R-Axion at Colliders

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We study the effective theory of a generic class of hidden sectors where supersymmetry is broken together with an approximate R-symmetry at low energy. The light spectrum contains the gravitino and the pseudo-Nambu-Goldstone boson of the R-symmetry, the R-axion. We derive new model-independent constraints on the R-axion decay constant for R-axion masses ranging from GeV to TeV, which are of relevance for hadron colliders, lepton colliders, and B factories. The current bounds allow for the exciting possibility that the first sign of supersymmetry will be the R-axion. We point out its most distinctive signals, providing a new experimental handle on the properties of the hidden sector and a solid motivation for searches of axionlike particles.

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In this Letter we argue that there are generic signs of supersymmetry (SUSY) to be looked for at colliders that have not yet been satisfactorily explored: those associated with the R-axion, the pseudo-Nambu-Goldstone boson (PNGB) of a spontaneously broken R-symmetry.

Although it is well known that supersymmetry must be broken in a "hidden sector," its dynamics is left unspecified in the vast majority of phenomenological studies, which instead focus on the "visible sector," e.g., the minimal supersymmetric standard model (MSSM). Here we point out that in an extensive class of models the hidden sector leaves its footprints in observables accessible to the current experimental program. In particular, we perform a thorough phenomenological study of the R-axion at high- and low-energy hadron and lepton colliders.

The $\mathcal{N}=1$ SUSY algebra contains a single $U(1)_R$ ("R-symmetry") under which supercharges transform, $[R, Q_\alpha] = -Q_\alpha$, such that components of a given supermultiplet have R charges r differing by one unit (e.g., gauge fields carry no R charge while gauginos have $r_\lambda = 1$). R-symmetry plays a crucial role in models of low-energy dynamical SUSY breaking. According to the general result of Nelson and Seiberg, an R-symmetry must exist in any generic, calculable model which breaks SUSY with F terms, and if the R-symmetry is spontaneously broken then SUSY is also broken [1]. Spontaneous R-symmetry breaking often occurs also in incalculable models like in [2,3]. If the SUSY-breaking vacuum is metastable, like in Intriligator-Seiberg-Shih (ISS) constructions [4], then an analogue of the Nelson-Seiberg result holds for an approximate R-symmetry [5]. When $U(1)_R$ is explicitly broken by a suitable deformation of the hidden sector, the R-axion gets a mass in addition to the irreducible contribution from supergravity [6] but can remain naturally lighter than the other hidden-sector resonances.

In light of the above observations, and contrary to what previously explored in the literature, we treat the R-axion mass m_a as a free parameter, together with its decay constant f_a. Our analysis shows that such a particle could well be the first sign of SUSY to show up in experiments. We also derive model-independent bounds on the scale of spontaneous R-symmetry breaking, opening a new observational window on the properties of the SUSY-breaking hidden sector.

Setup.—We parametrize with m_* the SUSY mass gap of the hidden sector and with g_* the coupling strength between hidden-sector states at m_*. The generic size of the SUSY-breaking vacuum expectation value (VEV) is then $F \sim m_*^2/g_*$, an outcome of naive dimensional analysis (NDA) with a single scale and coupling [7,8]. The R-axion decay constant is $f_a \sim m_*/g_*$, while the R-axion mass should satisfy $m_a \ll m_*$ in order for the R-symmetry axion to be a PNGB.

As a generic consequence of spontaneous SUSY breaking, a light gravitino is also present in the low-energy spectrum. In the rigid limit (i.e., $M_{\rm Pl} \to \infty$) the transverse degrees of freedom of the gravitino decouple, leaving a massless Goldstino in the spectrum. The effective action of the Goldstino and the R-axion can be written using the nonlinear superfield formalism of [9] and reads

$$\begin{aligned}\mathcal{L}_{\rm hid} &= \int d^4\theta\left(X^\dagger X + \frac{f_a^2}{2}\mathcal{R}^\dagger\mathcal{R}\right) \\ &\quad + \int d^2\theta(FX + w_R\mathcal{R}^2) + \text{c.c.} \\ &\supset -F^2 + i\bar{G}\bar{\sigma}^\mu\partial_\mu G + \frac{f_a^2}{2}(\partial_\mu a)^2 \\ &\quad - \frac{w_R}{F^2}(iG^2 e^{-2ia}\Box a + \text{c.c.}), \qquad (1)\end{aligned}$$

where X and $\mathcal{R} = e^{i\mathcal{A}}$ carry R charge 2 and 1, respectively (the R charge of a supermultiplet is identified with that of its bottom component) and satisfy the nonlinear constraints $X^2 = 0$ and $X(\mathcal{R}^\dagger\mathcal{R} - 1) = 0$. The above constraints, in combination with (1), give $X = G^2/2F^2 + \sqrt{2}G\theta - F\theta^2$, $\mathcal{A} = a + O(aG)$; see the Supplemental Material for more details [10].

Since the R charge of $\mathcal{R}$ is 1, its effective action differs from the one of a SUSY axion in that a superpotential term is allowed. This is controlled by the dimension-three parameter w_R, which is related to the VEV of the superpotential and satisfies the inequality $w_R < f_a F/2\sqrt{2}$, under the assumption of no extra light degrees of freedom other than the R-axion and the Goldstino [11,12]. The superpotential term induces cubic interactions between the R-axion and two Goldstini, proportional to m_a^2, that lead to an invisible decay channel for the R-axion. The corresponding decay rate of the R-axion into two Goldstini is

$$\Gamma(a \to GG) = \frac{1}{4\pi}\left(\frac{m_a^5 w_R^2}{f_a^2 F^4}\right) < \frac{1}{32\pi}\frac{m_a^5}{F^2} \qquad (2)$$

and it is bounded from above as a consequence of the upper bound on w_R, saturated only in free theories. Our power counting gives $w_R \sim F f_a$, making the width within an $O(1)$ factor of the upper limit in Eq. (2). For ordinary axions, w_R would instead break explicitly the associated global symmetry, resulting in a suppression of the decay width into Goldstini by extra powers of m_a^2/m_*^2. Hence, a sizable invisible decay width is a distinctive feature of the R-axion compared to other axionlike particles.

The R-axion mass is generated by sources of explicit R-symmetry breaking and can be parametrized as

$$m_a^2 \sim \frac{\epsilon_{\not R} F}{f_a^2} r_\epsilon^2 \ll m_*^2 \quad \text{from } \mathcal{L}_{\not R} = \int d^2\theta \frac{1}{2}\epsilon_{\not R} X R^{-r_\epsilon}, \qquad (3)$$

where r_ϵ is the R charge of the explicit-breaking spurion $\epsilon_{\not R}$, with $\epsilon_{\not R}/F \ll 1$ technically natural. Explicit examples of this mass hierarchy arise by adding suitable R-symmetry breaking deformations in calculable models of dynamical SUSY breaking like the 3-2 model [1,13,14] or in SUSY QCD at large N once the hidden gauginos and squarks get soft masses [15,16]. Moreover, in SUSY-breaking models like the one in [5], the explicit breaking of the R-symmetry is generically bounded from above ($\epsilon_{\not R}/F \ll 1$) by requiring the SUSY-breaking vacuum to be metastable.

The R-symmetry breaking contribution (3) can well be expected to dominate over the unavoidable supergravity (SUGRA) contribution arising from the tuning of the cosmological constant [6], which gives rise to $m_a^2 \sim (10\ \text{MeV})^2 \times m_*/10\ \text{TeV} \times m_{3/2}/\text{eV}$ (we refer to the Supplemental Material for its derivation [10]).

We now study the couplings of the R-axion with the visible-sector fields, which we take to be the MSSM (with matter or R parity). The superpartners get SUSY-breaking masses $m_{\rm soft}$ from their interactions with the hidden sector, which are controlled by a perturbative coupling g. This coupling is a proxy for the SM gauge coupling constants in gauge mediation models [17,18] or for Yukawa-type interactions in extended gauge mediation; see [19] for a review. The scaling of $m_{\rm soft}$ strongly depends on the type of mediation mechanism. We can estimate it as

$$m_{\rm soft} \sim \left(\frac{g}{g_*}\right)^n \times g\frac{F}{m_*} = \left(\frac{g}{g_*}\right)^{n+1} \times m_*. \qquad (4)$$

In this Letter we assume that gauginos get a mass via their coupling to the hidden-sector global current, so that $m_{\rm soft} \sim (g/g_*)^2 m_*$. Notice that if $g_* = 4\pi/\sqrt{N_{\rm mess}}$ we recover the ordinary gauge mediation scaling where $N_{\rm mess}$ is the number of messengers. Other scaling, e.g., the one of [20] for Dirac gauginos ($n = 0$), will be discussed elsewhere. Besides, whatever the scaling in Eq. (4), there is always a large portion of parameter space where the R-symmetry axion is lighter than the superpartners, which correspond to $r_\epsilon\sqrt{g_*\epsilon_{\not R}}/m_* \lesssim (g/g_*)^{n+1}$. It would also be interesting to depart from the NDA expectation for the scales F and f_a and explore models where a large separation between the two is realized.

We consider in the following a small SUSY-breaking scale $\sqrt{F}$ in the range from 1 to a few 100s of TeV. This regime is welcome for fine-tuning and Higgs mass considerations. The resulting gravitino mass lies in the window $10^{-4}\ \text{eV} \lesssim m_{3/2} \lesssim 5\ \text{eV}$, where the upper limit comes from cosmological and astrophysical bounds on gravitino abundance [21–23], while the lower limit comes from LEP [24] and LHC [25,26] bounds.

Since the visible sector feels the SUSY breaking only through g/g_* effects, we can treat the MSSM superfields linearly and "dress" the R-charged operators with appropriate powers of the R-axion. We also neglect subleading effects from explicit R-breaking terms, suppressed by powers of $\sim m_a/m_*$. The interactions of the R-axion with the MSSM gauge sector are then

$$\mathcal{L}_{\rm gauge} = \int d^2\theta \left(-i g_i^2 \frac{c_i^{\rm hid}}{16\pi^2} \mathcal{A} - \frac{m_{\lambda_i}}{2F} X \mathcal{R}^{-2} \right) \mathcal{W}_i^2 + {\rm c.c.}$$
$$\supset \frac{g_i^2 c_i^{\rm hid}}{16\pi^2} \frac{a}{f_a} F^i \tilde{F}^i - \frac{m_{\lambda_i}}{2} \lambda_i \lambda_i \left(e^{-2ia/f_a} + \frac{g_i^2 c_i^{\rm hid}}{4\pi^2} i \frac{a}{f_a} \right)$$
$$+ {\rm c.c.}, \tag{5}$$

where $\mathcal{W}$ is the field strength superfield ($r_{\mathcal{W}} = 1$) and i labels the SM gauge group, where we defined $g_1 = \sqrt{5/3} g_Y$ and $\tilde{F}^{i,\mu\nu} = 1/2 \epsilon^{\mu\nu\rho\sigma} F^i_{\rho\sigma}$. The Majorana gaugino masses are of order $m_{\lambda_i} \approx m_{\rm soft}$ by assumption. The coefficients $c_i^{\rm hid}$ encode the hidden-sector contributions to the mixed anomalies of the $U(1)_R$ with the SM gauge groups. For example, we get $c_i^{\rm hid} = -N_{\rm mess}$ for $i = 1, 2, 3$, for $N_{\rm mess}$ messengers chiral under $U(1)_R$ and in the $5 + \bar{5}$ of $SU(5)$ with zero R charge [in our NDA, $N_{\rm mess} \sim (4\pi/g_*)^2$]. We encode the contributions to the anomalies from the MSSM fields in the full loop functions.

The interactions in the Higgs sector can be written as

$$\mathcal{L}_{\rm Higgs} = \int d^4\theta \left(\frac{\mu}{F} X^\dagger \mathcal{R}^{2-r_H} - \frac{B_\mu}{F^2} |X|^2 \mathcal{R}^{-r_H} \right) H_u H_d + {\rm c.c.}$$
$$\supset \mu \tilde{h}_u \tilde{h}_d e^{i(2-r_H)a/f_a} - B_\mu h_u h_d e^{-i r_H a/f_a} + {\rm c.c.}, \tag{6}$$

where $\tilde{h}_{u,d}$ are the Higgsinos and $h_{u,d}$ the Higgs doublets. We have assumed the μ term, in addition to B_μ, is generated by the hidden dynamics. The actual value of $r_H = r_{H_u} + r_{H_d}$ thus depends on the charge assignments in the hidden sector. The charge assignment of the visible sector fields is modified by higher-dimensional operators in the Kahler potential like $|H_{u,d}|^2 |\mathcal{R}|^2$, etc., which lead to g/g_* suppressed effects that will be neglected in what follows.

The coupling to the MSSM Higgses proportional to B_μ induces, after electroweak symmetry breaking, a small mixing between a and the MSSM Higgs boson A [27]

$$\delta = r_H \frac{v}{f_a} \frac{s_{2\beta}}{2} \frac{1}{1 - m_a^2/m_A^2} \simeq r_H \frac{v}{f_a} \frac{s_{2\beta}}{2}. \tag{7}$$

If we assume the Yukawa interactions in the superpotential are allowed in the limit of exact $U(1)_R$ ($r_{H_u} + r_Q + r_U = 2$, etc.), the mixing δ is the only source of couplings between a and the SM fermions and we get

$$\mathcal{L}_{\rm f} = i r_H \frac{a}{f_a} [c_\beta^2 m_u \bar{u} \gamma_5 u + s_\beta^2 m_d \bar{d} \gamma_5 d + s_\beta^2 m_\ell \bar{\ell} \gamma_5 \ell]. \tag{8}$$

The same mixing induces

$$\mathcal{L}_{ahh} = \frac{\delta^2}{v} h (\partial_\mu a)^2, \tag{9}$$

where h is the SM-like Higgs, as well as extra interactions with the MSSM Higgses, whose phenomenological consequences we leave for future work [14]. Finally, the a couplings to sfermions also arise from its mixing with A and are proportional to the A terms. Since we assume all the sfermions to be heavy and the A terms to be small, these couplings do not play any role in the R-axion phenomenology discussed here.

Phenomenology.—We focus on R-axion masses in the range between 2 GeV and 2 TeV, and we refer to [28] for a LHC study for masses of $O(100)$ MeV. For definiteness we fix $r_H = 2$, which allows for an R-symmetric μ term and a B_μ term from spontaneous $U(1)_R$ breaking. We will comment on the phenomenological differences of the $r_H = 0$ case, where the role of μ and B_μ is reversed. The Majorana gaugino masses cannot be arbitrarily larger than the scale of spontaneous R breaking, so we take $f_a \gtrsim 0.3$ TeV and fix for illustrative purposes the gaugino masses to the grand unified theory (GUT) universal values $m_{\lambda_{1,2,3}} = 0.7$, 1.4, 3.6 TeV (different values do not change the R-axion phenomenology as long as $m_{\lambda_i} > m_a/2$). For $f_a \lesssim 1$ TeV, obtaining such heavy gauginos present model building challenges which are beyond the scope of this Letter.

We now discuss the different production modes of the R-axion. For the purposes of this Letter we ignore R-axion production from SUSY decay chains. As for any axionlike particle, the single production modes scale with $1/f_a^2$ and double production ones with $1/f_a^4$.

(i) At the LHC, the resonant a (+SM) production is dominated by gluon fusion, which we determine using the leading-order prediction multiplied by a constant K factor of 2.4 [29] (see the Supplemental Material for details [10]).

(ii) Also at the LHC, we have double a production from Higgs decays driven by the coupling in Eq. (9).

(iii) At lepton colliders the R-axion can be single produced via its coupling to the Z, which is dominated by the anomaly in Eq. (5).

(iv) At flavor factories we consider R-axion production via $B \to K^{(*)} a$ and $\Upsilon \to \gamma a$ decays. The ${\rm BR}(B \to K^{(*)} a)$ are computed from the general result of [30], accounting for the mixing of the R-axion with the CP-odd Higgs (7) [31], and choosing for reference $m_{H^\pm} = 1$ TeV (we take the form factor relevant for K^* from [32]). The ${\rm BR}(\Upsilon \to \gamma a)/{\rm BR}(\Upsilon \to ll)$ is computed using the standard Wilczek formula [33].

The R-axion decays to pairs of gauge bosons and of SM fermions, to Goldstini (invisible) and, if kinematically allowed, to supersymmetric particles. We refer the reader to the Supplemental Material for quantitative details [10].

In Fig. 1 we summarize the present constraints on the R-axion in the m_a-f_a plane as well as the most promising processes to search for it at future experiments. For $m_a \gtrsim m_h/2$, the most important bounds come from resonant a production at the LHC. The most distinctive feature of the R-axion is the invisible signal strength of Eq. (2), which is important for large m_a. This results in constraints from monojet searches at 8 and 13 TeV [34,35]. To draw them, we have determined the ratio of the a and a + jet(s) production cross sections via a MADGRAPH [56,57] simulation, for the different missing energy cuts for which the bounds are given in [34,35] (we believe this approximation to be sufficient for our purposes).

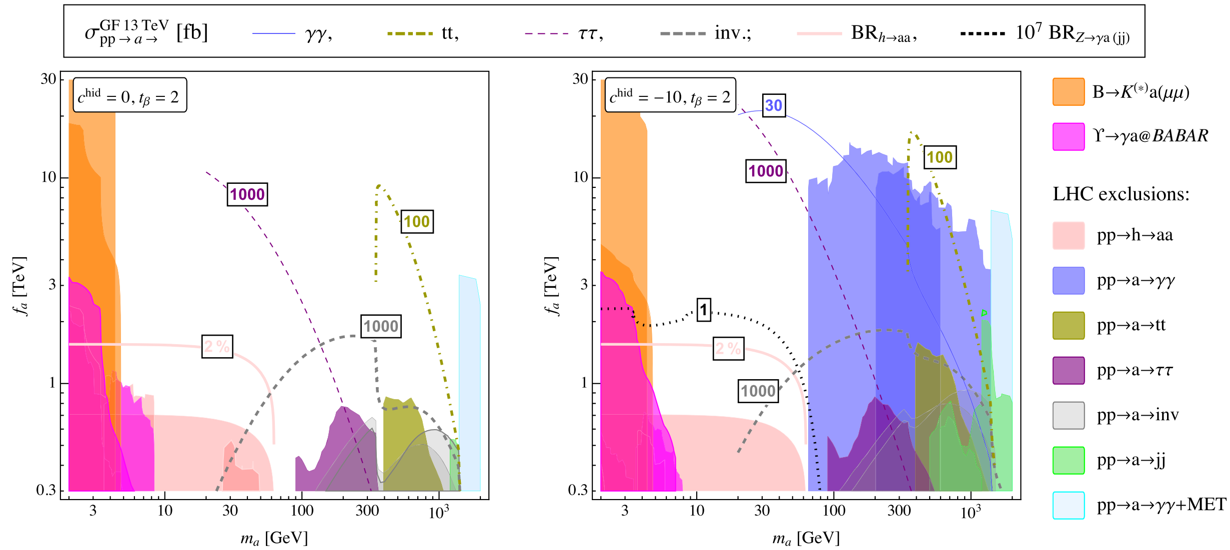

FIG. 1. Shaded areas show LHC8 and LHC13 [34–49], LHCb [50,51], *BABAR* [52–54], and Belle [55] exclusions. Contours show signal strengths at the LHC13 and Higgs and Z boson branching ratios. The tiny area in the lower right corner where $m_a \approx 4\pi f_a$ lies beyond the regime of validity of our effective description.

At $m_a > 1.4$ TeV the decay into bino pairs opens up, resulting in a $\gamma\gamma$ + MET final state via the prompt bino decay to $\gamma + G$. This is constrained by inclusive $\gamma\gamma$ + MET searches [36], which we translate in a bound $\sigma^{8\mathrm{TeV}}_{gg\to a\to\lambda_1\lambda_1} < 0.3$ fb.

For large anomalies, diphoton resonant searches [37–39] at 8 and 13 TeV dominate the collider phenomenology for $m_a > 60$ GeV [58]. Searches for a resonance decaying into dijet [40–42], into $t\bar{t}$ at 8 TeV [43,44], and into ditau at 13 TeV [45,46] give complementary bounds for $m_a > 500$ GeV, $m_a > 2m_t$, and 100 GeV $\lesssim m_a < 2m_t$, respectively.

For small anomalies the LHC constraints are sensibly weakened by the reduced production cross section and by the suppressed branching ratio in diphotons. A lower bound on f_a can still be derived from a combination of $t\bar{t}$, ditau, and monojet searches.

For $m_a \lesssim 2m_h$, the major constraint comes from the upper bound on $\mathrm{BR}(h \to \text{untagged}) < 32\%$ [47]. This bound depends only on the mixing in Eq. (7) and applies to both cases with $N_{\text{mess}} = 10$ and $N_{\text{mess}} = 0$. We also include constraints arising from exclusive Higgs decays (e.g., $h \to aa \to 4\mu$ or $h \to aa \to 2\tau 2\mu$) [48,49]. Finally, LEP constraints [59–64] are not relevant for the f_a considered in this study.

For $m_a \lesssim 9$ GeV, stringent constraints on f_a come from *BABAR* searches of $\Upsilon \to a\gamma$, with a decaying into tau or muon pairs [52,53] or hadrons [54]. For $m_a \lesssim 4$ GeV, we consider LHCb [50,51] and Belle [55] bounds from $B \to K^*\mu\mu$. In particular, the recent LHCb limit $\mathrm{BR}(B \to K^*\mu\mu) \lesssim 2\times 10^{-9}$ [50] puts the strongest bound on f_a, among all the searches we considered, reaching $f_a \gtrsim 200$ TeV for $N_{\text{mess}} = 0$.

At large t_β, δ becomes smaller [see Eq. (7)], reducing the bounds from Higgs branching ratio measurements and $B \to K$ transitions. The couplings to quarks and leptons are also t_β dependent; most importantly, the $t\bar{t}$ signal strength is reduced at large t_β. For $r_H = 0$ (and within our previous assumptions) the R-axion does not couple to SM fermions at the linear level. This makes it generically very difficult to be constrained for $m_a < m_h/2$. For larger values of m_a and irrespective of the values of t_β and r_H, diphoton constraints give $f_a \gtrsim 10$ TeV for large anomalies, while for small anomalies a milder bound on f_a is anyway given by monojet and dijet searches.

Promising signatures for future experimental programs are also shown in Fig. 1. The most distinctive one is the decay into two Goldstini, which gives a large invisible signal strength. This will be probed by monojet searches at the LHC (multijet + MET searches could also be relevant [65]) and constitutes a very good motivation for the high-luminosity LHC program. For large anomalies, the diphoton final state will be the most promising one at the LHC, while for small anomalies ditaus and $t\bar{t}$ will be more important. For $m_a < m_h/2$ we show how an improvement of the Higgs coupling measurements down to 1%–2% (which is within the reach of ILC [66]) would probe f_a up to 1.5 TeV. Even bigger values of f_a are within the reach of machines like CLIC, CEPC, and FCC-ee, which plan to probe Higgs coupling with a precision of roughly 10^{-3} [66]. For large anomalies, an important probe of a light R-axion would be $Z \to \gamma a$ measurements at future lepton colliders. A naive rescaling of the LEP I analysis [62] indicates that $Z \to \gamma a(jj)$ branching ratios in the ballpark of 10^{-7} could be probed at the FCC-ee, if $O(10^{12})$ Zs will be produced. We notice that the mass window $10 \lesssim m_a \lesssim 65$ GeV is less constrained by the searches we considered. This could be improved by extending the coverage of resonance searches, in particular $\gamma\gamma$, to lower invariant masses.

To distinguish the R-axion from other scalar resonances, a jet + MET signal would certainly help in combination with a pattern along the lines discussed above. Of course, to reinforce the R-axion interpretation of a possible signal, one would eventually need to find evidence for superpartners.

Conclusions.—The possibility that the R-axion could be the first sign of SUSY at colliders is well motivated from theoretical and phenomenological considerations, *a fortiori* given the strong LHC bounds on sparticles.

In this Letter we have investigated the low-energy dynamics of SUSY-breaking sectors with a light R-axion (and gravitino) coupled to the MSSM. Our results are summarized in Fig. 1, where we show how current and future colliders probe the space of R-axion masses and decay constants. We have also identified some promising signatures to cover the currently unconstrained part of the parameter space.

The R-axion constitutes a very interesting prototype of axionlike particles, with couplings that follow from well-defined selection rules of the theory, and whose mass can be safely considered a free parameter.

The rich phenomenology of the R-axion certainly deserves further investigation. The R-axion can give rise to nonstandard heavy Higgs decays or SUSY decay chains, it can be a further motivation for high-intensity experiments (in its light-mass window), and it could impact cosmological and astrophysical processes.

Finally, we wish to point out that other appealing features of SUSY, such as unification and dark matter, might find an interesting interplay with a light R-axion, opening new model-building avenues. We leave the exploration of this exciting physics for the future.

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