a)$ and $\lambda_a(m_a)$ given by Equations (7) and (9). Note that we do not consider the velocity dispersion effect given by Equation (10) for the case involving an axion mass distribution, since the dispersion at each halo mass is negligible compared to the dispersion coming from $P(\lambda)$, which has a much wider wavelength distribution.

Then the luminosity of axion decay at $\lambda$ is given by

$$L_{a,\lambda} = L_a F(\lambda) = L_a(M) F[\lambda_0/(1 + z)].$$

where $\lambda_0$ is the observed wavelength at $z = 0$. Now we can estimate the comoving emissivity of axion–photon decay by

$$\tilde{J}_a(z) = \frac{1}{4\pi} \int dM \frac{dn}{dM} L_{a,\lambda}(M, z).$$

Here $dn/dM(M, z)$ is the halo mass function (Cooray & Sheth 2002). The mean intensity of axion–photon decay at observed wavelength $\lambda_0$ is given by

$$I_\lambda(\lambda_0) = \int_{0}^{\lambda_{\text{max}}} \frac{c}{H(z)(1 + z)^2} \tilde{J}_a(z),$$

where $H(z)$ is the Hubble parameter, and we assume the flat $\Lambda$CDM model with $H(z) = \sqrt{\Omega_{\text{M}}(1 + z)^3 + \Omega_\Lambda}$ and $\Omega_\Lambda = 1 - \Omega_M$. The $\lambda_{\text{max}}$ is the maximum redshift we consider, and we take $\lambda_{\text{max}} = 30$ here. Note that in the case of a single axion mass, $I_\lambda$ at different wavelengths originates from different redshifts since the SED has a Gaussian profile with relatively small $\sigma_\lambda$. Since the structures at different redshift intervals are mainly uncorrelated, this suggests that the cross-correlation between intensity fluctuations at different wavelengths will be zero, contrary to the observations. Our motivation for using a mass distribution is exactly because of this reason: in order to preserve the observed cross-correlations between bands from 0.6 to 1.6 $\mu\text{m}$, we must allow different axion masses to contribute at the same redshift.

The one-halo and two-halo terms of the angular cross power spectrum for axion–photon decay at observed wavelengths $\lambda_0$ and $\lambda_0'$ can be evaluated by

$$C_{\ell,\text{1h}}^{\lambda_0\lambda_0'} = \frac{1}{(4\pi)^2} \int dz \frac{dN}{dz} \left(\frac{a}{\chi}\right)^2 \times \int dM \frac{dn}{dM} \tilde{a}^2(k|M, z)L_{a,\lambda} L_{a,\lambda'},$$

where $\tilde{a}^2(k|M, z)$ is the radial correlation function between axions at $k$ and $z$. Notice that $\tilde{a}^2(k|M, z)$ is a function of $|k|$ and $z$, and $L_{a,\lambda}$ and $L_{a,\lambda'}$ are the luminosity of axion decay at $\lambda$ and $\lambda'$, respectively.
\[
C_{\ell,2h}^{\lambda_0\lambda'_{0}} = \frac{1}{(4\pi)^2} \int \frac{d\chi}{dz} \left( \frac{a}{\chi} \right)^2 P_{\text{lin}}(k, z) \times \int dM \frac{dn}{dM} b(M, z) u(k|z) L_{a,\lambda} \times \int dM \frac{dn}{dM} b(M, z) u(k|z) L_{a,\lambda'}.
\]

Here \(\chi\) is the comoving distance, \(a = 1/(1 + z)\) is the scale factor, \(b(M, z)\) is the halo bias, and \(P_{\text{lin}}(k, z)\) is the linear power spectrum with \(k = \ell/\chi\) (Cooray & Sheth 2002). The term \(u(k|M, z)\) is the Fourier transform of the axion density profile in a halo, and we assume it follows the Navarro–Frenk–White (NFW) profile (Navarro et al. 1997). The total angular cross power spectrum for \(\lambda_0\) and \(\lambda'_{0}\) is then given by

\[
C_{\ell}^{\lambda_0\lambda'_{0}} = C_{\ell,1h}^{\lambda_0\lambda'_{0}} + C_{\ell,2h}^{\lambda_0\lambda'_{0}}.
\]

Then we use \(C_{\ell} = \lambda_0 \lambda'_{0} C_{\ell}^{\lambda_0\lambda'_{0}}\) to fit the data from the observations, as we show in the next section. We notice that the free-streaming effect of axions can suppress the angular power spectrum at small scales (e.g., Hannestad et al. 2005, 2010; Grin et al. 2008), but we find this effect would not change our constraints on \(m_a\) and \(\zeta\) significantly for \(m_a < 5.5\) eV. Hence, we ignore this effect in the data-fitting process for simplicity.

3. MODEL CONSTRAINTS

We primarily make use of two data sets in our fitting process. The first is from the \(HST\) observations in five optical and NIR bands centered at 0.606, 0.775, 0.850, 1.25, and 1.6 \(\mu m\). The data set is obtained by the Wide Field Camera 3 and Advanced Camera for Surveys, and it covers 120 square arcminutes in the Great Observatories Origins Deep Survey (Mitchell-Wynne et al. 2015). The second is from CIBER, a rocket-borne instrument that was designed to measure the spatial and spectral properties of the EBL, with fluctuation measurements in two NIR bands centered at 1.1 and 1.6 \(\mu m\). The CIBER data are from two flights in 2010 and 2012, using two 11 cm telescopes each with a four square degree field of view (Zemcov et al. 2014). We also included fluctuation data from \(Spitzer\) observations at 3.6 \(\mu m\) given by Cooray et al. (2012b), but we found that that cannot fit those measurements around \(\ell = 3 \times 10^3\) with an adequate chi-squared value, especially for an axion model involving a single mass. We give details and a discussion related to \(Spitzer\) model fits in the Appendix.

In order to fit the \(HST\) data, we consider four components: the axion decay model, the shot-noise term, the signal from high-\(z\) faint galaxies during the epoch of reionization, and a power-law component that accounts for Galactic foregrounds, such as diffuse Galactic light (DGL). At small scales, the shot noise dominates the power spectrum. Since it is scale-independent, its angular power spectrum is given by

\[
C_{\ell}^{\text{shot}} = A_{\text{shot}},
\]

where \(A_{\text{shot}}\) is the shot-noise amplitude factor, which is a constant for a given observed wavelength. Note that we do not consider the contribution of low-redshift galaxies to \(HST\) measurements, since it just contributes to the shot-noise term at multipole moments of \(\ell > 2 \times 10^3\) and can be treated as shot noise in these angular scales (Helgason et al. 2012).

For the \(HST\) bands centered at 1.25 and 1.6 \(\mu m\), the signal from high-\(z\) faint galaxies can be an important component of the total signal. The bands of 1.25 and 1.6 \(\mu m\) can cover the high-\(z\) signal from \(8 \lesssim z \lesssim 10.5\) and \(10.5 \lesssim z \lesssim 13\), respectively. We adopt an analytic model given by Cooray et al. (2012a) to estimate the angular power spectrum of the high-\(z\) galaxies. This model can provide the mean intensity and angular power spectra at different wavelengths with a few parameters, such as the star formation rate density, the escape fraction of ionizing photons, and star-formation efficiency. This model is also consistent with the observations of reionization history, optical depth of electron scattering, and UV luminosity functions at high redshifts (Cooray et al. 2012a). However, as claimed by Cooray et al. (2012a), this high-\(z\) model cannot explain the anisotropies of near-IR EBL since the amplitudes of high-\(z\) galaxy power spectra are at least two orders of magnitude smaller than the observations, and we need the other components to match the data. Here we add an amplitude factor in the model, and we have

\[
C_{\ell}^{\text{high}-z} = A_{\text{high}-z} C_{\ell,\text{model}}^{\text{high}-z},
\]

where \(A_{\text{high}-z}\) is the amplitude factor, which can be constrained by the data, and \(C_{\ell,\text{model}}^{\text{high}-z}\) is the high-\(z\) power spectrum from Cooray et al. (2012a). The detection of \(A_{\text{high}-z}\) from the \(HST\) data and the resulting implications in terms of the reionization model are discussed in Mitchell-Wynne et al. (2015).

At large scales, the angular power spectrum is dominated by \(C_{\ell}^{\text{f}}\) for foregrounds. We find it can be described by

\[
C_{\ell}^{\text{f}} = A_{\ell} \ell^{-3},
\]

where \(A_{\ell}\) is the amplitude factor for the foregrounds, which can be fitted by the data. One possible and dominant component in \(C_{\ell}^{\text{f}}\) could be the DGL, since the DGL term is proportional to \(\sim \ell^{-3}\), as shown in Zemcov et al. (2014) and Mitchell-Wynne et al. (2015).

Therefore, the model we use to fit the \(HST\) data at 1.25 and 1.6 \(\mu m\) is

\[
C_{\ell}^{\text{HST}} = C_{\ell}^{\text{axion}} + C_{\ell}^{\text{shot}} + C_{\ell}^{\text{high}-z} + C_{\ell}^{\text{f}}.
\]

For the \(HST\) bands of 0.606, 0.775, and 0.850 \(\mu m\), there is almost no high-\(z\) signal received (Mitchell-Wynne et al. 2015), so we fit the data by

\[
C_{\ell}^{\text{HST}} = C_{\ell}^{\text{axion}} + C_{\ell}^{\text{shot}} + C_{\ell}^{\text{f}}.
\]

For the CIBER data at 1.1 and 1.6 \(\mu m\), following Zemcov et al. (2014), we include the model of low-\(z\) residual galaxies given by Helgason et al. (2012) to calculate the low-\(z\) contribution for the CIBER data. This is due to the relatively shallow depth of the foreground mask in CIBER measurements relative to \(HST\) fluctuations. The shallow depth is such that the clustering of residual galaxies makes an appreciable contribution to CIBER fluctuations. Then, we fit the CIBER data by

\[
C_{\ell}^{\text{CIBER}} = C_{\ell}^{\text{axion}} + C_{\ell}^{\text{high}-z} + C_{\ell}^{\text{low}-z} + C_{\ell}^{\text{f}}.
\]

The \(C_{\ell}^{\text{low}-z}\) model is derived from the NIR luminosity function at different bands. We add a scale factor \(f_{\text{low}-z}\) in the model to vary the low-\(z\) angular power spectrum \(C_{\ell}^{\text{low}-z}\) in 1\(\sigma\) uncertainty, which means that \(f_{\text{low}-z} = 0, 1, \) and \(0.5\) are for 1\(\sigma\) lower, upper, and center values of \(C_{\ell}^{\text{low}-z}\), respectively. Note that the shot-noise term
has already been included in \( C_{\ell}^{\text{low-z}} \) (Zemcov et al. 2014). We use the same \( HST \, C_{\ell}^{\text{high-z}} \) and \( C_{\ell}^{\text{f}} \) at 1.25 and 1.6 \( \mu \text{m} \) to fit the CIBER \( C_{\ell}^{\text{high-z}} \) and \( C_{\ell}^{\text{f}} \) at 1.1 and 1.6 \( \mu \text{m} \), respectively, since these two \( HST \) bands have bandwidths similar to the two CIBER bands. This could help us to fit the two terms, given that the \( HST \) and CIBER data have different scale coverages from \( \ell = 3 \times 10^2 \) to \( 10^6 \). Note that the \( C_{\ell}^{\text{f}} \) is fixed to be the upper limit of DGL power spectra \( C_{\ell}^{\text{DGL}} \) at 1.1 and 1.6 \( \mu \text{m} \), measured in terms of the cross-correlation between CIBER and IRAS in Zemcov et al. (2014). In order to investigate this effect on the constraints of axion properties, we also explore the case with \( C_{\ell}^{\text{f}} \) fixed to be CIBER \( C_{\ell}^{\text{DGL}} \) at 1.1 and 1.6 \( \mu \text{m} \). As we can see in the next section, the fitting results with CIBER \( C_{\ell}^{\text{DGL}} \); the foreground at large angular scales is the maximum allowed given the cross-correlation results between CIBER and IRAS (Zemcov et al. 2014) for 1.1, 1.25, and 1.6 \( \mu \text{m} \). The fitting results for the case of a single axion mass are quite similar.

In the model-fitting process, we employ the MCMC method to constrain the free parameters in the model. The Metropolis–Hastings algorithm is adopted to determine the probability of

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**Figure 1.** The best-fit axion decay models for \( HST \) and CIBER anisotropy power spectra for the case with a mass distribution for axions given by Equation (7). The solid, dashed, and dash-dotted curves are the total, axion, and high-\( z \) power spectra, respectively. The shot noise and \( C_{\ell}^{\text{f}} \) are shown as blue dotted lines for 0.606 and 1.6 \( \mu \text{m} \) as an example (we do not show them for the other bands for clarity, but those terms are included in the model fits). The long dashed curves are the low-\( z \) power spectra from low-\( z \) galaxies for the CIBER data. Top: the fitting results with free \( C_{\ell}^{\text{f}} \) for all bands. Bottom: the fitting results with CIBER \( C_{\ell}^{\text{DGL}} \); the foreground at large angular scales is the maximum allowed given the cross-correlation results between CIBER and IRAS (Zemcov et al. 2014) for 1.1, 1.25, and 1.6 \( \mu \text{m} \). The fitting results for the case of a single axion mass are quite similar.
the acceptance of a new chain point (Metropolis et al. 1953; Hastings 1970). The likelihood function can be estimated by \( L \propto \exp(-\chi^2/2) \), and \( \chi^2 \) is given by

\[
\chi^2 = \sum_{i=1}^{N_d} \frac{(C_i^\text{obs} - C_i^\text{th})^2}{\sigma_i^2},
\]

where \( N_d \) is the number of data points, \( C_i^\text{obs} \) and \( C_i^\text{th} \) are the angular power spectra from observational data and theory, respectively, and \( \sigma_i \) is the error of the data at \( \ell \). In this work, we have the total \( \chi^2 = \chi^2_HST + \chi^2_{\text{CIBER}} \).

We assume the free parameters follow the flat prior probability distribution. For the case with a single axion mass,
and $C^\ell_{\text{CIBER}}$ for the case with an axion mass distribution. The shot-noise term and $C^\ell_{(i)}$ are also shown for comparison. We show a subset of the absolute intensity measurements from the literature (Overduin & Wesson 2004). The total intensity of the axion model is at most $1 \text{nW m}^{-2} \text{s}^{-1}$, which is roughly a factor of 10-20 smaller than the total galaxy contribution.

we assume $m_\ell \in (0, 20)$, and we set $m_0 \in (0, 20)$ and $\alpha \in (0, 5)$ for the case with an axion mass distribution. The solid, dashed, and dash-dotted curves are the total, axion, and high-$z$ power spectra, respectively. The shot-noise term and $C^\ell_{(i)}$ component are also shown for 0.606 $\mu$m as blue dotted lines as an example. The long dashed curves are the low-$z$ power spectra for the CIBER data at 1.1 and 1.6 $\mu$m. For 0.606, 0.775, and 0.850 $\mu$m, the anisotropies of axion emission (one-halo term) dominate the angular power spectra $\ell^4 < \ell < 2 \times 10^3$, while the high-$z$ power spectra (one-halo term) contribute significantly to the power spectrum in the 1.25 and 1.6 $\mu$m bands at these scales, especially around $\ell = 10^3$. At the smaller ($\ell > 2 \times 10^3$) and larger scales ($\ell < 10^4$), the angular power spectra are dominated by the shot noise and $C^\ell_{(i)}$, respectively.

For 1.1, 1.25, and 1.6 $\mu$m, we find the model with free $C^\ell_{(i)}$ can fit the HST and CIBER data well with $\chi^2_{\text{min}} = 264.5$ and reduced $\chi^2 = 2.9$. On the other hand, we find $\chi^2_{\text{min}} = 617.3$ and the reduced $\chi^2 = 6.6$ for the case where we set $C^\ell_{(i)} = C^\ell_{\text{DGL}}$, the residual foreground that is the maximum level allowed by the CIBER cross-correlation with IRAS maps, as described in Zemcov et al. (2014). For the case of a single axion mass, the results are similar, and we find $\chi^2_{\text{min}} = 289.2$ and the reduced $\chi^2 = 3.1$ for free $C^\ell_{(i)}$ and $\chi^2_{\text{min}} = 705.0$ and reduced $\chi^2 = 8.0$ for CIBER $C^\ell_{\text{DGL}}$. This indicates that there should be other additional components, besides the measured CIBER $C^\ell_{\text{DGL}}$, that contribute to this part of the power spectrum.

The best-fit values and 1$\sigma$ errors of all free parameters in our models for all cases are shown in Table 1. For the case of a single axion mass, we find that the fitting results of the cases with free $C^\ell_{(i)}$ and CIBER $C^\ell_{\text{DGL}}$ are quite similar to each other. The allowed parameter ranges of the axion model are consistent within 1$\sigma$ C.L., although the $\chi^2_{\text{min}}$ are much different between the two cases.

For the case of an axion mass distribution, we show the $P(m_\ell)$ with the best-fit values of $m_0$ and $\alpha$ in Figure 2. The blue solid and dashed lines are for the cases of free $C^\ell_{(i)}$ and CIBER $C^\ell_{\text{DGL}}$, respectively. We find they have shapes similar to the peaks located around 3.5 eV and wide distributions extending beyond 10 eV. The mean axion mass from these two $P(m_\ell)$ are $\bar{m}_\ell = 4.45$ and 3.90 for the free $C^\ell_{(i)}$ case and the CIBER $C^\ell_{\text{DGL}}$ case, respectively. Also, we derive the probability distribution for the mean axion mass $\bar{m}_\ell$ from the MCMC chains by integrating over the $P(m_\ell)$ calculated at each chain point with $m_0$ and $\alpha$. We show the two-dimensional (2D) contour maps of $\bar{m}_\ell$ versus $\zeta$ for both cases of free $C^\ell_{(i)}$ and CIBER $C^\ell_{\text{DGL}}$ in Figure 3. The 1D marginalized PDFs for $\bar{m}_\ell$ and $\zeta$ are also shown as blue curves. We find that the constraint results for these two cases are different but consistent within the 2$\sigma$ C.L., and the best-fit values and 1$\sigma$ errors are $\bar{m}_\ell = 4.39^{+0.18}_{-0.18}$ and $\zeta = 0.143^{+0.02}_{-0.02}$ for the free $C^\ell_{(i)}$ case, and $\bar{m}_\ell = 3.91^{+0.16}_{-0.12}$ and $\zeta = 0.293^{+0.03}_{-0.03}$ for the CIBER $C^\ell_{\text{DGL}}$. These values are also in a good agreement with the values of $\bar{m}_\ell$ derived from the $P(m_\ell)$ shown in Figure 2, which are obtained with the best-fit values of $m_0$ and $\alpha$.

Also, we find our fitting results for $A_{\text{high-}z}$, $f_{\text{low-}z}$, $A_t$, and $C^\ell_{\text{DGL}}$ are similar to the results of Mitchell-Wynne et al. (2015). The axion decay model can explain the EBL fluctuation as equally well as the IHL model by comparing the $\chi^2_{\text{min}}$ with the IHL model, especially for the free $C^\ell_{(i)}$ case (Zemcov et al. 2014; Mitchell-Wynne et al. 2015). Just based on a statistical comparison of the data and model fits, we cannot distinguish between the scenarios involving IHL and axion decays to explain the clustering excess seen in the intensity fluctuations of the IR background.

In Figure 4, by using Equation (14), we calculate the mean intensity spectra of the axion–photon decay with the best-fit values of $m_0$, $\alpha$, $m_\ell$, and $\zeta$ for the axion mass distribution case (blue curves) and single-mass case (red curves). The solid and dashed curves are for the free $C^\ell_{(i)}$ case and the CIBER $C^\ell_{\text{DGL}}$ case, respectively. The shaded regions denote the 1$\sigma$ uncertainties for the corresponding curves. The observed data are also shown for comparison (Overduin & Wesson 2004). As can be seen, the intensity spectra of the free $C^\ell_{(i)}$ and CIBER $C^\ell_{\text{DGL}}$ cases are consistent within 1$\sigma$ C.L., although the fitting

Figure 4. The axion–photon decay intensity spectrum with the best fits of $m_\ell$, $m_0$, $\alpha$, and $\zeta$. The blue and red curves denote the $C^\ell_{(i)}$ for axion mass distribution and single-mass cases, respectively. The solid and dashed curves are for the free $C^\ell_{(i)}$ and the CIBER $C^\ell_{\text{DGL}}$, respectively. The shaded regions are the 1$\sigma$ uncertainties for the relevant curves. For comparison, we show a subset of the absolute intensity measurements from the literature (Overduin & Wesson 2004). The total intensity of the axion model is at most 1 nW m$^{-2}$ s$^{-1}$, which is roughly a factor of 10-20 smaller than the total galaxy contribution.

4. RESULTS

In Figure 1, we show our best-fit results of the HST and CIBER data for the case with an axion mass distribution. The results for the case of a single axion mass are quite similar. The solid, dashed, and dash-dotted curves are the total, axion, and high-$z$ power spectra, respectively. The shot-noise term and $C^\ell_{(i)}$ component are also shown for 0.606 $\mu$m as blue dotted lines as an example. The long dashed curves are the low-$z$ power spectra for the CIBER data at 1.1 and 1.6 $\mu$m. For 0.606, 0.775, and 0.850 $\mu$m, the anisotropies of axion emission (one-halo term) dominate the angular power spectra $\ell^4 < \ell < 2 \times 10^3$, while the high-$z$ power spectra (one-halo term) contribute significantly to the power spectrum in the 1.25 and 1.6 $\mu$m bands at these scales, especially around $\ell = 10^3$. At the smaller ($\ell > 2 \times 10^3$) and larger scales ($\ell < 10^4$), the angular power spectra are dominated by the shot noise and $C^\ell_{(i)}$, respectively.
results of $m_0$ and $\alpha$ are not in very good agreement for these two cases (see Table 1). We find that our best-fit mean intensity spectra are much smaller than the measurements, by two orders of magnitude at least. On one hand, the observed data can be overestimated by underestimating the mean intensity of the foreground contamination from Galactic diffuse light and zodiacal light. On the other hand, it indicates that the axion–photon decay alone cannot provide enough intensity on the total NIR intensity, although it could offer a good interpretation for the excess of the NIR anisotropies. The main contribution of the NIR mean intensity might come from low-redshift ($z < 6$) galaxy emission (e.g., Gong & Cooray 2013).

We also compare our results with the other measurements and estimations in Figure 5. In the left panel, we show the $m_u$ versus $g_{\gamma\gamma}$ diagram with excluded regions from different measurements, such as the constraints from the evolution of horizontal branch stars (HB), the duration of the neutrino pulse of SN 1987a, Big Bang nucleosynthesis (BBN), cosmic microwave background (CMB), mean EBL intensity, DM, X-rays, and optical observations (Grin et al. 2007; Cadamuro & Redondo 2012). We also show the bound from cosmological constraints given by the axion smoothing effects on the power spectra of CMB and large scale structure (LSS) (as a gray vertical line with an arrow) (Hannestad et al. 2010). This bound excludes $m_u \gtrsim 1$ eV along the Kim–Shifman–Vainshtein–Zakharov (KSVZ) model (yellow region). In the right panel, the $m_u$ versus $T_{rh}$ relationship with excluded regions is shown (Grin et al. 2008). Here, $T_{rh}$ is the reheating temperature. The higher blue hatched region is excluded by the constraint $\Omega_a h^2 < 0.135$ assuming all DM is composed of axions. The lower green region shows the constraints from the WMAP1 and SDSS data.

We show our best-fit models as blue and red stars to denote the cases of axion mass distribution and single-mass cases, respectively. The results of the free $C_f^i$ case and CIBER $C_f^{\text{DGL}}$ case are shown by solid and dotted stars, respectively. In comparison, we convert the best fits of $m_u$ (or $\overline{m}_u$) and $\zeta$ to $g_{\gamma\gamma}$ using Equation (5), and we estimate $T_{rh}$ via the relation of $\Omega_a - T_{rh}$ given by Grin et al. (2008), with $\Omega_a$ obtained by Equation (3) using the best fit of $m_u$ or $\overline{m}_u$. As can be seen, the results of the single axion mass case are consistent with all
current measurements. For the case of an axion mass distribution, the results stay safely outside the excluded regions of \( m_a \) versus \( T_{\text{reh}} \), but they have some tensions with the HB stars and optical observations. Fortunately, however, we find the constraint result of the free \( C^{\text{QPL}}_f \) case with \( m_a \lesssim 4 \) eV is still consistent with the other observations (see the filled contours in Figure 3), while the result of the CIBER \( C^{\text{QPL}}_f \) case (blue dotted star) is almost discarded by the evolution of HB stars and LSS observations. Also, we need to note that the reheating temperature \( T_{\text{reh}} \) required by our model is only about 40 MeV, which is available in the low-temperature-reheating scenario (Grin et al. 2008); it is however lower than most of the theoretical expectations that are based on BBN constraints.

5. SUMMARY AND DISCUSSION

In this paper, we discuss the axion–photon decay as a potential origin of the NIR intensity fluctuations. According to the theoretical prediction, the axion particle can decay into two photons with wavelengths in the optical and NIR bands if the axion mass is within \( \sim 1–10 \) eV. This provides a possible solution for the excess of the clustered anisotropies at NIR wavelengths, as seen in measurements from *Spitzer*, CIBER, *Hubble*, and *AKARI* (Kashlinsky et al. 2005, 2007, 2012; Matsumoto et al. 2011; Cooray et al. 2012b; Zemcov et al. 2014; Mitchell-Wynne et al. 2015). We calculate the luminosity of the axion decay at wavelength \( \lambda \) in DM halos with mass \( M \) and assume the single axion mass case and the mass distribution case in the estimation. With the help of the halo model, we estimate the mean intensity spectrum and the angular power spectrum for the axion–photon decay.

In order to compare with the observational measurements on the anisotropy power spectrum, we adopt the MCMC method to fit our axion model with the data from all seven NIR bands of *HST* and CIBER. We find our models for axion decay can fit the data with a statistical accuracy equal to the existing IHL model (Cooray et al. 2012b; Zemcov et al. 2014). The best fits of the axion mass and photon coupling constant are \( m_a = 4.83 \) eV and \( \zeta = 0.02 \) for the single-mass case, and the mean axion mass \( \bar{m}_a = 4.39 \) and \( \zeta = 0.14 \) for the case where we invoke a mass distribution. The single-mass case, however, is ruled out by the lack of cross-correlation of fluctuations between different wavelengths, and the preferred model involves an axion mass distribution following Equation (7) with parameters given by the best-fit model.

We also compare the total mean intensity spectra of the axion decay and find that the axion decay is producing at most 1 nW m\(^{-2}\) sr\(^{-1}\), despite explaining the fluctuation measurements with *HST* and CIBER. Finally, we compare the best-fit values of \( m_a \) (or \( \bar{m}_a \)) and \( \zeta \) with the bounds given by different measurements and estimations, and we find our models basically do not conflict with these bounds, although there are some tensions between the constraints from the axion mass distribution case and the HB stellar evolution and optical observations. Therefore, our model can provide reasonable constraints on the axion properties and can offer a possible explanation for the excess of anisotropies of the NIR EBL.

We should also note that, besides the axion–photon decay, there are other possible models to explain the excess of the anisotropies of the NIR EBL, such as the IHL from diffuse stars in DM halos (Cooray et al. 2012b; Zemcov et al. 2014) and DCBHs formed in the early universe (Yue et al. 2013). Hence, the constraints on the axion mass given in this work should be taken as an upper limit, since we assume that axion decay alone provides the contribution to the EBL power spectra without considering the contributions from the models mentioned above. The axion mass and coupling strength can be smaller if other possible signals are included.

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APPENDIX

Besides the data from *HST* and CIBER discussed in the main paper, we also include the data from *Spitzer* fluctuation measurements at 3.6 \( \mu \)m given by Cooray et al. (2012b) to perform a joint model fit with *HST* and CIBER. We take the model similar to the *HST* case to calculate the angular power spectrum \( C^\text{Spitzer}_f \) at 3.6 \( \mu \)m such that

\[
C^\text{Spitzer}_f = C^\text{axion}_f + C^\text{shot}_f + C^\text{high-\zeta}_f + C^f. \tag{25}
\]

We use the same axion and high-\( \zeta \) faint galaxy model for all three data sets to estimate \( C^\text{axion}_f \) and \( C^\text{high-\zeta}_f \), and we use two parameters \( A^3 \text{hot}_f \) and \( A^3 \text{hot}_f \) to estimate \( C^\text{hot}_f \) and \( C^f_f \) for the Spitzer data, respectively.

In Figure 6, we show the fitting results for all three data sets including *Spitzer* at 3.6 \( \mu \)m for the case involving an axion mass distribution. The gray solid, dashed, and dash-dotted curves are the best-fit results for the total, axion decay, and high-\( \zeta \) galaxy power spectra at 3.6 \( \mu \)m, respectively. The best-fit values and 1\( \sigma \) errors of the free parameters are shown in Table 2.

**Table 2**

| Parameter | free \( C^f_f \{ m_a \} \) |
|-----------|-----------------------------|
| \( m_a \) | \( 4.40_{-0.22}^{+0.21} \) |
| \( \alpha \) | \( 3.10_{-0.66}^{+0.51} \) |
| \( \bar{m}_a \) | \( 4.53_{-0.50}^{+0.29} \) |
| \( \zeta \) | \( 0.09_{-0.10}^{+0.14} \) |
| \( \log_{10}(A_{\text{hot}}) \) | \( 1.62_{-0.16}^{+0.05} \) |
| \( \log_{10}(A_{\text{hot}}) \) | \( 0.47_{-0.04}^{+0.05} \) |
| \( \log_{10}(A_{\text{hot}}) \) | \( 6.19_{-3.50}^{+1.09} \) |

\( \chi^2_\text{min} \)

\( \chi^2_\text{min} (N - \text{dof}) \)

372.6

3.2
Table 2. We find that the fitting results of the free parameters are similar to the results from the HST+CIBER data discussed in the main paper (second column of Table 1). Although the best fits of $m_0$ and $\alpha$ are different for HST+CIBER and HST+CIBER+Spitzer, the derived mean axion mass $\bar{m}_a$ is consistent in the 1 $\sigma$ C.L.

We also find that the reduced $\chi^2$ is 3.2, which is comparable with the result of the HST+CIBER case shown in Table 1. While the model gives overall good fits for the HST and CIBER data, we find issues with the model-fit to Spitzer fluctuations around $\ell = 3 \times 10^3$. The power spectrum of axion decay at 3.6 $\mu$m is relatively small. Hence, we interpret this result to imply that axion decay alone cannot explain all of the NIR intensity fluctuations, and other sources, such as IHL (Cooray et al. 2012b) and DCBHs (Yue et al. 2013), are probably needed. However, since the reduced $\chi^2$ of HST+CIBER+Spitzer is comparable to that for HST+CIBER, the axion model still has the potential to fit the data of all bands. We will try to improve this model and explore the possibilities in the future work.

REFERENCES

Arvanitaki, A., Dimopoulos, S., Dubovsky, S., et al. 2010, PhRvD, 81, 123530
Baldry, I. K., & Glazebrook, K. 2003, ApJ, 593, 258
Cadamuro, D., & Redondo, J. 2012, JCAP, 02, 032
Cooray, A., Gong, Y., Smidt, J., & Santos, M. 2012a, ApJ, 756, 92
Cooray, A., & Sheth, R. 2002, PhR, 372, 1
Cooray, A., Smidt, J., De Bernardis, F., et al. 2012b, Natur, 494, 514
Finke, J. D., Razzaque, S., & Dermer, C. D. 2010, ApJ, 712, 238
Franceschini, A., Rodighiero, G., & Vaccari, M. 2008, A&A, 487, 837
Fukugita, M., & Peebles, P. J. E. 2004, ApJ, 616, 643
Gong, Y., & Chen, X. 2007, PhRvD, 76, 123007
Gong, Y., & Cooray, A. 2013, ApJL, 772, L12

Grin, D., Covone, G., Kneib, J.-P., et al. 2007, PhRvD, 75, 105018
Grin, D., Smith, T. L., & Kamionkowski, M. 2008, PhRvD, 77, 083020
Hannestad, S., Mirizzi, A., & Raffelt, G. G. 2005, JCAP, 07, 002
Hannestad, S., Mirizzi, A., Raffelt, G. G., & Wong, Y. Y. Y. 2010, JCAP, 1008, 001
Hastings, W. K. 1970, Biometrika, 57, 97
Hauser, M. G., & Dwek, E. 2001, ARA&A, 39, 249
Helgason, K., Ricotti, M., & Kashlinsky, A. 2012, ApJ, 752, 113
Inoue, Y., Inoue, S., Kobayashi, M. A. N., et al. 2013, ApJ, 768, 197
Kaplan, D. B. 1985, NuPhB, 260, 215
Kashlinsky, A., Arendt, R. G., Mather, J., & Moseley, S. H. 2005, Natur, 438, 45
Kashlinsky, A., Arendt, R. G., Mather, J., & Moseley, S. H. 2007, ApJL, 654, L5
Kashlinsky, A., Arendt, R. G., Ashby, M. L. N., et al. 2012, ApJ, 753, 63
Kephart, T. W., & Weiler, T. J. 1987, PhRvL, 58, 171
Madsen, J. 1990, PhRvL, 64, 2744
Madsen, J. 1991, PhRvD, 44, 999
Matsumoto, T., Seo, H. J., Jeong, W.-S., et al. 2011, ApJ, 742, 124
Metropolis, N., Rosenbluth, A. W., Rosenbluth, M. N., Teller, A. H., & Teller, E. 1953, JCP, 21, 1087
Mitchell-Wynne, K., Cooray, A., Gong, Y., et al. 2015, NatCo, 6, 7945
Navarro, J. F., Frenk, C. S., & White, S. D. M. 1997, ApJ, 490, 493
Overduin, J. M., & Wesson, P. S. 2004, PhR, 402, 267
Pececi, R., & Quinn, H. 1977, PhRvL, 38, 1440
Ressell, M. T. 1991, PhRvD, 44, 3001
Stecker, F. W., Malkan, M. A., & Scully, S. T. 2012, ApJ, 761, 128
Svrcek, P., & Witten, E. 2006, JHEP, 06, 051
Thompson, R. I., Eisenstein, D., Fan, X., Rieke, M., & Kennicutt, R. C. 2007, ApJ, 666, 658
Turner, M. S. 1987, PhRvL, 59, 2489
Turner, M. S. 1990, PhR, 197, 67
Weinberg, S. 1978, PhRvL, 40, 223
Wilczek, F. 1977, PhRvL, 40, 279
Yue, B., Ferrara, A., Salvaterra, R., Xu, Y., & Chen, X. 2013, MNRAS, 433, 1556
Zemcov, M., Smidt, J., Arai, T., et al. 2014, Sci, 346, 732