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Axinos as dark matter

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Abstract: Supersymmetric extensions of the Standard Model that incorporate the axion solution to the strong CP problem necessarily contain also the axino, the fermionic partner of the axion. In contrast to the neutralino and the gravitino, the axino mass is generically not of the order of the supersymmetry-breaking scale and can be much smaller. The axino is therefore an intriguing candidate for a stable superpartner. In a previous letter [1] it was shown that axinos are a natural candidate for cold dark matter in the Universe when they are generated non-thermally through out-of-equilibrium neutralino decays. Here, we extend the study of non-thermal production and include a competing thermal production mechanism through scatterings and decays of particles in the plasma. We identify axino masses in the range of tens of MeV to several GeV (depending on the scenario) as corresponding to cold axino relics if the reheating temperature $T_{\rm R}$ is less than about 5×10^4 GeV. At higher $T_{\rm R}$ and lower mass, axinos could constitute warm dark matter. In the scenario with axinos as relics, the gravitino problem finds a natural solution. The lightest superpartner of the Standard Model spectrum remains effectively stable in high-energy detectors but may be either neutral or charged. The usual highly restrictive constraint $\Omega_\chi h^2 \lesssim 1$ on the relic abundance of the lightest neutralino becomes void.

Keywords: Cosmology of Theories beyond the SM, Dark Matter, Supersymmetric Standard Model.

Contents

1. Introduction

Among several possible ways of addressing theoretical puzzles of the Standard Model, two are generally believed, for quite independent reasons, to be very compelling. Softly-broken low-energy supersymmetry (SUSY) seems to be the most attractive way of solving the hierarchy problem and of linking the electroweak scale with physics around the Planck scale. Moreover, it has the generic property of contributing only small corrections to electroweak parameters and also of predicting a light Higgs boson, both in agreement with the outcome of current precision measurements.

The most compelling way of resolving the strong CP problem of QCD seems to be provided by invoking the Peccei and Quinn (PQ) mechanism [2]. There, the CP-violating θ-term in the QCD Lagrangian, which is experimentally required to be excessively tiny, is replaced by a term involving a new fundamental pseudo-scalar field, the axion. This is achieved by introducing a global, chiral U(1) symmetry group, which is spontaneously broken at some high energy scale $f_a \sim 10^{11}$ GeV. The QCD anomaly breaks this chiral U(1) symmetry at the one-loop level, and hence the axion becomes not a true Goldstone boson but a pseudo-Goldstone boson with a tiny mass of order $\Lambda^2_{\rm QCD}/f_a$ [3]. As the lightest pseudo-scalar consistent with the known particle phenomenology, the axion can be very important in cosmology and astrophysical processes. Axion physics and cosmology are very rich fields which have been thoroughly explored in the literature [4].

JHEP05(2001)033

Axions can also be the dark matter (DM) in the Universe. Indeed, being extremely light and feebly interacting, and having a lifetime of order $> 10^{40}$ years, axions can be considered as stable for all practical purposes. Therefore, if axions are copiously produced in the early Universe, they can contribute substantially to the relic mass density. Since they are spin-zero particles, we can consider hot axions and cold axions separately. Hot axions are produced when the PQ symmetry is broken at $T \leq f_a$. Such hot axions are diluted by inflation if the reheating temperature after inflation is less than about $10^9\,\mathrm{GeV}$. At low temperature it is difficult to produce hot axions because of their tiny interaction strength. Since the axion potential is extremely flat, the axion vacuum expectation value (VEV) $\langle a \rangle$ will not move until the temperature of the Universe falls below $\sim 1\,\mathrm{GeV}$. These coherent axions produced around the QCD phase transition have very small kinetic energy and are therefore cold. They can constitute the DM of the Universe if $f_a \sim 10^{12}\,\mathrm{GeV}$ [5].

An inevitable prediction of combining the two well-motivated and independent hypotheses of axions and supersymmetry is the existence of the SUSY partner of the axion, the axino. Being massive and both electrically and color neutral, axinos are an intriguing possibility for the lightest supersymmetric particle (LSP) and a WIMP. In the presence of R-parity conservation there may be dramatic cosmological implications of the existence of such stable massive particles. In particular, they may contribute substantially to the relic mass density and may constitute the main component of the DM in the Universe. This is the possibility that we will explore in this study.

There are two other SUSY particles with well-known, and much-studied, cosmological properties which, ever since the early days of low-energy SUSY, have been considered for DM candidates. Perhaps the most popular of them is the lightest neutralino, χ, which, if stable, often provides the desired amount of relic abundance for natural ranges of other superpartner masses. Furthermore, since its characteristic interaction strengths are often of a sizeable fraction of the electroweak interactions, the neutralino has a good chance of being detected in high energy colliders and, if it indeed constitutes the dark matter, in WIMP dark matter searches.

The other traditional candidate, the gravitino, $\widetilde{G}$, arises by coupling SUSY to gravity in supergravity or superstring scenarios. It has been long known that gravitinos could be copiously produced in thermal processes in the early Universe and often suffer from the well-known "gravitino problem" [7] of excessive destruction of light elements and of contributing too much to the energy density. These problems can be avoided if the reheating temperature after inflation is sufficiently low, $T_{\mathrm{R}} \lesssim 10^9\,\mathrm{GeV}$ [6]–[8]. Therefore, the gravitino remains a possibility for dark matter. (For recent work see refs. [9]–[11].)

In contrast to the neutralino and gravitino cases, axinos have, undeservedly, attracted much less attention in the literature. This is even more surprising in view of the fact that axinos possess some distinctive properties that lead to important

JHEP05(2001)033

cosmological implications. In particular, in contrast to the case of the gravitino as well as the neutralino and other ordinary superpartners, the mass of the axino is generically *not* of the order of the SUSY-breaking scale and, as we will see later, can be much lower [12]–[14]. This alone makes the axino an intriguing candidate for the LSP and dark matter.

One important exception was provided by a comprehensive study of Rajagopal, Turner and Wilczek [13]. These authors considered both light ($\sim$ keV range) and more massive ($\sim$ GeV) axinos and studied their cosmological implications as either the LSP or the next-to-LSP (NLSP). They derived an upper bound $m_{\tilde{a}} < 2\,\mathrm{keV}$ on the mass of stable, "primordial" axinos which would hold in the absence of inflation. They also concluded that, in a class of interesting models, the axino could have a mass satisfying this bound and would constitute so-called *'warm'* dark matter. Such light axinos were studied also by other early papers [15, 16].

While warm DM has some interesting properties (see, e.g., ref. [17]), it has been the standard cosmological lore of the last several years that the invisible dynamical component of the mass–energy density of the Universe is probably predominantly in the form of *cold* dark matter (CDM) [18]. The paradigm has been based on good agreement of numerical simulations of the formation of large scale structures with observations [18]. More recent studies have revealed some possible problems with predictions of the standard CDM theory on sub-galactic scales, which still need to be clarified. Nevertheless, we believe that it is still well-justified to assume that most of the DM in the Universe is cold.

In a previous paper [1], we pointed out that axinos can naturally form cold dark matter. One way of producing such CDM axinos is through out-of-equilibrium decays of a heavy enough superpartner. Because axino couplings to other particles are suppressed by $1/f_a$, as the Universe cools down, all heavier SUSY partners first cascade-decay to the lightest ordinary SUSY particle (LOSP), e.g., the lightest neutralino. (By "ordinary" we mean any of the superpartners of the Standard Model particles.) The LOSPs then freeze out of thermal equilibrium and subsequently decay into axinos. If LOSP is a neutralino of tens of GeV or heavier, its lifetime, which scales like $1/m_\chi^3$, is often much shorter than $1\,\mathrm{sec}$ and the decay takes place before Big Bang Nucleosynthesis (BBN) [1].

It is worth mentioning that the neutralino will nevertheless appear to be stable in high-energy colliders since its lifetime will typically be significantly larger than $\sim 10^{-7}\,\mathrm{sec}$. The same would be true if the LOSP carried electric charge. In this case the LOSP would appear in a detector as a stable, massive, charged state [13].

Furthermore, since after freeze-out all neutralinos convert into axinos, one obtains a simple expression for the axino relic abundance $\Omega_{\tilde{a}} h^2 = (m_{\tilde{a}}/m_\chi)\,\Omega_\chi h^2$ [1]. (Here $\Omega_{\tilde{a}} = \rho_{\tilde{a}}/\rho_{\mathrm{crit}}$, where $\rho_{\tilde{a}}$ is the relic density of axinos and ρ_{crit} is the critical density, h is the dimensionless Hubble parameter, $m_{\tilde{a}}$ and m_χ are the respective masses of the axino and the neutralino, and $\Omega_\chi h^2$ is the abundance the neutralinos

JHEP05(2001)033

would have had today.) It follows that cases where $\Omega_{\tilde{a}}h^2 \sim 1$ will typically correspond to supersymmetric configurations for which $\Omega_\chi h^2 \gg 1$. In other words, cases traditionally thought to be excluded by the condition $\Omega_\chi h^2 \lesssim 1$ will now be readily allowed (actually, even favored!). They would also typically correspond to larger superpartner masses, which may be unwelcome news fo collider searches, but appears to be favored by gauge coupling unification [19, 20] as well as flavor and proton decay constraints [21, 20].

In addition to the non-thermal production (NTP) discussed above, axinos can also be generated through thermal production (TP), namely via two-body scattering and decay processes of ordinary particles still in thermal-bath. (We stress that, despite the name of the process, the resulting axinos will typically be already out of thermal equilibrium because of their exceedingly tiny couplings to ordinary matter, except at very large $T_{\rm R}$. They will nevertheless be produced in kinetic equilibrium with the thermal-bath and thus their momenta will have a thermal spectrum inherited from the scattering particles in the plasma.) While the analogous process of gravitino production has been studied in the literature [8, 9], this has not yet been done in the case of axinos.

In this study we aim to provide a detailed exploration of both TP and NTP mechanisms of producing axinos through processes involving ordinary superpartners, and to compare their relative effectiveness. Both mechanisms are quite generic, but specific results will be somewhat model-dependent. For definiteness, we will be working in the context of the Minimal Supersymmetric Standard Model (MSSM). Also, while any ordinary superpartner (neutral or not) could in principle be the LOSP, we will concentrate on the case of the neutralino. We will also, for reference, often assume $f_a = 10^{11}\,\mathrm{GeV}$, although it will be relatively straightforward to rescale our results for other values of f_a. Furthermore, since the axions and axinos are produced through different mechanisms and both can be CDM, one could consider the possibility that both contribute sizeably to the CDM relic abundance. While this scenario remains a viable possibility, we will, for definitness, explicitly assume that it is the axino that dominates the Universe.

We will show that, unsurprisingly, TP is more important at larger values of the reheating temperature $T_{\rm R}$. On the other hand, NTP will dominate at rather low (but not unreasonably low) values $T_{\rm R} \lesssim 5 \times 10^4\,\mathrm{GeV}$. This is also, broadly, the range of $T_{\rm R}$ for which other constraints will be satisfied.

Clearly we are most interested in the cases where the relic density of axinos $\rho_{\tilde{a}}$ will be close to the critical density $\rho_{\rm crit}$. More precisely, we will require $0.1 \lesssim \Omega_{\tilde{a}}h^2 \lesssim 0.3$. This condition will provide a strong constraint on the scenario. In particular, in the NTP case of axino production through $\sim \mathcal{O}(100\,\mathrm{GeV})$ bino decay, it will imply $m_{\tilde{a}} \gtrsim \mathcal{O}(10\,\mathrm{MeV})$ to $\mathcal{O}(1\,\mathrm{GeV})$, depending on whether we allow the SUSY-breaking scale $M_{\rm SUSY}$ to take very large values (tens of TeV) or require $M_{\rm SUSY} \lesssim 1\,\mathrm{TeV}$, respectively.

JHEP05(2001)033

Axinos will initially be nearly always relativistic (unless they are nearly degenerate with the parent neutralinos), hence they have to be heavy enough to "cool down" and become non-relativistic, or cold, by the epoch of matter dominance. We find that this leads to a somewhat weaker condition $m_{\tilde{a}} \gtrsim \mathcal{O}(100\,\mathrm{keV})$.

A similar bound arises from requiring that the relativistic axinos do not contribute too much to the energy density of radiation during BBN, since their production will take place near the time of BBN. Furthermore, a significant fraction of neutralino decays will proceed into $q\bar{q}$-pairs (+axino) through an intermediate photon or Z-exchange. The resulting hadronic showers may cause excessive destruction of light elements. We find that this can be easily avoided but that the specific bounds are strongly model-dependent. It would suffice to simply assume a large enough neutralino mass m_χ for the decay to take place before BBN. We find that such a crude requirement would be unnecessarily over-constraining. Nevertheless, for smaller m_χ the resulting lower bound on $m_{\tilde{a}}$ will indeed be typically much stronger than the above bounds. For example, for a bino of 60 GeV one finds $m_{\tilde{a}} \gtrsim 360\,\mathrm{MeV}$. However, increasing m_χ to 150 GeV removes the bound altogether.

We should mention other non-thermal mechanisms for re-populating the Universe with axinos and/or other relics such as gravitinos and moduli. First, it has recently been claimed [22] that gravitationally interacting particles can be copiously produced through inflaton field oscillations in the reheating or preheating process, and that non-thermal production of gravitinos can be much more efficient than through thermal processes.[1] This would further aggravate the gravitino problems with nucleosynthesis, thus leading to a much more severe constraint on the reheating temperature, sometimes as low as $T_{\mathrm{R}} \lesssim 10^5\,\mathrm{GeV}$, depending on the model. We do not expect the mechanism to contribute to the production of axinos in any significant way, especially in the regime of low T_{R} mentioned above. Furthermore, such mechanism would strongly depend on the inflationary model and on the interactions between the inflaton, and the axino and it would seem difficult, if not impossible, to obtain the required abundance of DM.

Second, axinos can be produced in gravitino decays, along with their non-SUSY partner, the axion [25]. It is worth emphasizing that the gravitino problem can be easily resolved if the axino is the LSP and the hierarchy of masses $m_{\mathrm{LOSP}} > m_{\widetilde{G}} > m_{\tilde{a}}$ is assumed. The process will take place long after nucleosynthesis ($t \sim 10^8\,\mathrm{sec}$) but decay products will now be completely harmless. The lightest ordinary superpartner will also decay directly to axinos, thus by-passing the dangerous late decays to gravitinos and energetic photons. In the relatively low $T_{\mathrm{R}} \lesssim 10^5\,\mathrm{GeV}$ regime, where NTP of axinos dominates, there is not even the gravitino problem anymore [22]. However, the mechanism resolves the gravitino problem even at much larger $T_{\mathrm{R}} \lesssim 10^{15}\,\mathrm{GeV}$ studied in ref. [25]. In other words, it seems that the gravitino

[1] This claim has very recently been disputed [23]. See also ref. [24].

JHEP05(2001)033

problem can be resolved, assuming the above hierarchy, in a model-independent way. At large $T_{\rm R}$ the axino would not contribute to cold DM. Instead, it would be warm DM while cold DM could be provided by, for instance, the axion.

It is clear that in discussing cosmological implications of axinos one quantity of crucial importance is the axino mass itself. Unfortunately, unlike the axino coupling, this quantity is rather poorly determined and is strongly model-dependent. As mentioned above, a distinctive feature of axinos is that typically their mass is not set by the SUSY-breaking scale and therefore can be much lower. In fact, it can span a wide range from $\sim$ eV to $\sim$ GeV [14], depending on the model. In this study we will therefore treat $m_{\widetilde{a}}$ as a basically free parameter.

Finally, we briefly mention the saxion — the $R = 1$ spin-zero component of the axion supermultiplet.[2] As with other scalars, the mass of the saxion arises from a soft term and is typically of the order of the SUSY-breaking scale. Cosmological properties of saxions can be quite important [26, 27]. In particular, saxions decay relatively fast which may cause significant entropy production which would reheat the Universe and lead to other consequences. However, these effects will not be relevant to us if the saxion mass is below 1 TeV which we will assume here. We will occasionally comment on the saxion below when relevant for axino cosmology.

The plan of the paper is as follows. In section 2 we discuss relevant axion and axino properties. In particular, we concentrate on the axino mass and couplings. In section 3 the thermal production of axinos is analyzed in detail. section 4 summarizes and extends the previously obtained results in the case of NTP [1]. Astrophysical and cosmological constraints are discussed in section 5. In section 6 the two production mechanisms are compared and further discussed. Finally, in section 7, implications for cosmology and for collider phenomenology are briefly discussed and concluding remarks are made.

2. The axino

First, let us briefly summarize the main properties of axions. The axion can be defined as a pseudo-scalar particle with the effective interaction given by

$$\mathcal{L}_a^{\rm eff} = \frac{a}{M} F\widetilde{F}\,, \tag{2.1}$$

where M is a model-dependent mass scale, F is the field strength of the gluon field and $\widetilde{F}$ is its dual. The potential arising from the above interaction settles $\langle a\rangle$ to zero. In general, there can arise some other small terms in the axion potential in addition to the one arising from (2.1). (When CP violation is considered $\langle a\rangle = 0$ is shifted,

[2] In the literature s is called saxion ($\equiv$ scalar partner of axion) [13], or saxino ($\equiv$ the scalar partner of axino) [27]. Since superpartner names ending with "-ino" seem to be reserved for (Majorana) fermions, here we will use the name 'saxion'.

but the shift is extremely small since a square of the weak interaction coupling and the pion mass appear in the resulting expression [28].) At low energy, one may not question how M is generated. It can arise in renormalizable or non-renormalizable theories, or in composite models. If a fundamental theory such as superstrings allows the above non-renormalizable interaction, then M is the compactification scale [29]. If a confining force at a scale Λ_a produces the Goldstone boson with the above coupling at low energy, then M is of order Λ_a and the axion is composite [30]. If the axion resides in the complex Higgs multiplet(s), it can be derived through the PQ mechanism with the spontaneously broken chiral U(1) global symmetry.

In phenomenologically acceptable axion models, the mass scale M is given by $M = 8\pi f_a/\alpha_s$, where f_a is the PQ scale, $10^9\,\mathrm{GeV} \lesssim f_a \lesssim 10^{12}\,\mathrm{GeV}$. In these models the axion mass is given by $m_a \sim \Lambda_{QCD}^2/f_a \sim 10^{-2}$–$10^{-5}$ eV; in other words the axion is very light. These models can be classified as hadronic axion models (usually called the KSVZ models [31]) and DFSZ axion models [32]. Let us briefly recall some basic features of these two most popular implementations of the PQ mechanism. (Detailed presentations can be found in several excellent reviews, e.g. in [4].) In both one assumes the existence of a complex scalar field ϕ, which is a singlet under the SM gauge group but carries a PQ charge. When ϕ develops a VEV, $\langle\phi\rangle = f_a$, the PQ symmetry gets broken and the phase of ϕ becomes the axion field. To relate the global charge of the complex scalar ϕ to the PQ charge (i.e. to couple the axion to the gluon anomaly), in the KSVZ scheme [31] one further introduces at least one heavy quark Q[3] which couples to ϕ through the term

$$\mathcal{L}_{PQ} = f_Q \phi \bar{Q}_R Q_L + \mathrm{h.c.}\,, \tag{2.2}$$

where f_Q is the Yukawa coupling. The PQ charge of ϕ is $+2$, while those of $Q_{R,L}$ are ± 1. The breaking of the PQ symmetry gives large mass to the heavy quark, $m_Q = f_Q f_a$. The axion interacts with ordinary matter through loops involving the exchange of the heavy quark (i.e. through the anomaly term).

This leads, below the PQ scale, to effective axion interactions terms; among these the most important one is its interaction with gluons

$$\mathcal{L}_{agg} = \frac{\alpha_s}{8\pi f_a} a F\widetilde{F}\,. \tag{2.3}$$

In the DFSZ model [32], instead of the heavy quark, one introduces two Higgs doublets [2, 33] $\phi_{u,d}$ of $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ which couple the SM quark sector to ϕ. The doublets carry hypercharges $\pm 1/2$ and PQ charges $-Q_{u,d}$, respectively. The ϕ_u field couples only to up-type quarks, while ϕ_d couples only to down-type quarks and leptons which carry PQ charges so as to respect the PQ symmetry.

[3] In general one can introduce, instead of a single colored fermion, a multiplet, e.g. a non-singlet vector-like quark representation of $\mathrm{SU}(2) \times \mathrm{U}(1)_Y$.

After the breaking of the PQ symmetry, again effective axion interactions with ordinary matter are generated, including the term (2.3). In both models there appears an additional quantity, the number of different vacua (or, equivalently, domain walls) N: $N = 1(6)$ for the KSVZ (DFSZ) model. For our purpose it is enough to note that its effect will be to replace f_a by f_a/N.

Both models can be readily supersymmetrized [34, 35]. The scalar field ϕ becomes promoted to the superfield, and accordingly the other fields as well. In particular, the axion, being the phase of ϕ, is also promoted to an axion supermultiplet Φ consisting of the pseudo-scalar axion a, its fermionic partner, the axino $\widetilde{a}$, and the scalar partner, the saxion s

$$\Phi = \frac{1}{\sqrt{2}}\,(s + ia) + \sqrt{2}\tilde{a}\theta + F_\Phi\theta\theta\,. \tag{2.4}$$

The axino is thus a neutral, $R = -1$, Majorana chiral fermion. Adding supersymmetry opens up a wide range of choices for the implementation and breaking of the PQ symmetry through a choice of different superpotentials and SUSY-breaking schemes, as will be illustrated below.

Now let us move on to the discussion of axino masses. We will briefly summarize the relevant results of several previous studies [12]–[15], [36, 37]. The overall conclusion will be that the mass of the axino is strongly model-dependent [13, 14]. It can be very small ($\sim$ eV), or large ($\sim$ GeV), depending on the model.

It is worth emphasizing that, unlike the case of the gravitino and ordinary superpartners, the axino mass does not have to be of the order of the SUSY-breaking scale in the visible sector, $M_{\rm SUSY}$ [12]–[14]. In global SUSY models, one sets $M_{\rm SUSY} = \mathcal{O}(1\,{\rm TeV})$ on the basis of naturalness. In SUGRA models $M_{\rm SUSY}$ is set by the gravitino mass $m_{\widetilde{G}} \sim M_{\rm S}^2/M_{\rm P}$, where $M_{\rm S} \sim 10^{11}\,{\rm GeV}$ is the scale of local SUSY-breaking in the hidden sector and $M_{\rm P} \simeq 2.4 \times 10^{18}\,{\rm GeV}$ denotes the Planck mass.

It is easy to see why $m_{\widetilde{a}}$ is not generically of the order of $M_{\rm SUSY}$. In the case of unbroken SUSY, all members of the axion supermultiplet remain degenerate and equal to the tiny mass of the axion given by the QCD anomaly. Once SUSY is broken, the saxion, being a scalar, receives a soft-mass term $m_s^2 s^2$ [12, 38], where $m_s \sim M_{\rm SUSY}$, similarly to the other scalar superpartners.

Of course, the axion, being instead a phase of the fields whose VEVs break the PQ symmetry, does not receive soft mass. Likewise, one cannot write a *soft* mass term for the axino since it is a chiral fermion (for the same reason there are no soft terms for, for instance, the MSSM higgsinos) and a superpotential term $W \sim$ (mass parameter) $\cdot\, \Phi\Phi$ is absent due to the PQ symmetry. The lowest-order term one can write will be a non-renormalizable term of dimension-5. The axino mass will then be of order $M_{\rm SUSY}^2/f_a \sim 1\,{\rm keV}$ [12]. (A more detailed, and more sophisticated, explanation can be found, e.g., in ref. [14].)

However, in addition to this inevitable source of axino mass, one can easily generate much larger contributions to $m_{\tilde{a}}$ at one-loop or even at tree-level. One-loop terms will always contribute but will typically be $\lesssim M_{\rm SUSY}$ (KSVZ) or even $\ll M_{\rm SUSY}$ (DFSZ). Furthermore, in non-minimal models where the axino mass eigenstate comes from more than one superfield, $m_{\tilde{a}}$ arises even at the tree-level. In this case $m_{\tilde{a}}$ can be of order $M_{\rm SUSY}$ but can also be much smaller. These points are illustrated by the following examples.

First let us see how tree-level contributions can be generated from superpotentials involving singlet fields. For example, consider the case where the PQ symmetry is assumed to be broken by the renormalizable superpotential term (in KSVZ or DFSZ models) [39]

$$W = f Z(S_1 S_2 - f_a^2)\,, \tag{2.5}$$

where f is a coupling, and Z, S_1 and S_2 are chiral superfields with PQ charges of $0, +1$ and -1, respectively. In this case the axino mass can be at the soft SUSY-breaking mass scale [36]. It arises from diagonalizing the mass matrix of the fermionic partners $\tilde{S}_1, \tilde{S}_2$, and $\tilde{Z}$,

$$\begin{pmatrix} 0 & m_{\tilde{a}} & f f_a \\ m_{\tilde{a}} & 0 & f f_a \\ f f_a & f f_a & 0 \end{pmatrix}, \tag{2.6}$$

where $m_{\tilde{a}} = f\langle Z\rangle$ and $f f_a \sim 10^{11}\,$GeV. The mass matrix in eq. (2.6) is for three two-component neutral fermions. These three will split into one Dirac fermion, and one Majorana fermion which is interpreted as the axino. The corresponding eigenvalues are $\lambda = -m_{\tilde{a}}$ and $\lambda = \pm\sqrt{2} f f_a + \mathcal{O}(m_{\tilde{a}})$. In the global SUSY limit, $\langle Z\rangle = 0$ and the tree-level axino mass would be zero. However, when S_1 and S_2 acquire VEVs and soft terms are included, $V = |f|^2(|S_1|^2 + |S_2|^2)|Z|^2 + (A_1 f S_1 S_2 Z - A_2 f f_a^2 Z + \text{h.c.})$, a linear term in Z is generated which induces $\langle Z\rangle$ of order $(A_1 - A_2)/f$ where $A_{1,2}$ are the soft trilinear mass parameters. The axino mass thus arises at the soft mass scale [36].

However, by choosing a more complicated superpotential, one can significantly lower the axino mass [14] even at the tree-level. Consider a supergravity superpotential consistent with the PQ symmetry as

$$W' = f Z(S_1 S_2 - X^2) + \frac{1}{3}\lambda(X - M)^3\,, \tag{2.7}$$

where X carries a zero PQ charge. In this case a minimization of the potential resulting from W' is much more complicated, and an approximate solution gives the lightest eigenvalue of the fermion mass matrix with $\langle V\rangle = 0$ which is $m_{\tilde{a}} = \mathcal{O}(A - 2B + C) + \mathcal{O}(m_{3/2}^2/f_a)$, where A, B and C are the respective trilinear, bilinear and linear soft breaking parameters. For the standard pattern of soft breaking terms, $B = A - m_{3/2}$ and $C = A - 2m_{3/2}$, the leading contribution $A - 2B + C$ vanishes and the tree-level axino mass becomes of order $m_{3/2}^2/f_a \sim 1\,$keV.

JHEP05(2001)033

The above example shows that in calculating even the tree-level axino mass one must carefully consider the full potential V generated in supergravity models. In general, the tree-level axino mass either can be of order $m_{3/2}$ or, depending on the PQ sector of the model and the pattern of soft breaking parameters, can be much less, as shown in the above examples. The detailed conditions for this to happen were analyzed by Chun and Lukas [37]. In particular one can also have a more complicated PQ sector and break the PQ symmetry by non-renormalizable terms in the superpotential leading to interesting cosmological implications [40].

If the tree-level axino mass is either zero or of order $m_{3/2}^2/f_a$, the contribution from loop diagrams can become more important. In the global SUSY version of the KSVZ model, axino mass will arise at the one-loop level with a SUSY-breaking A-term insertion at the intermediate heavy squark line. Then one finds $m_{\widetilde{a}} \sim (f_Q^2/8\pi^2)A$, where f_Q is the Yukawa coupling of the heavy quark to a singlet field containing the axion (compare eq. (2.2)), which gives $m_{\widetilde{a}}$ in the range of $\lesssim$ a few tens of GeV [41, 36]. In the DFSZ model [32] (in either global SUSY or supergravity), where such A-term contribution is absent, the axino mass remains in the keV range, as was pointed out in [15, 13]. Thus, in some models the axino mass can be rather small.

In gauge-mediated SUSY-breaking models (GMSB), the pattern of the axino mass is completely different. In the GMSB approach, the tree-level effect of supergravity is not important since the gravitino mass is much smaller than soft breaking masses, typically of order $\sim$ eV to $\sim$ keV. These models attracted much interest in the recent past since the gravitino production rate is quite large which leads to important implications for cosmology [42] and accelerator experiments [43]. In the GMSB models, the axino mass is further reduced by one more power of f_a and can range from 10^{-9} eV to 1 keV [44]. It remains model-dependent but again can be smaller than the gravitino mass. We will not discuss GMSB cases here any more.

In conclusion, we see that a complete knowledge of the superpotential and supersymmetry-breaking mechanism is necessary to pin down the axino mass. In general it can range from $\sim$ eV to $\sim$ tens of GeV. For the sake of generality, in discussing cosmological properties of axinos, we will use $m_{\widetilde{a}}$ as a free parameter.

In addition to its mass, the axino couplings to gauge and matter fields are of crucial importance for the study of its cosmological abundance. In general, the couplings of the members of the axion/axino supermultiplet will be determined by the PQ symmetry and supersymmetry. The most important term for our purpose will be the coupling to gauge multiplets. It is given by the same interaction term as the one which gives rise to the "θ" term in the QCD lagrangian. In supersymmetric notation it can be written as

$$\mathcal{L} = \frac{\alpha_Y C_{aYY}}{4\sqrt{2}\pi\,(f_a/N)}(\Phi B_\alpha B^\alpha)_{\theta\theta} + \frac{\alpha_2 C_{aWW}}{4\sqrt{2}\pi\,(f_a/N)}[B \to W] + \frac{\alpha_s C_{\rm agg}}{4\sqrt{2}\pi\,(f_a/N)}[B \to F] + \text{h.c.}, \tag{2.8}$$

where α_i $(i = Y, 2, s)$ are the coupling constant strengths of the SM subgroups, and B, W_3 and F are the respective gauge supermultiplets. The coefficient $C_{\text{agg}} = 1$ is universal, while C_{aYY} and C_{aWW} are model-dependent. For processes that involve the electroweak gauge bosons and superpartners, it is always possible to rotate away the interaction of the axion with the $\mathrm{SU}(2)_L$ gauge multiplet through some anomalous global chiral $\mathrm{U}(1)_X$ rotation since the PQ symmetry is global. Then the only interaction that remains to be specified will be the one between the axino and the $\mathrm{U}(1)_Y$ gauge multiplet. This will be equivalent to assigning PQ charges to left-handed doublets in such a way as to obtain $C_{aWW} = 0$. One will then also find $C_{aYY} = C_{a\gamma\gamma}$.

The rationale for removing the coupling C_{aWW} is as follows. We would like to keep the gauge couplings and top quark Yukawa coupling but neglect the other Yukawa couplings. This is because the neglected Yukawa couplings will not be important in the production cross section. Moreover we expect that the lightest neutralino is likely to contain a significant bino component [45, 46, 19], so that phenomenologically the interaction between the $\mathrm{U}(1)_Y$ field strength and the axion multiplet will be important. Keeping this in mind, we will work with the simplest interactions and use the basis where the C_{aWW} coupling is absent.

The procedure of global $\mathrm{U}(1)_X$ rotation and charge redefinition is somewhat subtle and requires a word of clarification on which we digress here. This rotation applies only to the high energies of interest to us, but not to the low-energy phenomena below the QCD scale. We can use a global $\mathrm{U}(1)_X$ rotation to redefine the PQ charges of the quarks like $Q' = Q - Q(\text{quark doublet})$. This will cancel the PQ charge of the quark doublets, but will also shift the PQ charges of the other fields. Since the Q' charges of the quark doublets are absent, one would expect that there is no C_{aWW} coupling either. However, there will be (in DFSZ models) lepton and Higgs doublets whose Q' charges will in general be non-zero. Thus the Q' charge should instead be defined by the rotation removing C_{aWW}, in which case Q' charges of quark doublets and lepton and Higgs doublets may in general be different from zero. Because the top quark Yukawa coupling is kept, the global $\mathrm{U}(1)_X$ transformation must either leave the quark fields unchanged or operate on both the left-handed and right-handed quarks and the Higgs field, similarly to the Standard Model $\mathrm{U}(1)_Y$.

The rotation will not change the C_{agg} coupling since $\mathrm{U}(1)_Y$ is anomaly-free. On the other hand, for leptons the Yukawa coupling can be neglected and we can therefore use any global $\mathrm{U}(1)_X$ such that the final C_{aWW} coupling is absent. The above rotation amounts to changing the PQ charges of the lepton doublets and hence to modifying all the couplings of the a-field with the field strengths of W_μ and B_μ. If we could neglect also the top quark Yukawa coupling, we could consider instead a basis without the C_{agg} coupling where the interaction (2.1) would be absent and no axion field could be defined. But that scenario would correspond to the decoupling of the axion from the SM sector which takes place in the massless quark case.

JHEP05(2001)033

In conclusion, the axion coupling must be defined rigorously at the axion vacuum. However, above the TeV scale we can choose any basis for calculating the axino production if the contribution for $E \lesssim 1$ TeV is not important.

After the $\mathrm{U}(1)_X$ rotation, the SU(2) singlets (quarks and leptons) carry Q' charges and hence the aYY coupling will not be absent:

$$\frac{\alpha_Y C_{aYY}}{f_a/N} a B \widetilde{B}\,, \tag{2.9}$$

where B denotes here the field strength of the $\mathrm{U}(1)_Y$ gauge boson B_μ and $\widetilde{B}$ denotes its dual. Since $B_\mu = A_\mu \cos\theta_W + Z_\mu \sin\theta_W$, we obtain

$$\frac{\alpha_Y \cos^2\theta_W C_{aYY}}{f_a/N} a F_{\mathrm{em}} \widetilde{F}_{\mathrm{em}}\,, \tag{2.10}$$

where, F_{em} is the field strength of the electromagnetic gauge boson and $\widetilde{F}_{\mathrm{em}}$ is its dual. This implies $C_{aYY} = C_{a\gamma\gamma}$.

In the DFSZ model with (d^c, e)-unification, $C_{aYY} = 8/3$. In the KSVZ model for $e_Q = 0, -1/3$, and $2/3$, $C_{aYY} = 0, 2/3$ and $8/3$, respectively [47]. Below the QCD chiral symmetry-breaking scale, $C_{a\gamma\gamma}$ and C_{aYY} are reduced by 1.92.

For our purpose the most important coupling will be that of axino-gaugino-gauge boson interactions which can be derived from eq. (2.8) and written in a more conventional way as a dimension-5 term in the lagrangian:

$$\mathcal{L}_{\tilde{a}\lambda A} = i\frac{\alpha_Y C_{aYY}}{16\pi\,(f_a/N)} \bar{\tilde{a}}\gamma_5[\gamma^\mu,\gamma^\nu]\widetilde{B}B_{\mu\nu} + i\frac{\alpha_s}{16\pi\,(f_a/N)}\bar{\tilde{a}}\gamma_5[\gamma^\mu,\gamma^\nu]\tilde{g}^b F^b_{\mu\nu}\,. \tag{2.11}$$

Here and below $\widetilde{B}$ denotes the bino, the fermionic partner of the $\mathrm{U}(1)_Y$ gauge boson B and $\tilde{g}$ stands for the gluino.

One can also think of terms involving dimension-4 operators coming, e.g., from the *effective* superpotential $\Phi\Psi\Psi$ where Ψ is one of the MSSM matter (super)fields. However, axino production processes coming from such terms will be suppressed at high energies with respect to processes involving eq. (2.11) by a factor m_Ψ^2/s, where s is the square of the center of mass energy. We will comment on the role of dimension-4 operators again below but, for the most part, mostly concentrate on the processes involving axino interactions with gauginos and gauge bosons, eq. (2.11), which are both model-independent and dominant.

3. Thermal production

As stated in the Introduction, thermal production proceeds through collisions and decays of particles that still remain in the thermal-bath. As we will see, its efficiency will strongly depend on effective interaction strengths of the processes involved and on characteristic temperatures of the bath. Particles like axinos and gravitinos are

JHEP05(2001)033

somewhat special, in the sense that their interactions with other particles are very strongly suppressed with respect to the Standard Model interaction strengths. Therefore such particles remain in thermal equilibrium only at very high temperatures. In the particular case of axinos (as well as axions and saxions), their initial thermal populations decouple at [13]

$$T_{\rm D} \sim 10^{10}\,{\rm GeV}\left(\frac{f_a/N}{10^{11}\,{\rm GeV}}\right)\left(\frac{\alpha_s}{0.1}\right)^{-3}. \tag{3.1}$$

At such high temperatures, the axino number density is the same as the one of photons and other relativistic species. In other words, such primordial axinos freeze out as *relativistic* particles. Rajagopal, Turner and Wilczek (RTW) [13] pointed out that, in the absence of a subsequent period of inflation, the requirement that the axino energy density be not too large ($\Omega_{\widetilde a} \lesssim 1$) leads to

$$m_{\widetilde a} < 12.8\,{\rm eV}\left(\frac{g_*(T_{\rm D})}{g_{\rm eff}}\right), \tag{3.2}$$

where $g_{\rm eff} = 1.5$ and g_* is the number of effectively massless degrees of freedom (particles with mass much smaller than temperature). In the MSSM at temperatures much higher than $M_{\rm SUSY}$, one has $g_* = 915/4$. The RTW bound (3.2) then takes a more well-known form [13]

$$m_{\widetilde a} < 2\,{\rm keV}, \tag{3.3}$$

and the corresponding axinos would be light and would provide warm or even hot dark matter [13]. We will not consider this case in the following, primarily because we are interested in cold DM axinos. We will therefore assume that the initial population of axinos (and other relics, such as gravitinos), which were present in the early Universe, was subsequently diluted away by an intervening inflationary stage and that the reheating temperature after inflation was smaller than $T_{\rm D}$. It also had to be less than f_a, otherwise the PQ symmetry would have been restored, thus leading to the well-known domain-wall problem associated with global symmetries.

At lower temperatures, $T < T_{\rm D}$, the Universe can be re-populated with axinos (and gravitinos) through scattering and decay processes involving superpartners in the thermal-bath. As long as the axino number density $n_{\tilde a}$ is much smaller than n_γ, the number density of photons in thermal equilibrium, its time evolution will be adequately described by the Boltzmann equation:

$$\frac{d n_{\tilde a}}{d t} + 3Hn_{\tilde a} = \sum_{i,j}\langle\sigma(i+j\to\widetilde a+\cdots)v_{\rm rel}\rangle n_i n_j + \sum_i\langle\Gamma(i\to\widetilde a+\cdots)\rangle n_i\,. \tag{3.4}$$

Here H is the Hubble parameter, $H(T) = \sqrt{(\pi^2 g_*)\,/\,(90M_{\rm P}^2)}\,T^2$, where g_* has been defined above and $\sigma(i+j\to\widetilde a+\cdots)$ is the scattering cross section for particles i,j into final states involving axinos, $v_{\rm rel}$ is their relative velocity, n_i is the ith particle number density in the thermal-bath, $\Gamma(i \to \widetilde{a} + \cdots)$ is the decay width of the ith particle, $\langle \cdots \rangle$ stands for thermal averaging. (Averaging over initial spins and summing over final spins is understood.) Note that on the *r.h.s.* we have neglected inverse processes, since they are suppressed by $n_{\widetilde{a}}$.

Solving the Boltzmann equation. In order to solve the Boltzmann eq. (3.4), it is convenient to introduce the axino TP yield

$$Y_{\widetilde{a}}^{\rm TP} = \frac{n_{\widetilde{a}}^{\rm TP}}{s}\,, \tag{3.5}$$

where $s = (2\pi^2/45) g_{s*} T^3$ is the entropy density, and normally $g_{s*} = g_*$ in the early Universe. We also change variables from the cosmic time t to the temperature T by $-dt/dT = 1/HT$ for the radiation dominated era. One integrates the Boltzmann equation from the reheating temperature $T_{\rm R}$ after inflation down to zero. The yield can be written as

$$Y_{\widetilde{a}}^{\rm TP} = \sum_{i,j} Y_{i,j}^{\rm scat} + \sum_i Y_i^{\rm dec}\,, \tag{3.6}$$

where the summation over indices i, j is the same as in eq. (3.4). The expressions for $Y_{i,j}^{\rm scat}$ and $Y_i^{\rm dec}$ are given by

$$Y_{i,j}^{\rm scat} = \int_0^{T_{\rm R}} dT\, \frac{\langle \sigma(i + j \to \widetilde{a} + \cdots) \rangle n_i n_j}{sHT} \tag{3.7}$$

and

$$Y_i^{\rm dec} = \int_0^{T_{\rm R}} dT\, \frac{\langle \Gamma(i \to \widetilde{a} + \cdots) \rangle n_i}{sHT}\,. \tag{3.8}$$

Explicit formulae for $Y_{i,j}^{\rm scat}$ and $Y_i^{\rm dec}$ can be found, for instance, in ref. [48].

Scatterings. In the case of axinos the main productions channels are the scatterings of (s)particles described by a dimension-5 axino-gaugino-gauge boson term in the lagrangian (2.11). Because of the relative strength of α_s, the most important contributions will come from strongly interacting processes. We will discuss them first.

The relevant 2-body processes for strongly interacting particles i, j into several final states involving axinos are listed in table 1. The corresponding cross sections $\sigma_n = \sigma(i + j \to \widetilde{a} + \cdots)$, where the label $n = A, \ldots, J$ counts the different allowed combinations of the initial and final-state particles, can be written as

$$\sigma_n(s) = \frac{\alpha_s^3}{4\pi^2 \left(f_a/N\right)^2} \overline{\sigma}_n(s)\,, \tag{3.9}$$

where $\sqrt{s}$ is the energy in the center-of-mass frame. Also listed there are the respective: spin factor $n_{\rm spin}$ (the number of spin combinations in the initial state), flavor

n	Process	$\overline{\sigma}_N$	$n_{\rm spin}$	$n_{\rm F}$	$\eta_1\eta_2$
A	$g^a + g^b \to \tilde{a} + \tilde{g}^c$	$\frac{1}{8}\|f^{abc}\|^2$	4	1	1
B	$g^a + \tilde{g}^b \to \tilde{a} + g^c$	$\frac{5}{16}\|f^{abc}\|^2\left[\log(s/m_{\rm eff}^2) - \frac{15}{8}\right]$	4	1	$\frac{3}{4}$
C	$g^a + \tilde{q}_k \to \tilde{a} + q_j$	$\frac{1}{8}\|T^a_{jk}\|^2$	2	$N_F \times 2$	1
D	$g^a + q_k \to \tilde{a} + \tilde{q}_j$	$\frac{1}{32}\|T^a_{jk}\|^2$	4	$N_F \times 2$	$\frac{3}{4}$
E	$\tilde{q}_j + q_k \to \tilde{a} + g^a$	$\frac{1}{16}\|T^a_{jk}\|^2$	2	$N_F \times 2$	$\frac{3}{4}$
F	$\tilde{g}^a + \tilde{g}^b \to \tilde{a} + \tilde{g}^c$	$\frac{1}{2}\|f^{abc}\|^2\left[\log(s/m_{\rm eff}^2) - \frac{29}{12}\right]$	4	1	$\frac{3}{4}\,\frac{3}{4}$
G	$\tilde{g}^a + q_k \to \tilde{a} + q_j$	$\frac{1}{4}\|T^a_{jk}\|^2\left[\log(s/m_{\rm eff}^2) - 2\right]$	4	N_F	$\frac{3}{4}\,\frac{3}{4}$
H	$\tilde{g}^a + \tilde{q}_k \to \tilde{a} + \tilde{q}_j$	$\frac{1}{4}\|T^a_{jk}\|^2\left[\log(s/m_{\rm eff}^2) - \frac{15}{8}\right]$	2	$N_F \times 2$	$\frac{3}{4}$
I	$q_k + \bar{q}_j \to \tilde{a} + \tilde{g}^a$	$\frac{1}{24}\|T^a_{jk}\|^2$	4	N_F	$\frac{3}{4}\,\frac{3}{4}$
J	$\tilde{q}_k + \tilde{\bar{q}}_j \to \tilde{a} + \tilde{g}^a$	$\frac{1}{24}\|T^a_{jk}\|^2$	1	$N_F \times 2$	1

Table 1: The cross sections for the different axino thermal-production channels involving strong interactions. Masses of particles have been neglected, except for the plasmon mass $m_{\rm eff}$. See text for explanation of the different symbols.

factor $n_{\rm F}$ (number of color-triplet chiral multiplets), and number density factor η_i, $i = 1, 2$ ($\eta_i = 3/4$ (1) for each initial-state fermion (boson). The group factors f^{abc} and T^a_{jk} of the gauge group SU(N) satisfy the usual relations $\sum_{a,b,c} |f^{abc}|^2 = N(N^2-1)$ and $\sum_a \sum_{jk} |T^a_{jk}|^2 = (N^2-1)/2$.

The diagrams listed in the table 1 are analogous to those involving gravitino production considered by Moroi et al. [8], and later by Bolz et al. [9], and we follow their classification. This analogy should not be surprising, since both particles are neutral Majorana superpartners. (We will return to the gravitino–axino analogy in section 6.) In figure 1 we present the Feynman diagrams corresponding to channels A and G. The diagrams for the other channels can be drawn in a similar fashion.

The diagrams in channels B, F, G and H involve t-channel exchange of massless gluons and are therefore divergent. In order to provide a mass regulator, one can introduce a plasmon mass $m_{\rm eff}$ representing the effective gluon mass due to plasma effect. We follow the prescription of ref. [7], which has also been used in refs. [8, 9], and assume $m_{\rm eff}^2 = g_s^2 T^2$ where g_s is the strong coupling constant.

Note that we have neglected in the gluon thermal mass a potentially sizeable factor which counts the number of colored degrees of freedom of the plasma but kept a constant term in the logarithmically divergent cross sections, which depends on the procedure used to regulate the divergence. While a more proper treatment exists in the literature [49, 11], based on Thermal Field Theory and taking into account also the Fermi or Bose nature of the particles in the thermal-bath, we believe that the above approximation is good enough for our purposes. This belief is supported in particular by the fact that in the case of gravitinos, the full treatment agrees with an approximate one within a factor 3 [11]. For simplicity, in computing $\overline{\sigma}_n(s)$ we neglected all mass terms (involving gauginos, axinos and scalars) and kept only the plasmon mass $m_{\rm eff}^2$.

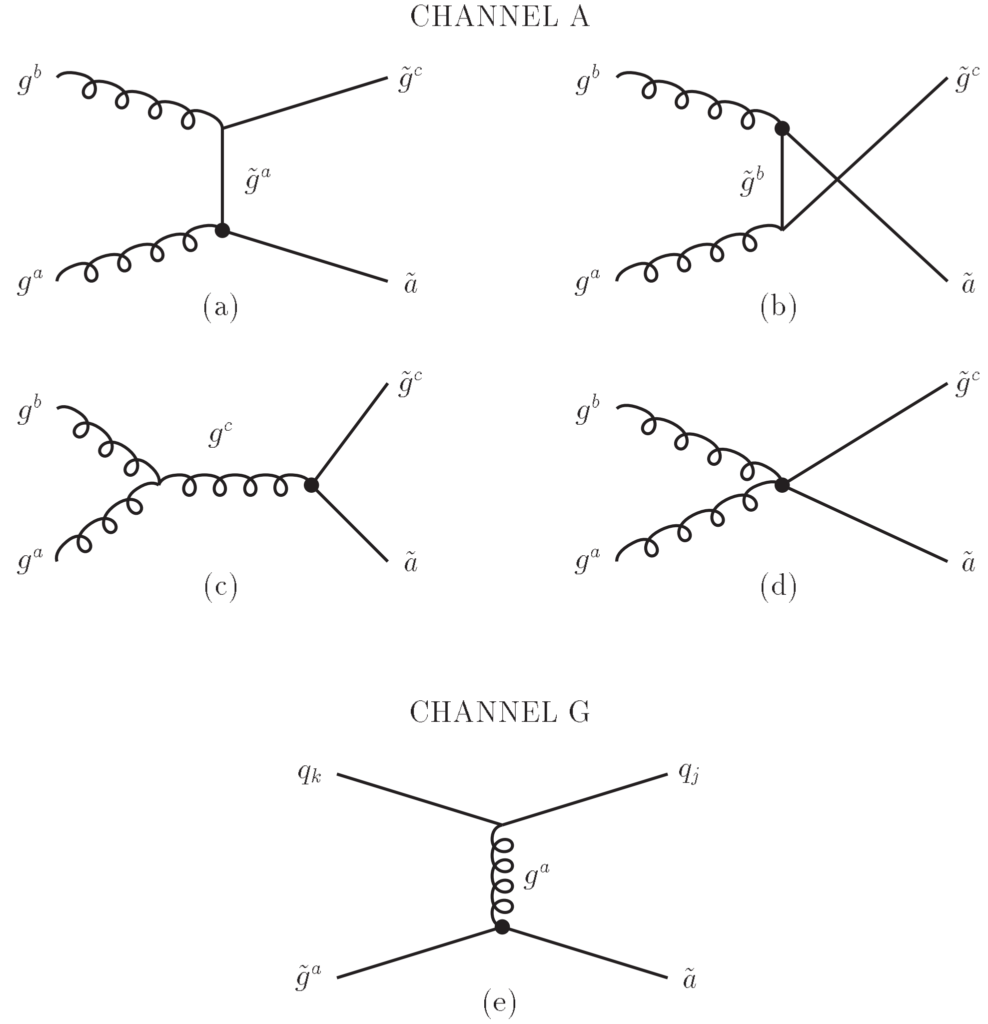

Figure 1: Feynman diagrams contributing to channels A and G in table 1. Thick dot denotes effective axino couplings to ordinary colored particles.

Finally, we comment on the scattering processes involving the hypercharge multiplet. It is clear that these will always be subdominant. This is not only a result of a weaker interaction strength relative to gluons and gluinos (so long as C_{aYY} is not too large), but is also caused by a smaller number of production channels. Indeed, since $\mathrm{U}(1)_Y$ is abelian, in this case many partial cross sections (all channel A, B and F contributions) vanish automatically. The others are given by the same expressions as in Table 1, with the substitution $\alpha_s^3 \to \alpha_Y^3 C_{aYY}^2$ and $|T_{ij}^a|^2 \to Y_i^2$, where Y_i is the hypercharge of the initial (and final) state.

Decays. In addition to scattering processes, axinos can also be produced through decays of heavier superpartners in thermal plasma. At temperatures $T \gtrsim m_{\tilde{g}}$, these are dominated by the decays of gluinos into LSP axinos and gluons. The relevant decay width is given by

$$\Gamma(\widetilde{g}^a \to \widetilde{a} + g^b) = \delta^{ab} \frac{\alpha_s^2}{128\pi^3} \frac{m_{\widetilde{g}}^3}{(f_a/N)^2} \left(1 - \frac{m_{\widetilde{a}}^2}{m_{\widetilde{g}}^2}\right)^3 \tag{3.10}$$

and one should sum over the color indices $a, b = 1, \ldots, 8$.

In addition, at temperatures $m_\chi \lesssim T_{\mathrm{R}} \lesssim m_{\widetilde{g}}$, neutralino decays to axinos also contribute, while at higher temperatures they are sub-dominant. The relevant contribution is given by the decay of the bino component:

$$\Gamma(\chi_i \to \widetilde{a} + B) = \frac{\alpha_{\mathrm{em}}^2 C_{a\chi_i B}^2}{128\pi^3} \frac{m_{\chi_i}^3}{(f_a/N)^2} \left(1 - \frac{m_{\widetilde{a}}^2}{m_{\chi_i}^2}\right)^3 . \tag{3.11}$$

Here α_{em} is the electromagnetic coupling strength, $C_{a\chi_i B} = Z_{iB} C_{aYY} / \cos^2\theta_W$ where Z_{iB} is the bino component of the ith neutralino χ_i ($i = 1, 2, 3, 4$). We use the basis $\chi_i = Z_{i1}\widetilde{B} + Z_{i2}\widetilde{W}_3 + Z_{i3}\widetilde{H}_b^0 + Z_{i4}\widetilde{H}_t^0$ of the respective fermionic partners (denoted by a tilde) of the electrically neutral gauge bosons B and W_3, and the MSSM Higgs bosons H_b and H_t.

Discussion and results. We have evaluated the integrals (3.7) and (3.8) numerically using expressions (3.9), (3.10) and (3.11). The results are presented in figure 2 for representative values of $f_a = 10^{11}\,\mathrm{GeV}$ and $m_{\widetilde{q}} = m_{\widetilde{g}} = 1\,\mathrm{TeV}$. The respective contributions due to scattering as well as gluino and neutralino decays are marked by dashed, dash-dotted and dotted lines.

It is clear that at high enough T_{R}, much above $m_{\widetilde{q}}$ and $m_{\widetilde{g}}$, scattering processes involving such particles dominate the axino production. For $T_{\mathrm{R}} \gg m_{\widetilde{q}}, m_{\widetilde{g}}$, $\langle\sigma_n v\rangle$ is almost constant and thus $Y^{\mathrm{scat}} \simeq 2 \times 10^{-5}\, M_{\mathrm{P}} T_{\mathrm{R}} \times (\sum_n \sigma_n)$. In other words, Y^{scat} grows linearly as T_{R} becomes larger. (In the numerical calculation we have also taken into account the running of the strong coupling constant. Then Y^{scat} is given, to a very good approximation, by the same formula with α_s replaced by $\alpha_s(T_{\mathrm{R}})$, and it grows like $\alpha_s^3(T_{\mathrm{R}})T_{\mathrm{R}}$. Using $\alpha_s(T_{\mathrm{R}})$ instead of its value at M_Z gives a correction of up to a factor 10. For example, $\alpha_s(M_Z)^3/\alpha_s(10^8\,\mathrm{GeV})^3 = 5.75$.) A similar result was estimated and used previously for the saxion in ref. [26]. In contrast, the decay contribution above the gluino mass threshold, $Y^{\mathrm{dec}} \simeq 5 \times 10^{-4}\,(M_{\mathrm{P}}\Gamma_{\widetilde{g}}/m_{\widetilde{g}}^2)$, remains independent of T_{R}. This is not surprising since the scattering term in eq. (3.4) behaves like const $\times\, n_i n_j \sim$ const $\times\, T^6$ relative to the decay term which scales like $m_{\widetilde{g}}^3 n_{\widetilde{g}} \sim m_{\widetilde{g}}^3 T^3$ plus additional suppression $\langle\Gamma\rangle \sim T^{-1}$. In the region of very high $T_{\mathrm{R}} \sim 10^9\,\mathrm{GeV}$, $Y_{\widetilde{a}}^{\mathrm{TP}}$ becomes comparable with the axino yield in thermal equilibrium $Y_{\widetilde{a}}^{\mathrm{EQ}} \simeq 2 \times 10^{-3}$. Since, as mentioned below eq. (3.4), we have neglected processes of axino re-annihilation to thermal-bath particles, there appears an artificial "edge" between the sloping solid curve representing $Y_{\widetilde{a}}^{\mathrm{TP}}$ and the dashed horizontal line corresponding to the axino equilibrium value $Y_{\widetilde{a}}^{\mathrm{EQ}} \approx 2 \times 10^{-3}$. Including axino re-annihilation would have the effect of rounding it. We leave it as is because we are not interested in very large values of T_{R}.

JHEP05(2001)033

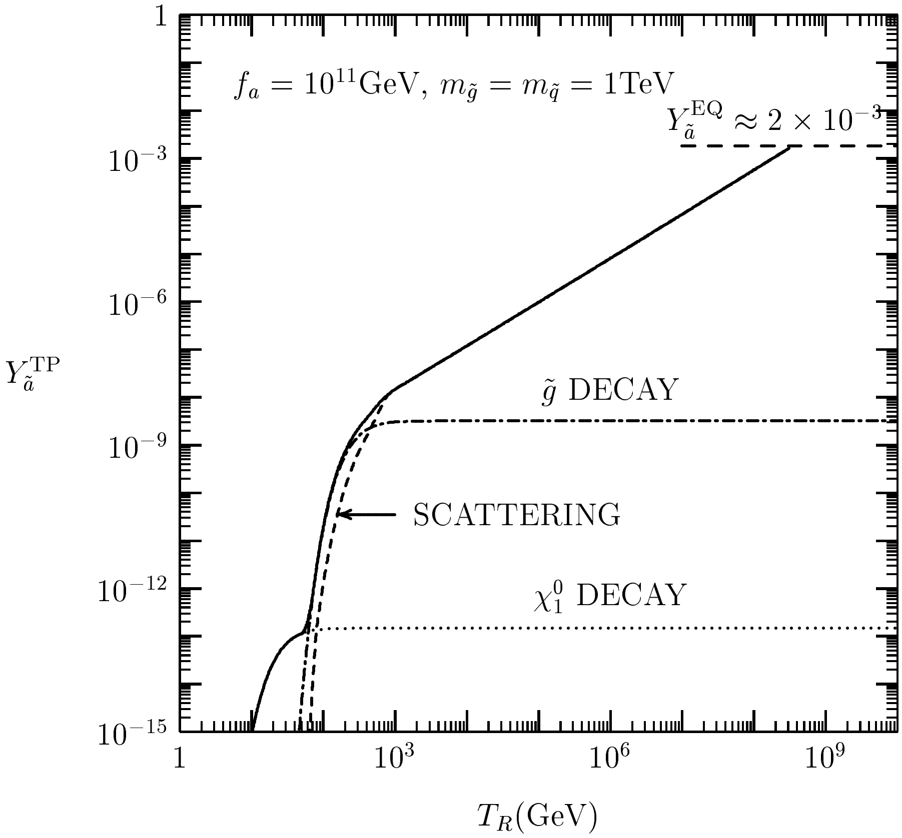

Figure 2: $Y_{\tilde{a}}^{\rm TP}$ as a function of $T_{\rm R}$ for representative values of $f_a = 10^{11}$ GeV and $m_{\tilde{q}} = m_{\tilde{g}} = 1$ TeV.

At $T_{\rm R}$ roughly below the mass of the squarks and gluinos, their thermal population starts to become strongly suppressed by the Boltzmann factor $e^{-m/T}$, hence causing a distinct knee in the scattering contribution in figure 2. It is in this region that gluino decays (dash-dotted line) given by eq. (3.10) become dominant, before they also become suppressed by the Boltzmann factor due to the gluino mass. For $m_\chi \lesssim T_{\rm R} \lesssim m_{\tilde{q}}, m_{\tilde{g}}$, the axino yield is well approximated by $Y^{\rm TP} \approx Y^{\rm dec} \simeq 5 \times 10^{-4} (M_{\rm P}\Gamma_{\tilde{g}}/T_{\rm R}^2)\, e^{-m_{\tilde{g}}/T_{\rm R}}$, and depends sensitively on the reheating temperature.

At still lower temperatures the population of strongly interacting sparticles becomes so tiny that at $T_{\rm R} \sim m_\chi$ neutralino decays given by eq. (3.11) start playing some role, until they too become suppressed by the Boltzmann factor. We indicate this by plotting in figure 2 the contribution of the lightest neutralino (dotted line). By comparing eqs. (3.10) and (3.11) we can easily estimate that the bino-like neutralino contribution is suppressed by $(\alpha_{\rm em}/8\alpha_3)^2\,(m_\chi/m_{\tilde{g}})^3 \sim 10^{-4}$ (with the color factor of 8 included), assuming the usual gaugino mass relations of the MSSM. It is clear that the values of $Y_{\tilde{a}}^{\rm TP}$ in this region are so small that, as we will see later, they will play no role in further discussions. We therefore do not present the effect of the decay of the heavier neutralinos.

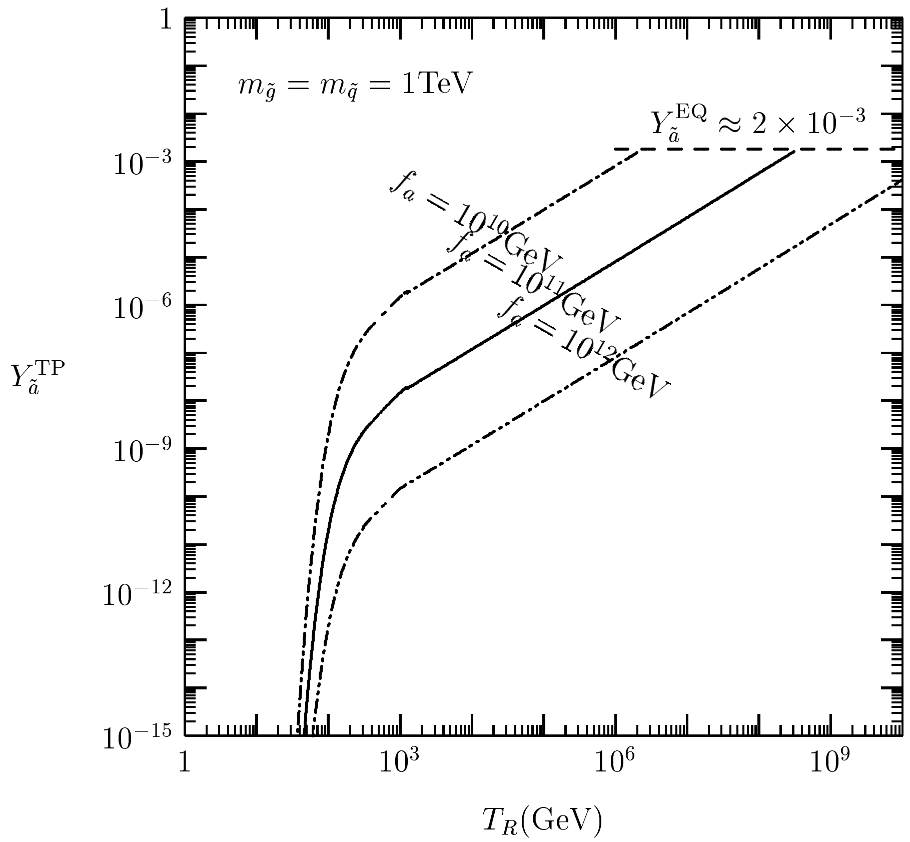

Figure 3: $Y_{\tilde{a}}^{\rm TP}$ as a function of $T_{\rm R}$ for representative values of f_a and $m_{\tilde{q}} = m_{\tilde{g}} = 1\,\text{TeV}$.

As noted at the end of section 2, there are also dimension-4 operators contributing to axino production processes. One such operator is given by $C_{\tilde{a}q\tilde{q}}(m_{\tilde{q}}/f_a)\bar{\tilde{a}}\gamma_5 q\tilde{q}^*$, where $C_{\tilde{a}q\tilde{q}}$ is an effective axino-quark-squark coupling whose size is model-dependent, but which usually arises at the two-loop level. It contributes to processes involving quarks in Table 1. For example, in channel G, in addition to the t-channel gluon exchange diagram drawn in figure 1, there will now be a diagram with a squark exchange in the s-channel. It will contribute $\alpha_s(C_{\tilde{a}q\tilde{q}}m_{\tilde{q}}/f_a)^2(1/s)$ compared to α_s^3/f_a^2 of dimension-5 operators. Hence it is suppressed at high energies, $s \gg m_{\tilde{q}}^2$, by a factor $(C_{\tilde{a}q\tilde{q}}/\alpha_s)^2(m_{\tilde{q}}^2/s)$, and it can be non-negligible only around the mass threshold $s \sim m_{\tilde{q}}^2$. This will, however, have a relatively small effect on the integration of the Boltzmann equation, unless $T_R \sim m_{\tilde{q}}$. For the DFSZ case, there is an additional contribution to the axino production coming from diagrams involving the Higgs supermultiplet in the initial or final state; we neglect such part in the present study since it is model-dependent and we expect it to be subdominant.

The sensitivity of $Y_{\tilde{a}}^{\rm TP}$ to the Peccei–Quinn scale f_a is presented in figure 3. As is clear from eq. (3.9), the rate of axino production is inversely proportional to f_a^2.

We emphasize that axinos produced in this way are already out of equilibrium. Their number density is very much smaller than n_γ (except $T_{\rm R} \sim 10^9\,$GeV and above) and cross sections for axino re-annihilation into other particles are greatly suppressed. This is why in eq. (3.4) we have neglected such processes. Nevertheless, even though axinos never reach equilibrium, their number density may be large enough to give $\Omega_{\widetilde{a}} \sim 1$ for large enough axino masses (keV to GeV range), as we will see later.

Before closing this section, it is worth discussing how the presence of saxions could influence the above results. The saxions will be produced thermally in a way analogous to the axinos but next they will decay with a relatively short lifetime $\tau_{\rm s}$ given by

$$\tau_{\rm s} = 2.65 \times 10^{-6}\,{\rm sec}\, \left(\frac{f_a/N}{10^{11}\,{\rm GeV}} \cdot \frac{0.1}{\alpha_s}\right)^2 \left(\frac{m_s}{1\,{\rm TeV}}\right)^{-3} . \tag{3.12}$$

This time needs to be compared with the time $\tilde{t}_{\rm s} = 2.9 \times 10^{-5}\,(1\,{\rm TeV}/m_s)^2$ when the saxions dominate the energy density. If $\tau_{\rm s} > \tilde{t}_{\rm s}$, the saxion decay generates significant entropy and hence dilutes the particle species decoupled before the decay, such as the axions and the axinos. In other words, the present axion density is lowered and therefore the upper bound on f_a is relaxed to values well beyond $10^{12}\,$GeV. Furthermore, the photon temperature now decreases more slowly, as pointed out in ref. [26]. These effects become important once the saxion mass is greater than $5\,$TeV [27] and f_a is large. In this study, however, we will neglect the reheating effect due to saxion decay since we will assume that the saxion mass is below $1\,$TeV. Finally, in case the saxion mass is much smaller than the SUSY-breaking scale, as for instance in no-scale models, then the saxion decays very late and can play the role of a late-decaying particle [26] instead of generating entropy.

4. Non-thermal production

As discussed in the introduction, axinos may also be produced in decay processes of particles which themselves are out of equilibrium, the decaying particle being one of the ordinary superpartners, the gravitino or the inflaton field. Below we will concentrate on the first possibility.

Let one of the ordinary superpartners be the LOSP and the NLSP (next-to-lightest supersymmetric particle). (They do not have to be the same. The case when the role of the NLSP will be taken by the gravitino will be discussed below.) A natural, albeit not unique, candidate for the LOSP is the lightest neutralino. This is because its mass is often well approximated by the bino mass parameter M_1 (nearly pure bino case) or by the μ-parameter (higgsino limit), neither of which grows much when evolved from the unification scale down to the electroweak scale. It is therefore natural to expect the neutralino to be the LOSP. Furthermore, in models employing full unification of superpartner masses (such as the CMSSM/mSUGRA), a mechanism of electroweak symmetry breaking by radiative corrections typically implies $\mu^2 \gg M_1^2$. As a result, the bino-like neutralino often emerges as the lightest ordinary superpartner [46, 19].

Axino production from bino-like neutralino decay was analyzed in the previous paper [1]. For the sake of completeness, we summarize here the relevant results as well as make additional points regarding the case of a general type of neutralinos. The process involves two steps. The LOSPs first freeze out of thermal equilibrium and next decay into axinos. Since the LOSPs are no longer in equilibrium, the Boltzmann suppression factor in this case does not depend on the thermal-bath temperature and has to be evaluated at the LOSP freeze-out temperature, which is well approximated by $T_{\rm f} \simeq m_\chi/20$. At $T < T_{\rm f}$ and for $\Gamma_\chi \ll H$ but large enough for the decay to occur in the radiation-dominated era, one finds that Y_χ is roughly given by [1]

$$Y_\chi(T) \simeq Y_\chi^{\rm EQ}(T_{\rm f}) \exp\left[-\int_T^{T_{\rm f}} \frac{dT'}{T'^3} \frac{m_\chi^2 \langle\Gamma_\chi\rangle_{T'}}{H(m_\chi)}\right] , \tag{4.1}$$

where $Y_\chi^{\rm EQ}(T_{\rm f})$ contains the Boltzmann suppression factor evaluated at $T_{\rm f}$, $\langle\Gamma_\chi\rangle_T$ is the thermally averaged decay rate for the neutralino at temperature T, and $H(m_\chi)$ is given below eq. (3.4). While we have used here the neutralino as the LOSP, the same mechanism would work for other ordinary superpartners as well.

It is clear that, in order for a two-step process to occur, the decay width of the LOSP has to be sufficiently small to allow for the freeze-out in the first place. Fortunately this is what usually happens since the interactions between the LOSP and the axino are suppressed by the large scale f_a. On the other hand, the lifetime of the LOSP must not be too large, otherwise the decay into axinos and ordinary particles would take place too late, during or after nucleosynthesis, and could destroy successful predictions for the abundance of light elements. This will be discussed in more detail in the next section. Here we note that it is truly remarkable that the relative strength of the axino interaction, in comparison with SM interactions, is such that the decay width falls naturally between these two limits.

It is worth mentioning that in principle the LOSP may decay into axinos even in the absence of freeze-out. It has also been recently pointed out in ref. [50] that even in the case of very low reheating temperatures $T_{\rm R}$, below the LOSP freeze-out temperature, a significant population of them will be generated during the reheating phase. Such LOSPs would then also decay into axinos as above. The process will be non-thermal in the sense that the decaying LOSPs will not be (yet) in thermal equilibrium. We will not pursue this possibility here but comment on it again when we discuss our results.

In ref. [1] we considered the non-equilibrium process

$$\chi \to \widetilde{a}\gamma\,. \tag{4.2}$$

This decay channel is always allowed. The decay rate for process (4.2) can be easily derived from eq. (3.11):

$$\Gamma(\chi \to \widetilde{a}\gamma) = \frac{\alpha_{\rm em}^2 C_{a\chi\gamma}^2}{128\pi^3} \frac{{m_\chi}^3}{(f_a/N)^2} \left(1 - \frac{m_{\widetilde{a}}^2}{m_\chi^2}\right)^3 , \tag{4.3}$$

where $C_{a\chi\gamma} = (C_{aYY}/\cos\theta_W) Z_{11}$, with Z_{11} standing for the bino part of the lightest neutralino.[4]

The neutralino lifetime can be written as

$$\begin{aligned} &\tau(\chi \to \widetilde{a}\gamma) = \\ &= 0.33\text{sec}\frac{1}{C_{aYY}^2 Z_{11}^2} \left(\frac{\alpha_{\rm em}^2}{1/128}\right)^{-2} \left(\frac{f_a/N}{10^{11}\,\text{GeV}}\right)^2 \left(\frac{100\,\text{GeV}}{m_\chi}\right)^3 \left(1 - \frac{m_{\widetilde{a}}^2}{m_\chi^2}\right)^{-3} . \end{aligned} \tag{4.4}$$

For large enough neutralino masses, an additional decay channel into axino and Z opens up,

$$\Gamma(\chi \to \widetilde{a}Z) = \frac{\alpha_{\rm em}^2 C_{a\chi\gamma}^2}{128\pi^3} \tan^2\theta_W \frac{{m_\chi}^3}{(f_a/N)^2} \times \text{PS}\left(\frac{m_Z^2}{m_\chi^2}, \frac{m_{\widetilde{a}}^2}{m_\chi^2}\right), \tag{4.5}$$

where the phase-space factor is given by

$$\text{PS}(x, y) = \sqrt{1 - 2(x+y) + (x-y)^2} \left[(1-y)^2 - \frac{x}{2}(1 - 6\sqrt{y} + y) - \frac{x^2}{2}\right]. \tag{4.6}$$

Note that this channel is always subdominant relative to $\chi \to \widetilde{a}\gamma$ because of both the phase-space suppression and the additional factor of $\tan^2\theta_W$. As a result, even at $m_\chi \gg m_Z, m_{\widetilde{a}}$, $\tau(\chi \to \widetilde{a}Z) \simeq 3.35\,\tau(\chi \to \widetilde{a}\gamma)$. It is also clear that the neutralino lifetime rapidly decreases with its mass ($\sim 1/m_\chi^3$). On the other hand, if the neutralino is not mostly a bino, its decay will be suppressed by the Z_{11}-factor in $C_{a\chi\gamma}$.

Other decay channels are the decay into axino and Standard Model fermion pairs through virtual photon or Z, but they are negligible with respect to the previous ones. We will discuss them later since, for a low neutralino mass, i.e. long lifetime, they can, even if subdominant, produce dangerous hadronic showers during and after nucleosynthesis.

In the DFSZ type of models, there exists an additional Higgs-higgsino-axino coupling, usually related to the MSSM μ-term in simplest realizations. For example, by using the fields S_1 and S_2, which break the PQ symmetry in the superpotential given by (2.5), one can generate the MSSM Higgs mass term through the renormalizable interactions

$$W_\mu = \lambda_i S_i H_u H_d \tag{4.7}$$

[4]In eq. (4.3) we have corrected an overall numerical factor with respect to [1, eq. (6)].

JHEP05(2001)033

for $i = 1, 2$, with a very small $\lambda_i \simeq m_W/f_a$ or through a non-renormalizable interaction like

$$W_\mu = \lambda'_i \frac{S_i^2 H_u H_d}{M_{\rm P}} \,. \tag{4.8}$$

After the PQ symmetry breaking, such superpotentials not only generate the MSSM μ-term [35, 51], but in general also give rise to a Higgs-higgsino-axino coupling of order μ/f_a. In the non-renormalizable case a four-fermion coupling of order μ/f_a^2 is generated. This coupling can be neglected since it does not contribute to the neutralino decay but the other can if the higgsino component of the neutralino is non-negligible. Note that after electroweak symmetry breaking such couplings produce a mixing between the axino and the higgsinos but in general these are of the order of $\mu v/f_a$, and therefore much less than μ, so that we can continue to consider the axino as an approximate mass eigenstate which is nearly decoupled from the MSSM.

An example of such an effective MSSM+axino model has been recently analyzed in ref. [52] for an effective potential of the type

$$W_\mu^{\rm eff} = \mu \left(1 + \frac{\epsilon}{v}\Phi\right) H_u H_d \,, \tag{4.9}$$

where $\epsilon \simeq v/f_a \simeq 10^{-8}$ and v is the Higgs VEV of the order of the weak scale. One needs to perform a full diagonalization of the 5×5 neutralino+axino mass matrix. Since in this case the higgsinos couple more strongly (or rather, less weakly) to the axino, the non-negligible mixing between axino and higgsino introduces additional decay channels, e.g. the one with intermediate virtual right-handed sleptons, that are absent in the KSVZ case.

By rescaling the results of [52, figure 2] to our central value $f_a = 10^{11}\,\text{GeV}$, one finds that for low bino mass, $m_\chi \leq 120\,\text{GeV}$, $m_{\tilde{a}} = 50\,\text{GeV}$, and specific choices of other supersymmetric parameters, in particular masses of the right-handed sleptons degenerate and larger than m_χ, the decay time of the neutralino into axino and lepton or quarks pairs is of the order of

$$\tau(\chi \to \tilde{a} l \bar{l} \text{ and } \tilde{a} q \bar{q}) \simeq 0.02\,\text{sec} \,; \tag{4.10}$$

it decreases at higher masses because of the dependence on m_χ^3 and of the opening of new channels ($b\bar{b}$ and WW^*), down to 3×10^{-6} sec at $m_\chi = 200\,\text{GeV}$. We then see that even in the DFSZ case the lifetime is of the right order of magnitude to allow the freeze-out and decay process since $H(m_\chi) \simeq 2 \times 10^7\,\text{sec}^{-1}(m_\chi/100\,\text{GeV})^2$.

Of course this result strongly depends on the type of DFSZ model considered: for example, the coupling ϵ between the axino and the Higgses can be suppressed or enhanced by mixing angles in the PQ sector since the axion multiplet is in general a combination of different multiplets (S_1 and S_2 in the above example, compare eq. (2.5)). Such angles are assumed to be of order 1 in the simple estimate above.

We therefore conclude that DFSZ-type models have to be analyzed case by case, but in general we expect them to give an even more efficient implementation of the non-thermal production through LOSP decay. A scenario of this type has been recently studied in detail in [40] with the conclusion that an axino (or better in this case flatino) LSP and CDM candidate is still possible and that its population is produced naturally after a phase of thermal inflation, with very low reheat temperature.

While in the discussion of the non-thermal production of axinos, we have concentrated on the neutralino as, in some sense, the most natural choice for a parent LOSP, we reiterate that in principle one could also consider other choices for the LOSP, including charged particles, which we will not do here.

Finally, we note for completeness that another way of producing axinos non-thermally is to consider gravitino decays [25]. In this case it is the gravitino, and not the LOSP, that is the NLSP. Such axinos are produced very late, at 10^8 sec. As we have stated in the Introduction, this channel may solve the gravitino problem [25]. We will come back to this point in the next section. Furthermore, axinos may be produced during (p)reheating in the decay of the inflaton field. Such processes are strongly model-dependent and we will not consider them here.

5. Constraints

Several non-trivial conditions have to be satisfied in order for axinos to be a viable CDM candidate. First, we will expect their relic abundance to be large enough, $\Omega_{\tilde{a}} h^2 \simeq 0.2$. This obvious condition will have a strong impact on other bounds. Next, the axinos generated through both TP and NTP will in most cases be initially relativistic. We will therefore require that they become non-relativistic, or cold, much before the era of matter dominance. Furthermore, since NTP axinos will be produced near the time of BBN, we will require that they do not contribute too much relativistic energy density to radiation during BBN. Finally, associated decay products of axino production will often result in electromagnetic and hadronic showers which, if too large, would cause too much destruction of light elements. In deriving all of these conditions, except for the first one, the lifetime of the parent LOSP will be of crucial importance.

First, from now on we will assume that the axinos give a dominant contribution to the matter density at the present time. This can be expressed as

$$\Omega_{\tilde{a}} h^2 = m_{\tilde{a}} Y_{\tilde{a}} \frac{s(T_{\rm now})}{\rho_{\rm crit}/h^2} \simeq 0.2\,, \tag{5.1}$$

where s is given under eq. (3.5) and we have assumed no significant entropy production. Since $\rho_{\rm crit}/h^2 = 0.8 \times 10^{-46}\,{\rm GeV}^4$ and $g_{s*}(T_{\rm now}) = 3.91$, this can be re-expressed as

$$m_{\tilde{a}} Y_{\tilde{a}} \simeq 0.72\,{\rm eV} \left(\frac{\Omega_{\tilde{a}} h^2}{0.2} \right), \tag{5.2}$$

which readily applies to both TP and NTP relics.

We note in passing that, for the initial population of axinos, the yield at decoupling is approximately $Y_{\tilde{a}} \simeq Y_{\tilde{a}}^{\rm EQ} \simeq 2 \times 10^{-3}$, which gives

$$m_{\tilde{a}} \simeq 0.36\,\mathrm{keV}\left(\frac{\Omega_{\tilde{a}}h^2}{0.2}\right). \tag{5.3}$$

This is an updated value for the Rajagopal-Turner-Wilczek bound (3.3).

For a fixed value of $\Omega_{\tilde{a}}h^2$, lower values of $Y_{\tilde{a}}$ give correspondingly larger $m_{\tilde{a}}$. For example, the turn-over region of $Y_{\tilde{a}} \simeq 10^{-10}$ in figure 2, where gluino decay becomes more important than scattering, gives $m_{\tilde{a}} \simeq 7.2\,\mathrm{GeV}\,(\Omega_{\tilde{a}}h^2/0.2)$.

Cold axinos. Next, we want to determine the temperature of the Universe at which the axinos will become non-relativistic. In nearly all cases axinos are initially relativistic and, due to expansion, will become non-relativistic at some later epoch, which depends on their mass and production mechanism. As mentioned in the Introduction, in the case of thermal production, even though the axinos are not in thermal equilibrium, they are produced in kinetic equilibrium with the thermal-bath. Their momenta therefore have a thermal spectrum, derived from the scattering particles in the plasma. The axinos become non-relativistic when the thermal-bath temperature reaches the axino mass:

$$T_{\rm NR} \simeq m_{\tilde{a}}\,, \tag{5.4}$$

even though they are not in thermal equilibrium.

In the case of non-thermal production, the momentum of the axinos depends strongly on the production mechanism. Here we will consider, as before, only the production through out-of-equilibrium neutralinos. In that case, axinos will be produced basically monochromatically, all with the same energy, roughly given by $m_\chi/2$, unless they are nearly mass-degenerate with the neutralinos. This is so because the neutralinos, when they decay, are themselves already non-relativistic. Thus the axinos will normally become non-relativistic only at later times, through momentum red-shift. However, in this case the simple estimate (5.4) used above for $T_{\rm NR}$ does not hold, since the axinos are not in kinetic equilibrium with the thermal-bath.

First we need to determine the temperature $T_{\rm dec}$ corresponding to the time of neutralino decay $t_{\rm dec}$. We will make the approximation that most axinos are produced in neutralino decays at the time equal to the neutralino lifetime $\tau_\chi = 1/\Gamma_\chi$, $t_{\rm dec} \simeq \tau_\chi$ (a sudden-decay approximation), which is a very reasonable one.

Because not all neutralinos decay at the same time as assumed in the sudden-decay approximation, at a given time the momenta of produced axinos will actually have a distribution that will be limited from above by $m_\chi/2$. Let us find this momentum distribution $f_{\tilde{a}}(t_0, p)$ of axinos coming from neutralino decays. When the axino was produced by the decay of neutralino at time t, it had the momentum $p = m_\chi/2$, and would be next red-shifted while the axino was relativistic. Since the scale factor

JHEP05(2001)033

R grows as $R \propto t^{1/2}$ during the RD epoch, at a later time t_0, the axino produced at a time t would have the momentum

$$p(t, t_0) = \frac{m_\chi}{2} \left(\frac{R(t)}{R(t_0)} \right) = \frac{m_\chi}{2} \left(\frac{t}{t_0} \right)^{1/2} . \qquad (5.5)$$

By counting the number of produced axinos during the time interval dt at t, we get

$$N_\chi \frac{dt}{\tau_\chi} = f_{\tilde{a}}(t_0, p) dp \,, \qquad (5.6)$$

where $N_\chi(t) = N_{\chi 0} e^{-t/\tau_\chi}$ is the number of neutralinos at t which is obtained from the decay equation $dN_\chi/dt = (1/\tau_\chi) N_\chi$. Thus,

$$f_{\tilde{a}}(t_0, p) = \frac{N_\chi}{\tau_\chi} \frac{dt}{dp} = N_{\chi 0} \frac{2p}{p(\tau_\chi, t_0)^2} e^{-(p/p(\tau_\chi, t_0))^2} . \qquad (5.7)$$

From this, we can find the total energy of axinos

$$E_{\tilde{a}}(t_0) = \int_0^{m_\chi/2} p f_{\tilde{a}}(t_0, p) dp = N_0 p(\tau_\chi, t_0) \left[\frac{\sqrt{\pi}}{2} \mathrm{Erf} \left[\left(\frac{t_0}{\tau_\chi} \right)^{1/2} \right] - \left(\frac{t_0}{\tau_\chi} \right)^{1/2} e^{-t_0/\tau_\chi} \right] . \qquad (5.8)$$

For $t_0 \gg \tau_\chi$, the energy density of axinos becomes

$$\rho_{\tilde{a}} = \frac{E_{\tilde{a}}(t_0)}{R(t_0)^3} = \frac{\sqrt{\pi}}{2} \frac{m_\chi}{2} \left(\frac{\tau_\chi}{t_0} \right)^{1/2} \frac{N_{\chi 0}}{R(t_0)^3} . \qquad (5.9)$$

This differs from the sudden-decay approximation only by a factor $\sqrt{\pi}/2 \approx 0.89$.

Using the sudden-decay approximation we obtain

$$\begin{aligned} T_{\mathrm{dec}} &= \left(\frac{90\, \Gamma_\chi^2 M_{\mathrm{Pl}}^2}{4\pi^2 g_*(T_{\mathrm{dec}})} \right)^{1/4} \\ &= 1.08 \times 10^9 \,\mathrm{GeV} \left(\frac{\Gamma_\chi}{\mathrm{GeV}} \right)^{1/2} \simeq 0.9 \,\mathrm{MeV} \left(\frac{\mathrm{sec}}{\tau_\chi} \right)^{1/2} , \end{aligned} \qquad (5.10)$$

where we have taken $g_*(T_{\mathrm{dec}}) \simeq 10$.

Axino momenta will red-shift with temperature as

$$p(T) = \frac{m_\chi}{2} \left(\frac{g_{s*}(T)}{g_{s*}(T_{\mathrm{dec}})} \right)^{1/3} \frac{T}{T_{\mathrm{dec}}} . \qquad (5.11)$$

as long as the axinos remain relativistic. (The factor involving g_{s*} at different times accounts for a possible reheating of the thermal-bath when some species become non-relativistic.) Axinos become non-relativistic at the epoch when $p(T_{\mathrm{NR}}) \simeq m_{\tilde{a}}$, which gives

$$\begin{aligned} T_{\mathrm{NR}} &= 2 \frac{m_{\tilde{a}}}{m_\chi} \left(\frac{g_{s*}(T_{\mathrm{dec}})}{g_{s*}(T_{\mathrm{NR}})} \right)^{1/3} T_{\mathrm{dec}} \\ &= 2.7 \times 10^{-5} m_{\tilde{a}} \left(\frac{100\,\mathrm{GeV}}{m_\chi} \right) \left(\frac{T_{\mathrm{dec}}}{1\,\mathrm{MeV}} \right) , \end{aligned} \qquad (5.12)$$

where we have taken $g_{s*}(T_{\mathrm{dec}}) \simeq g_*(T_{\mathrm{dec}}) \simeq 10$ and $g_{s*}(T_{\mathrm{NR}}) = 3.91$.

Using an explicit expression for $T_{\rm dec}$ (eq. (5.10)) then leads to

$$T_{\rm NR} = 2\,\frac{m_{\widetilde a}}{m_\chi}\left(\frac{g_{s*}(T_{\rm dec})}{g_{s*}(T_{\rm NR})}\right)^{1/3}\left(\frac{90\Gamma_\chi^2 M_{\rm Pl}^2}{\pi^2 g_*(T_{\rm dec})}\right)^{1/4} \tag{5.13}$$
$$= 53.8\,{\rm keV}\left(\frac{g_{s*}(T_{\rm dec})}{g_{s*}(T_{\rm NR})}\right)^{1/3}\frac{C_{aYY}Z_{11}}{g_*^{1/4}(T_{\rm dec})}\left(\frac{m_{\widetilde a}}{1\,{\rm GeV}}\right)\left(\frac{m_\chi}{100\,{\rm GeV}}\right)^{1/2}\left(\frac{10^{11}\,{\rm GeV}}{f_a/N}\right).$$

For $T_{\rm dec} < 200\,{\rm MeV}$, $g_{s*}(T_{\rm dec}) \simeq g_*(T_{\rm dec}) \simeq 10$ and for $T < 500\,{\rm keV}$, $g_{s*}(T) \simeq 3.91$, which finally leads to

$$T_{\rm NR} = 4.2\times 10^{-5} m_{\widetilde a} C_{aYY} Z_{11}\left(\frac{m_\chi}{100\,{\rm GeV}}\right)^{1/2}\left(\frac{10^{11}\,{\rm GeV}}{f_a/N}\right). \tag{5.14}$$

This epoch has to be compared with the matter-radiation equality epoch. Assuming that axinos constitute the largest part of the matter density, $\rho_{\rm matter}(T_{\rm eq}) = \rho_{\widetilde a}(T_{\rm eq})$, and since

$$\rho_{\widetilde a} = m_{\widetilde a} Y_{\widetilde a} s(T) = m_{\widetilde a} Y_{\widetilde a}\frac{2\pi^2}{45} g_{s*}(T)\,T^3\,, \tag{5.15}$$

the temperature at matter-radiation equality is given by

$$T_{\rm eq} = \frac{4 g_{s*}(T_{\rm eq})}{3 g_*(T_{\rm eq})} m_{\widetilde a} Y_{\widetilde a}\,. \tag{5.16}$$

Using eq. (5.2) and assuming $g_{s*}(T_{\rm eq}) = 3.91$ and $g_*(T_{\rm eq}) = 3.36$ allows us to express eq. (5.16) as

$$T_{\rm eq} = 1.1\,{\rm eV}\left(\frac{\Omega_{\widetilde a}h^2}{0.2}\right), \tag{5.17}$$

which holds for both thermal and non-thermal production.

In the former case, by comparing with eq. (5.4), one can easily see that $T_{\rm NR} > T_{\rm eq}$ is satisfied for any interesting range of $m_{\widetilde a}$. In the case of NTP, in eq. (5.16) we make a substitution $Y_{\widetilde a} = Y_\chi(T_{\rm f})$, the neutralino yield at freeze-out. The condition $T_{\rm NR} \gg T_{\rm eq}$ is satisfied for

$$m_{\widetilde a} \gg 41\,{\rm keV}\left(\frac{m_\chi}{100\,{\rm GeV}}\right)\left(\frac{1\,{\rm MeV}}{T_{\rm dec}}\right)\left(\frac{\Omega_{\widetilde a}h^2}{0.2}\right)$$
$$\gg 27\,{\rm keV}\frac{1}{C_{aYY}Z_{11}}\left(\frac{100\,{\rm GeV}}{m_\chi}\right)^{1/2}\left(\frac{f_a/N}{10^{11}\,{\rm GeV}}\right)\left(\frac{\Omega_{\widetilde a}h^2}{0.2}\right), \tag{5.18}$$

where the $T_{\rm dec}$ dependence on the neutralino lifetime is given in eq. (5.10). If axinos were lighter than the bound (5.18), then the point of radiation-matter equality would be shifted to a later time around $T_{\rm NR}$. Note that in this case the axino would not constitute cold, but warm or hot dark matter. We will see, however, that in the NTP case discussed here other constraints would require the axino mass to be larger than the above bound, so that we can discard this possibility.

BBN constraint on NTP axinos. In the case of non-thermal production, most axinos will be produced only shortly before nucleosynthesis and, being still relativistic, may dump too much to the energy density during the formation of light elements. The axino energy density now reads

$$\begin{aligned}
\rho_{\widetilde{a}}(T) &= Y_{\widetilde{a}} s(T) p(T) \\
&= \left(\frac{2\pi^2}{45}\right)\left(\frac{g_{s*}(T)^{4/3}}{g_{s*}(T_{\rm dec})^{1/3}}\right)\left(\frac{m_\chi Y_{\widetilde{a}}}{2T_{\rm dec}}\right) T^4 \\
&= 1.6\times 10^{-7}\left(\frac{g_{s*}(T)^{4/3}}{g_{s*}(T_{\rm dec})^{1/3}}\right)\left(\frac{m_\chi}{m_{\widetilde{a}}}\right)\left(\frac{1\,{\rm MeV}}{T_{\rm dec}}\right)\left(\frac{\Omega_{\widetilde{a}} h^2}{0.2}\right) T^4\,, \qquad (5.19)
\end{aligned}$$

where eqs. (5.2), (5.11), and an expression for the entropy density s given under eq. (3.5) have been used. In order not to affect the Universe's expansion during BBN, the axino contribution to the energy density should satisfy

$$\frac{\rho_{\widetilde{a}}}{\rho_\nu} \leq \delta N_\nu\,, \qquad (5.20)$$

where $\rho_\nu = (\pi^2/30)(7/4)T^4$ is the energy density of one neutrino species. Agreement with observations of light elements requires [53] $\delta N_\nu = 0.2$ to 1. Using eq. (5.10) leads, after some simple algebra, to

$$\begin{aligned}
m_{\widetilde{a}} &\gtrsim 274\,{\rm keV}\frac{1}{\delta N_\nu}\left(\frac{m_\chi}{100\,{\rm GeV}}\right)\left(\frac{1\,{\rm MeV}}{T_{\rm dec}}\right)\left(\frac{\Omega_{\widetilde{a}} h^2}{0.2}\right) \\
&\gtrsim 181\,{\rm keV}\frac{1}{\delta N_\nu}\frac{1}{C_{aYY} Z_{11}}\left(\frac{100\,{\rm GeV}}{m_\chi}\right)^{1/2}\left(\frac{f_a/N}{10^{11}\,{\rm GeV}}\right)\left(\frac{\Omega_{\widetilde{a}} h^2}{0.2}\right), \qquad (5.21)
\end{aligned}$$

where we have chosen the most conservative case $g_{s*}(T) = g_{s*}(T_{\rm dec}) \simeq 10$. (For $g_{s*}(T) = 3.91$ the bounds become respectively 78 keV and 55 keV.) As before, $T_{\rm dec}$ depends on the neutralino lifetime via eq. (5.10).

Showers. In the NTP case, if neutralino decays take place during or after BBN, produced bosons may lead to a significant depletion of primordial elements [18]. One often applies the crude constraint that the lifetime be less than about 1 sec. This in our case would provide a lower bound on m_χ. A detailed analysis [53] provides limits on the abundance of the decaying particle versus its lifetime. (See, e.g. [53, figure 3].)

First, the photons produced in reaction (4.2) carry a large amount of energy, roughly $m_\chi/2$. If the decay takes place before BBN, the photon will rapidly thermalize via multiple scatterings from background electrons and positrons ($\gamma + e \to \gamma + \gamma + e$) [54, 15]. The process will be particularly efficient at plasma temperatures above 1 MeV, which is the threshold for background $e\bar{e}$ pair annihilation, and which, incidentally, coincides with a time of about 1 sec. But a closer examination shows that also scattering with the high-energy tail of the CMBR thermalize photons very efficiently. As a result, the decay lifetime into photons can be as large as 10^4 sec. By comparing this with eq. (4.4) we find that, in the gaugino regime, this can be easily satisfied for $m_\chi < m_Z$. It is only in a nearly pure higgsino case and for a mass of tens of GeV that the bound would become constraining. We are not interested in such light higgsinos for other reasons, as will be explained later.

A much more stringent constraint comes from considering hadronic showers from $q\bar{q}$-pairs. These will be produced through a virtual photon and Z exchange, and, above the kinematic threshold for $\chi \to \widetilde{a}Z$, also through the exchange of a real Z-boson. We will now discuss this in some detail.

By comparing with Fig. 3 of ref. [53], we can see that the bound on hadronic showers can be written as

$$\frac{m_\chi n_\chi}{n_\gamma} \times BR(\chi \to q\bar{q}) < F(\tau_\chi), \tag{5.22}$$

where $BR(\chi \to q\bar{q})$ is the branching ratio of neutralinos into axinos plus $q\bar{q}$-pairs and $F(\tau_\chi)$ can be read out from ref. [53, figure 3].

Since $Y_{\widetilde{a}} \simeq Y_\chi(T_{\rm f})$, and remembering that

$$\frac{n_\chi}{n_\gamma} = \frac{\pi^4 g_{s*}}{45\zeta(3)} Y_\chi \simeq 7.04\, Y_\chi\,, \tag{5.23}$$

one can express the condition (5.22) as

$$\frac{m_{\widetilde{a}}}{m_\chi} > 5.07 \left(\frac{10^{-9}\,\mathrm{GeV}}{F(\tau_\chi)}\right)\left(\frac{\Omega_{\widetilde{a}}h^2}{0.2}\right) \times BR(\chi \to q\bar{q})\,. \tag{5.24}$$

As stated above, one can write

$$BR(\chi \to q\bar{q}) = BR(\chi \to \gamma^* \to q\bar{q}) + BR(\chi \to Z/Z^* \to q\bar{q}) + \text{interference term}\,. \tag{5.25}$$

Assuming $m_{\widetilde{a}} \ll m_\chi, m_Z$, in the region $m_\chi < m_Z$, we find that $BR(\chi \to \gamma^* \to q\bar{q})$ dominates and (summing over all light quark species up to the bottom) is of the order of 0.03–0.04, slowly increasing with m_χ. Above the Z threshold the intermediate Z channel also becomes sizeable, and the branching ratio grows faster with m_χ, up to $\simeq 0.06$ at $m_\chi = 150\,\mathrm{GeV}$. The interference part is always negligible.

We have computed $BR(\chi \to q\bar{q})$ both below and above the threshold for the process $\chi \to \widetilde{a}Z$ for negligible axino mass and used it in our numerical analysis. Our results are presented in Fig. 4 in the bino case $Z_{11}C_{aYY} \simeq 1$ and for $f_a/N = 10^{11}\,\mathrm{GeV}$. For the general case, one can still read out the lower bound on the ratio $m_{\widetilde{a}}/m_\chi$ from Fig. 4 by replacing m_χ with $(C_{aYY}Z_{11})^{2/3}(10^{11}\,\mathrm{GeV}/f_a/N)^{2/3}m_\chi$. In fact the main dependence on m_χ in eq. (5.24) comes from the factor m_χ and the neutralino lifetime in $F(\tau_\chi)$, while $BR(\chi \to q\bar{q})$ only slowly increases with m_χ.

These bounds are clearly much more stringent than those in (5.18) and (5.21), but they are, at the same time, strongly sensitive to the neutralino mass and composition.

JHEP05(2001)033

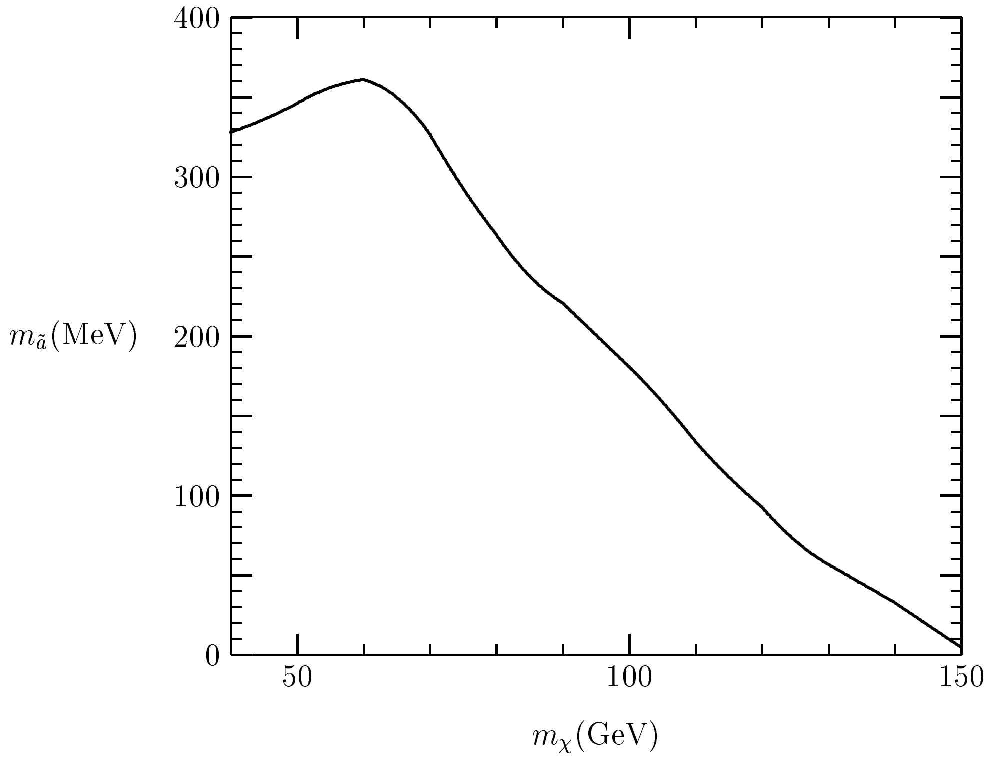

Figure 4: Lower bound on the axino mass from considering hadronic showers according to the condition (5.24), for $C_{aYY}Z_{11} = 1$ and $f_a/N = 10^{11}$ GeV. The bound disappears for $m_\chi = 150$ GeV when the lifetime drops below 0.1 sec.

Notice that, in the region of low m_χ, the bound on $m_{\tilde{a}}$ gradually strengthens with increasing m_χ (and therefore decreasing neutralino lifetime) because, for lifetimes of order 2–20 sec, the function $F(\tau_\chi)$ is shallow and the branching ratio and χ both increase. However, as the lifetime drops below 1 sec, $F(\tau_\chi)$ increases steeply and so the $m_{\tilde{a}}$ bound decreases almost linearly in m_χ before disappearing altogether for $m_\chi \gtrsim 150$ GeV (in the bino case) when $\tau_\chi \lesssim 0.1$ sec.

If the higgsino component of the decaying neutralino increases, so does the lifetime. There are two points to note here. One is that the neutralino yield now becomes much smaller, owing to co-annihilation with the next to lightest neutralino and lightest chargino until $m_\chi \gtrsim 500$ GeV, when the lifetime will be suppressed again. Furthermore, in the DFSZ model, new channels are present, as we discussed in the previous section, for which a typical lifetime will be very much smaller than 1 sec [52]. Only in the KSVZ model in the higgsino-like case, in a relatively small region $60 \lesssim m_\chi \lesssim 150$ GeV (the lower number being the current rough experimental bound), will there exist some restriction on the combination of neutralino mass and higgsino purity, which would further depend on the neutralino yield at freeze-out. Thus we conclude that overall we find no significant restriction on the neutralino mass from the lifetime constraint.

In summary, a lower bound $m_{\tilde{a}} \gtrsim \mathcal{O}(300\,\mathrm{keV})$ arises from requiring either the axinos to be cold at the time of matter dominance or that they do not contribute too much to the relativistic energy density during BBN. The constraint from hadronic destruction of light elements can be as strong as $m_{\tilde{a}} \gtrsim 360\,\mathrm{MeV}$ (in the relatively light bino case), but it is highly model-dependent and disappears for larger m_χ.

6. Thermal vs. non-thermal production

We now proceed to comparing the relic abundance of axinos produced non-thermally in neutralino decays with the thermal production case analyzed in section 3. Clearly, in the TP case the axino yield is primarily determined by the reheating temperature (for a fixed f_a). For large enough T_{R} ($T_{\mathrm{R}} \gg m_{\tilde{g}}, m_{\tilde{q}}$), it is proportional to T_{R}/f_a^2. (Compare figures 2 and 3) For the reheating temperature below the mass threshold of strongly interacting sparticles, it becomes Boltzmann suppressed by the factor $\sim \exp^{-m_{\tilde{g}}/T_{\mathrm{R}}}$ (or $\sim \exp^{-m_{\tilde{q}}/T_{\mathrm{R}}}$).

If the axino is the LSP, the present fraction of the axino energy density to the critical density is given by eq. (5.2). This relation allows us to redisplay the results of figure 2 in the plane $(m_{\tilde{a}}, T_{\mathrm{R}})$. For thermally regenerated axinos, eq. (5.2) gives an upper bound on T_{R} as a function of the axino mass from the requirement $\Omega_{\tilde{a}}^{\mathrm{TP}} h^2 \lesssim 1$. This bound is depicted by a thick solid line in figure 5.

A digression is in order here. An expert reader will have noticed that, while in the axino case, the bound $\Omega_{\tilde{a}} h^2 \lesssim 1$ gives $T_{\mathrm{R}} \sim 1/m_{\tilde{a}}$, in the gravitino case the analogous bound gives $T_{\mathrm{R}} \sim m_{\tilde{G}}$. (See, e.g., ref. [8, figure 1].) This difference is caused by the fact that the crucial effective lagrangian dimension-5 operator, which is responsible for the bound, is proportional to $1/m_{\tilde{G}}$ in the gravitino case, but exhibits no $m_{\tilde{a}}$ dependence in the case studied here. (Compare eq. (2) or ref. [8] with eq. (2.8).)

Since axinos couple like $1/f_a$, one would naively expect that T_{R} for which TP becomes important (or when too much relic abundance is generated) would be just f_a^2/M_{P}^2 of that for gravitinos and thus hopelessly low, $T_{\mathrm{R}} \lesssim 10^{-7}\,\mathrm{GeV}$. This is, however, not the case, since the gravitino production is dominated by the goldstino component, whose interaction is suppressed by the supersymmetry-breaking scale M_S, rather than the Planck scale: for example, the coupling to the gluino is $\propto m_{\tilde{g}}/M_S^2$. Thus T_{R} at which TP becomes significant will be of order 10^3–$10^4\,\mathrm{GeV}$, which is on the low side but still acceptable.

In the NTP case, the yield of axinos is just the same as that of the decaying neutralinos. This leads to the following simple relation [1]

$$\Omega_{\tilde{a}} h^2 = \frac{m_{\tilde{a}}}{m_\chi}\, \Omega_\chi h^2\,. \tag{6.1}$$

We remind the reader that $\Omega_\chi h^2$ stands for the abundance that the neutralinos would have had today, had they not decayed into axinos. It is related to the neutralino yield

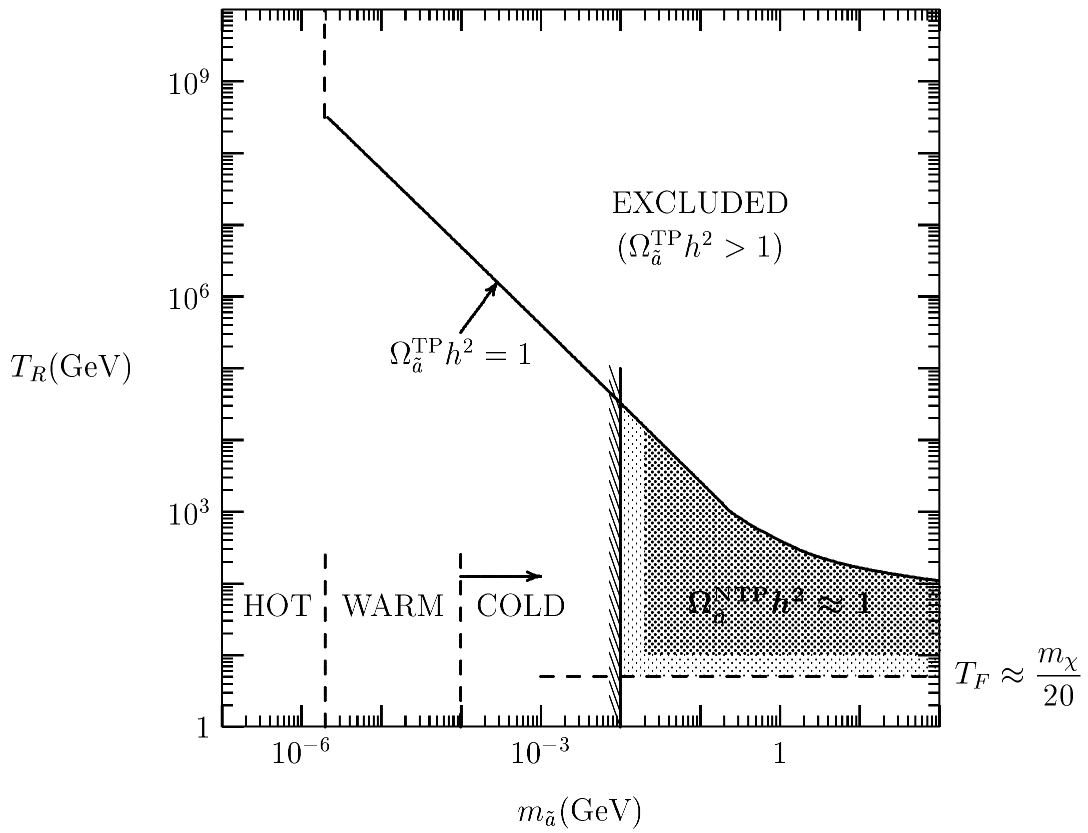

Figure 5: The solid line gives the upper bound from thermal production on the reheating temperature as a function of the axino mass. The dark region is the region where non-thermal production can give cosmologically interesting results ($\Omega_{\tilde a}^{\rm NTP} h^2 \simeq 1$) as explained in the text. We assume a bino-like neutralino with $m_\chi = 100\,\text{GeV}$ and $f_a = 10^{11}\,\text{GeV}$. The region of $T_{\rm R} \gtrsim T_{\rm f}$ is somewhat uncertain and is shown with light-grey color. A sizeable abundance of neutralinos (and therefore axinos) is expected also for $T_{\rm R} \lesssim T_{\rm f}$ [50] but has not been calculated. The vertical light-grey band indicates that a low range of $m_{\tilde a}$ corresponds to allowing SM superpartner masses in the multi-TeV range, as discussed in the text. The division of hot, warm and cold dark matter as a function of the axino mass shown in the lower left part is for axinos from non-thermal production.

through an analogue of eq. (5.2) (which actually applies to any stable massive species) and is determined by the effective cross section of neutralino pair-annihilation (as well as co-annihilation) into ordinary particles. When the rate of this process becomes less than the rate of the expansion of the Universe, the neutralinos freeze out. Typically this happens at freeze-out temperatures of $T_{\rm f} \simeq m_\chi/20$.

In contrast to the TP case, the NTP axino yield will also be independent of the reheating temperature (so long as $T_{\rm R} \gg T_{\rm f}$). In order to be able to compare the two production mechanisms, we will therefore fix the neutralino mass at some typical value. Furthermore we will map out a cosmologically interesting range of axino masses for which $\Omega_{\tilde{a}}^{\rm NTP} \sim 1$.

Our results are presented in Fig. 5 in the case of a nearly pure bino. We also fix $m_\chi = 100\,{\rm GeV}$ and $f_a = 10^{11}\,{\rm GeV}$. The dark region is derived in the following way. It is well known that $\Omega_\chi h^2$, the relic abundance of neutralinos, can take a wide range of values spanning several orders of magnitude. In the framework of the MSSM, which we have adopted, global scans give $\Omega_\chi h^2 \lesssim 10^4$ in the bino region at $m_\chi \lesssim 100\,{\rm GeV}$. (This limit decreases roughly linearly (on a log-log scale) down to $\sim 10^3$ at $m_\chi \simeq 400\,{\rm GeV}$.) For $m_\chi = 100\,{\rm GeV}$, by using eq. (6.1) we find that the expectation $\Omega_{\tilde{a}}^{\rm NTP} h^2 \simeq 1$ gives

$$10\,{\rm MeV} \lesssim m_{\tilde{a}} \lesssim m_\chi. \tag{6.2}$$

We note, however, that the upper bound $\Omega_\chi h^2 \lesssim 10^4$ comes from allowing very large $M_{\rm SUSY}$ (i.e. sfermion and heavy Higgs masses) in the range of tens of TeV. Restricting all SUSY mass parameter below about 1 TeV reduces $\Omega_\chi h^2$ below 10^2 and, accordingly, increases the lower bound $m_{\tilde{a}} \gtrsim 1\,{\rm GeV}$. For the sake of generality, in Fig. 5 we have kept the much more generous bound (6.2) but we marked a low range of $m_{\tilde{a}}$ with a light grey band to indicate the above point.

Likewise, for reheating temperatures just above $T_{\rm f}$, standard estimates of $\Omega_\chi h^2$ become questionable. We have therefore indicated this range of $T_{\rm R}$ with again light grey color. It has also been recently pointed out in ref. [50] that a significant population of LOSPs will be generated during the reheating phase even at $T_{\rm R}$ below the LOSP freeze-out temperature. Such LOSPs would then also decay into axinos as above. We have not considered such cases in our analysis and accordingly left the region $T_{\rm R} < T_{\rm f}$ blank, even though in principle we would expect some sizeable range of $\Omega_{\tilde{a}} h^2$ there.

We can see that, for large $T_{\rm R}$, the TP mechanism is more important than the NTP one, as expected. Note also that in the TP case the cosmologically favored region ($0.2 \lesssim \Omega_{\tilde{a}} h^2 \lesssim 0.4$) would form a very narrow strip (not indicated in figure 5) just below the $\Omega_{\tilde{a}}^{\rm TP} = 1$ boundary. In contrast, the NTP mechanism can give the cosmologically interesting range of axino's relic abundance for a relatively wide range of $m_{\tilde{a}}$ so long as $T_{\rm R} \lesssim 5 \times 10^4\,{\rm GeV}$. Perhaps in this sense, the NTP mechanism can be considered as somewhat more robust.

We have also marked in figure 5 some of the bounds discussed in section 5. They are normally not as restrictive as the shaded region in the figure. For example, the ranges of $m_{\tilde{a}}$ where non-thermally produced axinos would be hot/warm/cold dark matter are denoted, following eqs. (3.3) and (5.18). The potentially most stringent bound from hadronic showers would require $m_{\tilde{a}} \gtrsim 360\,{\rm MeV}$, but it is very sensitive to m_χ and disappears for $m_\chi > 150\,{\rm GeV}$ (compare figure 4).

JHEP05(2001)033

At larger bino mass, the lower bound on $m_{\tilde{a}}$ also increases, following eq. (6.1) and because $\Omega_\chi h^2$ decreases with m_χ. Other bounds do not give any additional constraints. In the higgsino case, one typically has $\Omega_\chi h^2 \ll 1$ (or, more properly, the higgsino number density at freeze-out is very much smaller than that of the bino with the same mass) owing mostly to co-annihilation, as mentioned at the end of the last section. A heavy higgsino with $m_\chi \gtrsim 500\,\mathrm{GeV}$ nevertheless remains an option for a LOSP. In this case, however, additional model-dependent (dimension-4) operators will contribute to the TP mechanism, and the upper curve in figure 5 will probably be moved upwards. Overall we find the bino case to be much more natural and robust.

7. Implications and conclusions

The intriguing possibility that the axino is the LSP and the dark matter WIMP possesses a number of very distinct features. This makes this case very different from those of both the neutralino and the gravitino. In particular, the axino can be a cold DM WIMP for a rather wide range of masses in the MeV to GeV range and for relatively low reheating temperatures $T_{\mathrm{R}} \lesssim 5 \times 10^4\,\mathrm{GeV}$. As T_{R} increases, thermal production of axinos starts dominating over non-thermal production and the axino typically becomes a warm DM relic with a mass broadly in the keV range. In contrast, the neutralino is typically a cold DM WIMP (although see ref. [55]).

Low reheating temperatures would favor baryogenesis at the electroweak scale. It would also alleviate the nagging “gravitino problem”. If additionally it is the axino that is the LSP and the gravitino is the NLSP, the gravitino problem is resolved altogether for both low and high T_{R}.

Phenomenologically, one faces a well-justified possibility that the bound $\Omega_\chi h^2 < 1$, which is often imposed in constraining a SUSY parameter space, may be readily avoided. In fact, the range $\Omega_\chi h^2 \gg 1$ (and with it typically large masses of superpartners) would now be favored if the axino is to be a dominant component of DM in the Universe. Furthermore, the lightest ordinary superpartner could either be neutral or charged but would appear stable in collider searches.

The axino, with its exceedingly tiny coupling to other matter, will be a real challenge to experimentalists. It is much more plausible that a supersymmetric particle and the axion will be found first. Unless the neutralino (or some other WIMP) is detected in DM searches, the axino will remain an attractive and robust candidate for solving the outstanding puzzle of the nature of dark matter in the Universe.

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JHEP05(2001)033

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JHEP05(2001)033