Dark and visible matter with broken R-parity and the axion multiplet

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Abstract

A small breaking of R-parity reconciles thermal leptogenesis, gravitino dark matter and primordial nucleosynthesis. We find that the same breaking relaxes cosmological bounds on the axion multiplet. Naturally expected spectra become allowed and bounds from late particle decays become so weak that they are superseded by bounds from non-thermal axion production. In this sense, the strong CP problem serves as an additional motivation for broken R-parity.
1 Introduction

A consistent cosmology has to explain the observed matter composition of the Universe. Standard thermal leptogenesis [1] provides a simple and elegant explanation for the origin of matter. However, thermal leptogenesis cannot be reconciled with an unstable gravitino in supergravity [2, 3] unless it is very heavy. Fortunately, the gravitino is one of the best-motivated candidates for dark matter. If it is the lightest superparticle, the condition that relic gravitinos do not overclose the Universe yields an upper bound on the reheating temperature [4]. It is remarkable that a reheating temperature of $\mathcal{O}(10^{10} \text{ GeV})$ can account for the observed matter from leptogenesis and at the same time for cold dark matter in the form of thermally produced relic gravitinos with a mass of $\mathcal{O}(100 \text{ GeV})$. Interestingly, the strong CP problem can be solved by the Peccei-Quinn mechanism [5, 6] in this cosmological setting. However, the next-to-lightest superparticle (NLSP) becomes so long-lived that the strongest conflict between both notions arises from NLSP decays that spoil the success of primordial nucleosynthesis [7]. This and axion physics lead to tighter constraints on the axion multiplet [8].

It has been shown that in the case of small R-parity breaking thermal leptogenesis, gravitino dark matter and primordial nucleosynthesis are naturally consistent [9]. Since a consistent cosmology should also enable a solution to the strong CP problem, we study in this letter how far the restrictions on the axion multiplet are softened, if R-parity is broken.

The small R-parity breaking couplings allow the lightest ordinary superparticle (LOSP) to decay into pairs of Standard Model particles. These decays happen instantaneously compared to the Hubble time after LOSP freeze-out or the R-parity-conserving decay into the gravitino. Thus they do not endanger the success of Big Bang nucleosynthesis (BBN). In this way the aforementioned NLSP decay problem is circumvented, which at the same time relaxes the constraint on the axino lifetime and thus the allowed range of the axino mass and possibly of other parameters. It is allowed to decay right before BBN, instead of the requirement with conserved R-parity to decay before NLSP freeze-out. Thus its decay temperature can be lowered by three orders of magnitude. In addition, saxion decays are allowed to produce NLSPs at any time before BBN.

The effect on the LOSP decay is the crucial impact of broken R-parity on the constraints on the axion multiplet. All members of the axion multiplet (axion $a$, axino $\tilde{a}$ and saxion $\phi_{sax}$) obtain additional R-parity breaking couplings, but R-parity breaking interactions and decays are suppressed by the Peccei-Quinn scale and additionally by the small R-parity breaking couplings. Thus they are produced and decay in the same way as if R-parity were conserved. One exception to this statement occurs, if the axino is the NLSP. In the following we fix the gravitino mass to be of $\mathcal{O}(100 \text{ GeV})$ and the reheating temperature at a rather large value of $\mathcal{O}(10^{10} \text{ GeV})$. This is a worst case scenario for the axion multiplet. We will comment on the DFSZ [10, 11] and KSVZ [12, 13] invisible axion models, especially regarding the possibility of an axino next-to-NLSP. Nevertheless our analysis is general and can be applied to axions in superstring models [14]. Our main results are comprised in Table 1. Since, for example, naturally expected spectra become allowed, our results serve as an additional motivation for broken R-parity.
2 Axino and saxion with conserved R-parity

**Axino**  Even if axinos do not enter thermal equilibrium after inflation, they are, nevertheless, regenerated by thermal scatterings and decays in the thermal plasma. The resulting density can be estimated in units of today’s critical density as

\[ \Omega_{\tilde{a}} h^2 \approx 7.8 \times 10^6 \left( \frac{m_{\tilde{a}}}{1 \text{ TeV}} \right) \left( \frac{T_R}{10^9 \text{ GeV}} \right) \left( \frac{10^{12} \text{ GeV}}{f_a} \right)^2 , \]  

(1)

where \( m_{\tilde{a}} \) and \( T_R \) denote the axino mass and the reheating temperature, respectively. The axion decay constant \( f_a = f_{PQ}/N \) with \( f_{PQ} \) denoting the Peccei-Quinn scale if \( N \) denotes the number of different vacua. For the KSVZ (DFSZ) model \( N = 1 \) (6). The produced LOSP\(^1\) and/or gravitino abundances from axino decay are orders of magnitude larger than the thermal abundances even if re-annihilation is taken into account [18, 19]. The produced LOSP abundance is inconsistent with the scenario, because the NLSP decay problem is worsened. Therefore, the decay is demanded to happen before LOSP freeze-out, so that one does not have to worry about the produced number of LOSPs since they thermalise normally. In this case the lower bound on the axino mass reads

\[ m_{\tilde{a}} \gtrsim 6 \times 10^2 \text{ GeV} \left( 1 - \frac{m_{\tilde{a}}^2}{m_{\tilde{g}}^2} \right)^{-1} \left( \frac{m_{\text{losp}}}{10^2 \text{ GeV}} \right)^2 \left( \frac{f_a}{10^{10} \text{ GeV}} \right)^2 \left( \frac{g_*(T_{\text{dec}}^{\tilde{a}})}{100} \right)^{1/4} , \]  

(2)

where \( m_{\tilde{g}} \) denotes the gluino mass, \( g_\star \) the effective number of relativistic degrees of freedom of the Universe and we have used that the freeze-out temperature \( T_{\text{losp}}^{\tilde{a}} \approx m_{\text{losp}}/25 \). By \( T_{\text{dec}}^{\tilde{a}} \) we denote the temperature of the Universe after the axino decay. Since in supergravity the axino mass is—as the other superparticle masses—generically of the order of the gravitino mass [20], i.e., \( m_{\tilde{a}} \sim m_{\text{susy}} \sim m_{3/2} \), the mass bound (2) can be interpreted as an upper bound on the axion decay constant, i.e., \( f_a \lesssim 10^{10} \text{ GeV} \). In (2) the decay into a gluino-gluon pair with width [15]

\[ \Gamma_{\tilde{a}g} = \frac{\alpha_s^2}{16\pi^3} \left( \frac{m_{\tilde{a}}}{f_a} \right)^2 \left( 1 - \frac{m_{\tilde{g}}^2}{m_{\tilde{a}}^2} \right)^3 , \]  

(3)

is assumed to be dominant and the strong coupling constant \( \alpha_s(\mu) = \alpha_s(m_{\tilde{a}}) \approx 0.1 \). The involved operator exists for any axion model that is able to solve the strong CP problem. Investigating other axino decay channels in the KSVZ model, we find that they require \( m_{\tilde{a}} \) to be larger than the expected gluino mass, which keeps (3) to be the dominant decay channel. Independently of \( f_a \) the axino must be sufficiently heavier than the gluino, which is expected to be among the heavier superparticles due to the running of its mass. For instance, with a gluino mass \( m_{\tilde{g}} = 1 \text{ TeV} \) and the parameter values appearing in (2) the axino mass is required to be larger than about 1.35 TeV.

\(^1\)The lightest ordinary superparticle (LOSP) is the lightest superparticle of the minimal supersymmetric standard model (MSSM). In this sense, gravitino and axino are extraordinary superparticles.
The DFSZ superfield is directly coupled to the Higgs superfields. Therefore, the DFSZ axino may dominantly decay into a Higgsino-Higgs pair. Using \( L = \frac{\mu}{f_a} (H_u H_d^0 + H_d H_u^0) \tilde{a} + \text{h.c.} \), we obtain in the decoupling limit\(^2\)

\[
\Gamma(\tilde{a} \rightarrow \tilde{H} + h) \simeq \frac{m_{\tilde{a}}}{4\pi} \sum_i |N_{\chi^0_i H_d} \sin \beta + N_{\chi^0_i H_u} \cos \beta|^2 \left( \frac{\mu}{f_a} \right)^2 ,
\]

(4)

where \( \mu \) is the usual supersymmetric \( \mu \)-parameter, \( \tan \beta = v_u/v_d \) is the ratio of the Higgs vacuum expectation values and \( N_{\chi^0_i H_d} \) and \( N_{\chi^0_i H_u} \) denote the Higgsino fractions of the \( i \)-th neutralino. We consider the case of degenerate, maximally mixed Higgsinos, which are then decoupled from the gauginos. For sufficiently large \( \tan \beta \) the mixing factor in (4) then becomes about \( 1/2 \). Thus, the sum of both Higgsinos softens the mass bound to

\[
m_{\tilde{a}}^{\text{dFsZ}} \gtrsim 5.4 \times 10^2 \, \text{GeV} \left( \frac{f_a}{10^{11} \, \text{GeV}} \right)^2 \left( \frac{100 \, \text{GeV}}{\mu} \right)^2 \left( \frac{m_{\text{kep}}}{10^2 \, \text{GeV}} \right)^2 \left( \frac{g_s(T_{\text{dec}})}{100} \right)^{1/2} .
\]

(5)

The upper bound on the Peccei-Quinn scale from late axino decay may thus be softened to \( f_a \lesssim 10^{11} \, \text{GeV} \) and for \( f_a \lesssim 10^{10} \, \text{GeV} \) we observe that the open decay channel into Higgsino-Higgs suffices in the DFSZ model. Other decay channels are subdominant compared to (3) and (4).

Altogether, only spectra with \( m_{\tilde{a}} > m_{\tilde{H}} \) are allowed in all axion models. This can be viewed as a problem, because it requires the axino to be heavier than naturally expected. The bound may become softened to \( m_{\tilde{a}} > m_{\tilde{H}} + m_h \) in the DFSZ model for a small \( \mu \)-parameter.

**Saxion**

Due to supersymmetry the couplings of the saxion have the same strength as the axino couplings. Thus the saxion is produced as efficiently as the axino. Its dominant decay producing MSSM particles is into a pair of gluons with [22]

\[
\Gamma_{g\bar{g}}^{\text{sax}} = \frac{\alpha_s^2}{32\pi^3} \frac{m_{sax}^3}{f_a^2}
\]

(6)

or—for the DFSZ model—into a pair of light Higgses with [23, 21]

\[
\Gamma(\phi_{\text{sax}} \rightarrow 2h) = \frac{m_{sax}}{8\pi} \left( 1 - \frac{4 f_a^2}{m_{sax}^2} \right)^{1/2} \left( \frac{\mu}{f_a} \right)^2 \left( \frac{\mu}{m_{sax}} \right)^2 .
\]

(7)

Since in these decays no superparticles are produced, they may happen right before BBN. Since the saxion receives its mass from SUSY breaking, one expects \( m_{sax} \sim m_{\text{susy}} \). If

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\(^2\)The result seems quite insensitive to the possibility of the Guidice-Masiero mechanism to generate the \( \mu \)-term, but this might be different for the saxion decay into two light Higgses [21]. For the neglected phase space factor we find \((1 + x_h - x_h^2)(1 + x_h + x_h^2)((1 - x_h)^2 - x_h^2))^{1/2} \) with \( x_h = m_{\chi^0_i} / m_{\tilde{a}} \) and \( x_h = m_h / m_{\tilde{a}} \).
the saxion is heavy enough to produce superparticle pairs, its decay could lead to the same worsening of the NLSP decay problem as the axino decay. In this situation the lower bound on the saxion mass from the decay into a pair of gluons becomes in the end

\[ m_{\text{sax}} \gtrsim 7.6 \times 10^2 \text{ GeV} \left( \frac{m_{\text{hosp}}}{10^2 \text{ GeV}} \right)^{\frac{1}{2}} \left( \frac{f_a}{10^{10} \text{ GeV}} \right)^{\frac{3}{2}} \left( \frac{g_* (T_{\text{dec}})}{100} \right)^{\frac{1}{2}}. \tag{8} \]

If the DFSZ saxion decays dominantly into Higgs pairs, we can derive an upper bound on the saxion mass

\[ m_{\text{sax}} \lesssim 740 \text{ GeV} \left( 1 - \frac{4m_{\text{h}}^2}{m_{\text{sax}}^2} \right)^{\frac{1}{2}} \left( \frac{10^{11} \text{ GeV}}{f_a} \right)^{2} \left( \frac{\mu}{300 \text{ GeV}} \right)^{4} \left( \frac{10^2 \text{ GeV}}{m_{\text{nlsp}}} \right)^{2} \left( \frac{100}{g_* (T_{\text{dec}})} \right)^{\frac{1}{2}} \tag{9} \]

We see that for \( f_a \lesssim 10^{10} \text{ GeV} \) the DFSZ saxion mass is only required to allow for the decay into a pair of light Higgses.

In addition, the saxion may decay as well into two axions with

\[ \Gamma_{\text{aa}} \simeq \frac{x^2 m_{\text{sax}}^3}{64\pi f_a^2}, \tag{10} \]

where the self-coupling \( x \) can be of order 1. The produced axions represent a form of dark radiation, i.e., decoupled, relativistic particles not present in the Standard Model. During BBN the energy density of dark radiation \( \rho_{\text{dr}} \) is constrained to be less than the energy density of one additional neutrino species \[ \rho_{\text{SM}} \mid_{\text{BBN}} \lesssim 0.14. \tag{11} \]

We have approximated the decay width of the saxion as \[ \Gamma_{\text{sax}} \simeq \Gamma_{\text{gg}} + \Gamma_{\text{aa}}. \tag{12} \]

Thus our conclusion should hold qualitatively for any axion model. Then the branching ratio reduces to

\[ B_{\text{aa}} \simeq 0.4 (1 + 50\pi^2 x^2)^{\frac{1}{2}} \left( \frac{10^{10} \text{ GeV}}{f_a} \right)^{\frac{1}{2}} \left( \frac{m_{\text{sax}}}{10^2 \text{ GeV}} \right)^{\frac{1}{2}} \left( \frac{Y_{\text{eq}}}{Y_{\text{sax}}} \right) \left( \frac{g_* (T_{\text{dec}})}{10.75} \right)^{\frac{1}{2}}. \tag{13} \]

In this inequality we have approximated the decay width of the saxion as \( \Gamma_{\text{sax}} \simeq \Gamma_{\text{gg}} + \Gamma_{\text{aa}} \). Thus our conclusion should hold qualitatively for any axion model. Then the branching ratio reduces to

\[ B_{\text{aa}} \simeq \frac{\Gamma_{\text{aa}}}{\Gamma_{\text{sax}}} = \frac{x^2}{x^2 + 2\alpha_s^2/\pi^2}. \tag{14} \]

\[ \text{Our expression differs from the literature in the numerical value and the dependence on } g_* . \text{ This is because we have taken into account properly the scaling of } \rho_{\text{dr}} \propto g_*^{4/3} \rho_{\text{SM}}^{1/3}. \text{ The simplifying assumption of entropy conservation } (g_* = \text{const.}), \text{ which is usually made, does indeed not hold during these stages of the Universe.} \]
The value for $f_a$ appearing in (13) corresponds to the upper bound on $f_a$ from axino decay and the axion energy density $\Omega_a$ (see below). Since in our scenario the reheating temperature is fixed at rather large values, all members of the axion multiplet enter thermal equilibrium after inflation for such small values of $f_a$. Therefore the saxion yield $Y_{sax}$ cannot be smaller than the equilibrium value $Y_{sax}^{eq} \simeq 1.21 \times 10^{-3}$. The appearing values for $f_a$ and $m_{sax}$ are chosen to show the worst situation in the considered scenario. Like a smaller $f_a$ or a larger $m_{sax}$, also a larger $x$ leads to an earlier decay, which corresponds to a smaller $\Omega_{sax}$ at its decay. Thus there is a self-curing effect for large $x$. The bound (13) represents indeed an implicit equation for the self-coupling $x$.

Evaluating it for $x$ it turns out that there is no constraint on $x$ at all in the scenario under consideration.

We point out that the absence of any bound on $x$ is due to the expectation $m_{sax} \sim m_{susy}$ and the restriction of $f_a$ to small values appropriate for the considered scenario.

Furthermore, $Y_{sax}$ could be much larger than $Y_{sax}^{eq}$, if the saxion is produced from coherent oscillations after inflation as discussed below. The saxion might decay before LOSP freeze-out, if the self-coupling is strong enough. Neglecting conservatively the saxion decay into gluons and Higgses we find

$$x \gtrsim 0.9 \left( \frac{10^2 \text{GeV}}{m_{sax}} \right)^{\frac{2}{7}} \left( \frac{f_a}{10^{10} \text{ GeV}} \right) \left( \frac{m_{\tilde{\text{g}}}}{10^2 \text{ GeV}} \right) \left( \frac{g_\ast(T_{\text{dec}}^{\tilde{\text{g}}})}{61.75} \right)^{\frac{3}{7}}.$$  (15)

In this situation there is no additional constraint on the saxion mass. Note that the decay into axions does not produce a significant amount of entropy.

If a new best-fit value demands additional radiation energy in the Universe [26, 27, 28, 29], which is often parameterised by a change of the effective neutrino degrees of freedom $\Delta N_{\text{eff}} > 0$, we can determine parameter values from (13), such that the additional energy is formed by axions from saxion decay. For $f_a \leq 10^{10}$ GeV and $m_{sax} \geq 10^2$ GeV the maximal $\Delta N_{\text{eff}}$ is 0.6. However, when these requirements are relaxed also larger $\Delta N_{\text{eff}}$ are possible. For example, we obtain $\Delta N_{\text{eff}} \simeq 1$ with a rather small saxion mass $m_{sax} = 10$ GeV, $f_a = 10^{10}$ GeV and a self-coupling $x = 0.1$.

### 3 R-parity violating case

**Axino** Since in the case of broken R-parity the NLSP decay problem is absent, the axino may decay right before BBN. Produced superparticles decay promptly into particles of the Standard Model, which thermalise normally. Then the lower bound (2) becomes

$$m_{\tilde{a}} \gtrsim 89 \text{ GeV} \left( 1 - \frac{m_\tilde{\text{g}}^2}{m_\tilde{a}^2} \right)^{-1} \left( \frac{f_a}{10^{12} \text{ GeV}} \right)^{\frac{2}{7}} \left( \frac{T_{\text{dec}}^{\tilde{\text{g}}}}{4 \text{ MeV}} \right)^{\frac{2}{7}},$$  (16)

where we still assume the decay channel $\tilde{a} \to \tilde{\text{g}}g$ to dominate and take $T_{\text{dec}}^{\tilde{\text{g}}}$ as lower bound on the temperature after the particle decay. The bound $T_{\text{dec}}^{\tilde{\text{g}}} = 4$ MeV, corresponding to a lifetime $\tau_{\text{max}} \sim 0.05$ s, considers the neutrino energy density during the neutrino thermalisation process [30, 31, 32], if the particle dominated the energy density
of the Universe before its decay. This is likely to happen in the considered scenario (see below). Weaker bounds $T^\text{dec}_\text{min} \sim 0.7 \text{ MeV}$ arise from BBN calculations [33, 34]. As in the following we omit the dependence on $g_a(T^\text{dec}_\text{min} = T^\text{dec}_\text{min}) = 10.75$. We see that due to $T^\text{dec}_\text{min} \ll T^\text{lo}_\text{NLSP}$ the axino mass bound becomes so weak that masses much smaller than the expected gluino mass would become allowed.

In this situation the axino decays always depend on which channels are kinematically open and thus, in principle, on the full spectrum. This is a qualitative difference to the R-parity conserving case. Particularly interesting is the possibility of an axino next-to-NLSP, which we will assume in the following. Then its dominant decay is fixed into the NLSP, in other words, into the LOSP. From (16) we see that the NLSP were allowed to be a gluino. A recent analysis reports $m_{\tilde{g}} > 322 \text{ GeV}$ as experimental lower limit for such a long-lived gluino [35]. From (5) we see that the DFSZ model would in addition allow for a Higgsino NLSP, if the corresponding decay channel were kinematically open.

Since the lightest neutralino is likely one of the lightest superparticles in the spectrum, one interesting decay is into neutralino and photon with [15]

$$\Gamma(\tilde{a} \to \chi_1^0 + \gamma) = \frac{\alpha^2_{\text{em}} C^2_{a\chi_1^0, \gamma} m_\tilde{a}^3}{128\pi^3 \bar{f}_a^2} \left(1 - \frac{m_{\chi_1^0}^2}{m_{\tilde{a}}^2}\right)^3,$$  \hspace{1cm} (17)

where $C_{a\chi_1^0, \gamma} = (C_{aBB}/\cos \Theta_W) N_{\chi_1^0 B}$, while $N_{\chi_1^0 B}$ is the bino fraction of the $i$-th neutralino and $\Theta_W$ denotes the weak mixing angle. We take the electromagnetic coupling constant $\alpha_{\text{em}}(\mu) = \alpha_{\text{em}}(m_{\tilde{a}}) \simeq 1/128$. The axion to two $B$ bosons coupling $C_{aBB}$ varies for different implementations of different axion models.\footnote{For example, in the DFSZ model with $(d', e)$ unification, $C_{aBB} = 8/3$. In the KSVZ model, for different electromagnetic charges of the heavy quark $e_Q = 0, -1/3, 2/3$, $C_{aBB} = 0, 2/3, 8/3$, respectively. Below the QCD scale, $C_{aBB}$ is reduced by 1.92 [36].}

For simplicity we set $C_{aBB} = N_{\chi_1^0 B} = 1$. We find

$$m_{\tilde{a}} \gtrsim 45 \text{ GeV} \left(1 - \frac{m_{\chi_1^0}}{m_{\tilde{a}}}\right)^{-1} \left(\frac{\bar{f}_a}{10^{10} \text{ GeV}}\right)^{2/3} \left(\frac{T^\text{dec}_\text{min}}{4 \text{ MeV}}\right)^{2/3},$$  \hspace{1cm} (18)

and conclude that a neutralino with a substantial bino component can be the NLSP, if the axino is the next-to-NLSP. Thus a superparticle spectrum with $m_{\tilde{g}} > m_{\tilde{a}} > m_{\chi_1^0}$ is now easily possible without any BBN conflict.

If the axino decay is fixed into a sneutrino, it decays via an intermediate neutralino into a photon and a sneutrino-neutrino pair, i.e., $\tilde{a} \to \gamma \chi_1^{0*} \to \gamma \tilde{\nu}_{\nu}$. This process is suppressed compared to (17) by an additional power of $\alpha_{\text{em}}$ and further factors depending on the neutralino composition and the exact spectrum. Thus in case of a sneutrino NLSP an axino next-to-NLSP might be possible only for parameter values at the boundaries of the allowed region and a tuned spectrum. Consequently, this situation is disfavoured.

For other sfermion NLSPs the situation depends much more strongly on the axion model. In the DFSZ model we estimate the kinematically unsuppressed decay width of...
the axino into a stop-top pair using the Lagrangian of [37] as

$$\Gamma(\tilde{a} \rightarrow \tilde{t} + t) \simeq \frac{m_\tilde{a}}{16\pi} \left( \frac{m_t X_u}{f_\tilde{a}} \right)^2, \quad (19)$$

where $m_t$ denotes the fermion mass and $X_u = 1/(\tan^2 \beta + 1)$. This decay arises from a dimension-four operator that becomes important at low decay temperatures and accordingly small masses. If only this channel were open, the resulting lower mass bound on the axino would become

$$m_\tilde{a} \gtrsim 2.4 \times 10^2 \text{ GeV} \left( \frac{f_\tilde{a}}{10^{12} \text{ GeV}} \right)^2 \left( \frac{T_{\text{dec}}^{\text{min}}}{4 \text{ MeV}} \right)^2 \left( \frac{173 \text{ GeV}}{m_t} \right)^4. \quad (20)$$

The lower bound would practically be given by $f_\tilde{a}$ and the requirement to have the channel kinematically unsuppressed, since a long-lived stop with mass below 249 GeV is excluded according to the CDF experiment [38]. It might well happen that the decay of the next-to-NLSP ($\tilde{a}$) into the NLSP ($\tilde{t}$) and its superpartner ($t$) is kinematically forbidden. The axino would decay violating R-parity. This is the same if the $\tilde{a}$ is the NLSP. We comment on this excluded case below.

The decay width into other sfermions is given by (19), if $m_t$ is replaced by the corresponding fermion mass and $X_u \rightarrow X_d = \tan^2 \beta/(1 + \tan^2 \beta)$ if appropriate. For example, the decay width into stau and tau is (19) with the replacements $m_t \rightarrow m_\tau \simeq 1.78 \text{ GeV}$ and $X_u \rightarrow X_d$. The corresponding lower bound on the axino mass becomes

$$m_\tilde{a} \gtrsim 2.3 \times 10^2 \text{ GeV} \left( \frac{f_\tilde{a}}{10^{12} \text{ GeV}} \right)^2 \left( \frac{T_{\text{dec}}^{\text{min}}}{4 \text{ MeV}} \right)^2. \quad (21)$$

Thus, for $f_\tilde{a}$ below $10^{12}$ GeV also this channel is restricted by the requirement of kinematic accessibility only. Since $m_\tau/m_\tilde{a} \ll 1$, this is not a strong constraint, if the stau is not much heavier than the gravitino. Consequently, if the stau is the NLSP, the DFSZ axino may be the next-to-NLSP.

Since the other leptons are lighter, in these cases the bound becomes tighter. While the smuon stays possible, the selectron would require $m_\tilde{a} \gtrsim 1 \text{ TeV}$ even for parameter values at the boundaries. These considerations expand to the quarks lighter than the top. The superpartners of the bottom, charm, and strange quarks are possibilities. For the down squark the Peccei-Quinn scale needs to be close to the lower limit. The situation for the up squark is even worse than the one for the selectron.

In the KSVZ model axino-sfermion-fermion interactions are one-loop-suppressed in the low-energy effective theory [39, 40], which weakens axino to sfermion-fermion decays substantially relative to the DFSZ case. The tree-level decay via intermediate neutralino into other sfermion-fermion pairs is comparable to the decay into sneutrino-neutrino. Consequently, the KSVZ axino is disfavoured, if the axino decay is fixed into a sfermion-fermion pair.

If the axino were the NLSP, i.e., for a spectrum with $m_{\tilde{\chi}} > m_\tilde{a} > m_3/2$, it would decay either into the gravitino or via R-parity violation. Both decay modes are strongly
suppressed by the Planck scale or by the Peccei-Quinn scale and the R-parity violating coupling, respectively. Therefore, the lifetime of the axino would always become much larger than the time of BBN. This is independent of whether R-parity is broken by bilinear or trilinear couplings \[41, 42, 43\]. Since \(\Omega_\tilde{a} \gg 1\), such a late decay would spoil the predictions of BBN. We conclude that an axino NLSP stays excluded in the R-parity violating scenario.

**Saxion** If R-parity is violated also superparticle pairs from saxion decays are harmless. If its dominant decay channel is into a pair of (massless) gluons, the most severe bound on the saxion mass is found to be similar to \((16)\). So for \(f_a \sim 10^{10} \text{ GeV}\) the saxion mass is practically not constrained in the R-parity violating case. However, considerations concerning the self-coupling \(x\) are not affected by R-parity violation.

If late-decaying particles enter thermal equilibrium after inflation, they likely dominate the energy density of the Universe at their decay, since the energy density of matter \(\rho_{\text{mat}}\) grows relative to the energy density of radiation \(\rho_{\text{rad}}\) as the scale factor \(a\) of the Universe, i.e., \(\rho_{\text{mat}}/\rho_{\text{rad}} \propto a\). If a particle dominates the energy density of the Universe, it might produce considerable entropy at its decay, which could dilute cosmological abundances to too small values. Axino and saxion enter thermal equilibrium at \(T_R \sim 10^{10} \text{ GeV}\) if \(f_a \lesssim 10^{11.5} \text{ GeV}\). The dilution factor due to their decay can be estimated to be \([44, 8]\)

\[
\Delta \simeq 2 \left( \frac{f_a}{10^{11} \text{ GeV}} \right) \left( \frac{1 \text{ TeV}}{m_{\tilde{a}/sax}} \right)^{1/2} \left( \frac{g^*(T_{\text{dec}})}{61.75} \right)^{1/4},
\]

(22)

assuming the decay into a gluino-gluon pair \((3)\) and two gluons \((6)\), respectively, to be dominant. The numerical difference between axino and saxion is small. Here, the upper bound on the dilution factor \(\Delta \geq 1\) arises from leptogenesis itself and is \(\Delta_{\text{max}} \sim 10^4\). Thus the occurring dilution is so small that there arises no constraint from entropy production by the thermally produced axion multiplet.

The saxion may also be produced by coherent oscillations around its potential minimum after inflation. The temperature at the onset of saxion oscillations is

\[
T_{\text{osc}}^{\text{sax}} \simeq 2.2 \times 10^{10} \text{ GeV} \left( \frac{m_{\text{sax}}}{1 \text{ TeV}} \right)^{1/2} \left( \frac{228.75}{g^*(T_{\text{osc}}^{\text{sax}})} \right)^{1/4}.
\]

(23)

If the reheating temperature \(T_R < T_{\text{osc}}^{\text{sax}}\) as in the standard scenario, the produced saxion abundance is \([24]\)

\[
\frac{\rho_{\text{osc}}^{\text{sax}}}{s} = \frac{1}{8} T_R \left( \frac{\phi_{\text{sax}}^i}{M_{\text{Pl}}} \right)^2 \simeq 4.2 \times 10^{-9} \text{ GeV} \left( \frac{T_R}{2 \times 10^9 \text{ GeV}} \right)^2 \left( \frac{f_a}{10^{10} \text{ GeV}} \right)^{2} \left( \frac{\phi_{\text{sax}}^i}{f_a} \right)^2,
\]

(24)

where \(\phi_{\text{sax}}^i\) denotes the initial amplitude of the oscillations and where we have assumed the simplest saxion potential, \(V = \frac{1}{2} m_{\text{sax}}^2 \phi_{\text{sax}}^2\).
The dilution factor due to saxion decays $\Delta \simeq 0.75 \frac{T_{\text{sax}}}{T_{\text{dec}}}$ with $T_{\text{sax}} = \frac{4}{3} \frac{\rho_{\text{osc}}}{s}$ is given by

$$\Delta \simeq 4 \times 10^{-10} \left( \frac{T_{\text{R}}}{2 \times 10^9 \text{ GeV}} \right) \left( \frac{f_a}{10^{10} \text{ GeV}} \right)^3 \left( \frac{\phi_{\text{max}}^s}{f_a} \right)^2 \left( \frac{1 \text{ TeV}}{m_{\text{sax}}} \right)^3 \left( \frac{g_*(T_{\text{dec}})}{10.75} \right)^{\frac{1}{2}}.$$

Since the dilution factor is by definition $\geq 1$, values for $\Delta$ smaller than one lead to the standard scenario with $\Delta = 1$. In this discussion we fix the reheating temperature at its lower boundary from thermal leptogenesis, $T_{\text{R}}^\text{min} = 2 \times 10^9 \text{ GeV}$ [45] or $2 \Delta \times 10^9 \text{ GeV}$ with $\Delta > 1$. The tightest bound on the initial amplitude in the standard scenario is found for the maximal axion decay constant $f_a = 10^{10} \text{ GeV}$ and a small saxion mass, for concreteness $m_{\text{sax}} = 10^2 \text{ GeV}$. It is $\phi_{\text{max}}^s \lesssim 7 \times 10^{13} \text{ GeV}$. In comparison, the loosest bound is found for the minimal axion decay constant $f_a = 6 \times 10^8 \text{ GeV}$ and a rather large saxion mass, for concreteness $m_{\text{sax}} = 1 \text{ TeV}$. It is $\phi_{\text{max}}^s \lesssim 1.4 \times 10^{15} \text{ GeV}$. Thus both bounds are far below the Planck scale. Bounds on the initial amplitude of saxion oscillations are not affected by R-parity violation. Conservatively allowing for smaller saxion masses as well, we summarise these bounds in Table 1 as $\phi_{\text{sax}}^i \lesssim (10^{13} - 10^{15}) \text{ GeV}$.

4 Constraints from the axion

First, the lower limit on the axion decay constant from axion physics [46], $f_a \gtrsim 6 \times 10^8 \text{ GeV}$, is not changed by the additional R-parity violating interactions. Furthermore, the axion still decays harmlessly into two photons with a lifetime much larger than the age of the Universe. Due to its tiny mass, the thermal relic density of axions is negligible. In contrast, a large axion density may be produced non-thermally by vacuum misalignment and topological defects.

The density produced by vacuum misalignment [47, 48, 49] is usually given as

$$\Omega_{a}^{\text{mis}} h^2 \sim a_0^2 \left( \frac{N f_a}{10^{12} \text{ GeV}} \right)^{7/6},$$

where $a_0$ comprises model-dependent factors. In the presence of gravitino dark matter we require $\Omega_a/\Omega_{\text{DM}} = r \ll 1$, which gives an upper bound

$$f_a < 10^{10} \text{ GeV} \quad \text{for} \quad a_0 = N = 1, \ r = 0.04.$$

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Note that this bound becomes tighter for larger $N$.

Topological defects occur, if the Peccei-Quinn symmetry is restored after inflation, i.e., $T_R > T_{PQ} \sim f_{PQ} = N f_a$. This leads to the formation of disastrous domain walls [50], if $N > 1$, and cosmic strings [51]. The abundance of axions from cosmic strings $\Omega_a^{\text{str}}$ exceeds that from vacuum misalignment [52, 53]. Since here the reheating temperature is fixed at about $2\Delta \times 10^9$ GeV, the situation is particularly interesting. There are two possible cases: i) topological defects are not created after inflation, so we do not have to care about them or ii) they occur and we have to take them into consideration. To avoid topological defects completely,

$$f_a > 2 \times 10^9 \text{ GeV} \frac{\Delta}{N},$$  

which could lead, depending on $N$, to a stronger lower bound on $f_a$ than given above. In the case without entropy production ($\Delta = 1$), this favours models with $N \geq 4$ such as the DFSZ model. If (28) is violated, $N = 1$—fulfilled by the KSVZ model—still avoids domain walls and the axion density from strings, $\Omega_a^{\text{str}} \sim 10 \times \Omega_a^{\text{mis}}$ [52, 53], gives a tighter upper bound than (27),

$$f_a < 1.3 \times 10^9 \text{ GeV} \quad \text{for} \quad a_0' = N = \Delta = 1, \ r = 0.04,$$  

where $a_0'$ comprises model-dependent factors at the production of axions from cosmic strings. Combining the above considerations, only the small interval $(1.3 - 2) \times 10^9$ GeV for $f_a$ might be excluded. In this sense, the allowed band for $f_a$ is not changed.

With entropy production ($\Delta > 1$) the reheating temperature is raised to compensate the dilution of the baryon asymmetry, so the Peccei-Quinn symmetry becomes restored for a larger range of values of $f_a$. We consider entropy production after the QCD phase transition by late particle decay and estimate the axion abundance from cosmic strings as $\Omega_a^{\text{str}} \sim 10 \times \Omega_a^{\text{mis}} / \Delta$ with $\Omega_a^{\text{mis}}$ as in (26). Then the upper bound (29) on $f_a$ is softened by a factor of $\Delta^{6/7}$, because the axions are diluted, while the amount produced remains the same. The maximal $\Delta^{\text{max}} \sim 10^4$ corresponds to $f_a \lesssim 4 \times 10^{12}$ GeV, so in this situation the Peccei-Quinn symmetry is probably restored, because $T_R \sim 2 \times 10^{13}$ GeV. We therefore list this upper bound on $f_a$ in Table 1. Larger values are possible, if the symmetry is not restored and if one accepts the axino to be heavier than the gluino.

Also the upper bound on $f_a$ from vacuum misalignment is affected by $\Delta > 1$ if the entropy is produced after the QCD phase transition. In this case the Universe is likely to be dominated by matter at the onset of axion oscillations at $T_a^{\text{dec}} \sim 1$ GeV and the upper bound on $f_a$ becomes [54]

$$\left( \frac{f_a}{10^{14} \text{ GeV}} \right)^2 \lesssim (r/0.02) \left( \frac{4 \text{ MeV}}{T_{\text{dec}}} \right),$$  

where $T_{\text{dec}}$ denotes as before the temperature of the Universe after the entropy-producing particle decayed. In this case the axion decay constant is constrained more strongly by $\Omega_a^{\text{str}}$ and too late axino and saxion decays. Altogether, a large dilution factor opens up
—at least for \( N = 1 \)—more parameter space, but the situation depends on the time of entropy production.

We would like to point out the appealing possibility that the saxion decay gives rise to the maximal dilution after the QCD phase transition. For the harmless value \( f_a = 10^{10} \) GeV this is the case for an initial amplitude of saxion oscillations equal to the geometric mean of the two scales involved, i.e., \( \phi_{\text{sax}}^i \sim \sqrt{f_a M_P} \), and a quite small saxion mass \( m_{\text{sax}} \sim 10 \) GeV [8]. However, it is understood that an adjusted initial amplitude might produce the desired amount of entropy for any of the allowed parameter values.

The given bounds from the different production mechanisms of axions refer in each case to “standard values” in parameter space. They can be relaxed or circumvented in “non-standard” scenarios. Constructing models that realise a small \( \Omega_a \) is beyond the scope of this letter.

5 Conclusions and outlook

We have presented how the constraints on the axion multiplet in scenarios of thermal leptogenesis with gravitino dark matter are relaxed, if R-parity is broken. Our results are comprised in Table 1. Obviously, the most important findings also hold in less restricted scenarios.

The saxion mass becomes practically unconstrained. An axino NLSP stays excluded, but the axino may be anywhere else in the superparticle mass spectrum as long as a decay channel into an ordinary superparticle of the MSSM is kinematically open and allowed at tree-level in the low-energy effective theory. This does not hold if the only possible decay is into a light fermion (up-quark, electron, neutrino) and its superpartner. In turn, for the DFSZ axion model, usually considered superparticle spectra are possible, including spectra with stau NLSP. The sufficient condition is that the axino decay into the NLSP is not too strongly kinematically suppressed. In the KSVZ model, decays into fermion-fermion pairs are loop-suppressed. However, in both models an open axino decay channel into a neutralino suffices. As a consequence, it will be interesting to construct concrete models with the naturally expected axino and saxion mass of order \( m_{\text{susy}} \), possibly along the lines of existing models like [21, 55, 56].

Constraints on the axion decay constant from axino and saxion decay are softened. Interestingly, the DFSZ model is—depending on the \( \mu \)-parameter—less restricted from particle decay already in the standard scenario. For example, the saxion mass could be unconstrained, if the Higgs boson is lighter than the NLSP. Since in the considered setting the reheating temperature is fixed at a large value, the constraints from axion overproduction on the decay constant are particularly interesting and depend on the axion model. However, they are not affected by R-parity violation. Therefore we have shown how late-time entropy production might soften the upper bounds on the axion decay constant and the initial amplitude of saxion oscillations after inflation.

There is no constraint on the saxion-axion-axion self-coupling as long as \( f_a < 10^{10} \) GeV and \( m_{\text{sax}} > 100 \) GeV. A large self-coupling of order one can remove constraints on the saxion mass, especially in the standard scenario. We remark that for
suitable values of the self-coupling and other model parameters one can obtain any desired amount of additional radiation energy in the Universe formed by axions. It will be interesting to identify (more) models realising a particular self-coupling and/or a small axion density \( \Omega_a \ll \Omega_{\text{DM}} \).

Furthermore, it might be interesting to investigate if other scenarios of dark and visible matter enable a solution of the strong CP problem. We conclude that broken R-parity does not only solve the NLSP decay problem but also makes it easier to solve the strong CP problem. Our results are an additional motivation for scenarios of thermal leptogenesis with gravitino dark matter and broken R-parity.

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| [GeV] | standard | \( \mathcal{R} \) | \( \mathcal{R} \land \Delta T_{<\Lambda_{\text{QCD}}} = \Delta_{\text{max}} \) |
|-------|----------|----------------|---------------------------------|
| \( f_a \) | \( \lesssim 10^{10} \) | \( \lesssim 10^{10} \) | \( \gtrsim 4 \times 10^{12} \) |
| \( m_\tilde{a} \) | \( > m_\tilde{g} \)  
(or \( > m_\tilde{\mu} + m_\tilde{b} \) in DFSZ) | \( > m_{\text{nlsp}} \)  
(or \( > m_{\text{nlsp}} \) in DFSZ) | \( \gtrsim \max\{m_{\text{nlsp}}^*, 300 \left( \frac{f_a}{10^{10} \text{ GeV}} \right)^2 \} \) |
| \( m_{\text{max}} \) | \( > 760 \left( \frac{f_a}{10^{10} \text{ GeV}} \right)^{\frac{2}{3}} \)  
(or \( > 2m_{\text{nlsp}} \) in DFSZ) | \( \gtrsim \frac{5}{3} \left( \frac{f_a}{10^{10} \text{ GeV}} \right)^{\frac{2}{3}} \)  
or \( \in [5 \left( \frac{f_a}{10^{10} \text{ GeV}} \right)^{\frac{2}{3}}, 2m_{\text{nlsp}}] \) | \( \gtrsim \frac{5}{3} \left( \frac{f_a}{10^{10} \text{ GeV}} \right)^{\frac{2}{3}} \) |
| \( \phi_{\text{max}} \) | \( \lesssim 10^{13} \ldots 10^{15} \) | \( \lesssim 10^{13} \ldots 10^{15} \) | \( \lesssim 5 \times (10^{14} \ldots 10^{16}) \) |

Table 1: Parameter constraints for the different scenarios (standard: R-parity conserved (\( \Delta = 1 \)), \( \mathcal{R} \): R-parity violated (\( \Delta = 1 \)) and violated R-parity with the maximal entropy dilution at the right time). Units are GeV where not written explicitly. Here, we assume the self-coupling \( x \ll 1 \). Only if a mass bound is sensitive to the actual value of \( f_a \), its dependence is given. By “nlsp” we indicate that the bound actually does not hold for any possible NLSP and depends on the axion model. The upper bound on \( f_a \) in the standard scenario is boldface to indicate that it arises from \( \Omega_a \) and the (KSVZ) axino decay. In the \( \mathcal{R} \) case it stems from \( \Omega_a^{\text{mis}} \) only. If, furthermore, matter dominates at \( T_{\text{osc}} \) it stems from \( \Omega_a^{\text{str}} \).
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