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Axionic Hot Dark Matter in the Hadronic Axion Window*

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Abstract

Mixed dark matter scenario can reconcile the COBE data and the observed large scale structure. So far the massive neutrino with a mass of a few eV has been the only discussed candidate for the hot dark matter component. We point out that the hadronic axion in the so-called hadronic axion window, $f_a \sim 10^6$ GeV, is a perfect candidate as hot dark matter within the mixed dark matter scenario. The current limits on the hadronic axion are summarized. The most promising methods to verify the hadronic axion in this window are the resonant absorption of almost-monochromatic solar axions from M1 transition of the thermally excited ^{57}Fe in the Sun, and the observation of the "axion burst" in water Čerenkov detectors from another supernova.

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1 Introduction

The cold dark matter (CDM) dominated universe with scale-invariant primordial density fluctuation has been the standard theory of structure formation. After COBE has found the finite density fluctuation in the cosmic microwave background radiation (CMBR), the standard CDM scenario was found to give too much power on smaller scales. Many modifications to the standard CDM scenario were proposed which solve the discrepancy: by introducing a small Hot Dark Matter (HDM) component [1], by "tilting" the primordial density fluctuation spectrum [2], by assuming a finite cosmological constant [3], or by introducing particles (such as ν_τ) whose decay changes the time of radiation-matter equality [4]. At this point, there is no clear winner among these possibilities.[1]

In this letter, we revisit the mixed dark matter (MDM) scenario from the particle physics point of view. This scenario has attracted strong interests because there has been a natural candidate for the HDM component: massive neutrino(s). A neutrino with a mass of a few eV can naturally contribute to a significant fraction of the current universe. However, it has not been easy to incorporate the HDM together with other neutrino "anomalies," unless all three generation neutrinos (possibly together with a sterile neutrino) are almost degenerate, and their small mass splittings explain various "anomalies." Such a scenario may be viewed as fine-tuned. Especially, the atmospheric neutrino anomaly is quite significant statistically now thanks to the SuperKamiokande experiment, which suggests the mass squared difference of $\Delta m^2 = 10^{-3}$–10^{-2} eV^2 between the muon and tau neutrinos. If we view the situation from the familiar hierarchical fermion mass matrices, it suggests the tau neutrino mass of 0.03 – 0.1 eV, and it appears difficult to accommodate the HDM based on massive neutrinos.

We point out that the hadronic axion [7] can be an alternative motivated candidate for the HDM component in the MDM model. Axion has been proposed as a solution to the strong CP problem in the QCD, and the hadronic axion (or KSVZ axion) is one version which predicts small coupling of the axion to the electron. There has been known a window of $f_a \sim 10^6$ GeV allowed by existent astrophysical and cosmological constraints if the axion coupling to photons is suppressed accidentally. This is referred to as the "hadronic axion window." Our main observation is that this window gives exactly the right mass of $m_a \sim$ a few eV and the number density of the axion appropriate for the HDM component in the MDM scenario.

2 Hadronic Axion

First, let us review the hadronic axion model [7]. The most important feature of the hadronic axion is that it does not have tree-level couplings to the ordinary quarks (u, d,

[1] However, a large "tilt" is difficult to obtain in many inflationary models. τCDM can be tested well by B-factory experiments in the near future [5]. The recent data from high-redshift supernovae prefer ΛCDM [6], but the possible evolution of supernovae needs to be excluded by more systematic comparison between nearby and high-z supernovae.

s, c, b, t) and leptons (e, ν_e, μ, ν_μ, τ, ν_τ). In this framework, we introduce new fermions which have Peccei-Quinn (PQ) charges, while ordinary fermions do not transform under $\mathrm{U(1)_{PQ}}$. Some of those new fermions, which we call PQ fermions hereafter, also have $\mathrm{SU(3)_C}$ quantum numbers. After the PQ symmetry is broken spontaneously, axion a appears as a pseudo-Nambu-Goldstone boson of the PQ symmetry.

The axion a couples to the photon with the operator

$$\mathcal{L}_{a\gamma\gamma} = \frac{1}{8} g_{a\gamma\gamma} a \epsilon^{\mu\nu\rho\sigma} F_{\mu\nu} F_{\rho\sigma} \equiv \frac{\alpha}{16\pi} \frac{C_{a\gamma\gamma}}{f_a} a \epsilon^{\mu\nu\rho\sigma} F_{\mu\nu} F_{\rho\sigma}, \tag{1}$$

where f_a is the axion decay constant. This interaction is induced by the mixing to the light mesons (π^0, η, η', and so on) as well as by the triangle anomaly of the PQ fermions. By using the chiral Lagrangian based on flavor $\mathrm{SU(2)}_L\times$ $\mathrm{SU(2)}_R$, we can estimate the coefficient $C_{a\gamma\gamma}$ as [8]

$$C_{a\gamma\gamma} = \frac{E_{\mathrm{PQ}}}{N} - \frac{2(4+z)}{3(1+z)}, \tag{2}$$

where $z = m_u/m_d$ which is estimated to be 0.56 by the leading order perturbation in quark masses in the chiral Lagrangian. (Hereafter, we use $z = 0.56$ for our estimation, unless we discuss quantities which are sensitive to the uncertainty in z.) In Eq. (2), the first term is from the $\mathrm{U(1)_{em}}$ anomaly of the PQ fermions, while the second term is due to the mixing between axion and light mesons. Simultaneously, we also obtain the formula for the axion mass as

$$m_a = \frac{\sqrt{z}}{1+z} \frac{f_\pi m_\pi}{f_a} \simeq 6.2\ \mathrm{eV} \times (f_a/10^6\ \mathrm{GeV})^{-1}, \tag{3}$$

where $f_\pi \simeq 93$ MeV is the pion decay constant, and m_π is the pion mass.

With this axion-photon-photon coupling, axion decays into two photons with the lifetime

$$\tau_a = \left[\frac{\alpha^2 C_{a\gamma\gamma}^2}{256\pi^3} \frac{m_a^3}{f_a^2} \right]^{-1} \simeq 1.2 \times 10^{12}\ \mathrm{yr} \times C_{a\gamma\gamma}^{-2} (m_a/10\ \mathrm{eV})^{-5}. \tag{4}$$

Notice that the lifetime of the axion is longer than the age of the Universe for $m_a \sim 10$ eV and $C_{a\gamma\gamma} \lesssim 1$, and hence primordial axions are still in the Universe. However, as we will see later, radiative decay of the axion may affect the background UV photons in spite of the long lifetime.

Here, we comment that $C_{a\gamma\gamma}$ is significantly affected by uncertainties in the chiral Lagrangian with which the mixing effect is usually calculated. First of all, the accuracy of the $\mathrm{SU(2)}_L\times$ $\mathrm{SU(2)}_R$ chiral Lagrangian is tested up to about 5 – 10 %. For example, by using the pion decay constant estimated from the leptonic decay width of $\pi^\pm$, $\Gamma(\pi^0 \to \gamma + \gamma)$ is calculated to be 7.73 eV [9], while experimentally, it is measured to be 7.7 ± 0.6 eV [10]. (Even though the center value given in Ref. [10] is in a good agreement, the single best measurement suggests the width to be 7.25 ± 0.23 eV [11], which is about 6 % off from the chiral Lagrangian prediction.) Furthermore, f_π estimated from the process $e^+ + e^- \to \pi^0 + e^+ + e^-$ [12] is about 10 % smaller than the one from the leptonic decay of $\pi^\pm$ [9]. Therefore, we may expect 5 – 10 % error in the calculation of the mixing effect from chiral Lagrangian. Another uncertainty is from the so-called Kaplan–Manohar ambiguity [13]. Within the lowest-order chiral perturbation theory, z is estimated to be 0.56. However, under the $\mathrm{SU}(2)_L\times\ \mathrm{SU}(2)_R$ flavor symmetry, the quark mass matrix $M = \mathrm{diag}(m_u, m_d)$ and its conjugate $(i\sigma^2)M^*(-i\sigma^2) = \mathrm{diag}(m_d^*, m_u^*)$ have the same transformation properties, and hence the following shifts are allowed: $m_u \to m_u' = m_u + \epsilon m_d^*$, $m_d \to m_d' = m_d + \epsilon m_u^*$, where ϵ is an unknown parameter [13]. Since the parameter ϵ is arbitrary, $z = m_u/m_d$ cannot be determined from the meson masses alone.[2] In particular, z much smaller than 0.56 (or even $z = 0$) may be allowed if we take this ambiguity into account [13]. This ambiguity cannot be resolved based on meson masses only, but can be by using the baryon masses to some extent. The uncertainty, however, remains large [14].[3] The mixing contribution to $C_{a\gamma\gamma}$ is affected by this uncertainty in z.

As we will see later, $C_{a\gamma\gamma}$ is constrained to be less than 0.01 – 0.1 from astrophysical arguments for the axion decay constant we are interested in. In general, $C_{a\gamma\gamma} \ll 1$ is possible if we adopt an accidental cancellation. With the lowest order chiral Lagrangian, cancellation occurs when $E_{\mathrm{PQ}}/N = 2(4+z)/3(1+z) \simeq 1.95$, but this estimation may not be so reliable. We believe that a better understanding of the quark masses is necessary to pin down the value of E_{PQ}/N for the accidental cancellation. With the current best knowledge, it is clear that the cancellation is quite possible for models with $E_{\mathrm{PQ}}/N \sim 2$ if we take the effects we discussed above into account. In particular, the possibility of the value obtained in grand-unified theories ($E_{\mathrm{PQ}}/N = 8/3$) may not be excluded.

The axion is also coupled to fermions: $\mathcal{L}_{aff} = g_{aff}\, a\bar{f}i\gamma_5 f$, which can again be estimated by using the chiral Lagrangian. Importantly, the hadronic axion does not couple to ordinary quarks and leptons at the tree level. Therefore, in particular, the axion-electron-electron coupling has an extra loop suppression factor [15]:

$$g_{aee} = \frac{3\alpha^2}{4\pi^2}\frac{m_e}{f_a}\left\{\frac{E_{\mathrm{PQ}}}{N}\ln(f_a/m_e) - \frac{2(4+z)}{3(1+z)}\ln(\Lambda_{\mathrm{QCD}}/m_e)\right\}. \tag{5}$$

On the other hand, mixing effects induce an axion-nucleon-nucleon coupling, even though the axion-quark-quark coupling vanishes at the tree level for a hadronic axion:

$$g_{aNN} = \frac{m_N}{f_a}\left\{(F_{A0} \mp F_{A3})\frac{1}{2(1+z)} + (F_{A0} \pm F_{A3})\frac{z}{2(1+z)}\right\}, \tag{6}$$

where $m_N \simeq 940$ MeV is the nucleon mass, and upper (lower) sign is for neutron (proton). The axial-vector isovector contribution has been quite well understood to be $F_{A3} \simeq -1.25$

[2] In the $\mathrm{SU}(3)_L\times\mathrm{SU}(3)_R$ chiral Lagrangian, the effect is formally higher order in quark masses, and hence $\epsilon \sim m_s/(4\pi f_\pi)$. Still, the ambiguity in z is rather large.

[3] However, if $z = 0$, strong CP problem is solved without introducing an axion. Therefore, we do not consider this possibility in this letter.

from the neutron β-decay. Isoscalar part F_{A0} used to be more ambiguous, since this quantity depends on the flavor-singlet axial-vector matrix element S (with $S \equiv \Delta u + \Delta d + \Delta s$ in Ref. [17]) as $F_{A0} \simeq -0.67S - 0.20$, where the constant piece is determined by the hyperon β-decay. In Ref. [17], however, S was estimated from experimental data including higher order QCD corrections, resulting in $S = 0.27 \pm 0.04$. Even though possible systematic uncertainties are not included in this calculation, we use this result as a reference when we discuss axion-nucleon-nucleon coupling. Because of these interactions, f_a is constrained by the axion emission from SN1987A.

3 Constraints on Hadronic Axion

Next, we summarize the constraints on the hadronic axion. In the later discussion, we will be interested in the case of $f_a \sim 10^6$ GeV so that hadronic axion becomes a good candidate of the HDM. Therefore, in this section, we pay an attention to this case.

Most importantly, the coupling of the hadronic axion to the electron is loop suppressed, as can be seen in Eq. (5). Therefore, the constraint on the axion-electron-electron coupling from the cooling of the red giant [18, 19] can be evaded. One can compare the current best upper limit ($g_{aee} \lesssim 2.5 \times 10^{-13}$ [19]) with Eq. (5), and see that g_{aee} for $f_a \sim 10^6$ GeV is smaller than the bound from the red giant for values of $E_{\rm PQ}/N \lesssim 7$.

A non-trivial constraint comes from the emission of the axion from a supernova. If an axion couples to nucleons strongly, the axion can be produced in the core of the supernova, and the axion emission may affect the cooling process of the supernova. In particular, the Kamiokande group and the IMB group measured the flux and duration time of the neutrino burst emitted from the SN1987A, and their results are consistent with the generally accepted theory of the core collapse. Therefore, they confirmed the idea that most of the energy released in the cooling process is carried off by neutrinos. If axion carries away too much energy from the supernova, it would conflict with those observations. The axion flux from the supernova can be suppressed enough in two parameter regions. If axion-nucleon-nucleon interaction is weak enough, the axion cannot be effectively produced in the core of the supernova. Quantitatively, for $f_a \gtrsim 10^9$ GeV, the axion flux can be small enough not to affect the cooling process [20]. On the contrary, if the axion interacts strongly enough, the mean free path of the axion becomes much shorter than the size of the core, and hence the axions cannot escape from the supernova. In this case, axion is trapped inside the so-called “axion sphere,” and the axion emission is also suppressed. (In this case, axions are emitted only from the surface of the axion sphere; this type of the axion emission is often called “axion burst.”) Quantitatively, for $f_a \lesssim 2 \times 10^6$ GeV (or equivalently, $m_a \gtrsim 3$ eV), the axion luminosity from SN1987A is suppressed enough [20].

For $f_a \lesssim 2 \times 10^6$ GeV suggested from the cooling of supernova, we have another constraint from the detection of axions in water Čerenkov detectors. In this parameter region, axion flux from the axion burst is quite sizable for its detection, even though it does not affect the cooling of SN1987A. If the axion-nucleon-nucleon coupling is strong enough, axions may excite the oxygen nuclei in the water Čerenkov detectors ($^{16}\mathrm{O} + a \to {}^{16}\mathrm{O}^*$), followed by radiative decay(s) of the excited state. If this process had happened, the Kamiokande detector should have observed the photon(s) emitted from the decay of $^{16}\mathrm{O}^*$. Due to the non-observation of this signal, $f_a \lesssim 3 \times 10^5$ GeV is excluded [21].

Another class of constraint is from the axion-photon-photon coupling. Because of this coupling, axion can be produced in Primakoff process in the presence of external electromagnetic field, and it also decays into two photons, which result in constraints on the (model-dependent) axion-photon-photon coupling.

One of the important constraints comes from the cooling of the horizontal-branch (HB) stars. If the axion-photon-photon coupling is too strong, axions are produced in the HB stars through the Primakoff process, and the emission of the axions affects the cooling of the HB stars. Then, the lifetime of the HB stars becomes shorter than the standard prediction, and the number of the HB stars are suppressed. However, the number of the HB stars are in a good agreement with theoretical expectations, and hence we obtain the upper bound on the axion-photon-photon coupling [22]:

$$g_{a\gamma\gamma} \lesssim 6 \times 10^{-11}\ \mathrm{GeV}^{-1}. \tag{7}$$

The important point is that $g_{a\gamma\gamma}$ has two sources: the electroweak anomaly of the PQ fermions and the mixing between the axion and light mesons (see Eqs. (1) and (2)). Furthermore, the mixing effect is usually calculated by using the chiral Lagrangian, and there is some uncertainty as discussed earlier. Therefore, it is difficult to convert the constraint (7) to the constraint on the PQ scale f_a. In fact, due to the model dependence, we only have an upper bound on the coefficient $C_{a\gamma\gamma}$:

$$C_{a\gamma\gamma} \lesssim 0.05 \times (f_a/10^6\ \mathrm{GeV}). \tag{8}$$

Notice that, in principle, any value of f_a can be viable with the cooling of the HB stars, if we adopt an accidental cancelation in $C_{a\gamma\gamma}$.

Another important constraint is from the effects of the radiative decay of the axion on the background UV photons. As noted in Eq. (1), axion is coupled to photons, and it decays into two photons with the lifetime given in Eq. (4). Even though the lifetime is longer than the age of the Universe, some fraction of the axion decays and we may see the emission line.

Constraint from the UV extragalactic light is discussed in Refs. [16, 23, 24]. Since the lifetime of the axion is longer than the age of the Universe, intensity of the photon is proportional to the inverse of the lifetime. Therefore, the intensity becomes smaller as the axion-photon-photon coupling gets weaker, and non-observation of the signal sets an upper bound on $C_{a\gamma\gamma}$. Overduin and Wesson looked for the emitted photon from the axion in the extra galactic light, and no signal of the axion was found. From their observation, they derived the upper bound on $C_{a\gamma\gamma}$ of 0.72 ($m_a = 3.8$ eV) to 0.014 ($m_a = 13.0$ eV) [24].

More stringent constraint may be obtained if we observe the photons emitted from the axions in clusters of galaxies. At the center of a cluster, axions are expected to be

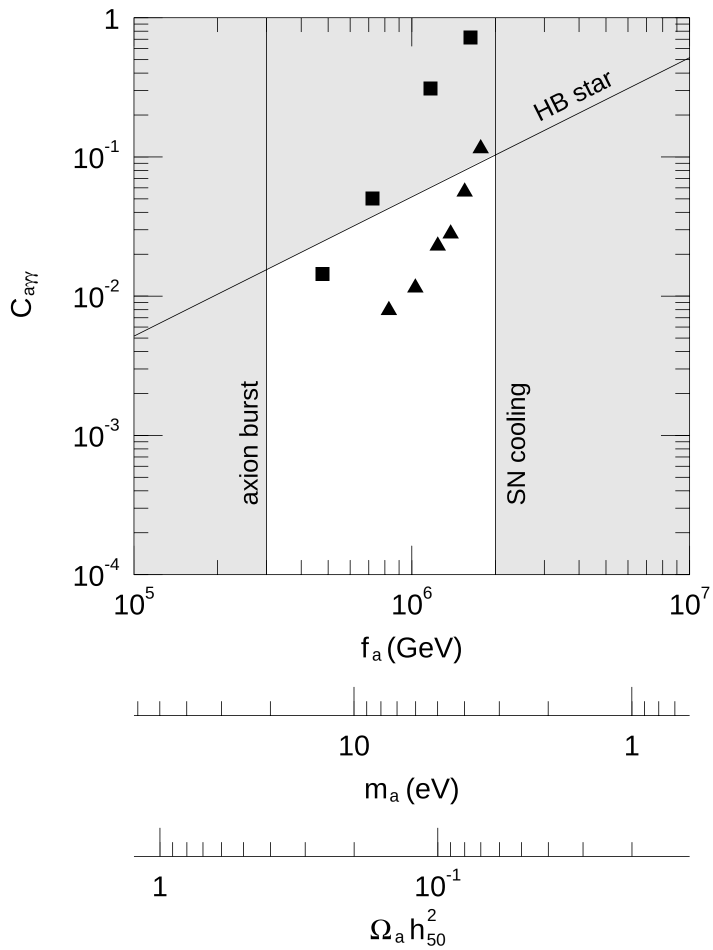

Figure 1: Astrophysical constraints on the hadronic axion model from the cooling of the supernova, axion burst, cooling of the HB stars, extragalactic light [24] (square), and emission line in clusters of galaxies [23] (triangle). Shaded region is excluded, and $C_{a\gamma\gamma}$ larger than squares and triangles are inconsistent with observations for fixed value of f_a.

gravitationally trapped, and its density is more enhanced than the cosmological density. Therefore, the emission lines may be more intense than the one from the extra galactic sources, and the constraint may be more stringent. With three samples of clusters, Ressell obtained the upper bound on $C_{a\gamma\gamma}$ of 0.12 ($m_a = 3.5$ eV) to 0.008 ($m_a = 7.5$ eV) [23], which is about one order of magnitude more stringent than the constraint from extra galactic background light. However, it is possible that the lines of sight of the particular galactic clusters are obscured by absorbing material, resulting in too stringent constraint [24]. If we adopt this argument, this constraint may be evaded.

All the constraints mentioned above are summarized in Fig. 1. As we have discussed, the hadronic axion with the axion decay constant in the following range is still viable with all the astrophysical constraints (if $C_{a\gamma\gamma}$ is small enough):

$$3 \times 10^5 \text{ GeV} \lesssim f_a \lesssim 2 \times 10^6 \text{ GeV} \quad (20 \text{ eV} \gtrsim m_a \gtrsim 3 \text{ eV}). \tag{9}$$

Notice that the constraints based on the axion emission from SN1987A is relatively model-independent. That is, in the hadronic axion model, the axion-nucleon-nucleon coupling is from the mixing between the axion and the light mesons, and hence it is independent of the $\mathrm{U(1)_{PQ}}$ charges of the PQ fermions.[4]

Finally, we comment on the constraint from the cooling of the red giants and the HB stars due to the axion-nucleon-nucleon coupling [25]. The axion-nucleon-nucleon coupling would allow an axion emission from red giants and the HB stars, and cause an additional energy loss rate which is proportional to m_a^2. This extra energy loss changes the brightness of these stars, and it also modifies the relative numbers of the red giants to the HB stars. Observed values of these quantities are in reasonable agreements with theoretical calculations, and hence we can obtain the upper bound on the axion emission rate. The constraint is quite sensitive to the flavor-singlet axial-vector matrix element $S \equiv \Delta u + \Delta d + \Delta s$, since axion-nucleon-nucleon coupling depends on S. For $S = 0.27$ as suggested in Ref. [17], axion mass smaller than about 12 eV is still allowed,[5] and larger axion mass is still viable if we adopt sizable uncertainty in S [25]. Therefore, we concluded that most of the parameter region for the axionic HDM is still alive.

4 Thermal Relic of Hadronic Axion

We have seen in the previous section that the hadronic axion with the decay constant in the window 3×10^5 GeV $\lesssim f_a \lesssim 2 \times 10^6$ GeV is astrophysically allowed as long as the axion-photon-photon coupling is sufficiently small. Now, we are in the position to discuss how the hadronic axion can be a good candidate for HDM. For this purpose, remember that the relevant mass range for the HDM is 1 eV – 10 eV, corresponding to the PQ scale of $f_a \sim 10^6$ GeV (see Eq. (3)), if the axion decouples around the same stage as when the neutrinos do.

For $f_a \sim 10^6$ GeV, the most important source of the primordial axions is the thermal production, rather than the coherent oscillation [16, 15]. Because of the couplings to nucleons (and to pions), axion are thermalized when $T \gtrsim 30 - 50$ MeV for $f_a \sim 10^6$ GeV. In the most recent calculation [15], the axion density is estimated as $[\rho_a/\rho_\nu]_{T\sim 1 \text{ MeV}} \simeq 0.4 - 0.5$, with ρ_a (ρ_ν) being the energy density of the axion (neutrino of one species), or equivalently,

$$\frac{n_a}{s} \simeq 0.02, \tag{10}$$

[4]It does suffer from the uncertainty in z mentioned earlier, however [15].

[5]The authors of Ref. [25] used $F_{A0} = -0.67S - 0.23$ from hyperon β decay without SU(3) breaking effects. A direct measurement, however, suggests $-0.67S - 0.20$ [26], and makes the S in their plot effectively smaller by 0.04.

where n_a is the number density of axion, and s is the total entropy density. (Here, we used $[\rho_a/\rho_\nu]_{T\sim 1\ \mathrm{MeV}} = 0.45$.) Then, the relic density of the axion is given by

$$\Omega_a = \frac{m_a n_a}{\rho_c} \simeq 0.2 \times h_{50}^{-2}(m_a/10\ \mathrm{eV}), \tag{11}$$

where h_{50} is the Hubble constant in units of 50 km/sec/Mpc. Thus, for $m_a \sim 10$ eV, Ω_a can be $0.1 - 0.2$ which is the requirement for the HDM in the MDM scenario. Note that the hadronic axion discussed here is a thermal relic with its mass of ~ 10 eV. Therefore, the axion here is a relativistic particle when the galactic scale crossed the horizon, and behaves as HDM.[6]

Comparing with Eq. (11), the window (9) is exactly where the axion has the right mass and number density to be the HDM component in the MDM scenario.

One may worry about the effect of the hadronic axion on the big-bang nucleosynthesis (BBN). At the time of the BBN, energy density of the axion is sizable ($[\rho_a/\rho_\nu]_{T\sim 1\ \mathrm{MeV}} \simeq 0.4 - 0.5$), and it raises the freeze out temperature of the neutron by speeding up the expansion rate of the Universe. As a result, in our case, more ^{4}He is synthesized than in the standard BBN case [15]. A few years ago, the observed value of the primordial ^{4}He abundance seemed to be unacceptably smaller than the theoretical prediction [29]. If this was true, a hadronic axion with $f_a \sim 10^6$ GeV could be extremely disfavored. However, the current situation is more controversial. Recently, both for D and ^{4}He, several new measurements have been done to determine their primordial abundances, but the results are not consistent with each other; some group reports low D abundance while the other results are much higher, and the same for ^{4}He. In particular, if we adopt a high value of the observed ^{4}He abundance [30], our scenario is consistent with the BBN. Since it is too premature to judge which measurements are reliable, we do not expect any solid argument based on the BBN which rules out the hadronic axion as the HDM component in the MDM scenario.

5 Prospect for Detecting Hadronic Axion

So far, we have seen that the hadronic axion in the current allowed parameter range almost automatically becomes appropriate for HDM. As discussed, this scenario is consistent with all the astrophysical constraints, if the axion-photon-photon coupling is suppressed enough, presumably by an accidental cancellation.

[6]It is interesting to note that the axion decay constant required in this scenario is rather close to the so-called messenger scale in models with gauge mediation of supersymmetry breaking [27], as well as the mass scale of the right-handed neutrino in the sneutrino CDM scenario [28]. It is conceivable that the field S which generates the supersymmetric and supersymmetry-breaking masses of the messengers carry the PQ charge and the messengers are the PQ fermions. The same field S can generate the required size of the right-handed neutrino mass in the sneutrino CDM scenario. The original scale of supersymmetry breaking, however, needs to be raised to make the gravitino heavier than the sneutrino, which can be achieved by making the messenger U(1) coupling constant somewhat small, ~ 0.03.

However, this scenario can be tested in the future in several observations. One possibility is to use the observation of the diffuse background UV photon. Accuracy of the current observation just excluded the axion-photon-photon coupling down to $C_{a\gamma\gamma} \lesssim 0.1 - 0.01$, as we have discussed. However, if the background photon spectrum will be well measured with a better resolution, the emission line from the axion decay may be found in the background photon spectrum. However, as we emphasized, $C_{a\gamma\gamma}$ is a model-dependent parameter. Therefore, a non-observation of the signal cannot exclude the possibility of hadronic axion HDM definitively, because of a possible accidental cancellation in $C_{a\gamma\gamma}$.

Therefore, a detection of hadronic axion which does not rely on axion-photon-photon coupling is strongly favored. One such possibility is to detect an axion burst from a future supernova at SuperKamiokande (or, in general, water Čerenkov detectors). An important point is that newer water Čerenkov detectors (like SuperKamiokande) have much larger fiducial volume than Kamiokande, and hence we can expect a larger event rate. Therefore, a hadronic axion with $f_a \sim 10^6$ GeV can be tested with a future supernova of the size and the distance of SN1987A, even though SN1987A could not exclude this possibility.

Calculation of the event rate suffers from the uncertainties in the axion-nucleon scattering cross section and modeling of supernovae. However, the detection of the signal appears plausible. For example, by rescaling the result given in Ref. [21], we expect a few events at SuperKamiokande for a supernova of the same size as SN1987A for $f_a \sim 10^6$ GeV. Of course, if a new supernova will be closer than SN1987A, we can expect larger number of events, and the hadronic axion HDM can be tested much easier.

Another interesting novel idea is due to Moriyama [31]. In the Sun, thermally excited ^{57}Fe nuclei can decay by emitting axions. Thanks to the Doppler broadening of the axion energy due to the thermal motion of ^{57}Fe, the same nuclide can resonantly absorb the axion. The detection rate was estimated and can be as high as 1 event/day/kg or more. A search was already performed along this line [32] even though they used a small target of 0.03 g to detect 14.4 keV gamma-ray escaping the target rather than the bolometric method suggested. They obtained an upper bound on the axion mass of 745 eV. Another experimental effort to detect solar axions is underway and may reach the axion mass as small as 3 eV in a few years [33].

6 Conclusions

In this letter, we have pointed out that the hadronic axion in the hadronic axion window ($f_a \sim 10^6$ GeV) can automatically be a good candidate of the Hot Dark Matter component in the mixed dark matter scenario. In order to evade an astrophysical constraint from the background UV light, axion-photon-photon coupling has to be suppressed in the hadronic axion window, probably by an accidental cancellation. This scenario may be tested by detecting the axion burst from a future supernova in water Čerenkov detectors, or detecting solar axions using resonant absorption.

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References

[1] See, for example, J.R. Primack, astro-ph/9707285.

[2] R. Davis *et al.*, *Phys. Rev. Lett.* **69**, 1856 (1992); A. Liddle and D. Lyth, *Phys. Lett.* **B291**, 391 (1992); F. Lucchin *et al.*, *Astrophys. J.* **401**, L49 (1992); J. Gelb *et al.*, *ibid.* **403**, L5 (1993); R. Cen *et al.*, *ibid.* **399**, L11 (1992).

[3] J. Ostriker and P. Steinhardt, *Nature* **377**, 600 (1995).

[4] S. Dodelson, G. Gyuk, and M. S. Turner, *Phys. Rev. Lett.* **72**, 3754 (1994); M. White, G. Gelmini, and J. Silk, *Phys. Rev.* **D51**, 2669 (1995); S. Bharadwaj and S. K. Kethi, astro-ph/9707143.

[5] J.L. Feng, T. Moroi, H. Murayama and E. Schnapka, hep-ph/9709411, to appear in *Phys. Rev.* **D**.

[6] S. Perlmutter *et al*, *Nature* **391**, 51 (1998).

[7] J.E. Kim, *Phys. Rev. Lett.* **43**, 103 (1979); M.A. Shifman, A.I. Vainstein, and V.I. Zakharov, *Nucl. Phys.* **B166**, 493 (1980).

[8] See, for example, E.W. Kolb and M.S. Turner, *The Early Universe* (Addison-Wesley, 1990).

[9] W.J. Marciano and A. Sarlin, *Phys. Rev. Lett.* **71**, 3629 (1993).

[10] Particle Data Group, R.M. Barnett *et al.*, *Phys. Rev.* **D54**, 1 (1996).

[11] H.W. Atherton *et al.*, *Phys. Lett.* **B158**, 81 (1985).

[12] H.-J. Behrend *et al.*, *Z. Phys.* **C49**, 401 (1991).

[13] D.B. Kaplan and A.V. Manohar, *Phys. Rev. Lett.* **56**, 2004 (1986).

[14] M.A. Luty and R. Sundrum, *Phys. Rev.* **D52**, 5202 (1995).

[15] S. Chang and K. Choi, *Phys. Lett.* **B316**, 51 (1993).

[16] M.S. Turner, *Phys. Rev. Lett.* **59**, 2489 (1987).

[17] J. Ellis and M. Karliner, hep-ph/9601280.

[18] G.G. Raffelt, *Phys. Rev.* **D33**, 897 (1986); D.S.P. Dearborn, D.N. Schramm, and G. Steigman, *Phys. Rev. Lett.* **56**, 26 (1986); T. Altherr, E. Petitgirard, and T. del Rio Gaztelurrutia, *Astropart. Phys.* **2**, 175 (1994).

[19] G. Raffelt and A. Weiss, *Phys. Rev.* **D51**, 1495 (1995).

[20] M.S. Turner, *Phys. Rev. Lett.* **60**, 1797 (1988); A. Burrows, M.S. Turner, and R.P. Brinkmann, *Phys. Rev.* **D39**, 1020 (1989); A. Burrows, M.T. Ressell, and M.S. Turner, *Phys. Rev.* **D42**, 3297 (1990).

[21] J. Engel, D. Seckel, and A.C. Hayes, *Phys. Rev. Lett.* **65**, 960 (1990).

[22] G.G. Raffelt, *Stars as Laboratories for Fundamental Physics: The Astrophysics of Neutrinos, Axions, and Other Weakly Interacting Particles* (University of Chicago Press, 1996).

[23] M.T. Ressell, *Phys. Rev.* **D44**, 3001 (1991).

[24] J.M. Overduin and P.S. Wesson, *Astrophys. J.* **414**, 449 (1993).

[25] W.C. Haxton and K.Y. Lee, *Phys. Rev. Lett.* **66**, 2557 (1991).

[26] S. Y. Hsueh *et al.*, *Phys. rev.* **D38**, 2056 (1988).

[27] M. Dine, A. Nelson, and Y. Shirman, *Phys. Rev.* **D51**, 1362 (1995); M. Dine, A. Nelson, Y. Nir, and Y. Shirman, *Phys. Rev.* **D53**, 2658 (1996).

[28] L.J. Hall, T. Moroi, and H. Murayama, hep-ph/9712515, to appear in *Phys. Lett.* **B**.

[29] N. Hata, R.J. Scherrer, G. Steigman, D. Thomas, T.P. Walker, S. Bludman, and P. Langacker, *Phys. Rev. Lett.* **75**, 3977 (1995); C.J. Copi, D.N. Schramm, and M.S. Turner, *Phys. Rev. Lett.* **75**, 3981 (1995).

[30] Y.I. Izotov, T.X. Thuan, and V.A. Lipovetsky, *Astrophys. J. Suppl. Series* **108**, 1 (1997).

[31] S. Moriyama, *Phys. Rev. Lett.* **75**, 3222 (1995).

[32] M. Krcmar, Z. Krecak, M. Stipcevic, A. Ljubicic, D. A. Bradley, IRB-ZEF-98/01, nucl-ex/9801005.

[33] S. Moriyama, private communications.