Axino LSP baryogenesis and dark matter

Angelo Monteux and Chang Sub Shin

New High Energy Theory Center, Department of Physics and Astronomy, Rutgers University, 136 Frelinghuysen Rd, Piscataway, NJ, 08854 U.S.A.

E-mail: amonteux@physics.rutgers.edu, changsub@physics.rutgers.edu

Received January 21, 2015
Revised March 31, 2015
Accepted April 27, 2015
Published May 20, 2015

Abstract. We discuss a new mechanism for baryogenesis, in which the baryon asymmetry is generated by the lightest supersymmetric particle (LSP) decay via baryonic $R$-parity-violating interactions. As a specific example, we use a supersymmetric axion model with an axino LSP. This scenario predicts large $R$-parity violation for the stop, and an upper limit on the squark masses between 15 and 130 TeV, for different choices of the Peccei-Quinn scale and the soft $X_i$ terms. We discuss the implications for the nature of dark matter in light of the axino baryogenesis mechanism, and find that both the axion and a metastable gravitino can provide the correct dark matter density. In the axion dark matter scenario, the initial misalignment angle is restricted to be $\mathcal{O}(1)$. On the other hand, the reheating temperature is linked to the PQ scale and should be higher than $10^4$–$10^5$ GeV in the gravitino dark matter scenario.

Keywords: baryon asymmetry, supersymmetry and cosmology, dark matter theory, axions

ArXiv ePrint: 1412.5586
1 Introduction

Two fundamental observations about our universe, namely the presence of dark matter and the prevalence of matter over antimatter, can individually find an explanation in theories that go beyond the Standard Model. However, it is harder to explain both in a unified framework. A simple argument based on symmetries can justify this: a stable dark matter candidate would be protected by a symmetry, while baryogenesis requires the violation of another symmetry, namely baryon number (or lepton number in high-energy leptogenesis). At the same time, flavor physics is sensitive to baryon-number violation, and proton decay is mediated by the simultaneous breaking of baryon and lepton number.

Although the symmetries involved in the two mechanisms need not to be related, they are in supersymmetric theories, where $R$-parity forbids renormalizable baryon-number-violating operators and at the same time provides a stable dark matter candidate, the Lightest Supersymmetric Particle (LSP). The LSP is also the source of the missing energy signature of supersymmetric events at colliders. Given the stringent LHC limits for light $R$-parity conserving supersymmetry (SUSY) and the null results in dark matter direct detection experiments, it is appealing to investigate phenomenological consequences of $R$-parity violation (RPV). Besides collider studies, its cosmological implication is very interesting because one can relate the baryogenesis mechanism and the nature of dark matter.

Sizable lepton and proton number violation are not consistent with bounds on the proton lifetime, thus we will focus on models allowing only for baryonic $R$-parity violation. In addition to the Minimal Supersymmetric Standard Model (MSSM) superpotential, we include the operator

$$\frac{1}{2} \lambda'_{ijk} U^c_i D_j^c D_k^c$$

(1.1)
where the RPV couplings $\lambda_i''_{jk}$ are antisymmetric under the exchange $j \leftrightarrow k$ and the sum over color indices is understood. Neutralinos decay away in the early universe. Without an additional symmetry to warrant dark matter stability, a natural candidate for dark matter should be super-weakly interacting: for example, the gravitino, which can be metastable and have a lifetime longer than the age of the universe, and the axion, which was originally introduced [1] to explain the absence of CP-violation in strong interactions (the strong CP problem). Introducing the axion supermultiplet gives a further implication that relates dark matter and baryogenesis; there is a fermionic superpartner, the axino, whose out-of-equilibrium decay through the operator (1.1) can generate a baryon asymmetry.

The idea of using the operator (1.1) for baryogenesis is not new. Refs. [2–4] discussed a baryon asymmetry generated at low temperatures (down to the MeV scale) from late decays of inflaton, gravitinos and axinos into superpartners, respectively. It is important that baryon asymmetry is generated at a temperature well below the superpartner mass scale so that it is not washed out by baryon-number-violating processes in the thermal bath [9]. For large RPV couplings, baryogenesis should happen at a temperature lower than about $m_{\tilde{q}}/20$. Otherwise, if the asymmetry is generated at or above the superpartner scale, all the $R$-parity violating couplings have to be smaller than $O(10^{-7})$ for baryons to survive.

We shortly review these baryogenesis mechanisms, which all involve $R$-parity violating $A$-terms. In [2], out-of-equilibrium decays of the inflaton generate a non-thermal population of squarks, which later decay to a quark and a gluino, provided that $m_{\tilde{q}} > m_{\tilde{g}} + m_q$. Interference between the tree level and two-loop diagrams gives a baryon asymmetry. In [3], the decay of a gravitino to a quark and a squark (or to a gluino and a gluon, with the subsequent gluino decay to a quark and a squark) was used to generate the asymmetry (given the hierarchies $m_{3/2} > m_{\tilde{q}} + m_q$ or $m_{3/2} > m_{\tilde{g}} > m_{\tilde{q}} + m_q$, respectively), through a one-loop diagram involving $A$-terms. In [4], the same diagrams were used to discuss the asymmetry generated by the decay of the axino (or saxion) to a gluino and finally to a quark and a squark.

For the cases described above, the parent particle decays to some on-shell superpartner via $R$-parity conserving interactions, and interference with a $R$-parity violating loop diagram generates the baryon asymmetry. A specific hierarchy is required in each case: the particle decaying out of equilibrium is always heavier than the superpartners that are on-shell in the interfering diagram. This automatically excludes the case of an LSP decay (in the following the LSP is defined as in $R$-parity conserving SUSY, and it is stable only when the RPV interactions are turned off). In fact, this reflects a more general statement: Nanopoulos and Weinberg proved in ref. [10] that an LSP defined in such a way cannot give a baryon asymmetry at first order in the baryon-number-violating interactions.

In this work, we will investigate how to generate both a baryon asymmetry at low energies and a correct dark matter density. We focus on models with baryonic RPV and show a new baryogenesis mechanism in which a late-decaying axino LSP gives rise to the observed baryon asymmetry (at second order in the $R$-parity violating interactions). Baryogenesis takes place below the weak scale and before Big Bang Nucleosynthesis. Both gravitino and axion dark matter are discussed, with the nature of dark matter constraining the parameter space for baryogenesis and vice versa.

This paper is organized as follows: we recall the Nanopoulos-Weinberg theorem in section 2, and discuss the implications for baryogenesis through an LSP decay. In section 3, we derive a new mechanism that generates substantial baryon asymmetry through late decays

\[1\] In another class of models that gives an asymmetry through the operator (1.1) [5–8], the decaying particles are thermally produced from freeze-out at a temperature well below the mass of the particles.
of an axino LSP. In section 4, we investigate the possibilities for dark matter candidates and discuss the range of parameter space in which both a baryon asymmetry and a correct dark matter density exist. We discuss the collider bounds on our model and conclude in section 5.

2 The Nanopoulos-Weinberg theorem and LSP baryogenesis

In ref. [10] it was shown that, given a particle \(X\) which is stable in the limit of no baryon-number-violating interactions, the decay rate of \(X\) into all final states with a given baryon number \(B\) equals the decay rate of its antiparticle \(\bar{X}\) into all states with baryon number \(-B\), at first order in the baryon-number-violating interactions. Because the net number of baryons is proportional to the difference of the two decay rates, no baryon asymmetry can be generated by an LSP decay at first order.

The Nanopoulos-Weinberg theorem was further investigated in ref. [11], where it was generally argued that a difference in the two decay rates exists only if the on-shell intermediate particles and the final particles have a different baryon number. In other words, the process to the right of the “cut” must violate baryon number. This results holds at all orders in the baryon-number-violating interactions. For an LSP decay, baryon-number-violating operators must also appear on the left of the “cut” to have an asymmetry at first order, with a baryon-number-conserving interaction to the left of the “cut”. The models of [2–4], where baryon asymmetry is generated by two-body decays, and of [12], where it comes from three-body decays, fall in this last category, as they used decays of heavier particles, and were indeed able to get an asymmetry at first order in the RPV couplings.

For the LSP case, it is natural to consider a low-energy effective theory containing the LSP \(\chi\), a Majorana fermion with mass \(m_\chi\), and the Standard Model fields, where all the heavier degrees of freedom have been integrated out. Baryon number violation is present in non-renormalizable operators which are suppressed by the heavier particles masses (in RPV SUSY, the squarks). The effective operators involving the quarks \(q\) and the LSP will be of the schematic form

\[
L_{|\Delta B|=1} \sim \frac{c_1}{\Lambda^2} (\chi q)(qq) + \text{h.c.}, \quad L_{|\Delta B|=2} \sim \frac{c_2}{\Lambda^5} (qq)(qq)(qq) + \text{h.c.},
\]

where the gauge indices are contracted to form gauge singlets and \(\Lambda\) is the mass scale of the heavier particles. Then the decay rate of \(\chi\) becomes

\[
\Gamma_{\chi \rightarrow qqq} \sim \frac{1}{(8\pi)^3} \frac{|c_1|^2 m_\chi^5}{\Lambda^4}.
\]

An asymmetry between \(\Gamma_{\chi \rightarrow qqq}\) and \(\Gamma_{\chi \rightarrow \bar{q}\bar{q}\bar{q}}\) will come from the interference between the tree level \(\Delta B = 1\) decay and the two-loop decay diagram obtained joining together the two operators in figure 1:

\[
\epsilon \equiv \frac{\Gamma_{\chi \rightarrow qqq} - \Gamma_{\chi \rightarrow \bar{q}\bar{q}\bar{q}}}{\Gamma_{\chi \rightarrow qqq} + \Gamma_{\chi \rightarrow \bar{q}\bar{q}\bar{q}}} \sim \frac{1}{(8\pi)^3} \frac{\text{Im}[c_1^2 c_2 m_\chi] m_\chi^2}{|c_1|^2 \Lambda^3}.
\]

\(\epsilon\) including the Majorana mass term, the theory has three complex parameters and two independent fields (\(\chi\) and \(qqq\)), resulting in one phase that cannot be removed by field redefinition of the \(q\)’s and \(\chi\). We can check that this is the phase appearing in \(\epsilon\) by considering the invariant combination of operators under the field redefinition, \((c_1^* \chi qqq)^2 (c_2 qqqqq)(m_\chi \chi \chi)\).
It is worth to note that the asymmetry is generated not only at second order in the baryon-number-violating interactions (as expected by the Nanopoulos-Weinberg theorem), but also at two-loops: a 1-loop diagram would require a dimension six, $\Delta B = 2$ operator in the effective theory, which is not allowed. Finally, the present baryon asymmetry depends on the LSP abundance at decay time,

$$Y_{\Delta B} \equiv \frac{n_B}{s} = \epsilon \left( \frac{n_\chi}{s} \right)_{t=1/\Gamma_\chi}. \quad (2.4)$$

This can be compared to its experimental value [13], $(Y_{\Delta B})_{\text{obs}} = (0.80 \pm 0.018) \times 10^{-10}$.

In the next section, we will see the axino LSP is a good example to realize this mechanism.

3 Baryogenesis from axino decays

In this section we present a concrete example of the LSP baryogenesis scenario just outlined. First, the particle has to decay out of thermal equilibrium. A typical candidate would be an LSP gravitino, produced in the reheating epoch [14, 15]: without R-parity, gravitinos decay to three quarks via the RPV operator of eq. (1.1). Compared to a non-LSP gravitino, this decay is much slower, as it is suppressed by the intermediate squark mass, by the RPV coupling $\lambda''$ and by the three-body kinematic factor. In order for the decay products not to interfere with Big Bang Nucleosynthesis (BBN), the gravitino LSP should be extremely heavy, $m_{3/2} \gtrsim 10^3 \text{TeV}$, with the other superpartners being even heavier [16]. Even so, all superpartners would have decayed in the early universe and there would be no dark matter candidate.

Thus, we examine a supersymmetric QCD axion model in which the axion ($a$), the pseudo-Goldstone boson associated with the spontaneous breaking of an anomalous Peccei-Quinn (PQ) U(1) symmetry at a scale $v_{\text{PQ}}$, solves the strong CP problem of QCD [17] and is a good dark matter candidate. Assuming that the axino ($\tilde{a}$), the fermionic superpartner of the axion, is the LSP, a baryon asymmetry can be generated by its decays in a more natural region of the parameter space than in the gravitino LSP case. The large value of $v_{\text{PQ}}$ ensures that axinos are out-of-equilibrium when they decay, and the lifetime is much longer than the period in which RPV interactions are in thermal equilibrium (thus, the asymmetry is not washed out). The chiral axion superfield can be written as

$$A = \frac{1}{\sqrt{2}}(s + ia) + \sqrt{2} \theta \tilde{a} + \theta^2 F_A, \quad (3.1)$$

where the saxion, the scalar superpartner of the axion, is denoted as $s$. We consider a following superpotential terms for $A$ to give interactions with the MSSM particles:

$$\Delta W = e^{\theta H_A/v_{\text{PQ}}} \mu H_u H_d + e^{\theta \Phi/v_{\text{PQ}}} M_{\Phi} \Phi \Phi^c, \quad (3.2)$$
where $H_u, H_d$ are the MSSM Higgs doublets, and $\Phi, \Phi^c$ are SM charged matter fields with $M_\Phi = \mathcal{O}(v_{PQ})$. The U(1)$_{PQ}$ symmetry is realized as $A \to A + i \theta v_{PQ}$, $H_u H_d \to e^{-iq_H \theta} H_u H_d$, and $\Phi \Phi^c \to e^{-iq_{\Phi^c} \theta} \Phi \Phi^c$. This is a hybrid of the DFSZ [18, 19] and KSVZ [20, 21] axion models, in which the axino decay is dominated by the first term as in the DFSZ case while at high temperature its thermal production is mostly given by that of the KSVZ model [22].

Because all the sparticles are heavier, we can consider a low-energy effective theory with SM quarks supplemented by the axino, a Majorana particle with a mass $m_{\tilde{a}}$. Non-renormalizable interactions for the quarks remain after integrating out the squarks in the diagrams of figure 2. The following effective axino interactions are given by the mixing of the axino with the higgsino, after integrating out the squarks:

$$
\mathcal{L}_{|\Delta B|=1} = \sum_{\alpha} \frac{\kappa_\alpha m_{u_\alpha}}{v_{PQ}} \lambda'_{\alpha} (\tilde{u}_{\alpha}^L)(d_\beta^c d_\gamma^c) + \text{h.c.}
$$

where $q_\alpha$ (respectively, $q^c_\alpha$) are the left-handed (right-handed) quarks. The holomorphic term in the second line comes from the left-right squark mixing, $X_{u_\alpha} \equiv A_{u_\alpha} + \mu \cot \beta$. $\kappa_\alpha$ is an $\mathcal{O}(1)$ coefficient given by the charges of the SM fields under the PQ symmetry.\footnote{In pure KSVZ models [20, 21], the SM fields are neutral under the PQ symmetry and the same operator would arise at 1-loop after integrating out the gluino; it would be further suppressed with respect to the DFSZ case.}

The six-fermion holomorphic $\Delta B = 2$ Lagrangian is obtained from the soft SUSY breaking $A$-terms, \( \Delta \mathcal{L}_{\text{soft}} = \lambda'_{ijk} A_{ijk} \tilde{u}_L \tilde{d}_R \tilde{d}_R \tilde{d}_R \), after integrating out the right-handed

\[ \Delta \mathcal{L}_{\text{soft}} = \lambda'_{ijk} A_{ijk} \tilde{u}_L \tilde{d}_R \tilde{d}_R \tilde{d}_R + \text{h.c.,} \]

where $\lambda'_{ijk}$ is the axino mass parameter.
squarks,

\[ \mathcal{L}_{\Delta B=2}^A = - \sum_{ijk,\alpha\beta\gamma,\delta\zeta,\lambda} \frac{\lambda^*_{ijk}A_{ijk}^*\lambda_{\delta\lambda\gamma}\lambda_{\alpha\beta\delta}}{m_{\lambda_R}^2m_{\lambda_R}^2m_{\lambda_R}^2} (d^c_i d^c_j)(u^c_i d^c_j) + \text{h.c.} \]  \tag{3.4}

and also from squark-quark-gaugino interactions, \( \Delta \mathcal{L}_{\text{susy}} = \sqrt{2} g_A \tilde{u}_{Ri} T^A u_i^c + \tilde{d}_{Ri} T^A d_i^c \lambda^A + \text{h.c.} \), after integrating out the gauginos:

\[
\mathcal{L}_{\Delta B=2}^{\tilde{A}} = - \sum_{\lambda,\alpha\beta\gamma,\delta\zeta} \frac{g_A^2 m_{\lambda R}^2 m_{\lambda R}^2 m_{\lambda R}^2}{4 m_{\lambda A}^2 m_{\lambda R}^2 m_{\lambda R}^2} (d^c_i d^c_j) (T^A u^c_i)(T^A u^c_j)(d^c_i d^c_j) + \text{h.c.}
\]

\[ + 4 \left( \tilde{u}_{Ra} \rightarrow \tilde{d}_{R\gamma}, u^c_\alpha \leftrightarrow d^c_\gamma, \tilde{u}_{R\delta} \leftrightarrow \tilde{d}_{R\delta}, u^c_\delta \leftrightarrow d^c_\delta \right) \]

\[ + 2 \left( \tilde{u}_{Ra} \leftrightarrow \tilde{d}_{R\gamma}, u^c_\alpha \leftrightarrow d^c_\gamma \right), \]  \tag{3.5}

where \( A \) indicates the gauge group index of the Standard Model, \( \tilde{\lambda}^A \) is the corresponding gaugino with Majorana mass \( m_{\lambda A}^2 \). In the denominator \( m_{\lambda A}^2 \) is used to imply \( |m_{\tilde{\lambda} A}|^2 \). If the gaugino masses are all of the same order, the gluino contribution is most important. If they follow the GUT relation as \( m_{3/2} : m_{\tilde{e}_R} : m_{\tilde{e}_L} \approx 1 : 4 : 36 \), the bino contribution would be also important. More detailed expressions are presented in the appendix. Exchange of neutral Higgsinos also generates \( \Delta B = 2 \) operator. However this contribution is proportional to \( y_t y_d \sim m_t m_d / v^2 \), so it is suppressed. For the sake of simplicity, we neglect it and decouple the Higgsinos in our discussion, so that the bino Majorana mass term is nearly the mass eigenvalue of the neutralino.

Because in eq. (3.3) the effective axino coupling to quarks is proportional to the up-type quark masses, the dominant one involves the top quark. The relevant \( R \)-parity violating couplings are \( \lambda''_{312}, \lambda''_{313}, \lambda''_{323} \). For simplicity, we take the assumption that \( \lambda''_{323} \) is the only dominant coupling.\(^4\)

In the limit of massless final states, the tree-level decay rate of the axino LSP is

\[
\Gamma_{\tilde{a} \rightarrow tbs} \approx \frac{\lambda''_{323}^2}{512 \pi^3} m_{\tilde{a}}^2 m_{t_R}^4 \left[ 1 + \left( \frac{X t m_t}{m_{t_L}^2} \right)^2 \right]. \tag{3.6}
\]

Thus the total decay rate is \( \Gamma_{\tilde{a}} = \Gamma_{\tilde{a} \rightarrow tbs} + \Gamma_{\tilde{a} \rightarrow \tilde{\nu}_\tau \tilde{\nu}_\tau} \approx 2 \Gamma_{\tilde{a} \rightarrow tbs} \). Corrections to this result are proportional to \( m_{\tilde{t}}^2 / m_{\tilde{a}}^2 \) and \( m_{\tilde{a}}^2 / m_{t_R}^2 \), and are shown in the appendix.

In order to obtain the baryon asymmetry as in eq. (2.4), we consider the axino thermal history and evaluate its abundance in section 3.1. The asymmetry parameter \( \epsilon \) is presented in section 3.2, and we discuss its implications for the sparticle spectrum in the same section.

3.1 Axino cosmology

Even if the interactions between the axinos and the MSSM particles are quite suppressed, axinos can be generated from the thermal bath in the early Universe. For a reheating temperature \((T_R)\) much higher than the MSSM sparticle masses and much lower than the

\(^4\)Note that we do not require all the other couplings to be zero, just to be small. This is justified in models where flavor symmetries determine a hierarchical structure of the \( R \)-parity violating couplings [23–25]. The other couplings would be suppressed by spurions of the flavor symmetry.
PQ breaking scale \( (\mu \ll T_R \ll M_\Phi \sim v_{\text{PQ}}) \), axinos are mainly produced by scattering processes mediated by gluinos [26–31]. The axino yield from this thermal production is

\[
\left( \frac{n_{\tilde{a}}}{s} \right)_{\text{TP}} = \min \left[ 0.001 g_3^6 \left( \frac{200}{g_\ast(T_R)} \right)^{3/2} \left( \frac{T_R}{10^{10} \text{GeV}} \right) \left( \frac{10^{12} \text{GeV}}{v_{\text{PQ}}} \right)^2 , 0.002 \left( \frac{200}{g_\ast(T_{\text{dec.}})} \right) \right],
\]

(3.7)

where \( n_{\tilde{a}} \) is the axino number density, \( s \) the entropy density of the Universe given by \( (2\pi^2/45) g_\ast s T^3 \) and \( g_\ast(T) \simeq g_\ast(T) \) is the number of relativistic degrees of freedom at the temperature \( T \). If the reheating temperature is high enough \( (T_R > T_{\text{dec.}}) \), the scattering processes can be in chemical equilibrium. In such a case, the axino number density is \( n_{\tilde{a}} = (3\zeta(3)/2\pi^2) T^3 \) before it decouples, and the corresponding yield after decoupling \( (T < T_{\text{dec.}}) \) is given by the second term of the r.h.s. of eq. (3.7).

In our cosmological consideration, axinos should decay before BBN, as the baryon asymmetry is generated by their decay. We denote by \( T_D \) the axino decay temperature, defined by \( \Gamma_{\tilde{a}} = H(T_D) \):

\[
T_D \simeq 28 \text{ MeV} \left( \frac{m_{\tilde{a}}}{\text{TeV}} \right)^{1/2} \left( \frac{m_{\tilde{a}}}{m_{\text{IR}}} \right)^2 \left( \frac{10^{12} \text{GeV}}{v_{\text{PQ}}} \right).
\]

(3.8)

The condition that the axino decays before BBN corresponds to \( T_D \gtrsim 10 \text{ MeV} \). In the absence of large axino-squark hierarchy,\(^5\) it gives an upper bound on \( v_{\text{PQ}} \) of about \( 10^{12} \text{GeV} \). Because \( T_D \lesssim m_{\tilde{a}} \), axinos are non-relativistic at decay times, and as such they will eventually dominate the energy density of the universe, unless they decay beforehand. The temperature at which the axino energy density equals the radiation energy density is

\[
T_{\text{eq}} = \frac{4}{3} m_{\tilde{a}} \left( \frac{n_{\tilde{a}}}{s} \right)_{\text{TP}}.
\]

(3.9)

For \( T_D > T_{\text{eq}} \), axinos decay before dominating the energy density; for \( T_D < T_{\text{eq}} \), they decay after, injecting a non-negligible amount of high-energy decay products in the thermal plasma. This has the effect of increasing the entropy, and the axino yield at decay is

\[
\left( \frac{n_{\tilde{a}}}{s} \right)_{T_D} = \min \left[ \frac{3 T_{\text{eq}}}{4 m_{\tilde{a}}} , \frac{3 T_D}{4 m_{\tilde{a}}} \right].
\]

(3.10)

One can check that, for \( T_D > T_{\text{eq}} \), the axino yield is given by eq. (3.7), while for lower decay temperatures it is given by \( (3/4)(T_D/m_{\tilde{a}}) \).

In the axino decay, a difference in the decay rate to quarks vs. antiquarks is needed to generate a baryon asymmetry: the parameter \( \epsilon = (\Gamma(\tilde{a} \to qqq) - \Gamma(\tilde{a} \to \bar{q}q\bar{q}))/\Gamma(\tilde{a} \to qqq + \Gamma(\tilde{a} \to \bar{q}q\bar{q})) \) gives the net asymmetry per axino decay, and the net baryon yield is

\[
Y_{\Delta B} \equiv \frac{n_B}{s} = \epsilon \left( \frac{n_{\tilde{a}}}{s} \right)_{T_D}.
\]

(3.11)

On a separate note, the saxion is also produced from the thermal plasma in a similar amount as the axino. Its decay rate is much larger than that of the axino, because the saxion can decay through the \( R \)-parity conserving interactions. The saxion always decays to axions.

\(^5\)We will see below that a large splitting would give too small of a baryon asymmetry anyway.
with $\Gamma_{s \rightarrow aa} = m_s^3/(64\pi v^2_{Q})$, where $m_s$ is the saxion mass. If kinematically allowed, the dominant decay mode of the saxion would be $s \rightarrow hh$ with

$$\Gamma_{s \rightarrow hh} = \frac{\mu^4}{4\pi m_s^2} \left(1 - \frac{4m_n^2}{m_s^2}\right)^{1/2} \approx 10^6 \left(\frac{\mu}{m_s}\right)^4 \left(\frac{1}{1\text{ TeV}}\right)^2 \Gamma_{\tilde{a}}. \quad (3.12)$$

Thus saxions decay much earlier than the axinos. Furthermore, because their decays do not produce any baryon asymmetry, the role of the saxion really is negligible.

### 3.2 Axino baryogenesis

As discussed in section 2, no contribution to $\epsilon$ comes at one-loop. The interference between the tree-level decay and the two-loop decay (involving the $\Delta B = 2$ interactions in figure 2(c), 2(d)) gives a non-zero asymmetry,

$$\epsilon = \epsilon_A + \epsilon_{\tilde{g}} + \epsilon_{\tilde{B}}, \quad (3.13)$$

where

$$\epsilon_A = \frac{|\lambda''_{323}|^2}{32\pi^3} \frac{m^2_{\tilde{t}}}{m^2_{h_{\tilde{R}}}} \frac{m^2_{h_{\tilde{B}}}}{m^2_{t_{\tilde{R}}}} \text{Im}[m_{\tilde{a}}A''_{323}] \left[1 + \frac{1}{2} \frac{m_{\tilde{a}}X_L}{m^2_{t_{\tilde{L}}}} + \frac{1}{16} \left(\frac{m_{\tilde{a}}X_L}{m^2_{t_{\tilde{L}}}}\right)^2\right], \quad (3.14)$$

$$\epsilon_{\tilde{g}} = \frac{\kappa_{\tilde{g}}}{32\pi^3} |g^2_{33}| \frac{m^2_{\tilde{t}}}{m^2_{h_{\tilde{B}}}} \frac{m^2_{h_{\tilde{B}}}}{m^2_{t_{\tilde{R}}}} \text{Im}[m_{\tilde{a}}m_{\tilde{g}}^*] \left[1 + \frac{1}{2} \frac{m_{\tilde{a}}X_L}{m^2_{t_{\tilde{L}}}} + \frac{1}{16} \left(\frac{m_{\tilde{a}}X_L}{m^2_{t_{\tilde{L}}}}\right)^2\right], \quad (3.15)$$

$$\epsilon_{\tilde{B}} = \frac{\kappa_{\tilde{B}} |\lambda''_{323}|^2 g^2_{33}}{32\pi^3} \frac{m^2_{\tilde{t}}}{m^2_{h_{\tilde{B}}}} \frac{m^2_{h_{\tilde{B}}}}{m^2_{t_{\tilde{R}}}} \text{Im}[m_{\tilde{a}}m_{\tilde{B}}^*] \left[1 + \frac{1}{2} \frac{m_{\tilde{a}}X_L}{m^2_{t_{\tilde{L}}}} + \frac{1}{16} \left(\frac{m_{\tilde{a}}X_L}{m^2_{t_{\tilde{L}}}}\right)^2\right], \quad (3.16)$$

where $\kappa_{\tilde{g}}, \kappa_{\tilde{B}}$ are squark mass-dependent dimensionless parameters defined in eq. (A.6). In the limit of universal squark masses, we get

$$\kappa_{\tilde{g}} = 2, \quad \kappa_{\tilde{B}} = \frac{14}{9}. \quad (3.17)$$

Here, all mass squared terms represent real values. We can check the bino contribution would be same order of that of gluino’s if the GUT relation ($m_\tilde{B} : m_\tilde{g} \simeq (3/5)g^2 : g^2_{33}$) is satisfied.

In the following discussion we consider the case with $m_\tilde{g} \sim m_\tilde{B}$ such that $\epsilon_\tilde{g} \gg \epsilon_{\tilde{B}}$. As for the decay rate $\Gamma_{\tilde{a}}$, this is computed in the limit of massless final states and heavy intermediate squarks, and the exact expression is discussed in the appendix. In the following we will denote by $\Phi_A$ the relative phase between the axino mass and the RPV $A$-term $A_{323}'$, $\Phi_A \equiv \text{Im}[m_{\tilde{a}}A''_{323}]/|m_{\tilde{a}}A''_{323}|$, and by $m_\tilde{q}$ the average squark mass scale, $m_\tilde{q} \equiv (m^2_{t_{\tilde{R}}} m^2_{b_{\tilde{R}}} m^2_{s_{\tilde{R}}})^{1/6}$. Similarly $\Phi_\lambda \equiv \text{Im}[m_{\tilde{a}}m_\lambda^*]/|m_{\tilde{a}}m_\lambda|$, and for the contribution of $\epsilon_{\tilde{g}}$, universal squark masses are taken so that $m_\tilde{q} = m_{t_{\tilde{R}}} = m_{b_{\tilde{R}}} = m_{s_{\tilde{R}}} = m$. As an example, we give three benchmark points that reproduce the correct baryon asymmetry: taking $m_{\tilde{a}} = 500\text{ GeV}$, $A''_{323} \simeq X_L \approx m_\tilde{q}$ and a common CP phase $\Phi_A = \Phi_\tilde{g} = \Phi_{\tilde{B}} = \Phi$, the other parameters are:

BP1: $m_\tilde{g} = 900\text{ GeV}, \quad m_\tilde{q} = 1.5\text{ TeV}, \quad |\lambda''_{323}| = 1, \quad \Phi = 1; \quad \epsilon = 3.4 \times 10^{-6}$;

BP2: $m_\tilde{g} = 1\text{ TeV}, \quad m_\tilde{q} = 2\text{ TeV}, \quad |\lambda''_{323}| = 0.5, \quad \Phi = 0.2; \quad \epsilon = 5.9 \times 10^{-8}$;

BP3: $m_\tilde{g} = 2\text{ TeV}, \quad m_\tilde{q} = 1\text{ TeV}, \quad |\lambda''_{323}| = 1, \quad \Phi = 1; \quad \epsilon = 1.5 \times 10^{-7}$. \quad (3.18)
At the benchmark point BP1, the asymmetry receives roughly equal contributions from the diagrams with the $A$-terms and the gluino-mediated processes, while at BP2 and BP3 the gluino contribution is the dominant source. This is easily explained by the large power of $\lambda^\mu_{323} = 0.5$ for $\epsilon_A$ at BP2 and by the small gluino mass boosting $\epsilon_\tilde{g}$ at BP3. For each benchmark point, we obtain the observed baryon asymmetry $Y_{\Delta B}$ by choosing appropriate values for $T_R$ and $v_{PQ}$. As these parameters also determine the dark matter abundance, we will discuss the benchmark points in relation to the nature of dark matter in the next section. But first, let us discuss our result, eq. (3.13) in more details:

- A large $\lambda^\mu_{323}$ is strongly preferred, especially for $\epsilon_A$. In particular, the lower limit on $\lambda^\mu_{323}$ is 0.03 for $\epsilon_\tilde{g}$ to generate the asymmetry (and 0.1 for $\epsilon_A$). An upper limit of $\lambda^\mu_{323} = 1.07$ was found from the condition that perturbativity is valid up to the GUT scale in [32], where the RG running of the $R$-parity violating couplings was considered. Although from the point of view of a low-energy effective theory this is not a problem, and one can just expect that new degrees of freedom appear around the Landau pole, we will assume $\lambda^\mu_{323} = 1$ in the rest of this paper.

- The baryon asymmetry is proportional to, roughly, $(m_{\tilde{\mu}}/m_{\tilde{q}})^3(m_t/m_{\tilde{q}})^2$. It is suppressed for a large hierarchy between the squark mass and either the axino mass or the weak scale. Even for $\lambda^\mu_{323} = 1$, $m_{\tilde{\mu}} \approx m_{\tilde{q}} \approx m_{\tilde{g}}$, there is a suppression by $m_t^2/m_{\tilde{q}}^2$, which points to an upper limit on the squark scale. In particular, in this limit we can write:

$$
\frac{Y_{\Delta B}}{(Y_{\Delta B})_{\text{obs}}} \sim \left( \frac{n_{\tilde{\mu}}/s}{10^{-3}} \right) T_\nu 10^4 N \frac{m_t^3}{m_{\tilde{q}}^3} \cdot c
$$

(3.19)

where $c$ is an $O(1)$ number determined by $A^\mu_{323}, \kappa_{\tilde{g}}, \frac{X_t}{m_{\tilde{q}}}$ and $N$ takes into account the ratio between the numerical and the analytical analysis (as shown in the appendix); for $m_{\tilde{\mu}} \sim m_{\tilde{q}} \sim m_{\tilde{g}} \sim 1$ TeV, we find $N \sim 6$, while for $m_{\tilde{\mu}} \sim m_{\tilde{q}} \sim m_{\tilde{g}} \sim 100$ TeV the value of $N$ increases to about 60 (this can be expected as in this limit the top mass is massless and the only corrections are due to the internal propagators, which increase the results). In any case, we find

$$
m_{\tilde{q}} \lesssim 10^2 \sqrt{N} cm_t \simeq 130 \text{ TeV}.
$$

(3.20)

We find an absolute upper bound on the squark masses at 130 TeV for large soft terms $A^\mu_{323} \simeq X_t \simeq 3m_{\tilde{q}}$; larger values for $A$-terms are potentially dangerous in that they can generate color-breaking vacua [33, 34].

- Additionally, the baryon asymmetry is proportional to the relative phases between the axino mass $m_{\tilde{\mu}}$ and the soft SUSY breaking parameters $A^\mu_{323}$ and $m_{\tilde{g}}$. There is no direct constraint on the CP phase of $A^\mu_{323}$ but an indirect constraint is provided by the null results in the measurement of the neutron Electric Dipole Moment (EDM). CP phases for the MSSM A-terms $A_{U,D}$ (in particular, the phase $\phi_{A_{Q\tilde{g}}}$ = Im[$m_{\tilde{g}}A_Q^\mu$]/$m_{\tilde{g}}^2$, where $m_{\tilde{g}}$ is the gluino mass and $Q = U, D$) contribute to the neutron EDM [35],

$$
|d_n| \simeq 2.5 \times 10^{-25} e \text{ cm } \left( \frac{\phi_{A_{Q\tilde{g}}}}{1/4} \right) \left( \frac{\text{TeV}}{m_{\tilde{g}}} \right)^2
$$

(3.21)

while experimentally the upper limit is $|d_n| < 2.9 \times 10^{-26} e \text{ cm}$ [36]. This implies either a small phase $\phi_{A_{Q\tilde{g}}}$ or superpartners in the multi-TeV range. In our model, we can...
have a large baryon asymmetry and a small contribution to the neutron EDM in two ways. First, unlike the models of [3] in which gluino decays contribute to the baryon asymmetry with the same CP phase as the neutron EDM, even with a common CP phase for all the $A$-terms the baryon asymmetry and the neutron EDM are proportional to different phases. Thus $\text{Arg}[m_A A'_{323}] = \text{Arg}[m_A A'_{Q}]$ could be maximal, while $\text{Arg}[m_A A'_{Q}]$ could be small. In this case, $\text{Arg}[m_A A'_{Q}]$ would also be maximal, and the baryon asymmetry receives contribution from both the $A$-term and the gluino. Second, the phase of $A_Q$ and $A''_{323}$ might be independent at the messenger scale so that $A''_{323}$ has a large phase while the MSSM $A$-terms could have small ones. The RG running can generates a non-zero (but small) $A_Q$ phase at low energies, that does not contribute too much to the neutron EDM.

Summarizing, the contributions to the neutron EDM depends on the SUSY breaking sector and on the phases generated at that scale. The CP phase needed for baryogenesis is not the same as the one contributing to neutron EDM.

To conclude this section, we have a mechanism for generating the right baryon asymmetry that points to large $R$-parity violation, not too large squark masses, and can be safe from the null experimental results for neutron EDM. Large $R$-parity violation is not a problem if it is confined in interactions involving heavy quarks, otherwise there are many potentially large baryon-number-violating contributions to low-energy flavor physics (see [37] for a review). Even if the only non-zero coupling is $\lambda''_{323}$, couplings involving light quarks are generated at 1-loop level [2]: for $\lambda''_{323} \simeq 1$, we find $\lambda''_{112} \simeq 10^{-8}, \lambda''_{223} \simeq 10^{-5}$, which are too small to significantly contribute to $K\bar{K}$ mixing or $n-\bar{n}$ oscillation. We can also revisit the assumption of single-coupling dominance in the decay of the axino and see if the presence of other couplings is consistent with flavor physics. An important bound for the case with a non-negligible $\lambda''_{313}$ coupling comes from contributions to $\Delta m_K$ [38],

$$|\lambda''_{323} \lambda''_{313}| < 3 \times 10^{-2} \left(\frac{m_t}{1 \text{ TeV}}\right)^2$$

which for $\lambda''_{323} \simeq 1$ and TeV-scale squarks, implies $|\lambda''_{313}| \lesssim 3 \times 10^{-2}$. Then, the single-coupling dominance assumption was justified and the $\lambda''_{313}$-mediated contribution to the axino decay is negligible.

### 4 Dark matter

We now turn our attention to the presence of dark matter. We first recall the dark matter density from the Planck satellite’s CMB measurements (combined with WMAP9 polarization maps) [13],

$$\Omega_{\text{DM}} h^2 = 0.1199 \pm 0.0027.$$  

Without $R$-parity, no supersymmetric particle is stable and indeed the axino, which can be a viable dark matter candidate in $R$-parity-conserving models [26, 39], decays and generates the baryon asymmetry. There are two natural candidates for dark matter that are already in the model: axions and gravitinos. Coherent oscillation of the axions can give rise to cold dark matter if the PQ symmetry breaking scale is properly taken. If the gravitino is the LSP, its lifetime can be long enough that it constitutes the dark matter at present times. The abundance of the gravitino can be sizable by taking proper values of its mass and the reheating temperature $T_R$.
4.1 Heavy gravitino scenario

When the gravitino is heavy enough to decay through the $R$-parity conserving interactions, the only possible candidate for dark matter is the axion. Axion cold dark matter is generated when the axion starts to oscillate coherently at the QCD phase transition. Its abundance is given as \[ \Omega_a h^2 = \frac{1}{\Delta_a} k_a \theta_a^2 \left( \frac{v_{\text{PQ}}}{10^{12} \text{GeV}} \right)^{7/6}, \] (4.2)
where $k_a$ is a numerical factor of $\mathcal{O}(1)$, $\theta_a$ is the axion misalignment angle, and $\Delta_a$ is the possible dilution factor from entropy release when axinos decay after the axion coherent oscillation has started. In viable parameter regions, we find that $\Delta_a$ is just $\mathcal{O}(1)$.

The initial angle $\theta_a$ is not averaged out because we assume the PQ symmetry is broken from the inflation epoch. There is no dark matter contribution from the axionic string decays for the same reason. With the natural value of the angle $\theta_a^2 = \langle \theta_a^2 \rangle \sim 3$, $v_{\text{PQ}} \sim 10^{11} \text{GeV}$ explains the present density of dark matter. From the dark matter constraint, a larger value of $v_{\text{PQ}}$ is allowed if we take a small value of $\theta_a$. However, $v_{\text{PQ}}$ cannot be too large, otherwise axinos will decay after the BBN era. Using eqs. (3.8) and (3.13), $v_{\text{PQ}}$ can be represented as

\[ v_{\text{PQ}} = 10^{11} \text{GeV} \left[ \frac{\lambda''_{323}}{(c\Phi)^1/2} \right]^{1/2} \left( \frac{m_{\tilde{a}}}{0.5 \text{ TeV}} \right)^{3/2} \left( \frac{0.8 \times 10^{-10}}{Y_{\Delta B}} \right)^{1/2} \left( \frac{1 \text{ TeV}}{M_{\text{SUSY}}} \right)^{5/2} \left( \frac{T_R}{10^5 \text{ GeV}} \right)^{1/2}, \] (4.3)
for $T_D > T_{\text{eq}}$. Otherwise, for $T_D < T_{\text{eq}},$

\[ v_{\text{PQ}} = 10^{12} \text{GeV} \left[ \frac{\lambda''_{323}}{(c\Phi)^1/2} \right]^{1/5} \left( \frac{m_{\tilde{a}}}{0.5 \text{ TeV}} \right)^{17/10} \left( \frac{Y_{\Delta B}}{0.8 \times 10^{-10}} \right)^{2/5} \left( \frac{10 \text{ MeV}}{T_D} \right)^{7/5}, \] (4.4)
where $c$ is $\mathcal{O}(1)$ coefficient. In these expressions, we set all SUSY breaking parameters as a common scale, $M_{\text{SUSY}}$, for simplicity. Thus, we get $v_{\text{PQ}} \lesssim 10^{12} \text{GeV}$ for reasonable parameter values. The allowed range is rather small,

\[ 10^{11} \text{GeV} \lesssim v_{\text{PQ}} \lesssim 10^{12} \text{GeV}. \] (4.5)

In order to produce sizable baryon asymmetry, the reheating temperature should be high enough, but it is notable that $T_R$ need not be as large as $v_{\text{PQ}}$. This is consistent with the assumption that the PQ symmetry is not restored in the reheating epoch. As an example, the observed dark matter abundance and baryon asymmetry are generated for $v_{\text{PQ}} = 10^{11} \text{GeV}$, $T_R = 1.5 \times 10^{7} \text{GeV}$ at the benchmark point BP1, where $m_{\tilde{a}} = 500 \text{GeV}$, $m_{\tilde{g}} = 900 \text{GeV}$, $m_{\tilde{q}} = 1.5 \text{TeV}$, $A'_{323} \simeq X_1 = m_{\tilde{q}}$, $\lambda''_{323} = 1$, $\Phi = 1$. Because the baryon asymmetry is inversely proportional to the squark mass $m_{\tilde{q}}$ and the PQ scale $v_{\text{PQ}}$ (through the axino abundance $n_{\tilde{a}}/s$), with such a high value of $v_{\text{PQ}}$ we can repeat the argument leading to eq. (3.19) and find an absolute upper bound of $m_{\tilde{q}} \lesssim 10\sqrt{N_c} \simeq 15 \text{TeV}$. Note that this upper limit is found taking $m_{\tilde{a}} \simeq m_{\tilde{g}}$ and large $A$-terms, $A'_{323} \simeq X_1 \simeq 3m_{\tilde{q}}$, so that it corresponds to a rather compressed region of the parameter space. For a more natural choice of parameters, the squark mass has to be below 8 TeV.

On the other hand, although the gravitino is not a present dark matter candidate, its lifetime can be long enough to cause problems. The decay rate of the gravitino is

\[ \Gamma_{3/2} = \frac{1}{32\pi} \left( n_{\nu} + \frac{n_C}{12} \right) \frac{m_{\tilde{a}3/2}^3}{M_P^2} = \left( 4 \times 10^5 \text{sec}^{-1} \right) \left( n_{\nu} + \frac{n_C}{12} \right) \left( \frac{m_{\tilde{a}3/2}}{\text{TeV}} \right)^3, \] (4.6)
where \( n_V \) \((n_C)\) is the number of vector \((\text{chiral})\) supermultiplets whose masses are smaller than the gravitino mass. When the gravitino is heavier than the MSSM sparticles, its decay products and their amounts are strongly constrained by successful prediction of the standard Big Bang nucleosynthesis \([41]\). For reheating temperatures around \(10^7\) GeV \(\) \(\text{(needed to generate enough axinos)}\), the gravitinos have to decay before the BBN era, i.e. \(\tau_{3/2} < \mathcal{O}(0.1)\) sec. This requires \(m_{3/2} \gtrsim 50\) TeV, corresponding to a spectrum typical of anomaly-mediation of SUSY breaking.

The late time decay of heavy gravitinos could also contribute to the baryon asymmetry, as in \([3]\). However, in \([3]\) the gravitinos were dominating the energy density of the universe at decay time, implying \(T_R \sim 10^{15}\) GeV, while in our case the yield of the gravitino is too small to contribute to a baryon asymmetry of \(n_B/s \sim 10^{-10}\).

### 4.2 Light gravitino scenario

When the gravitino is the true LSP, the axino is the NLSP and can decay to the gravitino and the axion with a decay rate \([42]\)

\[
\Gamma_{\tilde{a} \rightarrow a \psi_{3/2}} = \frac{1}{96\pi} \frac{m_{\tilde{a}}^5}{M_p^2 m_{3/2}^2} \left( 1 - \frac{m_{3/2}^2}{m_{\tilde{a}}^2} \right)^{1/2}.
\]

Because the branching ratio \(\text{Br}(\tilde{a} \rightarrow a \psi_{3/2})\) is quite small, the baryogenesis mechanism is effectively the same as for an axino LSP, and our previous discussion holds. However the non-thermal production of gravitinos by axino decays can provide a sizable abundance of dark matter as

\[
\Omega_{3/2}^{\text{NTP}} h^2 = 0.274 \times 10^9 \ \text{Br}(\tilde{a} \rightarrow a \psi_{3/2}) \left( \frac{m_{4/2}}{\text{GeV}} \right) \left( \frac{m_{\tilde{a}}}{s} \right)_{T_D} \left( \frac{1}{\lambda_{323}''} \right)^2 \left( \frac{m_{\tilde{q}}}{1\ \text{TeV}} \right)^4 \left( \frac{1\ \text{GeV}}{m_{3/2}} \right) \min \left[ 0.048 \left( \frac{T_R}{10^7\ \text{GeV}} \right), 0.015 \left( \frac{v_{\text{PQ}}}{10^{10}\ \text{GeV}} \right)^2 \right].
\]

The second line is evaluated for \(T_D > T_{\text{eq}}\). For \(T_D < T_{\text{eq}}\), there is a further dilution by the factor \(T_D/T_{\text{eq}}\). On the other hand, the thermal production at reheating reads \([15, 43, 44]\)

\[
\Omega_{3/2}^{\text{TP}} h^2 \simeq 0.07 \left( \frac{m_{3/2}}{1\ \text{TeV}} \right)^2 \left( \frac{1\ \text{GeV}}{m_{3/2}} \right) \left( \frac{T_R}{10^7\ \text{GeV}} \right).
\]

We note that for given \(m_{3/2}\) and \(m_{\tilde{q}} \gtrsim m_{\tilde{q}}\), the thermal production is always the dominant contribution. For relatively low \(T_R\), the non-thermal production also can be important when the gluino is lighter than squarks, also for small \(\lambda_{323}''\).

If light gravitinos are produced in the right amount, they can give the correct relic density, provided that their lifetime is longer than the age of the universe \(\) \(\text{(they decay via RPV interactions,} \psi_{3/2} \rightarrow qqq)\). As a matter of fact, the condition on the gravitino lifetime is stronger, as the hadronic decay products would contribute to the cosmic ray antiparticle population, which is looked at in experiments such as PAMELA or AMS-02 \([45–47]\). For example, in \([48]\) it was shown that the lifetime of a vanilla DM candidate decaying to \(b\bar{b}\) is constrained to be bigger than about \(5 \times 10^{27}\) sec from the non-observation by PAMELA of an excess in the \(p/\bar{p}\) fraction, for \(80\) GeV \(\lesssim m_{\text{DM}} \lesssim 500\) GeV \(\) \(\text{(future antideuterons experiments will do better in the lower mass range)}\). To translate these results to the case of a gravitino
decaying to three quarks, it is necessary to find how many antiprotons are generated and compare it to the case of a $b\bar{b}$ final state, for each value of the DM mass. This effort is being tackled by one of the authors in [49], and it is generally found that the number of antiprotons in the experimental energy range produced in the $\psi_{3/2} \rightarrow q\bar{q}q$ case is approximately the same as in the $\chi \rightarrow b\bar{b}$ case, with variations of around $\pm 30\%$, depending on the particle mass and the specific flavor structure of the final three-quarks state. It is then reasonable to take the lower bound $\tau_{\psi_{3/2}}^{\text{exp}} \gtrsim 10^{27}$ sec on the gravitino lifetime, when it makes up all of the dark matter. This bound is conservative enough to not be sensitive to the uncertainties in the precise number of antiparticles arising from the gravitino decay.

The gravitino lifetime can be computed as [47]

$$\tau_{\psi_{3/2} \rightarrow u_i d_j d_k} = 1.28 \times 10^{26} \text{sec} \left( \frac{1}{\lambda_{ijk}'} \right)^2 \left( \frac{3 \text{GeV}}{m_{3/2}} \right)^7 \left( \frac{m_q}{1 \text{TeV}} \right)^4. \quad (4.10)$$

In our model the biggest coupling is $\lambda''_{323}$, allowing the decay channel $\psi_{3/2} \rightarrow t\bar{b}s$ for $m_{3/2} \gtrsim m_t + m_b + m_s$. For a gravitino lighter than the top quark, decays would be mediated by the biggest $\lambda_{ijk}'$ coupling with $i \neq 3$. For example, if the next non-negligible coupling is $\lambda''_{323}$ the decay would go through the $cbs$ channel down to the bottom quark mass. Even with $\lambda''_{323} \simeq 1$, this coupling would not contribute to baryogenesis as the axino partial decay rate would be proportional to the charm quark mass (instead of the top quark mass). Remembering that $\lambda''_{323} \approx 10^{-5} \lambda''_{323}$ is generated at 1-loop anyway, we can consider the range $10^{-5} \lesssim \lambda''_{323} \lesssim 1$. The lower bound on the DM lifetime $\tau_{3/2} \gtrsim 10^{27}$ seconds implies an upper bound on the gravitino mass, $60 \text{GeV} \gtrsim m_{3/2} \gtrsim 4 \text{GeV}$ (for $m_q = 1 \text{TeV}$; these bounds scale as $m_q^{4/7}$).

Finally, there is also a lower bound on the gravitino mass, coming from the one-loop proton decay channel $p \rightarrow K^+ \psi_{3/2}$ setting a limit on $\lambda''_{323}$ [50]:

$$\lambda''_{323} \lesssim 5 \times 10^{-8} \left( \frac{m_q}{300 \text{GeV}} \right)^2 \left( \frac{m_{3/2}}{1 \text{TeV}} \right). \quad (4.11)$$

For $\lambda''_{323} = 1$, $m_q = 1 \text{TeV}$, the corresponding lower limit on the gravitino mass is $m_{3/2} \gtrsim 2 \text{MeV}$.

In figure 3, we fix the axino mass to $m_{\tilde{q}} = 0.5 \text{TeV}$ and vary the remaining parameters ($\lambda''_{323}, \lambda''_{323}, X_t, m_{3/2}$) in the $v_{PQ} - T_R$ plane. The black horizontal dashed lines are contours of different values of $m_{3/2}$ that give the correct dark matter relic abundance in the range $2 \text{MeV} \lesssim m_{3/2} \lesssim 4 \text{GeV}$. The ranges excluded by proton decay and cosmic ray observations are respectively shown at the bottom in green and at the top in red (taking $\lambda''_{323} \simeq 1$). The non-thermal production from axino decays contributes to a dip in the lines, more easily seen in the figure on the right. For higher reheating temperatures, a spike can be seen when axinos decay as they dominate the energy density of the universe ($T_D < T_{eq}$; at higher reheating temperatures axinos are produced more efficiently); this dilutes the gravitinos produced at reheating, allowing a higher reheating temperature than naively thought. The diagonal lines (becoming vertical around the center of the plot) are contours that give the correct baryon asymmetry, with different values of the soft terms and the CP phase $\Phi$. Their behavior can be understood in the following way: for low reheating temperatures, the axino yield depends on both $T_R$ and on the PQ scale $v_{PQ}$, see eq. (3.7). For high $T_R$, the yield is just given by the thermal scattering expression, independent of $v_{PQ}$. As we increase the variables $\Phi, A''_{ij}, X_i$ that determine the asymmetry parameter $\epsilon$, the correct baryon asymmetry can be generated at higher values of $v_{PQ}$, that is, with weaker axino interactions. In the shaded region to the left, not enough asymmetry can be generated, because the axino yield is too
Figure 3. Constraints in the $T_R$–$
u_{PQ}$ plane, keeping $m_3 = 500 \text{GeV}$ fixed. The black dashed lines are contours for different gravitino masses that give the observed dark matter abundance, while the diagonal/vertical continuous and dashed lines correspond to different values of $\Phi''_A = \Phi_A = \Phi, A''_{323}, X_t$ reproducing the correct baryon asymmetry. The shaded regions are excluded by the corresponding labelled constraints, which are further explained in the text. (a) Masses are fixed at $m_{\tilde{q}} = 1 \text{TeV}, m_{\tilde{g}} = 2 \text{TeV},$ corresponding to the benchmark point BP2. (b) for a light gluino, the baryon asymmetry is dominated by the gluino-mediated diagram; the fixed parameters are $m_{\tilde{q}} = 2 \text{TeV}, m_{\tilde{g}} = 1 \text{TeV},$ corresponding to BP3.

small to start with and $\epsilon$ cannot be too large. In the light-blue region on the right, a correct baryon asymmetry can be generated only by taking large values of the soft $A$-terms, which is dangerous from the point of view of color-breaking vacua (the stop squarks might acquire a vev). We excluded the region with $A''_{323} \simeq X_t \gtrsim 3m_{\tilde{q}}$. On the right, in the light yellow region the axion can be dark matter (depending on the precise value of the misalignment angle) and gravitinos can either have decayed already or be a sub-dominant dark matter component, while in the gray rightmost region the axino decays at $0.1\text{–}1 \text{sec}$, compromising the observed abundances for light nuclei produced during BBN.

We note that both baryogenesis and dark matter can be accounted for in most of the “axion window”, $10^9 \text{GeV} < \nu_{PQ} < 10^{12} \text{GeV}$, for reheating temperatures as low as $10^3 \text{TeV}$ and as high as $10^7\text{–}10^8 \text{GeV}$. It is interesting to point out that the choice of fixed parameters

\footnote{This is a conservative estimate, as slightly lower values of $X_t$ could also generate unstable/metastable vacua. See refs. [33, 34] for a more detailed discussion on the constraints on the MSSM $A$-terms. More recently, RPV $A$-terms were considered in ref. [51], under the assumption of CMSSM-like boundary conditions (universal $A$-terms at the GUT scale). Because the masses of $\tilde{b}_L$ and $\tilde{s}_L$ are also important for vacuum stability but do not appear in the baryon asymmetry parameter $\epsilon$ in eq. (3.13), we leave the general study of the effect of RPV $A$-terms for future work.}
in figure 3 is in some way optimal: for heavier squark masses (for fixed $m_{\tilde{a}}/m_{\tilde{q}}$) the asymmetry parameter $\epsilon$ becomes smaller. For an almost degenerate axino LSP, $m_{\tilde{a}} \approx m_{\tilde{q}}$, more parameter space opens up, as $\epsilon$ is bigger, and smaller $A$-terms (and phases) are allowed; in this case, the higher squark mass allowed is 3.5 TeV.

We finish this sub-section with a comment on the Higgs mass: large $A$-terms are needed to achieve a 125 GeV Higgs boson with light stops in the MSSM, and at the same time large $A$-terms increase the asymmetry parameter $\epsilon$. A 125 GeV Higgs with maximal mixing ($X_t \approx \sqrt{6} m_{\tilde{q}}$) allows non-maximal values for $\lambda''_{323}$ and the CP-violating phase $\Phi$, such as $\Phi = 0.03$.

5 Conclusions

We have discussed a new mechanism for baryogenesis through the $R$-parity violating decay of an axino LSP, at the two-loop level and at the second order in the baryon-number-violating couplings. A suitable dark matter density is also generated by related processes, namely by the coherent oscillation mechanism for axions and by thermal scatterings and the axino decays for gravitinos. The scenarios described are very predictive: for the case of axion dark matter, the allowed range for the squarks extends to about 15 TeV; additionally, the initial axion misalignment angle is large. For the case of gravitino dark matter, the gravitino mass is between a few MeV and a few GeV, with proton decay and cosmic rays experiments capable of narrowing this interval; in this case the upper limit on squark masses is higher, of order 130 TeV. The cited limits on the squark masses correspond to tuned regions of the parameter space, where the axino mass is very close to the squarks masses; requiring that the axino and the squarks masses are different by at least 20% brings down the upper squark mass limits to 8 TeV and 90 TeV. In both cases, the axino should be close to the squark mass, up to a factor of a few, and $R$-parity violation should be maximal, corresponding to prompt decays of superpartners.

At the LHC, the most important signatures of light RPV squarks are multijets, with at least two jets from each squark, and three jets from the decay of a gluino. The most relevant experimental searches are [52] from CMS and [53, 54] from ATLAS. In particular, ref. [53] studied the decay of pair-produced gluinos to six quarks, and used $b$-tagging to probe the flavor structure of the RPV couplings $\lambda''_{ijk}$. Gluino masses below 874 GeV are excluded for gluinos whose decay products include a top and a bottom (as it is the case for large $\lambda''_{323}$ coupling). This limits are close to excluding the benchmark point BP3, which had $m_{\tilde{g}} = 1$ TeV. Unfortunately the limits on the gluino masses, apart from the matter of naturalness, are of little importance for our baryogenesis model, which can be mediated by squarks only (even for the gluino-mediated process, the dependence on $m_{\tilde{g}}$ is weak, $\epsilon_{\tilde{g}} \propto 1/m_{\tilde{g}}$). In fact, because the cross section for pair-produced stops is smaller than for pair-produced gluinos, RPV squarks are relatively unprobed at the LHC; for example, LSP squarks are best probed at the Tevatron by the CDF experiment, excluding squark masses up to about 100 GeV [55]. With dedicated searches, the LHC at 14 TeV has the potential to exclude RPV squarks up to about a TeV [56, 57]. For our axion dark matter scenario, a big part of the natural region of the parameter space can be probed.

Acknowledgments

A.M. and C.S.S. are supported in part by DOE grants doe-sc0010008, DOE-ARRA- SC0003883, and DOE-DE-SC0007897.
A \mathcal{L}_{|\Delta B|=2} for non-universal right-handed squark masses

From the Lagrangian (3.5), we get the following interaction terms for the gluino exchange

$$
\mathcal{L}_{|\Delta B|=2}^g = - \sum_{\text{all indices}} \frac{g^2 m_B^4 \lambda_{\alpha \beta \gamma} \lambda_{\delta \zeta \epsilon \eta} \epsilon_{ij} \epsilon_{i'j'k'}}{4 |m_B|^2} \times \\
\left\{ 
\begin{array}{l}
\frac{((d_\alpha^c)^j (d_\gamma^c)^k)}{4 m_{\tilde{u}_{R\alpha}} m_{\tilde{u}_{R\gamma}}} ((u_\alpha^c)^{i'} (u_\gamma^c)^{i}) - \frac{1}{3} (u_\alpha^c)^{i'} (u_\gamma^c)^{i}) ((d_\epsilon^c)^j (d_\zeta^c)^k) \\
+ \frac{((u_\alpha^c)^i (d_\gamma^c)^j)}{m_{\tilde{d}_{R\alpha}} m_{\tilde{d}_{R\gamma}}} ((d_\epsilon^c)^j (d_\zeta^c)^k) - \frac{1}{3} (d_\epsilon^c)^j (d_\zeta^c)^k) ((u_\gamma^c)^{i'} (d_\epsilon^c)^{i'}) \\
+ \frac{((u_\alpha^c)^i (d_\gamma^c)^j)}{m_{\tilde{d}_{R\alpha}} m_{\tilde{d}_{R\gamma}}} ((d_\epsilon^c)^j) - \frac{1}{3} (d_\epsilon^c)^j ((u_\gamma^c)^{i'} (d_\epsilon^c)^{i'}) 
\end{array}
\right\}, \quad (A.1)
$$

and for the bino exchange

$$
\mathcal{L}_{|\Delta B|=2}^B = - \sum_{\text{all indices}} \frac{g^2 m_B^4 \lambda_{\alpha \beta \gamma} \lambda_{\delta \zeta \epsilon \eta} \epsilon_{ij} \epsilon_{i'j'k'}}{2 |m_B|^2} \times \\
\left\{ 
\begin{array}{l}
\frac{Y_u^2 ((d_\beta^c)^j (d_\gamma^c)^k) ((u_\alpha^c)^i (u_\gamma^c)^i) ((d_\epsilon^c)^j (d_\zeta^c)^k)}{4 m_{\tilde{u}_{R\alpha}} m_{\tilde{u}_{R\gamma}}} \\
+ \frac{Y_{\tilde{d}} ((u_\alpha^c)^i (d_\gamma^c)^j) ((d_\epsilon^c)^j (d_\zeta^c)^k) ((u_\gamma^c)^{i'} (d_\epsilon^c)^{i'})}{m_{\tilde{d}_{R\alpha}} m_{\tilde{d}_{R\gamma}}} \\
+ \frac{Y_{\tilde{d}} ((u_\alpha^c)^i (d_\gamma^c)^j) ((d_\epsilon^c)^j) ((u_\gamma^c)^{i'} (d_\epsilon^c)^{i'})}{m_{\tilde{d}_{R\alpha}} m_{\tilde{d}_{R\gamma}}} 
\end{array}
\right\}, \quad (A.2)
$$

where $ij \, i'j'k'$ are the indices of SU(3) anti-fundamental representation, and $Y_q$ is the U(1)$_Y$ hypercharge of $q$. We used

$$
\sum_A \left[ [T^A]^a_i (q_1)^b_j \right] \left[ [T^A]^c_d (q_2)^d_j \right] = \frac{1}{2} (q_1)^c (q_2)^a - \frac{1}{6} (q_1)^a (q_2)^c, \quad (A.3)
$$

for SU(3)$_c$ gauge group. If the coupling $\lambda_{323}^\mu = - \lambda_{332}^\mu$ is only nonzero, those interactions become

$$
\mathcal{L}_{|\Delta B|=2}^g = - \frac{g^2 m_B^4 \lambda_{\alpha \beta \gamma} \lambda_{\delta \zeta \epsilon \eta} \epsilon_{ij} \epsilon_{i'j'k'}}{3 |m_B|^2} \times \\
\left\{ 
\begin{array}{l}
\frac{1}{m_{\tilde{u}_{R\alpha}}^4} + \frac{1}{m_{\tilde{d}_{R\alpha}}^4} + \frac{1}{m_{\tilde{d}_{R\gamma}}^4} \right) \left((t^c)^i (t^c)^j)((s^c)^k (b^c)^{i'}) \right) ((s^c)^k (t^c)^{j'}) \\
+ \left( \frac{1}{m_{\tilde{d}_{R\alpha}}^2} + \frac{1}{m_{\tilde{d}_{R\gamma}}^2} \right) \left((b^c)^j (t^c)^i)((s^c)^k (t^c)^{i'}) \right) ((b^c)^j (s^c)^k) \\
+ \left( \frac{1}{m_{\tilde{d}_{R\alpha}}^2} + \frac{1}{m_{\tilde{d}_{R\gamma}}^2} \right) \left((t^c)^i (s^c)^k)((b^c)^j (t^c)^{i'}) \right) ((b^c)^j (s^c)^k) 
\end{array}
\right\}, \quad (A.4)
$$
and

\[ \mathcal{L}_{|\Delta B|=2} = - \frac{g^2 m_B^4 \lambda_{33}^2}{9 |m_B|^2} \left[ \frac{1}{m_{sR}^4} + \frac{1}{m_{bR}^4} - \frac{2}{m_{sR}^2 m_{bR}^2} \right] \left( (t_\epsilon^i)(b_\epsilon^j)(s^k_\epsilon^{(i})(b_\epsilon^j)^{k'})((s^k_\epsilon^{(e)})^{k'}(t_\epsilon^i)^{e'}) \right) \\
+ \left( \frac{4}{m_{tR}^2} + \frac{1}{m_{bR}^2} + \frac{2}{m_{sR}^2 m_{tR}^2} \right) \left( (b_\epsilon^j)(t_\epsilon^i)^{(e')}((s^k_\epsilon^{(e)})^{k'}((b_\epsilon^j)^{(e')}((s^k_\epsilon^{(e)})^{k'} \right) \\
+ \left( \frac{4}{m_{tR}^2} + \frac{1}{m_{sR}^2} + \frac{2}{m_{sR}^2 m_{tR}^2} \right) \left( (t_\epsilon^i)^{(e')}((s^k_\epsilon^{(e)})^{k'}((b_\epsilon^j)^{(e')}((s^k_\epsilon^{(e)})^{k'} \right) \right]. \quad (A.5) \]

We used the identity \((\psi_x \psi_y)(\psi_x \psi) + (\psi_y \psi_x)(\psi_y \psi) = 0\) for chiral fermions \(\psi_x, \psi_y, \psi_x, \psi\). Now we can easily evaluate \(\epsilon_\tilde{g}\) and \(\epsilon_\tilde{B}\), and obtain the result of (3.13) with

\[
\kappa_\tilde{g} = \frac{1}{3} + \frac{m_{tR}^4}{6 m_{sR}^4} + \frac{m_{tR}^4}{6 m_{bR}^4} + \frac{m_{sR}^2}{6 m_{bR}^4},
\]

\[
\kappa_\tilde{B} = \frac{8}{9} + \frac{m_{tR}^4}{9 m_{sR}^4} + \frac{m_{tR}^4}{9 m_{bR}^4} + \frac{2 m_{tR}^4}{9 m_{sR}^4} + \frac{2 m_{tR}^4}{9 m_{bR}^4}. \quad (A.6) \]

### B Exact expressions for decay rate and asymmetry parameter

In section 3 we presented the decay rate \(\Gamma_\tilde{g}\) and the asymmetry parameters \(\epsilon_\tilde{A}, \epsilon_\tilde{g}\) in the limit of heavy internal squarks and massless final states. From the Feynman diagrams in figure 2, the full expressions are of the form

\[
\Gamma_{\tilde{g} \to t^0 b^0 s^0} = \epsilon_{\alpha\beta\gamma} \epsilon_{\alpha\beta'} \frac{\epsilon_{\alpha}^\prime \epsilon_{\beta'}^\prime}{2m_\tilde{g}} \int \frac{d^3 p_{t}}{(2\pi)^3} \frac{d^3 p_{b}}{(2\pi)^3} \frac{d^3 p_{s}}{(2\pi)^3} (2\pi)^4 \delta^{(4)}(p_1 - p_t - p_b - p_s) \\
\times \left( \frac{(p_1 \cdot p_t)(p_b \cdot p_s) + 2m_{tR}^2 p_{tR} \cdot p_s}{m_{tL}^2} \right) \left( \frac{(p_1 \cdot p_t)(p_b \cdot p_s)}{m_{tL}^2} \right), \quad (B.1) \]

\[
\epsilon_\tilde{A} = -6 \epsilon_\tilde{g} m_{tR}^4 \frac{\epsilon_{\alpha}^\prime \epsilon_{\beta'}^\prime}{m_{tL}^2} \int \frac{d^3 p_{t}}{(2\pi)^3} \frac{d^3 p_{b}}{(2\pi)^3} \frac{d^3 p_{s}}{(2\pi)^3} (2\pi)^4 \delta^{(4)}(p_1 - p_t - p_b - p_s) \\
\int \frac{d^3 k_0}{(2\pi)^3} \frac{d^3 k_0}{(2\pi)^3} \frac{d^3 k_0}{(2\pi)^3} (2\pi)^4 \delta^{(4)}(p_1 - k_0 - k_0 - k_0) \\
\times \frac{4(p_b \cdot p_s)(k_0 \cdot k_0)}{((k_0 + k_0)^2 - m_{tR}^2)((p_b + p_s)^2 - m_{tR}^2)((k_0 + k_0)^2 - m_{tL}^2)((p_1 - k_0)^2 - m_{tL}^2)((k_0 + k_0)^2 - m_{tL}^2)} \quad (B.2) \]

where \(p_1\) is the initial momentum of the axino, \(p_i\)'s are the momenta of the final states \(t, b, s\) and \(k_i\)'s are the momenta of the on-shell intermediate states that generate an imaginary part for the integral. The expression for \(\epsilon_\tilde{A}\) has a similar form, where the denominator is changed to include the gaugino propagator.
Corrections for sizable top mass. The top quark mass is not negligible for axino masses below a TeV. For a non-zero top mass, the available phase space is reduced, and we write down the full dependence on \( m_t \):

\[
\Gamma_{\text{EFT}}^{\text{exact}} = \frac{\kappa_3^2 |\lambda_{323}|^2}{512\pi^3} \frac{m_A^5 m_t^2}{v_{1}\Pi m_{1}} \left[ f(x) + 8x g(x) \frac{m_t X_t}{m_t^2} + f(x) \left( \frac{m_t X_t}{m_t^2} \right)^2 \right], \tag{B.3}
\]

\[
\epsilon_{\text{EFT}}^{\text{exact}} = \frac{|\lambda_{323}|^4}{32\pi^3} m_{\tilde{a}}^2 m_{\tilde{g}}^2 m_{\tilde{g}} \frac{\text{Im}[m_{\tilde{a}} A_{323}^{\ast}]}{m_{\tilde{a}}^4 m_{\tilde{g}}^2} \left[ g(x)^2 + \frac{f(x) g(x) X_{m_{\tilde{a}}} + \frac{f(x)^2}{16} \left( \frac{m_{\tilde{a}} X_t}{m_t^2} \right)^2}{f(x) + 8x g(x) \frac{m_t X_t}{m_t^2} + f(x) \left( \frac{m_t X_t}{m_t^2} \right)^2} \right]. \tag{B.4}
\]

Here \( x = \frac{m_t}{m_{\tilde{a}}} \) and the functions \( f(x), g(x) \) are defined as follows:

\[
f(x) \equiv 1 - 8x^2 + 8x^6 - x^8 - 24x^4 \log x, \tag{B.4}
\]

\[
g(x) \equiv 1 + 9x^2 - 9x^4 - x^6 + 12(x^2 + x^4) \log x. \tag{B.5}
\]

As expected, both the decay rate and the asymmetry parameter decrease with respect to the massless limit, because the available phase space is smaller. For \( m_{\tilde{a}} = 500 \text{ GeV} \), the corresponding values are

\[
x_0 = 0.346, \quad f(x_0) = 0.42, \quad g(x_0) = 0.24. \tag{B.6}
\]

Corrections for effective theory breakdown. Our results were also computed with the assumption of large squark masses, where the effective field theory approximation is powerful. Because the baryon asymmetry is proportional to \( (m_{\tilde{a}}/m_{\tilde{q}})^3 \), the hierarchy between the axino and the squarks cannot be too big; for instance, we took \( m_{\tilde{a}} = \frac{2}{3}m_{\tilde{q}} = \frac{2}{3}m_{\tilde{q}}, m_{\tilde{a}} = \frac{1}{4}m_{\tilde{q}} = \frac{1}{4}m_{\tilde{q}} \) and \( m_{\tilde{a}} = \frac{1}{2}m_{\tilde{q}} = \frac{1}{2}m_{\tilde{q}} \) as concrete benchmark points in eq. (3.18) and in section 4. Thus, it should be verified that corrections to our previous results are small. In the limit of massless final states, the decay rate computed in the full theory is

\[
\Gamma_{\text{EFT}}^{\text{exact}} = \frac{3\kappa_3^2 |\lambda_{323}|^2}{512\pi^3} \frac{m_A^5 m_t^2}{v_{1}\Pi m_{1}} \left[ 1 + \left( \frac{m_t X_t}{m_t^2} \right)^2 \right] h(y), \tag{B.7}
\]

\[
h(y) \equiv \frac{6y^2 - 5y^4 + 2(3 - 4y^2 + y^4) \log(1 - y^2)}{y^8}, \tag{B.8}
\]

where \( y = \frac{m_{\tilde{a}}}{m_t} \) and at small \( y, h(y) = 1 + \frac{4}{5}y^2 + \frac{3}{2}y^4 + O(y^6) \). As expected, the decay rate increases as the axino mass approaches the squark masses. For the asymmetry parameter \( \epsilon \), the phase space integral is more complicated and we have to rely on numerical integration.
Figure 4. (a) Ratio of $\Gamma_{\text{full}}$, the decay rate computed numerically, and $\Gamma_0$, the analytical result of eq. (3.6), which was computed in the limit of zero $m_t$ and large squark masses, for $X_t \simeq m_{\tilde{q}}$. (b) Ratio of $\epsilon_{\tilde{g}}^{\text{full}}$ and $\epsilon_{\tilde{g}}^0$, the contribution to the asymmetry parameter due to $A$-terms. (c, d) Ratio of $\epsilon_{\tilde{g}}^{\text{full}}$ and $\epsilon_{\tilde{g}}^0$, the contribution to the asymmetry parameter due to gluino exchange. In (c), the gluino has been taken as heavier than the squark, $m_{\tilde{g}} = 2m_{\tilde{q}}$, as at the benchmark point BP2, while in (d) the gluino is lighter, $m_{\tilde{g}} = m_{\tilde{q}}/2$, as for BP3. The star markers indicate the benchmark points used in section 4. For all figures, the left-right squark mixing has been taken as moderate, $X_t \simeq m_{\tilde{q}}$.

**Exact numerical results.** Even in the limit of massless final states, we could not find a simple analytical expression for the asymmetry parameter $\epsilon$ when taking into account the squark and gluino propagators. We can integrate the phase space integrals numerically and check that the simple expressions given in the main part of the article do not introduce a large error. We show the full numerical results for $\Gamma$ and $\epsilon$ in the $m_t$–$m_{\tilde{g}}$ plane in figures 4. The benchmark points of eq. (3.18), used in section 4, are indicated by white star markers. We see that $\Gamma$ and $\epsilon$ decrease by about two and six, respectively. This is mainly due to
relaxing the approximation $m_t = 0$, reducing the phase space available for the decay. The exact numerical results have been used throughout the paper and for figure 3.

References

[1] R.D. Peccei and H.R. Quinn, *CP Conservation in the Presence of Instantons*, Phys. Rev. Lett. 38 (1977) 1440 [insPIRE].

[2] S. Dimopoulos and L.J. Hall, *Baryogenesis at the MeV Era*, Phys. Lett. B 196 (1987) 135 [insPIRE].

[3] J.M. Cline and S. Raby, *Gravitino induced baryogenesis: A Problem made a virtue*, Phys. Rev. D 43 (1991) 1781 [insPIRE].

[4] S. Mollerach and E. Roulet, *Axino induced baryogenesis*, Phys. Lett. B 281 (1992) 303 [insPIRE].

[5] Y. Cui and R. Sundrum, *Baryogenesis for weakly interacting massive particles*, Phys. Rev. D 87 (2013) 116013 [arXiv:1212.2973] [insPIRE].

[6] Y. Cui, *Natural Baryogenesis from Unnatural Supersymmetry*, JHEP 12 (2013) 067 [arXiv:1309.2952] [insPIRE].

[7] G. Arcadi, L. Covi and M. Nardecchia, *Out-of-equilibrium baryogenesis and superweakly interacting massive particle dark matter*, Phys. Rev. D 89 (2014) 095020 [arXiv:1312.5703] [insPIRE].

[8] F. Rompineve, *Weak Scale Baryogenesis in a Supersymmetric Scenario with R-parity violation*, JHEP 08 (2014) 014 [arXiv:1310.0840] [insPIRE].

[9] H.K. Dreiner and G.G. Ross, *Sphaleron erasure of primordial baryogenesis*, Nucl. Phys. B 410 (1993) 188 [hep-ph/9207221] [insPIRE].

[10] M. Bolz, A. Brandenburg and W. Buchmüller, *Thermal production of gravitinos*, Nucl. Phys. B 606 (2001) 518 [hep-ph/0012052] [insPIRE].

[11] M. Dine, W. Fischler and M. Srednicki, *A Simple Solution to the Strong CP Problem with a Harmless Axion*, Phys. Lett. B 104 (1981) 199 [insPIRE].

[12] A.R. Zhitnitsky, *On Possible Suppression of the Axion Hadron Interactions* (in Russian), Sov. J. Nucl. Phys. 31 (1980) 260 [Yad. Fiz. 31 (1980) 497] [insPIRE].
[20] J.E. Kim, Weak Interaction Singlet and Strong CP Invariance, Phys. Rev. Lett. 43 (1979) 103 [inspire].
[21] M.A. Shifman, A.I. Vainshtein and V.I. Zakharov, Can Confinement Ensure Natural CP Invariance of Strong Interactions?, Nucl. Phys. B 166 (1980) 493 [inspire].
[22] K.J. Bae, H. Baer, E.J. Chun and C.S. Shin, Mixed axion/gravitino dark matter from SUSY models with heavy axinos, Phys. Rev. D 91 (2015) 075011 [arXiv:1410.3857] [inspire].
[23] A. Monteux, Natural, R-parity violating supersymmetry and horizontal flavor symmetries, Phys. Rev. D 88 (2013) 045029 [arXiv:1305.2921] [inspire].
[24] E. Nikolidakis and C. Smith, Minimal Flavor Violation, Seesaw and R-parity, Phys. Rev. D 77 (2008) 015029 [arXiv:1305.2921] [inspire].
[25] C. Csáki, Y. Grossman and B. Heidenreich, MFV SUSY: A Natural Theory for R-Parity Violation, Phys. Rev. D 85 (2012) 095009 [arXiv:1111.1239] [inspire].
[26] L. Covi, H.-B. Kim, J.E. Kim and L. Roszkowski, Axinos as dark matter, JHEP 05 (2001) 033 [hep-ph/0101009] [inspire].
[27] A. Brandenburg and F.D. Steffen, Axino dark matter from thermal production, JCAP 08 (2004) 008 [hep-ph/0405158] [inspire].
[28] A. Strumia, Thermal production of axino Dark Matter, JHEP 06 (2010) 036 [arXiv:1003.5847] [inspire].
[29] E.J. Chun, Dark matter in the Kim-Nilles mechanism, Phys. Rev. D 84 (2011) 043509 [arXiv:1110.2219] [inspire].
[30] K.J. Bae, K. Choi and S.H. Im, Effective Interactions of Axion Supermultiplet and Thermal Production of Axino Dark Matter, JHEP 08 (2011) 065 [arXiv:1106.2452] [inspire].
[31] K.J. Bae, E.J. Chun and S.H. Im, Cosmology of the DFSZ axino, JCAP 03 (2012) 013 [arXiv:1111.5962] [inspire].
[32] B.C. Allanach, A. Dedes and H.K. Dreiner, Two loop supersymmetric renormalization group equations including R-parity violation and aspects of unification, Phys. Rev. D 60 (1999) 056002 [hep-ph/9902251] [inspire].
[33] J.E. Camargo-Molina, B. O’Leary, W. Porod and F. Staub, Stability of the CMSSM against sfermion VEVs, JHEP 12 (2013) 103 [arXiv:1309.7212] [inspire].
[34] N. Blinov and D.E. Morrissey, Vacuum Stability and the MSSM Higgs Mass, JHEP 03 (2014) 106 [arXiv:1310.4174] [inspire].
[35] J. Polchinski and M.B. Wise, The Electric Dipole Moment of the Neutron in Low-Energy Supergravity, Phys. Lett. B 125 (1983) 393 [inspire].
[36] C.A. Baker et al., An Improved experimental limit on the electric dipole moment of the neutron, Phys. Rev. Lett. 97 (2006) 131801 [hep-ex/0602020] [inspire].
[37] R. Barbieri et al., R-parity violating supersymmetry, Phys. Rept. 420 (2005) 1 [hep-ph/0406039] [inspire].
[38] P. Slavich, Constraints on R-parity violating stop couplings from flavor physics, Nucl. Phys. B 595 (2001) 33 [hep-ph/0008270] [inspire].
[39] L. Covi, J.E. Kim and L. Roszkowski, Axinos as cold dark matter, Phys. Rev. Lett. 82 (1999) 4180 [hep-ph/9905212] [inspire].
[40] Particle Data Group collaboration, J. Beringer et al., Review of Particle Physics (RPP), Phys. Rev. D 86 (2012) 010001 [inspire].
[41] M. Kawasaki, K. Kohri and T. Moroi, Big-Bang nucleosynthesis and hadronic decay of long-lived massive particles, Phys. Rev. D 71 (2005) 083502 [astro-ph/0408426] [inspire].
[42] E.J. Chun, H.B. Kim and J.E. Kim, *Dark matters in axino gravitino cosmology*, *Phys. Rev. Lett.* 72 (1994) 1956 [hep-ph/9305208] [inSPIRE].

[43] J. Pradler and F.D. Steffen, *Thermal gravitino production and collider tests of leptogenesis*, *Phys. Rev. D* 75 (2007) 023509 [hep-ph/0608344] [inSPIRE].

[44] V.S. Rychkov and A. Strumia, *Thermal production of gravitinos*, *Phys. Rev. D* 75 (2007) 075011 [hep-ph/0701104] [inSPIRE].

[45] N.-E. Bomark, S. Lola, P. Osland and A.R. Raklev, *Photon, Neutrino and Charged Particle Spectra from R-violating Gravitino Decays*, *Phys. Lett. B* 686 (2010) 152 [arXiv:0911.3376] [inSPIRE].

[46] L.A. Dal and A.R. Raklev, *Antideuteron Limits on Decaying Dark Matter with a Tuned Formation Model*, *Phys. Rev. D* 89 (2014) 103504 [arXiv:1402.6259] [inSPIRE].

[47] A. Monteux, E. Carlson and J. Cornell, *Gravitino Dark Matter and Flavor Symmetries*, *JHEP* 08 (2014) 047 [arXiv:1404.5952] [inSPIRE].

[48] A. Ibarra and S. Wild, *Prospects of antideuteron detection from dark matter annihilations or decays at AMS-02 and GAPS*, *JCAP* 02 (2013) 021 [arXiv:1209.5539] [inSPIRE].

[49] E. Carlson, J. Cornell and A. Monteux, *Astroparticle constraints on unstable gravitino dark matter*, in preparation.

[50] K. Choi, K. Hwang and J.S. Lee, *Constraints on R-parity and B violating couplings in gauge mediated supersymmetry breaking models*, *Phys. Lett. B* 428 (1998) 129 [hep-ph/9802323] [inSPIRE].

[51] N. Chamoun, H.K. Dreiner, F. Staub and T. Stefaniak, *Resurrecting light stops after the 125 GeV Higgs in the baryon number violating CMSSM*, *JHEP* 08 (2014) 142 [arXiv:1407.2248] [inSPIRE].

[52] CMS collaboration, *Search for three-jet resonances in pp collisions at √s = 7 TeV*, *Phys. Lett. B* 718 (2012) 329 [arXiv:1208.2931] [inSPIRE].

[53] ATLAS collaboration, *Search for massive particles in multijet signatures with the ATLAS detector in √s = 8 TeV pp collisions at the LHC*, ATLAS-CONF-2013-091 (2013) [ATLAS-COM-CONF-2013-081] [inSPIRE].

[54] ATLAS collaboration, *Search for pair production of massive particles decaying into three quarks with the ATLAS detector in √s = 7 TeV pp collisions at the LHC*, *JHEP* 12 (2012) 086 [arXiv:1210.4813] [inSPIRE].

[55] CDF collaboration, T. Aaltonen et al., *Search for Pair Production of Strongly Interacting Particles Decaying to Pairs of Jets in pp Collisions at √s = 1.96 TeV*, *Phys. Rev. Lett.* 111 (2013) 031802 [arXiv:1303.2699] [inSPIRE].

[56] Y. Bai, A. Katz and B. Tweedie, *Pulling Out All the Stops: Searching for RPV SUSY with Stop-Jets*, *JHEP* 01 (2014) 040 [arXiv:1309.6631] [inSPIRE].

[57] D. Duggan et al., *Sensitivity of an Upgraded LHC to R-Parity Violating Signatures of the MSSM*, arXiv:1308.3903 [inSPIRE].