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Review

Axino phenomenology

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Abstract. The strong CP problem is solved elegantly by the PQ mechanism which predicts the presence of a light pseudo Goldstone boson called the axion. In supersymmetric theories, the axion is accompanied by its fermionic partner called the axino. It can play an important role in collider, dark matter, and neutrino physics. We review general properties of the axino in relation to the standard axion models, and discuss various phenomenological and cosmological implications.

1 Strong CP problem and axion

Standard Model (SM) allows several CP violating parameters: CP phases in Yukawa (mass) matrices Y_f and θ parameters in gauge field strengths:

$$\begin{aligned}\mathcal{L}_{\rm Yuk} &= H_u\,\bar{q}_L Y_u u_R + H_d\,\bar{q}_L Y_d d_R + H_d\,\bar{l}_L Y_l l_R + h.c.\\ \mathcal{L}_\theta &= \frac{g_3^2}{32\pi^2}\theta_3 G^a_{\mu\nu}\tilde{G}^{a\mu\nu} + \frac{g_2^2}{32\pi^2}\theta_2 W^i_{\mu\nu}\tilde{W}^{i\mu\nu} + \frac{g_1^2}{16\pi^2}\theta_1 B_{\mu\nu}\tilde{B}^{\mu\nu}.\end{aligned} \tag{1}$$

Various CP violating phenomena observed in the electroweak interactions are well described by the CKM phase in the quark Yukawa (mass) matrices: the Jarlskog invariant [1]

$$\delta = {\rm ArgDet}[Y_u Y_u^\dagger, Y_d Y_d^\dagger] \tag{2}$$

which is measured to be $\delta = 1.19 \pm 0.05$ [2]. In the QCD sector, the CP violating parameters induce an electric-dipole moment (EDM) of a nucleon through the effective strong θ term:

$$\bar{\theta} = \theta_3 + {\rm ArgDet}(Y_u Y_d). \tag{3}$$

For the neutron, one finds [3]

$$d_n \sim e\frac{m_q}{m_n^2} \sim 2.5\times 10^{-16}\bar{\theta}\,e{\rm cm}, \tag{4}$$

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and thus the non-observation of the neutron EDM: $d_n < 3 \times 10^{-26}\, e\text{cm}$ [4] requires $|\bar{\theta}| < 10^{-10}$. This amounts to the strong CP problem: "Why is $\bar{\theta}$ vanishingly small, particularly, in contrast to δ?".

Let us remark that the weak θ parameter is unobservable due to the chiral nature of the electroweak symmetry, that is, θ_2 can be rotated away by the (anomalous) $B+L$ symmetry [5].

The strong CP problem is elegantly resolved by the Peccei-Quinn-Weinberg-Wilczek (PQWW) mechanism [6–9] introducing a spontaneously broken global $U(1)_{PQ}$ symmetry which has a QCD anomaly. The Goldstone boson of such a symmetry, called the axion a, has the anomaly coupling:

$$\mathcal{L}_a^{\text{anomaly}} = \frac{g_3^2}{32\pi^2} \frac{a}{f_a} G^a_{\mu\nu} \tilde{G}^{\mu\nu}_a, \tag{5}$$

where f_a is the axion decay constant proportional to the $U(1)_{PQ}$ symmetry breaking scale. The QCD condensatation generates a potential for the axion (redefined after absorbing $\bar{\theta}$):

$$V[a] \approx m_\pi^2 f_\pi^2 \left(1 - \cos\left[\frac{a}{f_a}\right]\right) \tag{6}$$

which sets $\langle a/f_a \rangle \equiv 0$ at the minimum and thus dynamically resolves the strong CP problem. The axion also becomes massive due to such a condensation potential: $m_a \approx m_\pi f_\pi / f_a$. The original PQWW axion model realized at the weak scale has been ruled out, but high-scale axion models can be realized by inroducing a heavy quark (KSVZ) [10,11] or two Higgs boublets (DFSZ) [12,13]. The allowed window of the axion scale is $10^9 \lesssim f_a/\text{GeV} \lesssim 10^{12}$ where the lower limit comes from star cooling processes and the upper limit from the axion cold dark matter contribution taking the initial mis-alignment angle θ_i of order one. One may allow higher f_a if $\theta_i \ll 1$ is assumed depending on cosmological scenarios. For recent reviews on axion physics and cosmology, we refer the readers to [14,15] and references therein.

2 Axino mass and couplings

In the supersymmetric extension of the Standard Model, the axion field is promoted to a chiral supermultiplet A containing its scalar s (saxion) and fermion $\tilde{a}$ (axino) partner:

$$A \sim (s + ia, \tilde{a}) \tag{7}$$

which is essentially a phase superfield of a PQ symmetry breaking superfield $S \sim f_a e^{A/f_a}$. The axion fraction in S and the precise relation between $\langle S \rangle$ and f_a are model-dependent. While the axion being a pseudo-Goldstone boson is very light, the saxion and axino masses are expected to be around the supersymmetry breaking scale or well below depending on the supersymmetry breaking mechanism and the PQ symmetry breaking sector [16–18].

The KSVZ axion model is realized in supersymmetry by introducing a $U(1)_{PQ}$-invariant superpotential

$$W_{\text{KSVZ}} = \lambda S Q Q^c \tag{8}$$

where (Q, Q^c) is a vector-like pair of a heavy quark with the mass $M_Q = \lambda\langle S\rangle$. Integrating out the heavy quark superfield, one obtains the conventional axion-gluon-gluon coupling (5) as well as it supersymmetric counterpart, the axino-gluino-gluon coupling:

$$\mathcal{L}_{\tilde{a}}^{\text{anomaly}} = \frac{g_s^2}{32\pi^2}\frac{\tilde{a}}{f_a}\sigma^{\mu\nu}\tilde{g}^a G^a_{\mu\nu} + h.c.. \tag{9}$$

In the supersymmetric realization of the DFSZ axion, the Higgsino mass term μ is generated naturally from the PQ symmetry breaking *à la* Kim-Nilles [19,20]:

$$W_{\text{DFSZ}} = \lambda\frac{S^2}{M_*}H_1H_2, \tag{10}$$

where M_* is an UV scale. Notice that the electroweak μ term is induced by $\mu \sim \lambda f_a^2/M_*$. For instance, one gets $\mu \sim 1$ TeV for $f_a = 10^{10}$ GeV, $M_* = 10^{16}$ GeV and $\lambda = 0.1$. Due to the superpotential (10), the DFSZ axino has the couplings with Higgsino-Higgs and fermion-sfermion in addition to the anomaly coupling (9):

$$\mathcal{L}_{\tilde{a}}^{\text{DFSZ}} = c_H\frac{\mu}{f_a}\tilde{a}[\tilde{H}_1H_2 + \tilde{H}_2H_1] + c_f\frac{m_f}{f_a}\tilde{a}[f\tilde{f}^c + f^c\tilde{f}] + h.c., \tag{11}$$

where c_H and c_f are model-dependent parameter of order one.

3 Axino abundance and dark matter

While the axion is a good dark matter candidate, the axino can be another candidate if it is the lightest supersymmetric particle (LSP). Due to its superweak interaction, it becomes a warm dark matter candidate if it is produced by freeze-out process [21]. It can be produced also by a resonant axino-neutrino conversion in R-parity and lepton number violating models and become a cool (not so cold) dark matter candidate [22]. However, the major production mechanism is the generation by thermal scattering [23–29]. The resulting axino aboundance follows the simple Boltzmann equation

$$\frac{dY_{\tilde{a}}}{dT} = -\frac{\gamma}{sHT} \tag{12}$$

in terms of the axino number density normalized by the entropy density s at a given temperature T: $Y_{\tilde{a}} \equiv n_{\tilde{a}}/s$. Here γ is the thermal scattering rate for the axino production and H is the Hubble rate. Notice that this is an classsic example of the process dubbed as "freeze-in" [30].

In the case of the KSVZ axino [31,32], its thermal generation is governed by the $2 \to 2$ processes such as $\tilde{g}\tilde{g} \to \tilde{g}\tilde{a}$ involving the axino-gluino-gluon coupling (9) [24–26] assuming $T \ll M_Q$ [28]. The corresponding scattering rate behaves as $\gamma \propto T^6/f_a^2$ and thus the final axino abundance is given by

$$Y_{\tilde{a}} = \frac{\gamma}{sH}\Big|_{T=T_R} \tag{13}$$

where T_R is the reheat temperature after inflation. As a result, the axino relic density normalized by the present critical energy density is found to be

$$\Omega_{\tilde{a}}^{\rm KSVZ} h^2 \approx 0.11 \left(\frac{m_{\tilde{a}}}{44\,{\rm MeV}}\right) \left(\frac{T_R}{1\,{\rm TeV}}\right) \left(\frac{10^{11}\,{\rm GeV}}{f_a}\right)^2 . \tag{14}$$

This puts an upper limit on the reheat temperature $T_R < 10^7$ GeV if the axino is stable and thus becomes a dark matter candidate heavier than a few keV gauranteeing sufficient density perturbations.

For the DFSZ axino [27,29], the thermal production could be dominated by the $1 \to 2$ processes like $\tilde{H} \to \tilde{a} H$ or the $2 \to 2$ processes such as $t\tilde{t}^* \to \tilde{a} H$ involving the axino-Higgsino-Higgs coupling (10). In the case of the Higgsino decay, the axino production rate goes like $\gamma \propto \Gamma_{\tilde{H}} T^3$ where

$$\Gamma_{\tilde{H}} \approx \frac{c_H^2}{32\pi} \frac{\mu^3}{f_a^2}. \tag{15}$$

Thus one finds

$$Y_{\tilde{a}} = 2 \frac{\gamma}{sH}\Big|_{T=\mu} \tag{16}$$

which gives the axino relic density:

$$\Omega_{\tilde{a}}^{\rm DFSZ} h^2 \approx 0.11 \left(\frac{m_{\tilde{a}}}{18\,{\rm keV}}\right) \left(\frac{\mu}{1\,{\rm TeV}}\right) \left(\frac{10^{11}\,{\rm GeV}}{f_a/c_H}\right)^2 . \tag{17}$$

A similar result can be drawn also from the axino-stop-top coupling (10).

The above relation (17) implies that the DFSZ axino heavier than about 20 keV cannot be stable. If the axion decay process $\tilde{a} \to H\tilde{H}$ ($\tilde{H}$ may further decay to the LSP χ), it has a very small decay rate:

$$\Gamma_{\tilde{a}} \approx \frac{c_H^2}{16\pi} \left(\frac{\mu}{f_a}\right)^2 m_{\tilde{a}}, \tag{18}$$

which corresponds to the decay temperature $T_D = g_*^{-1/4}\sqrt{3\Gamma_{\tilde{a}} M_P}$. Defining $x_D \equiv m_\chi/T_D$ where χ can be the lighter Higgsino, one gets

$$x_D \approx 24 \left(\frac{g_*}{70}\right)^{1/4} \left(\frac{1\,{\rm TeV}}{m_{\tilde{a}}}\right)^{1/2} \left(\frac{m_\chi}{\mu}\right) \left(\frac{f_a/c_H}{10^{11}\,{\rm GeV}}\right). \tag{19}$$

Thus, the decay temperature can be even smaller than the usual freeze-out temperature T_f of the LSP, that is, $x_D > x_f$ where $x_f \equiv m_\chi/T_f \approx 25$. In this case, the decay-produced LSPs re-annihilate and their final number density $Y_\chi \equiv n_\chi/s$ follows the Boltzmann equation:

$$\frac{dY_\chi}{dT} = \frac{s\langle\sigma_A v\rangle}{HT} Y_\chi^2 \tag{20}$$

where $\langle\sigma_A v\rangle$ is the LSP annihilation rate. When $\langle\sigma_A v\rangle$ is independent of T, one finds

$$Y_\chi \approx \left.\frac{H}{s\langle\sigma_A v\rangle}\right|_{T=T_D} \tag{21}$$

leading to the final dark matter relic density:

$$\Omega_\chi h^2 \approx 0.1\frac{10^{-10}\,\mathrm{GeV}^{-2}}{(g_*/70)^{1/2}}\frac{x_D}{\langle\sigma_A v\rangle}. \tag{22}$$

Thus, one can get the right dark matter density for $\langle\sigma_A v\rangle \approx x_D \times 10^{-10}\,\mathrm{GeV}^{-2}$ which can be even larger than the standard freeze-out value. This opens up a new parameter space for the neutralino dark matter and thus has an important impact on the scenarios of mixed axion/neutralino dark matter [33–36].

4 Light Higgsino and axino dark matter

The electroweak symmetry breaking in supersymmetry requires a potential minimization condition:

$$\frac{m_Z^2}{2} = \frac{m_{H_d}^2 - m_{H_u}^2\tan^2\beta}{\tan^2\beta - 1} - \mu^2 \tag{23}$$

where $m_{H_{u,d}}$ are the soft masses of the two Higgs doublets, $\tan\beta \equiv v_u/v_d$ is the ratio of their vacuum expectation values, and μ is the Higgs bilinear parameter in the superpotential. As LHC finds no supersymmetry signals, masses of supersymmeric particles are pushed above TeV range and thus the minimization condition (23) requires a certain degree of fine-tunning. Barring too much fine-tunning, one may arrange $m_{H_{u,d}}$ and μ not too larger than m_Z. Thus, it is conceivable to have light Higgsinos decaying to the axino dark matter in the context of supersymmetric DFSZ axion model [37]. Such light Higgsinos can be copiously produced at the LHC even through Drell-Yan processes and decay to axino plus the Higgs boson h or Z boson. Given the decay rate of the Higgsino (15), one obtains the typical decay length of

$$c\tau_{\widetilde{H}^0} \approx 135\,\mathrm{m}\left(\frac{200\,\mathrm{GeV}}{\mu}\right)^3\left(\frac{f_a/c_H}{10^{12}\,\mathrm{GeV}}\right)^2. \tag{24}$$

Therefore, the Higgsino decay signals:

$$\widetilde{H}^0 \to \tilde{a}Z,\, \tilde{a}h \to \text{displaced dilepton/dijet + MET} \tag{25}$$

can be detected by MATHUSLA [38]. At the LHC, the displaced 4ℓ + MET search is most sensitive, but MATHUSLA case can be different because of different background composition and particle identification. The branching ratios depend on the value of $\tan\beta$ and μ, but in general all decays modes above have similar branching ratios.

The signal rate is enhanced by the fact that heavier (charged/neutral) Higgsino states all decay promptly to the lightest Higgsino state $\widetilde{H}_1^0$ (with invisibly soft particles) rather than decaying directly to the axino LSP. Although heavier Higgsino states are mass-splitted $\mathcal{O}(100)$ MeV only by electroweak quantum correction and small gaugino mixings, the splitting is big enough for this to be true. Thus the signal

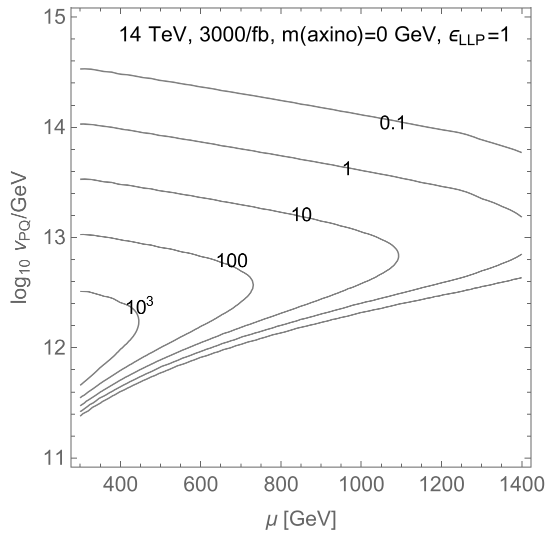

Fig. 1. Number of long-lived Higgsino decays to DFSZ axino+Z/h in MATHUSLA assuming Drell-Yan Higgsino production at the HL-LHC. A perfect LLP detection efficiency is assumed.

rate is a sum of all kinds of pair productions of Higgsino states. Figure 1 shows MATHUSLA prospect of detecting Higgsino LLP in the plane of $(\mu, f_a = v_{PQ})$ assuming an unknown detector efficiency to be maximal ($\epsilon_{LLP} = 1$) for the axino mass $m_{\tilde{a}} = 0$. The results will depends linearly on ϵ_{LLP}.

5 Axino as a sterile neutrino

The existence of a sterile neutrino is hinted by simultaneous explanation of diverse neutrino anomalies [39]. It has been suggested that the axino arising in the supersymmetric axion solution to the strong CP problem is a natural candidate of a sterile neutrino [40].

As mentioned eariler, the axino mass is model-dependent and can be made very light (e.g., $m_{\tilde{a}} \sim 1$ eV) particularily in the context of gauge mediated supersymmetry breaking [41]. The axino-neutrino mixing can arise naturally in a generalization of the Kim-Nilles mechanism with R-parity violation. To see this, let us realize the DFSZ axion model by assigning the following PQ charges [42]:

Superfields	S	H_u	H_d	L	E^c	Q	U^c	D^c
PQ charges	1	−1	−1	−2	3	1	0	0

(26)

This allows us to have the superpotential

$$W'_{\rm DFSZ} = a\frac{S^2}{M_*}H_u H_d + a'_i \frac{S^3}{M_*^2} H_u L_i + b_{ijk}\frac{S}{M_*} L_i L_j E^c_k + b'_{ijk}\frac{S}{M_*} L_i Q_j D^c_k \tag{27}$$

in addition to the usual quark and lepton Yukawa terms. After the PQ symmetry breaking, $\langle S\rangle \sim f_a$, one obtains an effective superpotential

$$\begin{aligned} W'_{\rm DFSZ} = &\mu e^{2A/f_a} H_u H_d + \mu'_i e^{3A/f_a} H_u L_i \\ &+ e^{A/f_a}(\lambda_{ijk} L_i L_j E^c_k + \lambda'_{ijk} L_i Q_j D^c_k) \end{aligned} \quad (28)$$

replacing S by $f_a e^{A/f_a}$ with the axion supermultiplet A. Notice that the leading order terms include the bilinear and trilinear R-parity (and lepton number) violating couplings which can generate the desired neutrino mass matrix [43]. The axino-neutrino mixing arises from the second term in (28):

$$\mathcal{L}_{\nu\tilde a} = m_{ia}\nu_i \tilde a \;\; \text{with} \;\; m_{ia} = 3\mu'_i \langle H^0_u\rangle / f_a \quad (29)$$

which determines the axino-neutrion mixing angle $\theta_{ia} \sim m_{ia}/m_{\tilde a} \ll 1$. Therefore, the DFSZ axino combined with R-parity violation can provide a $3+1$ neutrino oscillation scheme.

Our formulation also has interesting implications in dark matter physics. In variant models, the axino can be interpreted as a keV sterile neutrino dark matter [44]; a source of the 511 keV gamma-line via the decay mode $\tilde a \to e^+ e^- \nu$ [45]; and the 3.5 keV X-line [46–50] or the 130 GeV gamma-line via $\tilde a \to \gamma\nu$ [51].

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