Review of axino dark matter

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We review the status of axino dark matter. Two hierarchy problems, the strong CP problem and the gauge hierarchy problem, have led to introducing into particle physics a spontaneously broken global Peccei-Quinn symmetry and a softly broken supersymmetry, respectively. Combining them implies the presence of not only an axion, but also of its scalar component, saxion, and their fermionic partner, axino. Among these, the axion and the axino are attractive dark matter candidates. Various possibilities for the axino as dark matter are discussed.

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I. INTRODUCTION

The axino, stable or almost stable on cosmological time scales, is a well-motivated dark matter (DM) candidate. Axinos differ from ordinary weakly interacting massive particles (WIMPs) in that they are exceedingly weakly interacting (EWIMPs), which radically changes their cosmological properties and also ways of testing them in experiment.

Relic axinos can be produced in a hot plasma or in decays of heavy particles in the early universe. The most interesting case is when it can be cold dark matter (CDM) - the possibility that was first considered in Ref. 1 where axinos were generated in the decays of heavier particle after freezeout in the process that was dubbed non-thermal production (NTP). This was next extended in Ref. 2 to include thermal production (TP) through scatterings and decays of particles in thermal equilibrium. It was also pointed out there that the axino could as well be hot DM (HDM) or warm DM (WDM), or that, for instance, one of its populations could be warm while the other cold. Since then the axino as DM in cosmology, astrophysics and collider experiments has been studied in many papers [3-14].

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The axino as the superpartner of the axion was first considered right after it was recognized that softly broken supersymmetry (SUSY) was relevant for particle physics [15-17]. Therefore, an experimental confirmation of axino existence would validate two theoretical hypotheses designed to solve two respective hierarchy problems: the strong CP problem by a very light axion [18] and the gauge hierarchy problem by SUSY [19].

Similarly to the axion, axino interactions with Standard Model (SM) and minimal supersymmetric SM (MSSM) particles are suppressed by the axion decay constant \( f_a \). On the other hand, the quantity that is most relevant for the axino in astro-particle physics, and at the same time most poorly known, is its mass \( m_{\tilde{a}} \). Several theoretical calculations of the axino mass followed [20-22]. A method for calculating axino mass applies to any goldstino (the superpartner of a Goldstone boson). A goldstino related to the Goldstone boson has a root in a global U(1) symmetry and receives its mass below the SUSY breaking scale. SUSY breaking triggers the super-Higgs mechanism and is related to the gravitino mass \( m_{3/2} \) which has recently been clarified in [22]. Even though generically axino mass is of order \( m_{3/2} \), a theoretically allowed mass range encompasses all the range relevant for hot, warm and cold DM axinos. We will discuss this in more detail below.

In the early days, a very light HDM-like axino from the decay of a photino was considered [23] to constrain the photino mass dependence on the axion decay con-
A general low energy axion interaction Lagrangian can be written in terms of the effective couplings with the SM fields $c_1$, $c_2$, and $c_3$, which arise after integrating out all heavy PQ-charge carrying fields. The resulting effective axion interaction Lagrangian terms are 

$$\mathcal{L}_{\text{eff}}^{\text{int}} = c_1 \frac{(\partial_\mu a)}{f_a} \sum_q \bar{q}_q \gamma^\mu \gamma_5 q - \sum_q (\bar{q}_L m_q R e^{i a q_0/f_a} + \text{h.c.}) + \frac{c_3}{32\pi^2 f_a} a G G \nabla^2 \rho \rho$$

where $c_3$ can be set to one by rescaling $f_a$. The axion decay constant $f_S$, $\theta = a/f_S$, is defined up to the domain wall number, $f_S = N_{\text{DW}} f_a$. The derivative interaction term proportional to $c_1$ preserves the PQ symmetry. The $c_2$-term is related to the phase of the quark mass matrix, and the $c_3$-term represents the anomalous coupling. The axion-lepton interaction term $\mathcal{L}_{\text{leptons}}$ is analogous to the axion-quark interaction term.

Two prototype field theory models for very light axions have been considered in the literature. At the SM level, one considers the six SM quarks, $u, d, s, \ldots$, as strongly interacting matter fermions. Above the electroweak scale $v_{\text{EW}} \approx 247$ GeV one additionally introduces beyond the SM (BSM) heavy vector-like quarks ($Q_1$, $\bar{Q}_1$), which in the interaction Lagrangian are next integrated out.

At the field theory level the axion is present if there exist quarks carrying the net PQ charge $\Gamma$ of the global $U(1)_{\text{PQ}}$ symmetry. In the Kim-Shifman-Vainshtein-Zakharov (KSVZ) model one introduces only heavy quarks as PQ charge carrying quarks. This results in $c_1 = c_2 = 0$, and $c_3 = 1$ below the $v_{\text{EW}}$, or below the QCD scale $\Lambda_{\text{QCD}}$. The gluon anomaly term (the $c_3$ term), induced by an effective heavy quark loop, then solves the strong CP problem. The axion field is a component of the SM singlet scalar field $S$. In the Dine-Fischler-Srednicki-Zhitnitski (DFSZ) model instead one does not assume any net PQ charge in the BSM sector, and instead the SM quarks are assigned the net PQ charge, i.e., $c_1 = c_3 = 0$ and $c_2 \neq 0$ below the electroweak scale $v_{\text{EW}}$. Here also, the axion is predominantly a part of the SM singlet scalar field $S$. Several specific implementations of the KSVZ and DFSZ frameworks can be found in Refs. \[18, 28\] which, however, require a whole host of additional BSM fields. Therefore, any references to the properties of the KSVZ and the DFSZ models can serve at best as just guidelines. In this respect, it is unfortunate that there exists only one reference clarifying the axion-photon-photon coupling from a string-derived BSM framework [29].

SUSY models of very light axion can provide a clue about the magnitude on the axion decay constant $f_a$. In Ref. [30] it was speculated that it is related to the MSSM Higgs/higgsino mass parameter $\mu$ as $f_a \sim \sqrt{\mu M_P}$. Recently, a permutation symmetry $S_2 \times S_2$ has been used to relate $f_a$ and $\mu$, realizing the old hypothesis in terms

\[\text{II. AXION/AXINO MODELS}\]

The strong CP problem can most naturally be solved by introducing a very light axion field $a$ [12] which couples to the gluon anomaly

$$\mathcal{L} = \frac{\alpha_s a}{8\pi f_a} G^{\mu\nu} \tilde{G}_{\mu\nu}^a,$$

where $\alpha_s = g_s^2/4\pi$ is the strong coupling constant and $\tilde{G}^a_{\mu\nu} = \frac{1}{2} \varepsilon^{\mu\nu\rho\sigma} G^a_{\rho\sigma}$ is the dual of the field strength for eight gluons $G^a_{\mu\nu}$ ($a = 1, 2, \ldots, 8$). This interaction term can be obtained after integrating out colored heavy fields below the PQ symmetry breaking scale $f_a$ but above the electroweak scale $v_{\text{EW}}$. The axion decay constant $f_a$ is constrained from astrophysical and cosmological considerations to a narrow window $10^{10} \text{ GeV} \lesssim f_a \lesssim 10^{12} \text{ GeV}$ [18]. The upper bound is obtained here assuming that the initial misalignment angle is of order one, and can be lifted if the angle were assumed smaller than one [25].

\[\text{1} \quad \text{For other recent reviews on axinos as DM particles, see [1].}\]
of a discrete symmetry \[31\]. Using a discrete symmetry toward an approximate global symmetry is theoretically welcome in that it evades the wormhole breaking of the PQ symmetry \[32\].

Before one considers spontaneous symmetry breaking of \(U(1)_{\text{PQ}}\), the axion Lagrangian can be said to have the axion shift symmetry (which is just a phase rotation), \(a \rightarrow a + \text{const.}\), and the physical observables are invariant under the PQ phase rotation. Below \(f_a\), the PQ rotational symmetry is broken, which is explicitly reflected as a breaking of the axion shift symmetry through the appearance of the \(c_3\) term in Eq. (2). However, the \(c_2\) term enters into the phase and a discrete shift of the axion field can bring it back to the original value. The \(c_3\) term is the QCD vacuum angle term and if the vacuum angle is shifted by \(2\pi\) then it comes back to the original value. Thus, even though the \(U(1)_{\text{PQ}}\) is broken, one of its discrete subgroups, i.e., the one corresponding to the common intersection of the subgroups corresponding to the \(c_2\) and \(c_3\) terms, can never be broken. As a result, the combination \(c_2 + c_3\) is invariant under the axion shift symmetry, and \(c_2 + c_3\) is defined to be an integer signifying the unbroken discrete subgroup of \(U(1)_{\text{PQ}}\). It is called the domain wall number \(N_{\text{DW}} = |c_2 + c_3|\) \[33\].

When an axion model is supersymmetrized \[15–17\], there appears a fermionic SUSY partner of the axion field called the axino \(\tilde{a}\), as well as a real scalar field \(s\) named the saxion. Together with the axion they form an axion supermultiplet \(A\)

\[
A = \frac{1}{\sqrt{2}}(s + ia) + \sqrt{2} \tilde{a} \theta + F_A \vartheta \theta,
\]

where \(F_A\) stands for an auxiliary field of \(A\) and \(\vartheta\) for a Grassmann coordinate. The interaction of the axion supermultiplet is obtained by supersymmetrizing the axion interaction in Eq. (2). In particular, the interaction of the axion supermultiplet \(A\) with the vector multiplet \(V_a\), which is a SUSY version of the \(c_3\) term in Eq. (2), is given by

\[
\mathcal{L}_{\text{eff}}^{\text{ax}} = - \sum_V \frac{\alpha_V C_{AVV}}{2 \sqrt{2} \pi f_a} \int A \, \text{Tr} \left[ V_a V^a \right] + \text{h.c.},
\]

where \(\alpha_V\) denotes a gauge coupling, \(C_{AVV}\) is a model dependent constant and the sum goes over the SM gauge groups. From this the relevant axino-gaugino-gauge boson and axino-gaugino-sfermion-sfermion interaction terms can be derived and are given by \[34\]

\[
\mathcal{L}_a^{\text{eff}} = \frac{\alpha_a}{16 \pi f_a} \bar{a} \gamma_5 [\gamma^\mu, \gamma^\nu] \tilde{G}^b G^b_{\mu\nu} + \frac{\alpha_a}{4 \pi f_a} \bar{a} \tilde{g} a \sum_q g_q T^a \tilde{q} \tilde{q}
+ \frac{\alpha_2 C_{WY}}{16 \pi f_a} \bar{a} \gamma_5 [\gamma^\mu, \gamma^\nu] \tilde{W}^b W^b_{\mu\nu}
+ \frac{\alpha_2}{4 \pi f_a} \tilde{W}^a \sum_f g_f \tilde{f} B T^a \tilde{f} D
+ \frac{\alpha_Y C_{AYY}}{16 \pi f_a} \bar{a} \gamma_5 [\gamma^\mu, \gamma^\nu] \tilde{Y} Y_{\mu\nu}
+ \frac{\alpha_Y}{4 \pi f_a} \tilde{Y} \sum_f g_Y \tilde{f}^* Q Y \tilde{f},
\]

where the terms proportional to \(\alpha_2\) correspond to the \(SU(2)_L\) and the ones proportional to \(\alpha_Y\) to the \(U(1)_Y\) gauge groups, respectively. \(C_{WY}\) and \(C_{AYY}\) are model-dependent couplings for the \(SU(2)_L\) and \(U(1)_Y\) gauge group axino-gaugino-gauge boson anomaly interactions, respectively, which are defined after the standard normalization of \(f_a\), as in Eq. (1) for the \(SU(3)_c\) term. Here \(\alpha_2, \tilde{W}, W_{\mu\nu}\), and \(\alpha_Y, \tilde{Y}, Y_{\mu\nu}\) are, respectively, the gauge coupling, the gaugino field and the field strength of the \(SU(2)_L\) and \(U(1)_Y\) gauge groups. \(\tilde{f}_D\) represents the sfermions of the \(SU(2)_L\)-doublet and \(\tilde{f}\) denote the sfermions carrying the \(U(1)_Y\) charge.

Similarly, one can derive supersymmetrized interactions of the axion supermultiplet with matter multiplet as a generalization of the \(c_1\) and \(c_2\) terms in Eq. (2). Ref. \[35\] considered a generic form of effective interactions and clarified the issue of energy scale dependence of axino interactions. At some energy scale \(p\) which is larger than the mass of the PQ-charged and gauge-charged multiplet \(M_{\phi}\), the axino-gaugino-gauge boson interaction is suppressed by \(M_{\phi}^2/p^2\). This suppression is manifest in the DFSZ axion model or in the KSVZ model if the heavy quark mass is relatively low compared to the PQ scale in which case the heavy quark is of course not integrated out.

However, SUSY must be broken at low energy and thus a SUSY relation between the axino and the axion is modified. In fact, the most important axino parameter in cosmological considerations, the axino mass \(m_a\), does not even appear in Eq. (5). SUSY breaking generates the masses for the axino and the saxion and modifies their definitions. The saxion mass is set by the SUSY soft breaking mass scale, \(M_{\text{SUSY}}\) \[16, 36\]. The axino mass, on the other hand, is strongly model dependent. An explicit axino mass model with SUSY breaking was first constructed long time ago \[30\] with the superpotential for the PQ symmetry transformation \(S \rightarrow e^{i \alpha} S\) and \(\bar{S} \rightarrow e^{-i \alpha} \bar{S}\).

\[
W = \sum_{i=1}^{n_I} Z_i (S \bar{S} - f_i^2), \quad n_I \geq 2.
\]

With \(n_I = 1\), the \(U(1)_{\text{PQ}}\) symmetry is spontaneously
broken but SUSY remains unbroken. The case $n_I = 2$ was also considered in [30] but this model gives $m_3 = 0$.

As first pointed out by Tamvakis and Wyler [16], the axino mass is expected to receive at least a contribution at the order of $m_3 \sim O(M^2_{SUSY}/f_a)$ at tree level in the spontaneously broken global SUSY. In the literature, a whole range of the axino mass was considered, and in fact it can be even much smaller [17, 21, 22, 23, 37, 38], or much larger, than the $M_{SUSY}$ [21]. Because of this strong model dependence, in cosmological studies one usually assumes the axino interactions from the $U(1)$ dependence, in cosmological studies one usually assumes the axino mass as a free parameter.

Recently, the issue of a proper definition of the axion and the axino was studied in the most general framework, including non-minimal Kähler potential [22]. In that study, axino mass is given by $m_3 = m_3/2$ for $G_A = 0$, where $G = K + \ln |W|^2$ and $G_A \equiv \partial G / \partial A$. For $G_A \neq 0$, axino mass depends on the details of the Kähler potential, and it was shown that the case given by Eq. (9) belongs to one of these examples. In gauge mediation scenario, the gaugino mass is the dominant axino mass parameter. In the case of gravity mediation, the axino mass is likely to be greater than the gravitino mass but one cannot rule out lighter axinos [22].

One lucid but often overlooked aspect of the axino is that its definition must be given at a mass eigenstate level. The coupling to the QCD sector given in the first line of Eq. (3) can plausibly be that of the axino but it does not give the axino mass. This is because the axino is connected to two kinds of symmetry breaking, the PQ global symmetry breaking and the SUSY breaking, which in general are not orthogonal to each other. The PQ symmetry breaking produces an almost massless pseudo-Goldststone boson (the axion), while SUSY breaking produces a massless goldstino. The massless goldstino is then absorbed into the gravitino to make it heavy via the super-Higgs mechanism. This raises the question of what the axino really is. This issue is shown in Fig. 1 taken from Ref. [22]. The axino must be orthogonal to the massless goldstino component. Therefore, for the axino to be present in a spontaneously broken supergravity theory, one has to introduce at least two chiral fields [30]. Even though its name refers to the axion-related QCD anomaly, one must select the component which is orthogonal to the goldstino. If there are two SM singlet chiral fields, this is easy since there is only one component left beyond the goldstino. However, if more than two chiral fields are involved in SUSY breaking, more care is needed to identify the orthogonal mass eigenstate. Among the remaining mass eigenstates beyond the goldstino, a plausible choice for the axino field is the component whose coupling to the QCD anomaly term is the biggest. For two initial chiral fields in Fig. 1, $\tilde{a}'$ has the anomaly coupling of Eq. (3) and hence the $\tilde{a}$ coupling to the QCD sector is equal to or smaller than those given in Eq. (3). The remaining coupling is the one to the $s = \pm 4$ components of a massive gravitino. Therefore, for the two initial chiral fields, axino cosmology must include the gravitino, as well, if $\tilde{a}'$ is not identical to $\tilde{a}$. The “leakage” is parametrized by the $F$-term of the initial axion multiplet $A$. With more than two initial chiral fields, the situation involves more mass parameters. One notable corollary of Ref. [22] is that the axino CDM relic abundance for $m_3 < m_{3/2}$ is an over-estimation if $A$ obtains the $F$-term.

The saxion mass is most likely of order $m_{3/2}$, and therefore saxions decay to SM particles. A notable cosmological implications of the saxion decay is known to be entropy production and a dilution of cosmic particles and the cosmic energy density. When applied to axion cosmology, the effect leads to some increase of the cosmological upper bound on $f_a$ [39, 42]. Axions and axinos produced from the decays of saxions can also affect the cosmic microwave background temperature anisotropy by contributing an additional relativistic component [43–47].

III. AXINOS IN COSMOLOGY

In this section we proceed to review axinos as relics in cosmology. We will treat axino mass as a free parameter ranging from eV to multi-TeV scales. A schematic representation of the strength and mass of axino DM is shown in Fig. 2, which is an updated version of the figure originally included in Ref. [48] and next modified in [18]. Marked schematically in the figure are also other well motivated candidates for DM, with HDM, WDM and CDM classes of candidates shown in red, pink and blue colors, respectively. As one can see, depending on axino mass and production mechanism, cosmic axinos may as well fall into one or actually more than one (e.g., CDM and WDM) population of relics, as discussed below.

FIG. 1: Axion (blue) and goldstino (red) multiplets. The axion direction $a$ is defined by the PQ symmetry and the goldstino ($g_{1/2}$) and axino ($\tilde{a}$) directions are defined by the fermion mass eigenvalues. The primed fields are not mass eigenstates.
itino regeneration and decay [52, 53]. If the axino mass upper bound of Ref. [24] is relaxed, the calculation follows the same line of logic which was used for the gravitino regeneration and decay [52, 53]. If the axino mass is between around a MeV to several GeV, the correct axino CDM density is obtained with $T_R$ less than about $5 \times 10^4$ GeV.

Thermal production of axinos is described by the Boltzmann equation where the first term on the r.h.s. corresponds to scatterings and the second one to decays [2-5, 6, 51].

$$\frac{dn_\tilde{a}}{dt} + 3 H n_\tilde{a} = \sum_{i,j} \langle \sigma(i + j \to \tilde{a} + \ldots) v_{rel} \rangle n_i n_j$$

$$+ \sum_i \langle \Gamma(i \to \tilde{a} + \ldots) \rangle n_i,$$

where $H$ denotes the Hubble parameter, $\sigma(i + j \to \tilde{a} + \ldots)$ is the scattering cross section for particles $i, j$ into final states involving axinos and $n_i$ stands for the number density of the $i$th particle species, while $\Gamma(i \to \tilde{a} + \ldots)$ is the corresponding decay width into final states involving axinos. Approximate solutions for the number density of relic axinos are given in Ref. [53].

In Fig. 3 (taken from Ref. [34]) where it was updated from Refs. [2, 3], we show the axino yield $Y$ resulting

A. Production of relic axinos

As stated in Introduction, there are two generic ways of producing relic axinos in the early Universe: thermal production from scatterings and decays of particles in thermal equilibrium, and non-thermal production from the decays of heavier particles after their freezout.

1. Thermal production

Primordial axinos decouple from thermal equilibrium at the temperature [24]

$$T_{\text{dec}} = 10^{11} \text{ GeV} \left( \frac{f_\tilde{a}}{10^{12} \text{ GeV}} \right)^2 \left( \frac{0.1}{\alpha_s} \right)^3.$$

They overclose the Universe unless their mass is bounded to be smaller than keV [24]. In inflationary cosmology, the population of primordial axinos is strongly diluted by cosmic inflation; however axinos are re-generated during reheating. When the reheating temperature $T_R$ is below the decoupling temperature, axinos do not reach the equilibrium level. However, axinos can be produced from the scatterings in thermal plasma, and the number density is proportional to $T_R$, in which case the keV mass upper bound of Ref. [24] is relaxed. The calculation follows the same line of logic which was used for the gravitino regeneration and decay [52, 53]. If the axino mass

![FIG. 2: Several well-motivated candidates of DM are shown. $\sigma_{\text{int}}$ is the typical strength of the interaction with ordinary matter. The red, pink and blue colors represent HDM, WDM and CDM, respectively. We updated the previous figures [18, 18] by including the sterile neutrino DM [49, 51].](image1)

![FIG. 3: Thermal yield of axino, $Y_{\text{axino}} \equiv n_\tilde{a}/s$, versus $T_R$. For strong interactions, the effective thermal mass (ETM) approximation (black) is used. We use the representative values of $f_\tilde{a} = 10^{11}$ GeV and $m_\tilde{q} = m_\tilde{q} = 1$ TeV. For comparison, we also show the HTL approximation (dotted blue/dark grey) and that of Strumia (green/light grey). We also denote the yield from squark (solid green/light grey) and gluino decay (dashed red), as well as out-of-equilibrium bino-like neutralino decay (dashed black) with $C_{\tilde{a}YY} = 8/3$.](image2)
from scatterings and decays involving from strong interactions in the KSVZ model. For different values of \( f_a \), the curves move up or down proportional to \( 1/f_a^2 \). The contribution from SU(2)_L and U(1)_Y interactions are suppressed by the gauge coupling since the cross section \( \sigma \propto \alpha^3 \). (For comparison, in Fig. 3 we also show the yield from bino-like neutralino decay after freezeout which is subdominant at larger \( T_R \) but becomes important at low \( T_R \).)

In the case of scatterings, we compare three different prescriptions for treating the infrared (IR) divergence that have been used in the literature. In Ref. 2 an effective thermal mass (ETM) approximation was used to regulate the infrared divergence from massless gluon. A more consistent way using a hard thermal loop (HTL) approximation was used in Ref. 3. The technique is, however, valid only in the regime of a small gauge coupling, \( \alpha_s \ll 1 \), which corresponds to the reheating temperature \( T_R \gg 10^9 \) GeV where, as we shall see, axino as DM is too warm. In Ref. 3 full resummed finite-temperature propagators for gluons and gluinons were used which extended the validity of the procedure down to \( T_R \gtrsim 10^9 \) GeV. However, the gauge invariance in the next leading order is not maintained. We conclude that there currently remains a factor of a few uncertainty in the thermal yield of axinos at high \( T_R \).

As noted in Ref. 37, when the temperature is higher than the mass \( M_Q \) of the PQ-charged and gauge-charged matter in the model which induce the axino-gaugino-gauge boson interaction, the interaction amplitude is suppressed by \( M_Q^2/T^2 \), in addition to the suppression by the PQ scale \( f_a \). This is most notable in the DFSZ model where the higgsino mass \( \mu \) is around the weak scale and the temperature is higher than this scale.

Axinos can also be efficiently produced through decays of thermal particles via the second term in Eq. 2 when \( T_R \) is comparable to the mass of the decaying particles. At larger \( T_R \), the contribution from decays becomes independent of temperature and in any case strongly subdominant relative to that from scatterings. At lower \( T_R \) the production becomes exponentially reduced due to the Boltzmann suppression factor for the population of the decaying particles in the thermal plasma.

One of the decay channels is a two-body decay of a gaugino into an axino and a gauge boson 2. The gaugino-axino-sfermion-sfermion interaction in Eq. 5 generates three-body decays of a gaugino into an axino and two sfermions, which is subdominant to the two-body decay. In the KSVZ model, an effective dimension-4 axino-quark-squark coupling is generated at a one loop and the squark decay can produce important amount of axinos 3. Those axino production from thermal gluons, neutralinos and squarks are shown in Fig. 3.

In the DFSZ framework, the dominant production contribution comes from scatterings involving SU(2)_L interactions and also from the decays of a higgsino into an axino and a Higgs boson due to a tree-level axino-Higgs-higgsino interaction term 52, 56 which is proportional to the higgsino mass \( \mu \). Axino production from higgsino decays in thermal equilibrium is comparable to, or, for large \( \mu \), can even be larger than, that from squark decays for which a coupling exists already at a tree level due to the \( c_2 \) interaction term, which is proportional to the mass of the quark. Generally, in the DFSZ framework axino production from thermal decays dominates 59 over that from scatterings 2, 35, 50, 57 which is suppressed by the quark mass at higher temperature 52. Therefore, the axino abundance is independent of the reheating temperature, if it is high enough compared to the higgsino mass.

2. Non-thermal production

Axinos can be produced from out-of-equilibrium decays of heavier non-thermal particles. In this case, the axino abundance is independent of the reheating temperature and its relic number density simply depends on
the number density and the decay modes of the decaying mother particles.

A particularly interesting case is axino production from decays of the next-to-lightest SUSY particle (NLSP) as the lightest ordinary SUSY particle (LOSP). Other ordinary superpartners that are heavier than the NLSP first cascade-decay to the NLSP which next freezes out from thermal plasma. The NLSP then decays to the axino LSP, however, this occurs much later on the cosmic time scale, with the lifetime around the order of a second or less \( f_a \). The non-thermal production of axinos from the NLSP decay is given by

\[
\Omega_{\tilde{a}}^{\text{NTP}} = \frac{m_{\tilde{a}}}{m_{\text{NLSP}}} \Omega_{\text{NLSP}} \simeq 2.7 \times 10^{10} \left( \frac{m_{\tilde{a}}}{100 \text{ GeV}} \right) Y_{\text{NLSP}}. \tag{9}
\]

Cosmological axino production can also proceed from decays of other non-thermal relics, e.g., an inflaton, moduli, a saxion, Q-balls \[58\], etc, however, specific implementations are strongly model dependent and will not be considered here. Axino production from NLSP decay therefore provides a conservative estimate of the relic density of axinos.

CDM relic axino production from bino decays was first considered in Ref. \[1\], and next in more detail in Ref. \[2\] along with thermal production and with cosmological constraints on CDM, and also on HDM and WDM axinos. Additional processes of thermal axino production from squark decays in thermal plasma were subsequently considered in Ref. \[2\], and next applied in Ref. \[4\] in an analysis of the Constrained MSSM (CMSSM) with a neutralino and a stau as NLSP, where also analogous tau-stau-axino couplings were obtained and applied. More recently, these couplings were re-derived in Ref. \[8\] in a full two-loop calculation including four-body hadronic decays and also for \( f_a \) larger than \( 10^{12} \text{ GeV} \). These couplings are smaller and not important for thermal production, but they are important for the non-thermal production when the stau is the NLSP. Colored NLSP was considered in \[60, 61\], however, their contribution is negligible due to their late freezeout.

In Fig. 4 (taken from Ref. \[3\] where it was updated from Refs. \[2, 3\]) we consider the total abundance of axinos (the sum of thermal and non-thermal production) to show an upper limit on the reheating temperature for a given axino mass from the total CDM abundance in the KSVZ model. Here we fix \( f_a \) = \( 10^{11} \text{ GeV} \) and the cases (I), (II) and (III) denote the different assumed values of NLSP abundance: \( Y_{\text{NLSP}} = 0, 10^{-10}, \) and \( 10^{-8} \), respectively. For a small axino mass, less than some 10 MeV, thermal production is dominant and depends on the reheating temperature. However, for a larger mass, NTP provides the dominant contribution. The regions above/to the right of the curves are excluded due to the overabundance of DM.

Now we proceed to review cosmological implications of the axino depending on its mass.

### B. Axinos in cosmology

If axinos are produced very late, at the time of larger than one second after the Big Bang, from decays of an NLSP, the injection of high energetic hadronic and electromagnetic particles can affect the abundance of light elements produced during Big Bang Nucleosynthesis (BBN) \[62, 63\]. For charged NLSP the constraints are even stronger \[64\]. This provides an upper bound on the abundance or the lifetime of the NLSP; especially for large values of \( f_a \) the constraint is severe as discussed in \[2, 4, 8\]. However, as long as \( f_a \lesssim 10^{12} \text{ GeV} \), the lifetime of bino-like NLSP in a mass range of a few hundred GeV is less than 1 second, and for the stau is similar, which makes axino DM free from the BBN problem.

Constraints from BBN may also be applicable when the axino is heavy and unstable, in which case it decays into lighter MSSM particles and SM particles \[40, 65\].

All these scenarios depend primarily on one parameter, the axino mass, and we discuss them in the order of increasing mass.

#### 1. Axino as HDM

The axino with mass in the eV and sub-eV range produced in decays of an out-of-equilibrium sub-GeV photino was considered early on in Ref. \[23\]. This case of axino can be considered as HDM. A primordial thermal population of axinos decoupled at \( \sim 10^{11} \text{ GeV} \) and was diluted away by cosmic inflation. In R-parity conserving models the photino lifetime is a function of \( f_a \) and thus the HDM axino abundance depends on it \[23\]. This is relevant both in the standard Big Bang and the inflationary cosmology. This is because the photino abundance is calculated from the photino decoupling temperature which is below the reheating temperature after inflation and hence the photino abundance is independent of the cosmological scenarios.

A related sub-eV mass fermion useful for DM is gravitino for \( m_{3/2} \lesssim 1 \text{ keV} \) \[60\]. Since the decoupling temperature of gravitino is close to the Planck mass, primordial gravitinos were diluted out in the inflationary universe. However, axinos can decay to sub-eV gravitinos \[67\]. Sub-eV gravitinos are possible in the gauge mediated SUSY breaking scenario. In the unstable axino case sub-eV gravitinos can become HDM in the universe, which was called ‘axino-gravitino cosmology’ \[68\].

#### 2. Axino as WDM

Rajagopal, Turner and Wilczek \[24\] considered axinos in the keV range. They obtained the axino mass bound \( m_{\tilde{a}} \lesssim 2 \text{ keV} \) for axinos to be WDM. However, this bound can be relaxed if the reheating temperature \( T_R \) after inflation is much lower than the PQ symmetry breaking scale.
3. Axino as CDM

Covi, Kim and Roszkowski \cite{1} considered CDM axinos. Axinos with mass higher than 10 keV for TP or 10 MeV for NTP are non-relativistic enough to be CDM, as marked in Fig. 3. Even though they are relativistic at the time of production, their velocity is red-shifted with the expansion of the universe and they have small free streaming length at the time of structure formation.

For CDM axinos, relatively low reheating temperatures are preferred, as shown in Fig. 4. Therefore, axino is a good candidate for DM in models of thermal inflation which takes place at a late time and naturally predicts a low reheating temperature \cite{2,3,72,74}.

Allowing for $R$-parity violation, axino DM can decay with a lifetime much longer than the age of the universe. Photons from axino decay can be a DM signature and can explain some astrophysical anomalies \cite{77,81}.

Scenarios with a mixed axion-axino population of CDM have also been considered. In Fig. 3 we show a numerical result from a scan over the MSSM with 19 model parameters with a SUSY axion model \cite{11}. In the relic density versus PQ scale plane the dominant axion DM is shown in red and axino DM in blue.

4. Cosmology with superheavy axino

As a digression, we note that, for super-heavy axino, it is the neutralino that most likely is the LSP, in which case the neutralino population from heavy axino decays could constitute CDM \cite{63}. This case is particularly interesting if the axino mass is greater than the gravitino mass \cite{22} and the data from PAMELA \cite{82} and recently from AMS-02 \cite{83} which may imply a TeV scale of WIMP mass, if the WIMP is CDM. In \cite{84} it was shown that the superheavy axino can account for TeV-scale cosmic ray positrons produced as decay products of, for example, an NMSSM singlino $N$ to $\tilde{e}^+\tilde{e}$ plus $e^+\bar{e}$, but not to antiprotons. This is possible in a string derived flipped SU(5) grand unification model \cite{85}. In this case, $\tilde{e}$ eventually decays to LSP neutralino plus SM particles. Of course, the final population of the LSP is not enough to account for the present CDM density in the decaying DM scenario but the mother singlino density accounts for most of the CDM density, which is given by the non-thermal production from superheavy axino decays. Heavy axino decays to singlino and the singlino decays to positron and selection, and the selectron finally decays to the LSP neutralino. This $N$–WIMP decay scenario was proposed to explain the property that the PAMELA data does not contain any large excess of antiprotons but a significant flux of positrons \cite{86}. On the other hand, if the decaying DM scenario is ruled out in favor of the scattering production of positrons, the superheavy axino case can be predominantly decaying to the LSP plus SM particles, for example in the MSSM extended by the PQ symmetry with $S$ and $\bar{S}$ of Eq. (6), not introducing extra singlets $N$-type in the NMSSM.

The decay of heavy axinos provides a non-thermal population of LSP DM such as gravitinos or neutralinos \cite{63}. Therefore, the abundance of axino before decay is also constrained and gives a quite strong limit on the reheating temperature \cite{57}. The neutralino DM scenario from heavy axino decays was further studied in \cite{12,13} where numerical results for the production of mixed axion and neutralino DM were presented. The effect of saxion production and decay was also considered and the entropy production allowed the dilution of pre-existing particles. The saxion decay can also produce a sufficient population of relativistic axions which can be identified as dark radiation in CMB observations \cite{43,49}. 

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure5.png}
\caption{The relic density of mixed axion/axino dark matter versus the axino mass in the scan with 19 parameters in the MSSM with SUSY axion multiplet. Figure taken from Ref. \cite{11}.}
\end{figure}
IV. CDM AXINO PRODUCTION AT COLLIDERS

The exceedingly weak strength of axino interactions makes axino detection in direct DM search experiments as well as at collider experiments rather hopeless. However, at the LHC a signal of axino DM in $R$-parity conserving scenarios may show up in the properties of the LOSP.

With axino DM, the LOSP typically decays with the lifetime of less than one second. Therefore, even if it is unstable on cosmological time scales, it is practically stable inside a collider detector, although it will decay outside it. One spectacular signature of non-standard DM would be to detect an electrically charged (stau) or colored (stop) massive particle as a LOSP that would be seemingly stable in a detector.

However, this would still leave the question open whether the LSP is the axino or for example the gravitino. Measuring the LOSP lifetime would not be sufficient. If the LOSP is an electrically charged particle, such as stau, then this could be possible with the analysis of three-body decay of the charged NLSP slepton into the corresponding lepton, the LSP and a photon. One way would be to measure the spin of the LSP through the polarization of the final-state lepton and photon [88, 89]. Another is to measure the branching ratio and the angular distributions of the decay products in the three-body decays of the LOSP [90].

Even if a stable massive neutral particle, such as the neutralino, is detected at the LHC through a missing energy signature, this will still not guarantee its DM nature conclusively since it could decay outside of detectors. A signal indicating the same mass would therefore be necessary in direct and indirect DM searches. The absence of one could indicate the existence of lighter stable particle, such as the axino or the gravitino, playing the role of DM. In this way, with the axino as DM, a large region of the (C)MSSM which would be forbidden with neutralino DM becomes allowed [4].

With much luck, one could in principle measure at the LHC enough quantities to estimate, at least roughly, the relic abundance of the LOSP and compare it with the cosmological value. If the two were radically different, again this would indicate a non-standard choice for DM. A model-independent study of axino DM from collider information was performed in Ref. [91]. Based on the thermal and non-thermal production of axinos and the cold DM density,

$$\Omega_{\tilde{a}} h^2(T_R, m_{\tilde{a}}, m_{\tilde{\chi}^0_2}, \ldots) + \frac{m_{\tilde{a}}}{m_{\tilde{\chi}^0_2}} \Omega_{\text{LOS}} h^2 = \Omega_{\text{CDM}} h^2 \simeq 0.1,$$

and from the information inferred on the relic density of LOSP at colliders, one could attempt to determine a relation between the reheating temperature and the mass of axino. This is shown in Fig. 6 (taken from Ref. [91]).

V. CONCLUSIONS

Axino is an intriguing candidate for DM in SUSY with the axion solution of the strong $CP$ problem. The mass and interactions are highly dependent on the axion model and its coupling to SM particles are suppressed by the PQ scale $f_a$. However, SUSY breaking alters the connection so that the axino becomes massive and can also have mixing with the goldstino [22]. Axinos can be produced in the thermal plasma in the early Universe in scatterings and decays of heavier particles and in out-of-equilibrium decays of heavier particles. Axino relic density can be obtained to coincide with the correct density of DM. Axinos can be HDM, WDM or CDM depending on their mass range and production mechanism. There could also be two populations of axinos, for example one warm and one cold, from different production modes. In the case of heavy axinos, neutralinos can be produced from decays of axinos and extend the available parameter space and help to explain the positron anomaly in PAMELA and AMS-02. The identification of LOSP and its interactions at colliders might provide a relationship between
the reheating temperature and the axino mass, and also glimpses into the earliest moments of the Universe after inflation if the DM is made up of thermally produced axinos.

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[1] L. Covi, J. E. Kim and L. Roszkowski, Phys. Rev. Lett. 82 (1999) 4180 [hep-ph/9905212].
[2] L. Covi, H. B. Kim, J. E. Kim and L. Roszkowski, JHEP 0105 (2001) 033 [hep-ph/0101009].
[3] L. Covi, L. Roszkowski and M. Small, JHEP 0207 (2002) 023 [hep-ph/0206119].
[4] L. Covi, L. Roszkowski, R. Ruiz de Austri and M. Small, JHEP 0406 (2004) 003 [hep-ph/0402240].
[5] A. Brandenburg and F. D. Steffen, JCAP 0408 (2004) 008 [hep-ph/0405158].
[6] A. Strumia, JHEP 1006 (2010) 036 [arXiv:1003.5847 [hep-ph]].
[7] For a recent review, see, L. Covi and J. E. Kim, New J. of Phys. 11 (2009) 105003 [arXiv: 0902.0769 [astro-ph/CQ]].
[8] A. Freitas, F. D. Steffen, N. Tajuddin, and D. Wyler, Phys. Lett. B679 (2009) 270 [arXiv:0904.3218 [hep-ph]]; A. Freitas, F. D. Steffen, N. Tajuddin, and D. Wyler, Phys. Lett. B682 (2009) 193 [arXiv:0909.3293 [hep-ph]].
[9] E. J. Chun, H. B. Kim, K. Kohri, D. H. Lyth, JHEP 0803 (2008) 061 [arXiv:0801.4108 [hep-ph]]; S. Kim, W. J. Park and E. D. Steward, JHEP 0901 (2009) 015 [arXiv:0807.3607 [hep-ph]].
[10] H. Baer and A. D. Box, Eur. Phys. J. C 68 (2010) 523 [arXiv:0910.0333 [hep-ph]]; H. Baer, R. Dermisek, S. Rajagopalan, H. Summy, JCAP 1007 (2010) 014 [arXiv:1004.3297 [hep-ph]]; H. Baer, S. Kraml, A. Lessa, S. Sekmen, JCAP 1011 (2010) 040 [arXiv:1009.2950 [hep-ph]]; H. Baer, S. Kraml, A. Lessa, S. Sekmen, JCAP 1104 (2011) 039 [arXiv:1012.3769 [hep-ph]].
[11] H. Baer, A. D. Box, H. Summy, JHEP 1010 (2010) 023 [arXiv:1005.2215 [hep-ph]].
[12] H. Baer, A. Lessa, S. Rajagopalan and W. Sreeha-wong, JCAP 1106 (2011) 031 [arXiv:1105.5413 [hep-ph]]; H. Baer and A. Lessa, JHEP 1106 (2011) 027 [arXiv:1104.4807 [hep-ph]].
[13] H. Baer, A. Lessa and W. Sreeha-wong, JCAP 1201 (2012) 036 [arXiv:1110.2491 [hep-ph]].
[14] J. U Kang and G. Panotopoulos, JHEP 0805 (2008) 036 [arXiv:0805.0553 [hep-ph]].
[15] H. P. Nilles and S. Raby, Nucl. Phys. B198 (1982) 102.
[16] K. Tamvakis and D. Wyler, Phys. Lett. B112 (1982) 451.
[17] J. M. Frere and J. M. Gerard, Lett. Nuovo Cim. 37, (1983) 135.
[18] For a recent review, see, J. E. Kim and G. Carosi, Rev. Mod. Phys. 82 (2010) 557 [arXiv: 0807.3125 [hep-ph]].
[19] For a review, see, H. P. Nilles, Phys. Rep. 110 (1984) 1.
[20] E. J. Chun, J. E. Kim and H. P. Nilles, Phys. Lett. B287 (1992) 123.
[21] E. J. Chun and A. Lukas, Phys. Lett. B 357 (1995) 43 [hep-ph/9503223].
[22] J. E. Kim and M. -S. Seo, Nucl. Phys. B864 (2012) 296 [arXiv:1204.5495 [hep-ph]].
[23] J. E. Kim, A. Masiero and D. V. Nanopoulos, Phys. Lett. B139 (1984) 346.
[24] K. Rajagopal, M.S. Turner and F. Wilczek, Nucl. Phys. B358 (1991) 447.
[25] K.-J. Bae, J.-H. Huh and J. E. Kim, JCAP 0809 (2009) 005 [arXiv:0806.0497 [hep-ph]].
[26] J. E. Kim, Phys. Rev. Lett. 43 (1979) 103; M. A. Shifman, V. I. Vainstein, V. I. Zakharov, Nucl. Phys. B166 (1980) 4933.
[27] M. Dine, W. Fischler and M. Srednicki, Phys. Lett. B104 (1981) 199; A. P. Zhitnitskii, Sov. J. Nucl. Phys. 31 (1980) 260.
[28] J. E. Kim, Phys. Rev. D58 (1998) 055006 [arXiv:hep-ph/9802220].
[29] K.-S. Choi, I.-W. Kim and J. E. Kim, JHEP 0703 (2007) 116 [arXiv:hep-ph/0612107].
[30] J. E. Kim, Phys. Lett. B136 (1984) 78.
[31] J. E. Kim, Phys. Rev. Lett. 111 (2013) [arXiv:1303.1822 [hep-ph]].
[32] S. M. Barr and D. Seckel, Phys. Rev. D46, 1992 (539); M. Kamionkowski and J. March-Russell,Phys. Lett. B282 (1992) 137 [hep-th/9202003]; R. Holman, S. D. H. Hsu, T. W. Kephart, E. W. Kolb, R. Watkins, and L. M. Widrow, Phys. Lett. B282 (1992) 132 [hep-ph/9203206]; S. Ghigna, M. Lusignoli and M. Roncadelli, Phys. Lett. B283 (1992) 278; B. A. Dobrescu, Phys. Rev. D55, 1997 (5826) [hep-ph/9609221].
[33] P. Sikivie, Phys. Rev. Lett. 48 (1982) 1156.
[34] K.-Y. Choi, L. Covi, J. E. Kim and L. Roszkowski, JHEP 1204 (2012) 106 [arXiv:1108.2282 [hep-ph]].
[35] K. J. Bae, K. Choi and S. H. Im, JHEP 1108 (2011) 065 [arXiv:1106.2452 [hep-ph]].
[36] J. F. Nieves, Phys. Rev. D33 (1986) 1762.
[37] T. Goto and M. Yamaguchi, Phys. Lett. B276 (1992) 103.
[38] P. Moxhay and K. Yamamoto, Phys. Lett. B151 (1985) 363.
[39] J. E. Kim, Phys. Rev. Lett. 67 (1991) 3465.
[40] S. Chang and H. B. Kim, Phys. Rev. Lett. 77 (1996) 591.
[41] M. Kawasaki, K. Nakayama and M. Senami, JCAP **0803** (2008) 009 [arXiv:0711.3083 [hep-ph]].
[42] M. Kawasaki, N. Kitajima and K. Nakayama, Phys. Rev. **D83** (2011) 123521 [arXiv:1104.1262 [hep-ph]].
[43] K. Choi, K.-Y. Choi and C. S., Phys. Rev. **D86** (2012) 083529 [arXiv:1208.2496 [hep-ph]].
[44] J. Hasenkamp, Phys. Lett. **B707** (2012) 121 [arXiv:1107.4319 [hep-ph]].
[45] P. Graf and F. D. Steffen, JCAP **1302** (2013) 018 [arXiv:1208.2951 [hep-ph]].
[46] K. J. Bae, H. Baer and A. Lessa, JCAP **1304** (2013) 041 [arXiv:1301.7428 [hep-ph]].
[47] P. Graf and F. D. Steffen, arXiv:1302.2143 [hep-ph].
[48] L. Roszkowski, *Pramana* **62** (2004) 389 [arXiv:hep-ph/0404052].
[49] A. Boyarsky, A. Neronov, O. Ruchayskiy and M. Shaposhnikov, Mon. Not. Roy. Astron. Soc. **370** (2006) 213 [astro-ph/0512509].
[50] A. Boyarsky, O. Ruchayskiy and M. Shaposhnikov, Ann. Rev. Nucl. Part. Sci. **59** (2009) 191 [arXiv:0901.0011 [hep-ph]].
[51] K. N. Abazajian, M. A. Acero, S. K. Agarwalla, A. A. Aguilar-Arevalo, C. H. Albright, S. Antusch, C. A. Arguelles and A. B. Balantekin et al., arXiv:1204.3739 [hep-ph].
[52] J. R. Ellis, J. E. Kim, D. V. Nanopoulos, Phys. Lett. **B145** (1984) 181. E. Holtmann, M. Kawasaki, K. Kohri, and T. Moroi, Phys. Rev. **D60** (1999) [hep-ph/9805405].
[53] M. Kawasaki and T. Moroi, Prog. Theor. Phys. **93** (1995) 879 [hep-ph/9403061].
[54] M. E. Gomez, S. Lola, C. Pallis and J. Rodriguez-Quintero, JCAP **0901** (2009) 027 [arXiv:0809.1859 [hep-ph]].
[55] K. Choi, K. Hwang, H. B. Kim, T. Lee, Phys. Lett. **B467** (1999) 211-217. [hep-ph/9902291].
[56] E. J. Chun, Phys. Rev. **D84** (2011) 043509 [arXiv:1104.2219 [hep-ph]].
[57] K. J. Bae, E. J. Chun and S. H. Im, JCAP **1203** (2012) 013 [arXiv:1111.5962 [hep-ph]].
[58] L. Roszkowski, O. Seto, Phys. Rev. Lett. **98** (2007) 161304. [hep-ph/0608013].
[59] P. A. R. Ade et al. [Planck Collaboration], arXiv:1303.5076 [astro-ph.CO].
[60] C. F. Berger, L. Covi, S. Kraml and F. Palorini, JCAP **0810** (2008) 005 [arXiv:0807.0211 [hep-ph]].
[61] L. Covi, J. Hasenkamp, S. Pokorski and J. Roberts, JHEP **0911** (2009) 003 [arXiv:0908.3399 [hep-ph]].
[62] K. Jedamzik, Phys. Rev. **D70** (2004) 063524 [arXiv:astro-ph/0402344].
[63] M. Kawasaki, K. Kohri and T. Moroi, Phys. Lett. **B625** (2005) 0421 [arXiv:hep-ph/0501287].
[64] M. Pospelov, Phys. Rev. Lett. **98** (231301) 2007 [hep-ph/0605215].
[65] K.-Y. Choi, J. E. Kim, H. M. Lee, and O. Seto, Phys. Rev. **D77**, 123501 (2008) [arXiv:0801.0491[hep-ph]].
[66] H. Pagels and J. R. Primack, Phys. Rev. Lett. **48** (1982) 223.