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Axino light dark matter and neutrino masses with R-parity violation

Eung Jin Chun

Korea Institute for Advanced Study
Seoul 130-722, Korea
E-mail: ejchun@kias.re.kr

Hang Bae Kim

BK21 Division of Advanced Research & Education in Physics, Hanyang University
Seoul 133-791, Korea
E-mail: hbkim@hanyang.ac.kr

Abstract: Motivated by the recent observation of the 511 keV γ-ray line emissions from the galactic bulge and an explanation for it by the decays of light dark matter particles, we consider the light axino whose mass can be in the $1-10\,\mathrm{MeV}$ range, particularly, in the context of gauge-mediated supersymmetry breaking models. We discuss the production processes and cosmological constraints for the light axino dark matter. It is shown that the bilinear R-parity violating terms provide an appropriate mixing between the axino and neutrinos so that the light axino decays dominantly to $e^+e^-\nu$. We point out that the same bilinear R-parity violations consistently give both the lifetime of the axino required to explain the observed 511 keV γ-rays and the observed neutrino masses and mixing.

Keywords: Cosmology of Theories beyond the SM, Beyond Standard Model.

Contents

1. Introduction

The Peccei-Quinn (PQ) mechanism to solve the strong CP problem [1], when combined with supersymmetry (SUSY) which is the solution to the gauge hierarchy problem, predicts a singlet fermion called the axino. It can be light in certain supersymmetry breaking mechanism, and become the lightest supersymmetric particle (LSP) providing a good candidate for the particle dark matter (DM) in various mass ranges [2 – 6].

Phenomenologically viable supersymmetric models are implemented with the R-parity to assure the stability of the proton, which also implies the stability of the LSP. However, R-parity is not dictated from any deep theoretical principle. The small violation of R-parity is an attractive option for generating the neutrino masses and mixing [7]. Even with the R-parity violation, the LSP can be cosmologically stable and it may provide an indirect detection mechanism of the DM by leaving imprints in γ-rays from the galactic center and in the diffuse background [8].

Recent observation of 511 keV γ-rays by the SPI spectrometer aboard the INTEGRAL satellite not only confirmed the previously measured total flux but also revealed the morphology of the bulge emission, which is highly symmetric and centered on the galactic center with a full width half maximum of $\sim 8^\circ$ [9 – 11]. The observed emission of 511 keV γ-rays can be well explained by e^+e^- annihilations via positronium formation. But the origin of these galactic positrons remains a mystery. Many astrophysical sources have been suggested, including massive stars, neutron stars, black holes, supernovae, and X-ray binaries. The generic problem of astrophysical sources is that they have difficulty in explaining both the total flux and the high bulge-to-disk ratio of observed 511 keV γ-rays. Given this difficulty, suggested were alternative explanations that light dark matter (LDM) particles annihilating or decaying in the galactic bulge are the sources of the galactic positrons [12 – 18].

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In addition to positrons, annihilations or decays of LDM particles produce direct γ-rays via the internal bremsstrahlung processes. The observation of γ-rays from the galactic center in the energy range $1-100\,\mathrm{MeV}$ bounds the mass of LDM particles to be less than about $20\,\mathrm{MeV}$ [19]. It was also claimed that astrophysical sources are missing for the diffuse γ-ray background in the energy range $1-20\,\mathrm{MeV}$ from the observed spectrum, and that direct γ-rays from annihilations or decays of LDM particles can fit the spectrum when the produced positrons are normalized to fit the 511 keV γ-rays from the galactic bulge [21, 20]. Concerning the annihilating LDM, its mass less than $10\,\mathrm{MeV}$ is practically excluded because it leads to a much longer supernovae cooling time which makes impossible the emission of sufficiently energetic neutrinos observed in SN1987A [22].

In view of above observations, the axino in R-parity violating supersymmetric models is a well-motivated candidate for the MeV dark matter whose decay can explain the observed 511 keV line emission from the galactic bulge as suggested by Hooper and Wang [13]. Indeed, R-parity violation is required to make the axino decay and its lifetime can be very long since its interactions are suppressed by the PQ scale. An interesting question one may ask is whether the same R-parity violation can also generate the observed neutrino masses and mixing.

In this article, we show that the axino LDM scenario is consistent with the usual mechanism of generating the neutrino masses and mixing at tree-level through the small bilinear R-parity violating couplings $\sim 10^{-6}$ [23]. Such small bilinear terms turn out to induce an appropriate axino-neutrino mixing through which the light axinos decay to the positrons with the right range of the lifetime [13];

$$\tau_{\rm dm} \sim \frac{4\times 10^{26}}{m_{\rm dm}({\rm MeV})}\,{\rm sec}\,. \tag{1.1}$$

This has to be contrasted to the case of [13] where the trilinear couplings $\lambda_{i11} \sim 0.1$ were considered.

We also discuss how the MeV axino can arise, particularly, in gauge mediated supersymmetry breaking (GMSB) schemes where the saxion is predicted to get the mass in the range of $4-50\,\mathrm{GeV}$. Axinos are produced thermally or non-thermally in the early universe and the amount of axinos can be correctly adjusted for the appropriate reheat temperature and/or MSSM parameters. If the saxion abundance is comparable to the axino abundance as is the case of the thermal regeneration, the saxion decay to ordinary particles can cause a problem of upsetting the standard prediction of the big bang nucleosynthesis (BBN). Such a "saxion problem" puts another cosmological constraints on the axino LDM models.

2. The axino mass

The axion supermultiplet $A = (s + ia, \tilde{a})$ consists of the pseudo-scalar axion a, its scalar partner, the saxion s, and its fermionic partner, the axino $\tilde{a}$. It has the model-independent interactions with the gluon supermultiplet W_α

$$\mathcal{L}_A^{\rm eff} = \frac{\alpha_{\rm s}}{16\pi f_a} A W_\alpha W^\alpha \Big|_F\,, \tag{2.1}$$

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where f_a is the PQ symmetry breaking scale. At present particle phenomenology, astrophysical and cosmological observations restrict the range of f_a to be 10^9 GeV $\lesssim f_a \lesssim 10^{12}$ GeV. Then the axion mass is given by $m_a \sim \Lambda_{\rm QCD}^2/f_a \sim 10^{-2} - 10^{-5}$ eV.

The axino mass depends crucially on the way of supersymmetry breaking. In generic supergravity (SUGRA) models, it is expected to get the typical soft mass of order $m_{3/2} \sim 100\,\mathrm{GeV}$ and some special arrangement, e.g. no-scale model, is needed to allow the axino mass in the MeV scale [3, 24]. Light axino can arise naturally in GMSB models where SUSY breaking scale is lower than the PQ symmetry breaking scale [4]. Let us show how the MeV axino is predicted in GMSB models. Consider the DFSZ axion model [1] where the MSSM fields are charged under the PQ symmetry. Upon the PQ symmetry breaking, an effective Kähler potential between the axion supermultiplet A and the other fields Φ_i is generated as follows;

$$K_{\rm eff} = e^{x_i \frac{A+A^\dagger}{f_a}} \Phi_i^\dagger \Phi_i \tag{2.2}$$

where x_i is the PQ charge of Φ_i. Taking the terms of order A^2 and $\Phi_i = H_{1,2}$, one has a contribution to the axino mass; $m_{\tilde a} \approx F_{H_i} v/f_a^2 \approx \mu v^2/f_a^2 \ll \mathrm{MeV}$ which is negligible in our context. In GMSB models, Φ_i can be one of the hidden sector superfields, say $\hat X$, which is assumed to take the vacuum expectation value; $\langle \hat X \rangle = X + \theta^2 F_X$ leading to the effective supersymmetry breaking scale, $\Lambda \equiv F_X/X = 10^4 - 10^5\,\mathrm{GeV}$ [25]. Then, one obtains the axino and saxion mass as

$$m_{\tilde a} = x_X^2 \frac{X F_X}{f_a^2} \approx \left(\frac{X}{f_a}\right)^2 \Lambda\,, \tag{2.3}$$

$$m_s^2 = 2 x_X^2 \frac{F_X^2}{f_a^2}\,. \tag{2.4}$$

The axino mass in the range $1 - 10\,\mathrm{MeV}$ is obtained with $X/f_a \sim 10^{-3} - 10^{-4}$. These equations also give us the relation;

$$m_s^2 \approx 2\, m_{\tilde a} \Lambda \tag{2.5}$$

leading to the saxion mass $m_s \approx 4.5 - 45\,\mathrm{GeV}$.

3. The origin of cosmic axinos and cosmological constraints

There are two known ways in which axinos are produced in the early universe. One is the thermal production from the hot thermal bath after reheating. The other is the non-thermal production from decays of the lightest ordinary supersymmetric particles (LOSPs).

The decoupling temperature of axinos is estimated as [2]

$$T_D \sim 10^{10}\ \mathrm{GeV} \left(\frac{f_a}{10^{11}\,\mathrm{GeV}}\right) \left(\frac{\alpha_s}{0.1}\right)^{-3}, \tag{3.1}$$

where α_s is the strong coupling constant. If the reheat temperature T_R after inflation is higher than the decoupling temperature, the universe is overpopulated by axinos if the axino mass is larger than a few keV. Therefore, we only consider the case that the reheat temperature is lower than the decoupling temperature. In this case, axinos are produced from the thermal bath through scattering of quarks and gluons, though the number density of them do not reach the thermal equilibrium. The amount of axinos produced in this way, so called regeneration, is estimated to be [6]

$$\Omega_{\tilde{a}} h^2 \approx 0.28 \left(\frac{m_{\tilde{a}}}{\mathrm{MeV}}\right)\left(\frac{T_R}{10^5\,\mathrm{GeV}}\right)\left(\frac{f_a}{10^{11}\,\mathrm{GeV}}\right)^{-2}. \tag{3.2}$$

Thus, for the axino with mass $1-10\,\mathrm{MeV}$ to be the LDM, the relevant range of reheating temperature is $10-100\,\mathrm{TeV}$.

The axinos from decays of LOSPs can be cosmologically interesting when the axino mass is around the marginal value of order $10\,\mathrm{MeV}$. For this size of axino mass, the reheat temperature must be lower than $10\,\mathrm{TeV}$ to suppress the thermal production (regeneration). The amount of produced axinos is simply connected to that of LOSPs by

$$\Omega_{\tilde{a}} h^2 = \frac{m_{\tilde{a}}}{m_\chi}\Omega_\chi h^2, \tag{3.3}$$

and independent of the reheat temperature. When we take $m_\chi = 100\,\mathrm{GeV}$ and $m_{\tilde{a}} = 10\,\mathrm{MeV}$, the required value of $\Omega_\chi h^2$ is $\sim 10^4$. Such a high value is reached for very large M_{SUSY} in the range of tens of TeV. Thus, the non-thermal production of axinos for LDM could only be marginally relevant.

Even though relic axinos dominantly come from regeneration, the existence of LOSPs and their decay to axinos can produce radiative or hadronic cascades during or after the BBN, and alter its standard predictions on the light element abundances. To avoid this, the mass of LOSP needs to be large enough to make its lifetime much less than 1 section For example, in the case of the neutralino, one requires $m_\chi > 150\,\mathrm{GeV}$.

Let us discuss here how the accompanied saxion can also upset the standard prediction of the BBN, which is called "the saxion problem" [26]. Contrary to the axino, the saxion has the axion-like couplings to the quarks, $\frac{m_q}{f_a} s\bar{q}q$, or leptons, $\frac{m_l}{f_a} s\bar{l}l$, so that its life-time is much shorter than the axino LSP. On the other hand, during the axino regeneration (3.2), the saxions are also populated by the same amount and thus one finds

$$m_s Y_s \approx 10^{-9}\left(\frac{m_s}{\mathrm{MeV}}\right)\left(\frac{\mathrm{MeV}}{m_{\tilde{a}}}\right)\left(\frac{\Omega_{\tilde{a}} h^2}{0.28}\right)\,\mathrm{GeV}. \tag{3.4}$$

where Y_s is the saxion number density in unit of the entropy density. Note that this quantity is strongly constrained by the BBN. In the mass range $m_s \gtrsim \mathcal{O}(10)\,\mathrm{GeV}$, the above equation gives $m_s Y_s \gtrsim 10^{-5}\,\mathrm{GeV}$ for $m_{\tilde{a}} = 1\,\mathrm{MeV}$. Now that the saxion decays mainly to bottom and charm quarks, one finds a strong limit on the saxion lifetime: $\tau_s \lesssim 10^{-2}$ sec [27]. Specifically, the mass relation (2.5) gives us $m_s \approx 14\,\mathrm{GeV}$ for the axino mass $m_{\tilde{a}} \approx 1\,\mathrm{MeV}$ and $\Lambda = 10^5\,\mathrm{GeV}$. Then, the saxion lifetime,

$$\tau_s \approx \left[\frac{1}{8\pi}\frac{m_b^2}{f_a^2} m_s\right]^{-1} \lesssim 10^{-2}\ \mathrm{sec}\,, \tag{3.5}$$

becomes short enough to avoid the saxion problem for $f_a \lesssim 3\times 10^{11}\,\mathrm{GeV}$. In the case of supergravity models where one expects to get $m_s \approx 10^{2-3}\,\mathrm{GeV}$, the saxion is free of such a problem.

4. Axino-neutrino mixing and axino decay

Let us now assume the generation of the bilinear superpotential term, H_1H_2, and its R-parity and lepton number violating extension, L_iH_2 as a result of the PQ symmetry breaking;

$$W_{\rm eff} = \mu H_1H_2 + \epsilon_i\mu L_iH_2 \tag{4.1}$$

where μ and $\epsilon_i\mu$ carry PQ charges whose sizes are determined by the PQ charge assignments for $H_{1,2}$ and L_i. In eq. (2.2), the leading terms in A,

$$K_{\rm eff} = \frac{A}{f_a}[x_{H_i}H_i^\dagger H_i + x_{L_j}L_j^\dagger L_j] + \cdots\,, \tag{4.2}$$

give rise to the following axino-Higgsino and axino-neutrino mass terms;

$$\mathcal{L}_{\rm mixing} = x_{H_1}\frac{\mu v_2}{f_a}\,\tilde{a}\tilde{H}_1 + x_{H_2}\frac{\mu v_1}{f_a}\,\tilde{a}\tilde{H}_2 + x_{L_i}\frac{\epsilon_i\mu v_2}{f_a}\,\tilde{a}\nu_i + h.c.\,. \tag{4.3}$$

For $\mu v/f_a \ll m_{\tilde{H}}$ and $\epsilon_i\mu v_2/f_a \ll m_{\tilde{a}}$, one has the approximate mixing angles between the axino and Higgsino or neutrino as follows;

$$\theta_{\tilde{a}\tilde{H}} \sim \frac{v}{f_a} \quad \text{and} \quad \theta_{\tilde{a}\nu_i} = x_{L_i}\frac{\epsilon_i\mu v_2}{f_a m_{\tilde{a}}}\,. \tag{4.4}$$

The axino-neutrino mixing derived above induces the effective vertex of $\tilde{a}\nu_iZ$ and $\tilde{a}l_iW$ with the coupling $\sim g\theta_{\tilde{a}\nu_i}$. This gives rise to the four-quark operator as follows:

$$\mathcal{L}_{e^+e^-} \approx \frac{G_F}{\sqrt{2}}\,\theta_{\tilde{a}\nu_i}\,\bar{\nu}_i\gamma_\mu\gamma_5\tilde{a}\,\bar{e}\gamma^\mu(2\delta_{i1}-\gamma_5)e \tag{4.5}$$

where we omitted the small correction due to the vector part of the charged current.

Another important interaction to consider is the axino-photon-neutrino vertex arising from the photino-neutrino mixing. The bilinear term L_iH_2 induces the mixing between neutrinos and neutralinos of order ϵ_i. Then the supersymmetric anomaly coupling of axino-photon-photino leads to the axino-photon-neutrino coupling which is written down schematically as follows;

$$\mathcal{L}_\gamma = \frac{C_{a\gamma\gamma}\alpha_{\rm em}}{8\pi f_a}\epsilon_i\bar{\nu}_i\gamma_5\sigma_{\mu\nu}\tilde{a}F^{\mu\nu}, \tag{4.6}$$

where $C_{a\gamma\gamma}$ is an order-one parameter taking into account the precise values of the $\mathrm{U}(1)_{\rm em}$ anomaly and the photino-neutrino mixing. From the vertices (4.5) and (4.6), we get the following decay widths of the axino;

$$\begin{aligned}\Gamma_{\nu_ie^+e^-} &= \frac{G_F^2m_{\tilde{a}}^5}{192\pi^3}\,\theta_{\tilde{a}\nu_i}^2[\frac{1}{4}+\delta_{i1}]\\ \Gamma_{\nu_i\gamma} &= \frac{C_{a\gamma\gamma}^2\alpha_{\rm em}^2}{(16\pi)^3}\frac{m_{\tilde{a}}^3}{f_a^2}\,\epsilon_i^2\,.\end{aligned} \tag{4.7}$$

Let us first note that the photon mode is suppressed by $\alpha_{\rm em}^2$ compared to the e^+e^- mode;

$$\frac{\Gamma_{\nu\gamma}}{\Gamma_{\nu e^+e^-}} \approx \frac{3C_{a\gamma\gamma}^2\alpha_{\rm em}^2}{32G_F^2\mu^2v^2} \approx 10^{-4} \tag{4.8}$$

for $\mu/C_{a\gamma\gamma} = 100\,\mathrm{GeV}$. It is smaller than the internal bremsstrahlung process of e^+e^- mode which is suppressed by $\alpha_{\rm em}$ and also produces a direct γ-ray. This is enough to be consistent with the observations of the MeV γ-ray spectrum [19]. Then, the axino decay is determined by the process $\tilde{a} \to \nu e^+e^-$ whose lifetimes is given by

$$\tau_{\tilde{a}} \approx 10^{26}\,\mathrm{sec}\left(\frac{1\ \mathrm{MeV}}{m_{\tilde{a}}}\right)^3\left(\frac{f_a}{10^{11}\ \mathrm{GeV}}\right)^2\left(\frac{10^{-7}}{|x_L\epsilon|}\right)^2\left(\frac{100\ \mathrm{GeV}}{\mu}\right)^2 \tag{4.9}$$

which is in the right range to explain the observation (1.1) consistently with the neutrino data as will be shown in the following section.

5. Consistency with the neutrino data and experimental signatures

One of the interesting aspect of R-parity violation is that it can be the origin of the observed neutrino masses and mixing [23]. The general superpotential allowing R-parity and lepton number violation includes the following bilinear and trilinear terms;

$$W_{Rp} = \epsilon_i\mu L_i H_2 + \frac{1}{2}\lambda_{ijk}L_iL_jE^c_k + \lambda'_{ijk}L_iQ_jD^c_k\,. \tag{5.1}$$

According to the observation of ref. [13], an appropriate life time of the axino decay $\tilde{a} \to \nu_{\mu,\tau}e^+e^-$ can arise with trilinear R-parity violating couplings $\lambda_{211,311} \sim 0.1$. Such trilinear couplings can generate the 2-3 components of the neutrino mass matrix;

$$M^\nu_{ij} \approx \frac{1}{8\pi^2}\lambda_{i11}\lambda_{j11}\frac{m_e^2\mu\tan\beta}{m^2_{\tilde{e}}} \tag{5.2}$$

where $\tan\beta \equiv \langle H^0_2\rangle/\langle H^0_1\rangle$ and $m_{\tilde{e}}$ is the selectron soft mass. While the charged-current and e–μ–τ universality put the bound $\lambda_{i11} \lesssim 0.1(m_{\tilde{e}}/200\ \mathrm{GeV})$ [28], the above one-loop mass can reach the observed atmospheric neutrino mass scale $m_\nu \approx 0.05\,\mathrm{eV}$ only for an extreme value of $\mu\tan\beta \approx 50\,\mathrm{TeV}$ taking the boundary value of $\lambda_{i11} = 0.1\,(m_{\tilde{e}}/200\,\mathrm{GeV})$. In order to generate the other components of the neutrino mass matrix, one needs to introduce some other trilinear couplings such as $\lambda_{i22,j33}$ which induce M^ν_{11}, M^ν_{12} and M^ν_{13} through the combinations of $\lambda_{1jj}\lambda_{1jj}$, $\lambda_{133}\lambda_{233}$ and $\lambda_{122}\lambda_{322}$, respectively. Then, appropriate neutrino masses can be obtained for the trilinear couplings, $\lambda_{i22} \sim 10^{-4}$ and $\lambda_{i33} \sim 10^{-5}$, where the small ratios $\lambda_{i22}/\lambda_{i11}$ and $\lambda_{i33}/\lambda_{i11}$ are dictated by the factors of m_e/m_μ and m_e/m_τ, respectively. Such a hierarchy among λ_{ijj} appears *ad-hoc* considering the usual hierarchy in the quark and lepton Yukawa couplings.

Nevertheless, if there exits the trilinear coupling λ_{i11} of order 0.1, they leads to a remarkable experimental signature of resonant single sneutrino production in the future linear collider [29, 30], non-observation of which would rule out the axino LDM decaying through the trilinear couplings.

The observed neutrino masses and mixing can be more naturally explained if one invokes the presence of the bilinear term of the order 10^{-6} [23]. The bilinear R-parity violation generates neutrino masses at tree-level through the neutrino-neutralino mixing.

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In addition to the ϵ_i term in the superpotential (5.1), the scalar potential also contains the R-parity violating bilinear soft terms as follows;

$$V_0 = m^2_{L_i H_1} L_i H_1^\dagger + B_i L_i H_2 + h.c., \tag{5.3}$$

where B_i is the dimension-two soft parameter. Generically, one has $B_i = \epsilon_i \tilde{B}\mu$ with a dimension-one soft parameter $\tilde{B}$ for the μ term, and the soft mass-squared $m^2_{L_i H_1}$ contains the supersymmetric term $\epsilon_i \mu^2$. Upon the electroweak symmetry breaking, the sneutrino field gets nontrivial vacuum expectation value;

$$\frac{\langle \tilde{\nu}_i \rangle}{v_1} = -\frac{m^2_{L_i H_1} + B_i \tan\beta}{m^2_{\tilde{\nu}_i}}\,, \tag{5.4}$$

which is expected to be of order ϵ_i up to the soft mass dependence. These bilinear parameters induce mixing between neutrinos and neutralinos. For the small mixing mass, the week-scale seesaw with heavy neutralino mass scale $\sim 100\,\mathrm{GeV}$ leads to the well-known neutrino mass matrix at tree-level;

$$M^\nu_{ij} = -\frac{M_Z^2}{F_N} \xi_i \xi_j \cos^2\beta \tag{5.5}$$

where $\xi_i \equiv \epsilon_i - \langle \tilde{\nu}_i \rangle / v_1$ and $F_N = M_1 M_2 / M_{\tilde{\gamma}} + M_Z^2 \cos 2\beta/\mu$ with $M_{\tilde{\gamma}} = c_W^2 M_1 + s_W^2 M_2$. From eq. (5.5), one obtains the size of $|\xi| = \sqrt{\sum_i |\xi_i|^2}$ consistently with the atmospheric neutrino mass scale as follows;

$$|\xi| = 0.7 \times 10^{-6} \frac{1}{\cos\beta} \left(\frac{F_N}{M_Z}\right)^{1/2} \left(\frac{m_\nu}{0.05\,\mathrm{eV}}\right)^{1/2}. \tag{5.6}$$

This is compatible with the axino lifetime relation (4.9) for $\xi_i \sim \epsilon_i$. Note that the smaller neutrino mass explaining the solar neutrino oscillation can come from one-loop diagrams involving the trilinear couplings of order, $\lambda_{i33}, \lambda'_{i33} \sim 10^{-4:-5}$.

Let us finally remark that the bilinear R-parity violation leads to a distinct prediction on the lepton number violating decays of the lightest neutralino χ in the future colliders. The mass matrix of the form (5.5) enables us to determine the relation $5|\xi_1| \lesssim |\xi_2| = |\xi_3|$ from the neutrino data on the mixing angles. As the parameters ξ_i determine also the couplings of the R-parity violating processes; $\chi \to l_i^\pm W^\mp$, the above mixing angle relation can be tested in the decay of the neutralino whose branching ratios satisfies $\mathrm{Br}(eW) : \mathrm{Br}(\mu W) : \mathrm{Br}(\tau W) = |\xi_1|^2 : |\xi_2|^2 : |\xi_3|^2$ [23]. It is intriguing to note that future colliders can provide an indirect test for either scenario of the axino LDM decaying through λ_{i11} or ϵ_i.

6. Conclusion

The axino with the mass in the $1 - 10\,\mathrm{MeV}$ range is a good candidate for the LDM, which not only constitutes CDM but also explains the observed 511 keV γ-rays from the galactic bulge through its decay. The desired mass of the axino can be realized in certain supergravity models with some special arrangements, e.g., no-scale Kähler potential, or in gauge-mediated SUSY breaking models. The origin of relic axinos can be either the thermal production from the thermal bath after reheating or the non-thermal production from the LOSP decays. Both require a rather low reheat temperature $T_R \sim 10 - 100\,\mathrm{TeV}$.

The long lifetime of the axino is a result of the R-parity violation and the suppression of axino interactions with ordinary particles by the PQ scale. As is well-known, the small violation of R-parity by bilinear terms is an attractive option for generating the neutrino masses and mixing. We found an interesting fact that the same small R-parity violating bilinear terms can explain the observed 511 keV γ-rays as well as the observed neutrino mass matrix consistently within the current observational bounds and the reasonable choice of model parameters. This connection has a virtue that the explanation of neutrino masses and mixing by R-parity violating bilinear terms has testable predictions in the future colliders, thereby provides an indirect test of decaying axino LDM. The LDM is an very attractive idea in that if it turns out to be true, the morphology of 511 keV gamma-rays will serve as a good probe of the dark matter halo density profile. The decaying LDM models require more curspy density profile to fit the observed morphology of 511 keV γ-rays from the galactic bulge than the annihilation models. We expect this leads to interesting astrophysical implications [31, 32].

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