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Pure Gravity Mediation and Spontaneous $B{-}L$ Breaking from Strong Dynamics

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Abstract

In pure gravity mediation (PGM), the most minimal scheme for the mediation of supersymmetry (SUSY) breaking to the visible sector, soft masses for the standard model gauginos are generated at one loop rather than via direct couplings to the SUSY-breaking field. In any concrete implementation of PGM, the SUSY-breaking field is therefore required to carry nonzero charge under some global or local symmetry. As we point out in this note, a prime candidate for such a symmetry might be $B{-}L$, the Abelian gauge symmetry associated with the difference between baryon number B and lepton number L. The F-term of the SUSY-breaking field then not only breaks SUSY, but also $B{-}L$, which relates the respective spontaneous breaking of SUSY and $B{-}L$ at a fundamental level. As a particularly interesting consequence, we find that the heavy Majorana neutrino mass scale ends up being tied to the gravitino mass, $\Lambda_N \sim m_{3/2}$. We illustrate our idea by means of a minimal model of dynamical SUSY breaking, in which $B{-}L$ is identified as a weakly gauged flavor symmetry. We also discuss the effect of the $B{-}L$ gauge dynamics on the superparticle mass spectrum as well as the resulting constraints on the parameter space of our model. In particular, we comment on the role of the $B{-}L$ D-term.

Contents

1 Introduction: SUSY and $B-L$ breaking by the same chiral field

Pure gravity mediation (PGM) [1,2] is an attractive, viable and minimal scheme for the mediation of supersymmetry (SUSY) breaking to the visible sector.[1] The main idea behind this mediation scheme is that, given a rather high SUSY breaking scale of $\mathcal{O}\left(10^{11}\cdots10^{12}\right)$ GeV, soft SUSY breaking in the minimal supersymmetric standard model (MSSM) can be solely achieved by means of gravitational interactions. In PGM, squarks and sleptons receive large masses of the order of the gravitino mass, $m_{3/2} \sim 100\cdots1000\,\text{TeV}$, via the tree-level scalar potential in supergravity (SUGRA) [5]. Meanwhile, gauginos obtain one loop-suppressed masses around the TeV scale via anomaly mediation (AMSB) [6]. Because of the large sfermion mass scale, PGM easily accounts for a standard model (SM) Higgs boson mass of 126 GeV [7], while, at the same time, it is free of several notorious problems that other, low-scale realizations of gravity mediation are usually plagued with (such as the cosmological gravitino problem [8] or the SUSY flavor problem [9]).

In particular, PGM does not suffer from the cosmological Polonyi problem [10], which one typically encounters in ordinary gravity mediation. There, the SUSY-breaking (or "Polonyi") field X couples directly to the chiral field strength superfields belonging to the SM gauge interactions,

$$W \supset \frac{X}{M_{\text{Pl}}}\mathcal{W}^\alpha\mathcal{W}_\alpha\,, \tag{1}$$

with $M_{\text{Pl}} = (8\pi G)^{-1/2} \simeq 2.44\times10^{18}\,\text{GeV}$ denoting the reduced Planck mass and which results in gaugino masses of $\mathcal{O}\left(m_{3/2}\right)$. To be able to write down such couplings in the superpotential, one has to require that the field X be completely neutral. This, however, potentially leads to severe problems in the context of cosmology. Given a completely uncharged field X, the origin $X = 0$ does not have any special meaning in field space, which is why X is expected to acquire some vacuum expectation value (VEV) of $\mathcal{O}\left(M_{\text{Pl}}\right)$ during inflation. In this case, a huge amount of energy ends up being stored in the coherent oscillations of the Polonyi field after inflation. Once released in the perturbative decay of the Polonyi field at late times, this energy then results in dangerous entropy production as well as unacceptably large changes to the predictions of big bang nucleosynthesis. A number of solutions to this infamous Polonyi problem have been put forward

[1] For closely related mediation schemes, see [3,4].

over the years in the context of ordinary gravity mediation (see, e.g., [11]). At the same time, PGM resolves the Polonyi problem in the arguably simplest way, i.e., by requiring that there are no such couplings as in Eq. (1). More precisely, in PGM, SUSY is required to be broken by a *non-singlet field*, so that the SM gauginos are massless at tree level. This serves the purpose to lower the gaugino masses relative to the sfermion masses by a loop factor down to the TeV scale, so that the wino may eventually correspond to the lightest supersymmetric particle (LSP) and provide a viable candidate for dark matter (DM) in the form of weakly interacting massive particles (WIMPs) [2, 4, 12].

A crucial question, which needs to be addressed in any implementation of PGM, then is: Under which symmetry could the SUSY-breaking field X be possibly charged? Interesting candidates for such a symmetry are, e.g., a discrete R symmetry or a local $U(1)$ symmetry. Two of us have recently studied the former scenario in more detail in [13], which is why we will not pay any further attention to the possibility of a discrete R symmetry in the following. Instead, in this note, we shall focus on the possibility of a $U(1)$ symmetry being responsible for vanishing gaugino masses at tree level. A prime candidate for such a protective $U(1)$ symmetry is $B-L$, the Abelian gauge symmetry associated with the difference between baryon number B and lepton number L. This symmetry is essential to the seesaw mechanism [14] and may explain the origin of matter parity in the MSSM [15]. In addition, it may also play an important role in the early universe during the stages of reheating and leptogenesis (see, e.g., [16]). Furthermore, supposing that the field X is indeed charged under $B-L$, the auxiliary field F_X also needs to carry nonzero $B-L$ charge. In the SUSY-breaking vacuum at low energies, where $\langle |F_X| \rangle \neq 0$, the F-term of the SUSY-breaking field X therefore not only breaks SUSY, but also $B-L$. Assuming, within the framework of PGM, that the gaugino mass term in Eq. (1) is indeed forbidden by virtue of a local $B-L$ symmetry, thus, establishes a link between the spontaneous breaking of SUSY and the spontaneous breaking of $B-L$ at an elementary level. As we shall argue in this paper, this has several interesting phenomenological implications; most importantly, a direct connection between the heavy neutrino mass scale Λ_N in the seesaw extension of the MSSM and the gravitino mass $m_{3/2}$,

$$\Lambda_N \sim m_{3/2} \sim 100 \cdots 1000\,\text{TeV}\,. \tag{2}$$

Note that this relation nicely embodies the connection between the spontaneous breakings of $B-L$ and SUSY in our model, which is why it may be regarded as the hallmark signature of our scenario. As a consequence, the heavy Majorana neutrinos in the MSSM end up being much lighter than usually expected according to, e.g., the standard embedding of the seesaw mechanism into grand unified theories (GUTs). Our scenario is, hence, inconsistent with the notion of standard thermal leptogenesis (featuring a hierarchical heavy neutrino mass spectrum) [17][2] and, instead, requires

[2]Standard thermal leptogenesis requires the lightest sterile neutrino to have a mass of at least $M_{N_1} \sim 10^9$ GeV [18]. Besides that, simple alternatives to the paradigm of thermal leptogenesis may easily involve heavy Majorana neutrinos with masses almost as large as the scale of grand unification, $M_{N_1} \sim 10^{15} \cdots 10^{16}$ GeV; see [19] for a recent example.

some form of low-scale leptogenesis, such as resonant leptogenesis [20] (where the heavy neutrino masses are highly degenerate), in order to account for the baryon asymmetry of the universe.

The purpose of the present paper now is to illustrate our idea by means of a minimal example. More concretely, we shall demonstrate how to embed the spontaneous breaking of $B-L$ into one of the simplest models of dynamical SUSY breaking (DSB), i.e., the simplest realization of the vector-like DSB model à la IYIT [21], which is based on strongly coupled $Sp(1) \cong SU(2)$ gauge dynamics in combination with four fundamental matter fields. Here, following up on earlier work presented in [22], we shall identify $B-L$ as a weakly gauged flavor symmetry of the IYIT model (see Sec. 2). Next to the anticipated link between the spontaneous breaking of SUSY and $B-L$ and the prediction for the heavy neutrino mass scale in Eq. (2), this then provides us with important (partly tachyonic) corrections to the MSSM sfermion masses. These mass corrections consist, for one thing, of tree-level sfermion masses induced by the $B-L$ D-term and, for another thing, of effective sfermion masses induced by gauge mediation at the one-loop level [23] (see Sec. 3). Both corrections need to be sufficiently suppressed in order to ensure the stability of the low-energy vacuum. Fortunately, as we shall discuss in more detail in Sec. 3, the suppression of the $B-L$ D-term contributions to the MSSM sfermion masses turns out to be parametrically well controlled, thanks to the fact that we are able to derive an *explicit* expression for the $B-L$ D-term in terms of the underlying model parameters. In fact, owing to this calculability of the $B-L$ D-term, we are capable of tuning its magnitude to an arbitrarily small value by imposing an approximate flavor symmetry in the IYIT sector. Our set-up therefore features an interesting mechanism to maintain control over the $B-L$ D-term, which might otherwise spoil large parts of our construction.[3] Meanwhile, we find that the suppression of the gauge-mediated sfermion masses imposes an upper bound on the $B-L$ gauge coupling constant, $g \lesssim 10^{-3}$, which renders our model testable/falsifiable in a future multi-TeV collider experiment. Finally, in Sec. 4, we are going to conclude, giving a brief outlook as to how our study could possibly be continued.

2 Embedding $B-L$ into the IYIT SUSY breaking model

In its most general formulation, the IYIT model of dynamical SUSY breaking is based on a strongly coupled $Sp(N)$ gauge theory featuring $2N_f = 2(N+1)$ "quark" fields Ψ^i that transform in the fundamental representation of $Sp(N)$. At energies below the dynamical scale Λ, this theory is best described in terms of $(2N+1)(N+1)$ gauge-invariant composite "meson" operators $M^{ij} \simeq \left\langle \Psi^i \Psi^j \right\rangle / (\eta\Lambda)$, which are subject to a quantum mechanically deformed moduli constraint [24],

$$\mathrm{Pf}\left(M^{ij}\right) \simeq \left(\frac{\Lambda}{\eta}\right)^{N+1} , \quad \eta \sim 4\pi \, . \tag{3}$$

[3] We believe that the applicability of this technical result may extend well beyond the purposes of the present paper, which may make it also interesting from a more general perspective, i.e., if one is more interested in the general business of gauging global flavor symmetries of strongly coupled DSB models and perhaps less interested in the concrete phenomenology of a weakly gauged $B-L$ symmetry in the context of the IYIT model.

Here, $\text{Pf}\left(M^{ij}\right)$ denotes the Pfaffian of the antisymmetric meson matrix M^{ij} and η is a numerical factor that may be estimated based on naive dimensional analysis (NDA) [25]. In order to break SUSY in the IYIT model, one introduces Yukawa couplings between the quark fields Ψ^i, the fundamental degrees of freedom (DOFs) at energies above the dynamical scale Λ, and a set of $(2N+1)(N+1)$ singlet fields Z_{ij} in the tree-level superpotential,

$$W_{\text{tree}}^{\text{IYIT}} = \frac{1}{2}\lambda'_{ij}\, Z_{ij}\, \Psi^i \Psi^j \,. \tag{4}$$

In the effective theory at energies below the dynamical scale Λ, this gives rise to an effective superpotential for the meson fields M^{ij}, which lifts all flat direction in moduli space,

$$W_{\text{eff}}^{\text{IYIT}} \simeq \frac{1}{2}\lambda_{ij}\, \frac{\Lambda}{\eta}\, Z_{ij}\, M^{ij} \,. \tag{5}$$

This superpotential implies F-term conditions for the singlet fields, $M^{ij} = 0$, which cannot be satisfied while simultaneously fulfilling the moduli constraint in Eq. (3), $\text{Pf}\left(M^{ij}\right) \neq 0$. In the true vacuum of the IYIT model, SUSY is hence spontaneously broken because some of the singlet fields' F-terms are nonzero, i.e., SUSY is broken via the O'Raifeartaigh mechanism [26].

For all Yukawa couplings in Eq. (4) being equal, $\lambda'_{ij} \equiv \lambda$, the IYIT tree-level superpotential exhibits a global $SU(4) \times Z_4$ flavor symmetry.[4] Allowing for generic, numerically different Yukawa couplings, this symmetry is, however, broken down to an Abelian $U(1)_A \times Z_4$ flavor symmetry,

$$\lambda'_{ij} \text{ all different} \quad \Rightarrow \quad SU(4) \times Z_4 \rightarrow U(1)_A \times Z_4 \,, \tag{6}$$

with the axial $U(1)_A \subset SU(4)$ being associated with a global quark field rotation, $\Psi^i \rightarrow e^{iq\theta}\Psi^i$. In [22], this global $U(1)_A$ flavor symmetry has been promoted to a weakly gauged Fayet-Iliopoulos (FI) symmetry, $U(1)_A \rightarrow U(1)_{\text{FI}}$, in order to demonstrate how to generate a theoretically consistent and field-dependent FI D-term in the context of dynamical SUSY breaking. The advantage of such dynamically generated FI-terms is that they do not suffer from the usual problems that other FI models are plagued with [27–29]. Once coupled to SUGRA, constant, field-independent FI-terms, e.g., always require the presence of an *exact* continuous global symmetry [27], which is problematic from the perspective of quantum gravity [30]. On the other hand, field-dependent FI-terms in string theory [31], generated via the Green-Schwarz mechanism of anomaly cancellation [32], imply the existence of a shift-symmetric modulus field [29], which causes cosmological problems [10], as long as it is not properly stabilized (which is hard [33]). As shown in [22], dynamically generated and field-dependent FI-terms in field theory, by contrast, avoid all of these problems, rendering them the arguably best candidates for FI-terms with relevant implications for low-energy phenomenology.

In this paper, we shall now take the analysis of [22] one step further and promote the global $U(1)_A$ flavor symmetry of the IYIT model to a local $U(1)_{B-L}$ symmetry. For simplicity, we will

[4]Here, the discrete Z_4 symmetry corresponds to a phase shift of all quark fields by $\pi/2$, i.e., all quarks transform as $\Psi^i \rightarrow i\Psi^i$ under this Z_4 symmetry. In fact, this Z_4 is nothing but the anomaly-free subgroup of the anomalous $U(1)'$ symmetry that is contained in the full $U(4)$ flavor symmetry at the classical level, $U(4) \cong SU(4)\times U(1)' \supset SU(4)\times Z_4$.

restrict ourselves to the IYIT model in its simplest version from now on. That is, we will focus on the $Sp(1) \cong SU(2)$ case in combination with four quark flavors. If we assign $B-L$ charges $\pm q/2$ to these quark fields, the six meson fields at low energies end up carrying the following charges,

$$[M_+] = +q\,, \quad [M_-] = -q\,, \quad [M_0^a] = 0\,, \quad a = 1,2,3,4\,, \tag{7}$$

and similarly for the six singlet fields $Z_\pm$ and Z_0^a, which we also re-label according to their $B-L$ charges. The effective superpotential as well as the effective Kähler potential for these fields read[5]

$$W_{\text{eff}} \simeq \frac{\Lambda}{\eta}\left(\lambda_+\, M_+\, Z_- + \lambda_-\, M_-\, Z_+ + \lambda_0^a\, M_0^a\, Z_0^a\right)\,, \tag{8}$$

$$K_{\text{eff}} \simeq M_+^\dagger e^{2qgV} M_+ + M_-^\dagger e^{-2qgV} M_- + Z_+^\dagger e^{2qgV} Z_+ + Z_-^\dagger e^{-2qgV} Z_- + \sum_a |M_0^a|^2 + \sum_a |Z_0^a|^2\,. \tag{9}$$

Here, the vector field V stands for the $B-L$ vector multiplet, the auxiliary D component of which gives rise to the following D-term scalar potential,

$$V_D = \frac{1}{2} D^2 = \frac{q^2 g^2}{2}\left[|M_-|^2 - |M_+|^2 + |Z_-|^2 - |Z_+|^2\right]^2\,. \tag{10}$$

After imposing the quantum mechanically deformed moduli constraint in Eq. (3),

$$\text{Pf}\left(M^{ij}\right) = M_+\, M_- - M_0^1\, M_0^4 + M_0^2\, M_0^3 \simeq \left(\frac{\Lambda}{\eta}\right)^2\,, \tag{11}$$

one finds that the vacuum manifold of the low-energy theory exhibits exactly three local minima. In the limit of a vanishingly small gauge coupling constant g, these are respectively located at

$$\begin{aligned}
&\text{Vacuum I:} && \left<M_+ M_-\right> \simeq \left(\frac{\Lambda}{\eta}\right)^2, && \left<|M_+|\right> = \frac{\sqrt{\lambda_+\lambda_-}}{\lambda_+}\frac{\Lambda}{\eta}, && \left<|M_-|\right> = \frac{\sqrt{\lambda_+\lambda_-}}{\lambda_-}\frac{\Lambda}{\eta}, \\
&\text{Vacuum II:} && -\left<M_0^1 M_0^4\right> \simeq \left(\frac{\Lambda}{\eta}\right)^2, && \left<|M_0^1|\right> = \frac{\sqrt{\lambda_0^1\lambda_0^4}}{\lambda_0^1}\frac{\Lambda}{\eta}, && \left<|M_0^4|\right> = \frac{\sqrt{\lambda_0^1\lambda_0^4}}{\lambda_0^4}\frac{\Lambda}{\eta}, \\
&\text{Vacuum III:} && \left<M_0^2 M_0^3\right> \simeq \left(\frac{\Lambda}{\eta}\right)^2, && \left<|M_0^2|\right> = \frac{\sqrt{\lambda_0^2\lambda_0^3}}{\lambda_0^2}\frac{\Lambda}{\eta}, && \left<|M_0^3|\right> = \frac{\sqrt{\lambda_0^2\lambda_0^3}}{\lambda_0^3}\frac{\Lambda}{\eta},
\end{aligned} \tag{12}$$

with all other meson and singlet VEVs vanishing, respectively. The vacuum energies in these three vacua respectively scale with the geometric means of the corresponding pairs of Yukawa couplings,

$$V_{\text{I}} = 2\,\lambda_+\lambda_-\left(\frac{\Lambda}{\eta}\right)^4, \quad V_{\text{II}} = 2\,\lambda_0^1\lambda_0^4\left(\frac{\Lambda}{\eta}\right)^4, \quad V_{\text{III}} = 2\,\lambda_0^2\lambda_0^3\left(\frac{\Lambda}{\eta}\right)^4. \tag{13}$$

For $\lambda_+\lambda_- < \min\left\{\lambda_0^1\lambda_0^4, \lambda_0^2\lambda_0^3\right\}$, the lowest lying vacuum therefore corresponds to the one where $\left<M_+M_-\right> \simeq (\Lambda/\eta)^2$, i.e., the one in which $B-L$ is spontaneously broken by the nonvanishing VEVs

[5] Throughout the analysis in this paper, we will take the Kähler potential to be canonical for all fields and neglect all effects induced by higher-dimensional terms in the effective Kähler potential. These terms are uncalculable and, in principle, always present in the IYIT model. On the other hand, they are suppressed compared to the canonical terms in the Kähler potential by factors of $\mathcal{O}\left(\lambda^2/\eta^2\right)$ [34], which is why we can safely ignore them, as long as we stay in the perturbative regime, $\lambda \ll \eta$, and do not venture into the strongly coupled limit, where $\lambda \sim \eta$.

of the charged meson fields $M_\pm$. In the following, we shall assume that this condition is satisfied, so that in the low-energy vacuum of the IYIT model $B-L$ is indeed spontaneously broken.

In view of this result, two comments are in order: (i) First of all, we remark that it is actually an open question whether the deformed moduli constraint as stated in Eq. (3) really ends up being fulfilled *exactly* in the IYIT model or whether $\mathrm{Pf}\left(M^{ij}\right)$ could, in fact, also display a significant deviation from $(\Lambda/\eta)^2$ in the true vacuum. In the former case, only some of the singlet fields Z_{ij} acquire nonzero F-terms, while in the latter case also the $Sp(N)$ glueball field $T \propto \langle gg \rangle$ turns out to contribute to SUSY breaking with a nonzero F-term (see [13, 22] for an extended discussion of this issue). Our results will not be qualitatively affected by the choice between these two options, which is why, in this paper, we decide to neglect the possibility of a dynamical glueball field and work with $\mathrm{Pf}\left(M^{ij}\right) \equiv (\Lambda/\eta)^2$ for simplicity in the following. (ii) Our results in Eqs. (12) and (13) only hold in the weakly gauged limit, $g \rightarrow 0$. Once we turn on the $B-L$ gauge interactions, the vacuum manifold of the IYIT model becomes distorted. That is, while the loci of vacua II and III remain unchanged, vacuum I begins to shift in the $M_\pm$ plane, as soon as the coupling constant g is allowed to take a small, but nonzero value. More precisely, for small g, we find

$$\left<\left|M_\pm\right|\right> = \frac{\lambda}{\lambda_\pm}\frac{\Lambda}{\eta}\left[1 \pm \frac{\gamma^2}{\rho^4}\left(1-\rho^4\right)^{1/2} + \mathcal{O}\left(\gamma^4\right)\right] , \tag{14}$$

where we have introduced λ, ρ and γ as important combinations of the parameters $\lambda_\pm$ and g,

$$\lambda = \sqrt{\lambda_+\lambda_-} , \quad \rho = \left[\frac{1}{2}\left(\frac{\lambda_+}{\lambda_-} + \frac{\lambda_-}{\lambda_+}\right)\right]^{-1/2} , \quad \gamma = \frac{qg}{\lambda} . \tag{15}$$

Here, λ denotes the geometric mean of λ_+ and λ_-, the parameter $\rho \in [0,1]$ is a convenient measure for the amount of flavor symmetry violation in the charged meson sector,[6] and γ characterizes the strength of the $B-L$ gauge interactions relative to the strength of the IYIT Yukawa interactions. Eq. (14) illustrates that, while vacuum I always remains on the $M_+M_- = (\Lambda/\eta)^2$ hypersurface, its "flavor composition" in terms of M_+ and M_- begins to change in consequence of the $B-L$ gauge interactions, once the coupling strength g takes larger and larger values. A more detailed investigation of these next-to-leading order effects in the gauge coupling constant g is left for future work (especially a study of the dynamics in the large-g regime, where $\gamma \gg 1$). In this paper, we will, by contrast, content ourselves with a leading-order analysis, meaning that wherever possible we will simply neglect all effects of $\mathcal{O}(g)$.

So far, we have identified the condition under which the low-energy vacuum of the IYIT model not only breaks SUSY, but also $B-L$. Next, let us discuss the properties of this vacuum in a bit more detail. In doing so, we will mostly review some earlier results presented in [22], which is why we will be rather brief in what follows. The physical mass eigenstates at low energies are contained

[6] Note that, for equal Yukawa couplings, $\lambda_+ = \lambda_-$, the parameter ρ goes to $\rho = 1$, while, for drastically different Yukawa couplings in the charged meson sector, $\lambda_+ \ll \lambda_-$ or $\lambda_- \ll \lambda_+$, it approaches $\rho = 0$. Moreover, ρ^2 can also be interpreted as the ratio between the harmonic and geometric means of λ_+^2 and λ_-^2 (see [13] for details).

in the following two linear combinations of the singlet fields Z_+ and Z_-,

$$X = \frac{1}{\sqrt{2}} \left(Z_+ + Z_-\right) \,, \quad Y = \frac{1}{\sqrt{2}} \left(Z_+ - Z_-\right) \,, \tag{16}$$

the $B-L$ vector multiplet V as well as in the goldstone multiplet A of spontaneous $B-L$ breaking,[7]

$$M_\pm = \left<|M_\pm|\right> e^{\pm A/f_A} \,, \quad f_A = K_0^{1/2} \,, \quad K_0 = \left< |M_+|^2 \right> + \left< |M_-|^2 \right> . \tag{17}$$

Here, the decay constant f_A ensures the correct normalization of the goldstone field A and K_0 represents the VEV of the Kähler potential in global SUSY. While the actual goldstone phase $a \in A$ remains massless and is absorbed by the $B-L$ vector field $A_\mu \in V$ upon spontaneous $B-L$ breaking, all other DOFs contained in A obtain soft SUSY-breaking masses via the IYIT superpotential. The goldstone field A, hence, vanishes in the true vacuum, $\left<A\right> = 0$, which allows us to expand the effective superpotential for X, Y and A in powers of A. Up to $\mathcal{O}\left(A^2\right)$, we have

$$W_{\text{eff}} \simeq \mu^2 X - m\, Y A + \frac{m^2}{2\,\mu^2}\, X A^2 \,, \tag{18}$$

where μ and m denote the F-term SUSY breaking scale as well as the soft SUSY-breaking mass resulting from the IYIT superpotential, respectively, (see Eq. (15) for the definitions of λ and ρ)

$$\mu = 2^{1/4} \lambda^{1/2}\, \frac{\Lambda}{\eta} \,, \quad m = \frac{\mu^2}{f_A} = \rho\, \lambda\, \frac{\Lambda}{\eta} \,. \tag{19}$$

Correspondingly, the gravitino mass $m_{3/2}$ needs to take the following value in our set-up,

$$m_{3/2} = \frac{\mu^2}{\sqrt{3}\, M_{\text{Pl}}} = \lambda \left(\frac{2}{3}\right)^{1/2} \frac{\left(\Lambda/\eta\right)^2}{M_{\text{Pl}}} \,, \tag{20}$$

in order to ensure that the cosmological constant (almost) vanishes in the low-energy vacuum. Requiring the gravitino mass to take a certain value, say, $m_{3/2} = 1000\,\text{TeV}$, thus allows us to eliminate either λ or the dynamical scale Λ/η from our analysis. We opt for the latter, so that

$$\frac{\Lambda}{\eta} \simeq 1.7 \times 10^{12}\,\text{GeV} \left(\frac{1}{\lambda}\right)^{1/2} \left(\frac{m_{3/2}}{1000\,\text{TeV}}\right)^{1/2} . \tag{21}$$

As evident from Eq. (18), X corresponds to the SUSY-breaking goldstino field, while Y represents the Dirac mass partner of the $B-L$ goldstone field A. In terms of the charged meson fields $M_\pm$, the F component of the goldstino field X is given as (see Eqs. (8) and (16))

$$-F_X^* = \frac{1}{\sqrt{2}} \left(\lambda_+ M_+ + \lambda_- M_-\right) \frac{\Lambda}{\eta} \,, \tag{22}$$

which acquires a VEV $\left<|F_X|\right> = \mu^2$ in the true vacuum. Since F_X does *not* transform as a singlet under $B-L$, its nonzero VEV not only breaks SUSY, but also $B-L$. We emphasize that this is

[7]In [13] (making use of some earlier results presented in [35]), the $U(1)_A$ flavor symmetry of the IYIT model has been identified with the global Peccei-Quinn (PQ) symmetry appearing in the axion solution to the strong CP problem, $U(1)_A \to U(1)_{\text{PQ}}$, rather than with a local $B-L$ symmetry. In this case, the field A then turns out to correspond to the chiral axion superfield in a supersymmetric version of the KSVZ axion model [36].

one of the key features of the set-up considered in this paper. Furthermore, we note that X is massless at tree level (see Eq. (18)). At the classical level, the complex scalar contained in X, hence, corresponds to a flat (or modulus) direction of the scalar potential. This vacuum degeneracy is, however, lifted at the loop level [34], which renders the "sgoldstino" a pseudomodulus, after all. The effective sgoldstino mass m_X has recently been re-evaluated in [13] (see Eq. (116) therein). As it turns out, m_X ends up being a complicated function of the Yukawa couplings $\lambda_\pm$ and λ_0^a. For this reason, we will not state the full expression here, but merely restrict ourselves to the result in the flavor-symmetric limit, in which $\lambda_0^a \equiv \lambda$ for all $a = 1, 2, 3, 4$,[8]

$$m_X^2 = \frac{2\ln 2 - 1}{16\pi^2}\left(1 + \frac{4}{\rho^6}\right)\left(\frac{m}{\mu}\right)^4 m^2 \,. \tag{23}$$

Last but not least, it is instructive to examine the effective Kähler potential for the charged meson fields $M_\pm$ as a function of V and A (see Eq. (9)). Again expanding in powers of A, we find

$$K_{\text{eff}} = K_0 - 2qg\,\xi\,V_A + m_V^2 V_A^2 + \mathcal{O}\left(V_A^3\right) \,, \quad V_A = V + \frac{1}{\sqrt{2}\,m_V}\left(A + A^\dagger\right) , \tag{24}$$

where ξ denotes the $B{-}L$ FI parameter, $\xi \equiv \langle D\rangle/(qg)$, and m_V is the $B{-}L$ vector boson mass,

$$\xi = \left\langle\, |M_-|^2 \,\right\rangle - \left\langle\, |M_+|^2 \,\right\rangle , \quad m_V = \sqrt{2}qgf_A \,. \tag{25}$$

Eq. (24) nicely illustrates how the goldstone field A is eaten by the $B{-}L$ vector multiplet V upon spontaneous $B{-}L$ breaking. In terms of the parameters of our model, ξ and f_A are given as

$$\xi = \left(\frac{\lambda_+}{\lambda_-} - \frac{\lambda_-}{\lambda_+}\right)\left(\frac{\Lambda}{\eta}\right)^2 = \frac{2\left(1-\rho^4\right)^{1/2}}{\rho^2}\left(\frac{\Lambda}{\eta}\right)^2 , \quad f_A = \left(\frac{\lambda_+}{\lambda_-} + \frac{\lambda_-}{\lambda_+}\right)^{1/2}\frac{\Lambda}{\eta} = \frac{\sqrt{2}}{\rho}\frac{\Lambda}{\eta} \,. \tag{26}$$

3 Phenomenological consequences for neutrinos and sparticles

In the previous section, we have shown how the spontaneous breaking of $B{-}L$ may be accommodated in the IYIT model of dynamical SUSY breaking. Let us now study the phenomenological implications of this embedding. First of all, we note that our set-up offers an intriguing possibility to generate Majorana masses for the right-handed neutrinos in the seesaw extension of the MSSM. Suppose that the charge q of the meson fields $M_\pm$ is actually given as $q = 2$. Then, gravitational interactions at the Planck scale will result in the following operators in the effective theory (above and below the dynamical scale Λ, respectively),

$$W \supset \frac{1}{2}\frac{c_i}{M_{\text{Pl}}}\Psi^3\Psi^4\,N_iN_i \,, \quad W_{\text{eff}} \supset \frac{1}{2}\frac{c_i}{M_{\text{Pl}}}\frac{\Lambda}{\eta}\,M_-\,N_iN_i \,, \tag{27}$$

[8] In this limit, the vacua I, II and III become degenerate (see Eq. (13)). The breaking of the $SU(4) \cong SO(6)$ flavor symmetry down to $SO(5)$ then results in five massless particles: the $B{-}L$ goldstone phase plus four genuine goldstone bosons, which may cause trouble at low energies. Therefore, in order to avoid such massless particles, the global $SU(4)$ symmetry should actually never be *exactly* realized. Instead, it should be at most realized as an *approximate* symmetry, so that all Yukawa couplings merely end up taking values *close* to each other.

where the c_i are dimensionless coefficients of $\mathcal{O}(1)$ and with the fields N_i denoting the left-handed superfields the fermionic components of which correspond to the hermitian conjugates of the right-handed neutrinos needed for the seesaw mechanism. Upon spontaneous $B-L$ breaking, these couplings then turn into Majorana mass terms for the neutrino fields N_i,

$$W \supset \frac{1}{2} M_i \, N_i N_i \,, \quad M_i = c_i \, \Lambda_N \,, \quad \Lambda_N = \frac{1}{\rho} \left[1 \pm \left(1 - \rho^4\right)^{1/2}\right]^{1/2} \frac{\left(\Lambda/\eta\right)^2}{M_{\rm Pl}} \,, \tag{28}$$

where the sign in the square brackets depends on whether λ_- is smaller (+) or larger (−) than λ_+. The heavy neutrino mass scale therefore turns out to be tied to the gravitino mass (see Eq. (20))!

$$\Lambda_N = \frac{(3/2)^{1/2}}{\rho\,\lambda} \left[1 \pm \left(1 - \rho^4\right)^{1/2}\right]^{1/2} m_{3/2} \,. \tag{29}$$

In the flavor-symmetric limit, $\rho \rightarrow 1$, we find in particular,

$$\Lambda_N \simeq 1200\,\textrm{TeV} \left(\frac{1}{\lambda}\right) \left(\frac{m_{3/2}}{1000\,\textrm{TeV}}\right) . \tag{30}$$

We emphasize that this relation between the heavy neutrino mass scale Λ_N and the gravitino mass $m_{3/2}$ is one of the most important phenomenological consequences of our model.

Next, before turning to the phenomenological implications of our model for the MSSM sparticle spectrum, we mention in passing that a coupling of the neutrino fields N_i to the singlet field Z_- would, by contrast, *not* allow for a successful generation of the heavy neutrino mass scale Λ_N. In SUGRA, the field Z_- acquires a VEV of $\mathcal{O}\left(m_{3/2}\right)$ [13, 22], which is why one might naively think that a coupling of the form $Z_- NN$ in the superpotential may also result in neutrino masses of $\mathcal{O}\left(m_{3/2}\right)$. This is, however, not so because of the large F-term of the field Z_-, which results in additional mass terms for the scalar neutrino fields of $\mathcal{O}\left(\mu\right)$. After diagonalizing the sneutrino mass matrix, one then finds that some of the sneutrinos end up being tachyonic with masses of $\mathcal{O}\left(-\mu\right)$, which renders the coupling $Z_- NN$ unfeasible. On the contrary, we actually have to make sure that the coupling $Z_- NN$ is forbidden, since it will otherwise interfere with our mechanism to generate the mass scale Λ_N. This is best done by invoking R symmetry, under which the neutrino fields carry charge 1, the meson fields charge 0 and the singlet fields charge 2 (see also [13, 22, 35]). R symmetry then allows the couplings in Eq. (27), but forbids couplings of the form $Z_- NN$.

A second important consequence of our set-up for low-energy phenomenology are tree-level as well as loop-induced corrections to the masses of the MSSM sfermions. Here, the tree-level mass corrections originate from the nonvanishing VEV of the auxiliary $B-L$ D field, $\left<D\right> = qg\,\xi$. To see this, recall that the total tree-level scalar potential in SUGRA takes the following form,

$$V = V_F + V_D = e^{K/M_{\rm Pl}^2} \left[\left(W_i + \frac{W}{M_{\rm Pl}^2} K_i\right) K^{i\bar{\jmath}} \left(\overline{W}_{\bar{\jmath}} + \frac{\overline{W}}{M_{\rm Pl}^2} K_{\bar{\jmath}}\right) - 3 \, \frac{\left|W\right|^2}{M_{\rm Pl}^2} + \frac{1}{2} \, e^{-K/M_{\rm Pl}^2} D^2 \right] , \tag{31}$$

where the indices i and $\bar{\jmath}$ refer to differentiation w.r.t. to the complex scalars ϕ_i and $\phi_{\bar{\jmath}}^*$, respectively, and where $K^{i\bar{\jmath}}$ denotes the inverse of the Kähler metric, $K^{i\bar{\jmath}} \equiv \left(K_{i\bar{\jmath}}\right)^{-1}$. The superpotential W, the Kähler potential K and the $B-L$ D-term are all nonvanishing in the true vacuum,

$$\left<W\right> \equiv W_0 \equiv e^{-K_0/M_{\rm Pl}^2/2} \, m_{3/2} \, M_{\rm Pl}^2 \,, \quad \left<K\right> \equiv K_0 \,, \quad \left<D\right> \equiv D_0 = qg\,\xi \,. \tag{32}$$

For one reason or another, these VEVs are fine-tuned such that the cosmological constant (almost) vanishes. This is to say that, in the low-energy vacuum, the total scalar potential is (almost) zero,

$$\langle V \rangle = \left\langle K^{i\bar{\jmath}} \mathcal{F}_i \mathcal{F}_{\bar{\jmath}}^* \right\rangle + \frac{1}{2} D_0^2 - 3\, e^{K_0/M_{\rm Pl}^2} \frac{|W_0|^2}{M_{\rm Pl}^2} = 0 \,, \quad \mathcal{F}_i = e^{K/M_{\rm Pl}^2/2} \left(W_i + \frac{W}{M_{\rm Pl}^2} K_i \right) . \tag{33}$$

Together, Eqs. (32) and (33) allow us to solve for the gravitino mass in terms of the total SUSY breaking scale $\Lambda_{\rm SUSY}$ (which reduces to μ in the global SUSY limit and for small g, see Eq. (20)),

$$m_{3/2}^2 = \frac{\Lambda_{\rm SUSY}^4}{3 M_{\rm Pl}^2} \,, \quad \Lambda_{\rm SUSY}^4 = F_0^2 + \frac{1}{2} D_0^2 \,, \quad F_0^2 = \left\langle K^{i\bar{\jmath}} \mathcal{F}_i \mathcal{F}_{\bar{\jmath}}^* \right\rangle . \tag{34}$$

Each MSSM sfermion $\tilde{f}$ now appears with a canonically normalized term in the Kähler potential,

$$K = K_0 + \tilde{f}^\dagger e^{2 q_f g\, V} \tilde{f} + \cdots = K_0 + \tilde{f}^\dagger \tilde{f} + \cdots . \tag{35}$$

Inserting this expansion into Eq. (31) yields the universal tree-level MSSM sfermion mass in PGM,

$$V = \exp\left(\left|\tilde{f}\right|^2 / M_{\rm Pl}^2 \right) V_0 + \left(e^{K_0/M_{\rm Pl}^2} \frac{|W_0|^2}{M_{\rm Pl}^4} - \frac{D_0^2}{2 M_{\rm Pl}^2} \right) \left|\tilde{f}\right|^2 + \cdots = V_0 + m_0^2 \left|\tilde{f}\right|^2 + \cdots \tag{36}$$

where $V_0 \equiv \langle V \rangle = 0$. Making use of the definition of the gravitino mass in Eq. (32), we find

$$m_0^2 = \frac{V_0}{M_{\rm Pl}^2} + e^{K_0/M_{\rm Pl}^2} \frac{|W_0|^2}{M_{\rm Pl}^4} - \frac{D_0^2}{2 M_{\rm Pl}^2} = m_{3/2}^2 - \frac{D_0^2}{2 M_{\rm Pl}^2} = m_{3/2}^2 + \Delta m_0^2 \,, \quad \Delta m_0^2 = - \frac{q^2 g^2 \xi^2}{2 M_{\rm Pl}^2} . \tag{37}$$

Here, the first contribution to m_0, given by the gravitino mass $m_{3/2}$, corresponds to the universal soft mass for all sfermions in PGM *in absence* of a nonzero D-term, while the second contribution to m_0 represents a universal shift in m_0 induced by the nonzero FI parameter ξ. In the context of our SUSY breaking model and assuming that $qg \sim 1$, one naively expects $D_0 \sim \xi \sim \Lambda^2$, so that

$$m_0 \sim m_{3/2} \sim \Delta m_0 \sim \frac{\Lambda}{M_{\rm Pl}} \Lambda \,. \tag{38}$$

This means that the ξ-induced shift in the soft sfermion mass, Δm_0, may, under certain circumstances, become roughly as large as the "bare" soft mass in absence of a nonzero FI term, $\Delta m_0 / m_{3/2} \sim 1$. Since the shift Δm_0 represents a tachyonic mass correction, it is, however, important that Δm_0 never exceeds $m_{3/2}$. This results in an upper bound on the ratio D_0/F_0,

$$m_0^2 = \frac{1}{3 M_{\rm Pl}^2} \left(F_0^2 - D_0^2 \right) \geq 0 \quad \Rightarrow \quad \frac{D_0}{F_0} \leq 1 \,. \tag{39}$$

Note that this bound on the magnitude of the D-term applies independently of the fact that the sfermion $\tilde{f}$ carries nonzero $B{-}L$ charge. Instead, it holds universally for any $U(1)$ symmetry that may contribute to the total vacuum energy with a nonvanishing D-term. In the context of our DSB model, the VEV of the D-term is always trivially smaller than the VEV of the IYIT F-term, at least as long as we stay in the weakly gauged regime, where $\gamma \ll 1$, (see Eqs. (19) and (26))

$$\frac{D_0}{F_0} = \gamma \left[\frac{2^{1/2}}{\rho^2} \left(1 - \rho^4 \right)^{1/2} + \mathcal{O}\left(\gamma^2\right) \right] \ll 1 \,. \tag{40}$$

Whether or not D_0 always remains smaller than F_0 also in the strongly gauged regime, i.e., for $\gamma \gg 1$, is an open question, which we leave for future work. While general SUGRA theorems suggest that this may very well be the case [37], it would still be interesting to determine the precise upper bound $D_0/F_0|_{\max}$ on the ratio D_0/F_0 in the context of the IYIT model.

Next to the universal soft mass m_0 in Eq. (37), each sfermion receives a further tree-level mass correction m_D, which depends on its respective $B{-}L$ charge q_f. Because of the interaction with the $B{-}L$ D field in the Kähler potential (see Eq. (35)), each sfermion explicitly appears in V_D,

$$V_D = \frac{q^2 g^2}{2}\left[\xi - \frac{q_f}{q}\big|\tilde{f}\big|^2 + \cdots\right]^2 = \frac{1}{2}D_0^2 + m_D^2\big|\tilde{f}\big|^2 + \cdots\,, \quad m_D^2 = -q_f g\, D_0 = -q\, q_f g^2 \xi\,, \qquad (41)$$

so that we eventually obtain for the total tree-level mass $m_{\tilde{f}}^{\rm tree}$ of an MSSM sfermion,

$$\left(m_{\tilde{f}}^{\rm tree}\right)^2 = m_0^2 + m_D^2\,, \quad m_0^2 = m_{3/2}^2 - \frac{q^2 g^2 \xi^2}{2M_{\rm Pl}^2}\,, \quad m_D^2 = -q\, q_f g^2 \xi\,. \qquad (42)$$

We hence see that sfermions $\tilde{f}$ with charge q_f such that $q_f\xi > 0$ acquire a negative mass squared as long as the FI parameter ξ is not substantially suppressed w.r.t. the dynamical scale Λ,

$$m_0 \sim \frac{\Lambda^2}{M_{\rm Pl}}\,, \quad \xi \sim \Lambda^2\,, \quad q\,\xi > 0 \quad \Rightarrow \quad \left(m_{\tilde{f}}^{\rm tree}\right)^2 \sim -\Lambda^2\left[1 + \mathcal{O}\left(\frac{\Lambda^2}{M_{\rm Pl}^2}\right)\right]. \qquad (43)$$

This poses a serious problem, which, in general, may be regarded as a fundamental obstacle to identifying any gauged $U(1)$ flavor symmetry featuring a nonzero D-term with $B{-}L$. Of course, a trivial way out of this problem is to assume an extremely small $B{-}L$ gauge coupling, $g \lesssim \Lambda/M_{\rm Pl}$, so as to suppress m_D^2 by a factor $g^2 \lesssim (\Lambda/M_{\rm Pl})^2$. In the case of PGM, where one typically has $\Lambda \sim 10^{12}\,\textrm{GeV}$ (see Eq. (21)), this would mean that g should take at most a value of $\mathcal{O}\left(10^{-6}\right)$. Such a tiny gauge coupling is certainly rather unusual, which leads one to wonder whether there is not a possibility to somehow lift the upper bound on g by means of another suppression mechanism.

One of the main conceptual achievements in the present paper is the realization that this is indeed possible! Our main observation is that the FI parameter ξ itself could be parametrically suppressed, $\xi/\Lambda^2 \lesssim (\Lambda/M_{\rm Pl})^2$, in consequence of an enhanced flavor symmetry. The logic behind this idea is the following: Under generic circumstances, all we could say about ξ is that it arises from a combination of scalar VEVs in the D-term scalar potential. If, e.g., two scalar fields $\phi_\pm$ with charges ±1 were involved in the generation of ξ, we would write

$$V_D = \frac{g^2}{2}\left[\left\langle\, |\phi_-|^2 \,\right\rangle - \left\langle\, |\phi_+|^2 \,\right\rangle - q_f\big|\tilde{f}\big|^2 + \cdots\right]^2, \quad \xi = \left\langle\, |\phi_-|^2 \,\right\rangle - \left\langle\, |\phi_+|^2 \,\right\rangle. \qquad (44)$$

At this level of the description, a suppressed value of ξ would merely correspond to a fine-tuning among the VEVs of ϕ_+ and ϕ_-, which might appear very unnatural at first sight. In order to explain why ξ should be much smaller than one would naively expect, $|\xi| \ll \left\langle\, |\phi_\pm|^2 \,\right\rangle$, we therefore require a more detailed description of how ξ is actually generated in the course of spontaneous SUSY breaking—which is exactly the case in the DSB model studied in the present paper. Within the IYIT model supplemented by a weakly gauged flavor symmetry, we are able to derive an explicit expression for ξ in terms of the underlying model parameters (see Eq. (26)). The question as to whether or not ξ has a chance of ending up suppressed is then no longer a question pertaining to scalar VEVs, but rather to the Yukawa couplings in the Lagrangian. This opens up the possibility to render ξ arbitrarily small by imposing an approximate flavor symmetry among these couplings.

Recall that the parameter ρ in Eq. (26) characterizes the quality of the "exchange symmetry" "+" $\leftrightarrow$ "−" in the charged meson sector (see Eq. (15) and Footnote 6). In the limit of an exact exchange symmetry, ρ goes to 1 and the FI parameter ξ trivially vanishes altogether,

$$\xi = \frac{2\left(1-\rho^4\right)^{1/2}}{\rho^2}\left(\frac{\Lambda}{\eta}\right)^2 \quad \overset{\rho\rightarrow 1}{\longrightarrow} \quad 0\,. \tag{45}$$

However, before we put too much trust in this limit, we first have to clarify the actual meaning of this exchange symmetry. To do so, note that the exchange symmetry "+" $\leftrightarrow$ "−" can, in fact, be re-formulated as a Z_2 parity acting on the following linear combinations of the fields $M_\pm$ and $Z_\pm$,

$$\frac{1}{\sqrt{2}}\left(M_+ + M_-\right)\,, \quad \frac{1}{\sqrt{2}}\left(Z_+ + Z_-\right)\,, \quad \frac{1}{\sqrt{2}}\left(M_+ - M_-\right)\,, \quad \frac{1}{\sqrt{2}}\left(Z_+ - Z_-\right)\,, \tag{46}$$

where the first two linear combinations transform even and the last two linear combinations transform odd under this Z_2 parity. For generic Yukawa couplings λ_0^a (see Eq. (8)), this Z_2 parity can, however, *not* be realized at the level of the fundamental quark fields above the dynamical scale Λ. For instance, if we tried to realize the Z_2 exchange symmetry by assigning the following transformation behavior to the four fundamental quark fields,

$$\Psi^1 \leftrightarrow \Psi^3\,, \quad \Psi^2 \leftrightarrow \Psi^4\,, \tag{47}$$

the gauge-invariant composite meson fields at low energies, $M_+ \propto \left<\Psi^1\Psi^2\right>$, $M_- \propto \left<\Psi^3\Psi^4\right>$, $M_0^1 \propto \left<\Psi^1\Psi^3\right>$, $M_0^2 \propto \left<\Psi^1\Psi^4\right>$, $M_0^3 \propto \left<\Psi^2\Psi^3\right>$, $M_0^4 \propto \left<\Psi^2\Psi^4\right>$, would transform as follows,

$$M_+ \leftrightarrow M_-\,, \quad M_0^1 \leftrightarrow M_0^1\,, \quad M_0^2 \leftrightarrow M_0^3\,, \quad M_0^4 \leftrightarrow M_0^4\,. \tag{48}$$

In this case, it would not be sufficient to simply set $\lambda_+ = \lambda_-$ in order to realize the exchange symmetry in the superpotential; we also would have to require that $\lambda_0^2 = \lambda_0^3$. This tells us that it is, in general, not possible to identify the Z_2 symmetry as a subgroup of the global $SU(4)$ flavor symmetry, which we obtain in the limit of equal Yukawa couplings, $\lambda_{ij} \equiv \lambda$. For generic Yukawa couplings in the neutral meson sector, the exchange symmetry in the charged meson sector should rather be regarded as an accidental symmetry of the low-energy effective theory, which we happen to encounter once we set $\lambda_+ = \lambda_-$. As nothing but an accidental symmetry of the effective superpotential, the exchange symmetry is then expected to be explicitly broken by higher-order terms in the effective Kähler potential, so that we basically loose all control over its quality.

The lesson from these considerations is that it is not enough to simply send the parameter ρ to 1 in order to suppress the FI parameter ξ. Instead, we have to impose a larger (approximate) global flavor symmetry, not merely a Z_2 exchange symmetry in the charged meson sector. Here, an obvious choice is to require the full $SU(4)$ flavor symmetry to be approximately realized in the IYIT sector, so that $Z_2 \subset SU(4)$. In this case, it is then possible to identify the Z_2 exchange symmetry with a global flavor symmetry of the fundamental theory at high energies and it is conceivable that the parameter ρ indeed takes a value very close to 1. Meanwhile, we caution that the $SU(4)$ symmetry of the IYIT superpotential should not attain an arbitrarily good quality, as this would render the three low-energy vacua of the IYIT model degenerate (see Eq. (13)). In fact, in the limit of an exact $SU(4)$ symmetry, the vacua I, II and III become connected to each other via, in total, four flat directions, which may be regarded as coordinates of the compact space $SO(6)/SO(5)$ [34] and which might cause serious problems at low energies. On the other hand, as long as the $SU(4)$ symmetry is only approximately realized, these four directions in field space have masses that scale with the differences between the geometric means $\lambda = (\lambda_+\lambda_-)^{1/2}$, $\lambda_{14} = \left(\lambda_0^1\lambda_0^4\right)^{1/2}$ and $\lambda_{23} = \left(\lambda_0^2\lambda_0^3\right)^{1/2}$. For $\lambda_0^1 = \lambda_0^4$ and $\lambda_0^2 = \lambda_0^3$, e.g., we find that the neutral mesons M_0^a give rise to two complex mass eigenstates, m_{14}^- and m_{23}^-, with almost vanishing masses [13],

$$m_{m_{14}^-}^2 = \left(\lambda_{14}^2 - \lambda^2\right)\left(\frac{\Lambda}{\eta}\right)^2 , \quad m_{m_{23}^-}^2 = \left(\lambda_{23}^2 - \lambda^2\right)\left(\frac{\Lambda}{\eta}\right)^2 . \tag{49}$$

Requiring that the masses squared of these complex scalars remain positive, $\lambda < \min\left\{\lambda_{14}, \lambda_{23}\right\}$, is then equivalent to the condition that the $B{-}L$-breaking vacuum should be the lowest-lying among the three low-energy vacua of the IYIT sector (see the discussion below Eq. (13)).

Another caveat applying to the quality of the global flavor symmetry in the IYIT sector pertains to the anomaly-free Z_4 symmetry which is realized even for all Yukawa couplings λ_{ij} being different (see Eq. (6)). This symmetry is broken by the VEVs of the charged meson fields, together with SUSY and $B{-}L$, down to a Z_2 parity (under which all quarks transform odd, $\Psi^i \rightarrow -\Psi^i$). If this symmetry was exact, its spontaneous breaking would result in the formation of *stable* domain walls, which might have disastrous cosmological consequences [38].[9] Thus, also the Z_4 symmetry of the IYIT superpotential should only be approximately realized, so that its breaking leads at most to the formation of *unstable* domain walls, which quickly annihilate after their production. This is, e.g., achieved by higher-order terms in the Kähler potential that explicitly break Z_4.

After these qualitative remarks, we are now ready to study the suppression of the FI parameter ξ in more quantitative terms. To do so, let us first expand $m_{\tilde{f}}^{\text{tree}}$ in Eq. (42) around $\rho = 1$,

$$\left(m_{\tilde{f}}^{\text{tree}}\right)^2 = m_{3/2}^2\left[1 - 6^{1/2}\lambda\,\gamma^2\,\frac{q_f}{q}\,\frac{M_{\text{Pl}}}{m_{3/2}}\,\epsilon + \mathcal{O}\left(\epsilon^2\right)\right] , \quad \epsilon \equiv \left(1-\rho^4\right)^{1/2} . \tag{50}$$

[9]Whether or not stable Z_4 domain walls would lead to cosmological problems depends on the scale of inflation: If the Hubble scale during inflation, H_0, is low, $H_0 \lesssim \Lambda$, the Z_4-breaking phase transition takes place during inflation and all dangerous domain walls are inflated away. However, in the case of large-scale inflation, $H_0 \gtrsim \Lambda$, the Z_4 symmetry is only broken after the end of inflation, so that the associated formation of domain walls would pose a problem. Note that similar considerations may also help to explain why the SUSY breaking scale needs, in fact, to be much higher than one would naively expect according to arguments based on the idea of electroweak naturalness [39].

Here, we have introduced the parameter $\epsilon \in [0,1]$ to describe small deviations from the flavor-symmetric limit, $\epsilon \to 0$. Note that ϵ not only directly parametrizes the suppression of ξ, it also corresponds to the relative difference between the Yukawa couplings λ_+ and λ_- squared,

$$\xi = 2\,\epsilon \left(1-\epsilon^2\right)^{-1/2} \left(\frac{\Lambda}{\eta}\right)^2 , \quad \epsilon = \left|\frac{\lambda_+^2 - \lambda_-^2}{\lambda_+^2 + \lambda_-^2}\right| . \tag{51}$$

Our philosophy in the following will now be that the parameter ϵ can, in principle, take arbitrarily small values, so that the exchange symmetry in the charged meson sector becomes arbitrarily good. We emphasize that this is not in contradiction with our above remarks regarding the quality of the $SU(4)$ or Z_4 flavor symmetries, as it only pertains to the relation between the Yukawa couplings λ_+ and λ_-. We can always render the total flavor symmetry sufficiently broken by retaining a (small) hierarchy among λ, λ_{14} and λ_{23}, irrespectively of how close λ_+ and λ_- are to each other. Given the sfermion mass $m_{\tilde{f}}^{\rm tree}$ in Eq. (50), we then find that, in the small-ϵ regime, the tree-level bound on the gauge coupling constant g scales as follows with the suppression factor ϵ,

$$m_{\tilde{f}}^{\rm tree} \geq 0 \quad \Rightarrow \quad g \leq g_{\rm max}^{\rm tree} \approx \left(\frac{\lambda}{6^{1/2}\,q\,q_f}\frac{m_{3/2}}{M_{\rm Pl}}\frac{1}{\epsilon}\right)^{1/2} . \tag{52}$$

For $q=2$ and $q_f=1$, we can, hence, lift the bound on g to some $\mathcal{O}(1)$ value, if ϵ is of $\mathcal{O}\left(10^{-13}\right)$,

$$g_{\rm max}^{\rm tree} \simeq 0.9 \left(\frac{\lambda}{1}\right)^{1/2} \left(\frac{m_{3/2}}{1000\,\text{TeV}}\right)^{1/2} \left(\frac{10^{-13}}{\epsilon}\right)^{1/2} , \tag{53}$$

so that the magnitude of the FI parameter ξ is pushed just below the gravitino mass squared,

$$\xi \simeq \frac{m_{3/2}^2}{q\,q_f\,(g_{\rm max}^{\rm tree})^2} \simeq 0.6\,m_{3/2}^2 . \tag{54}$$

We, thus, find that an approximate flavor symmetry among the couplings of the IYIT sector allows us to sufficiently suppress the $B{-}L$ D-term. Here, the key feature of our analysis has been the calculability of the D-term in the context of the IYIT model, due to which we were able to compute an explicit expression for ξ in terms of the underlying model parameters (see Eq. (26)). We believe that this feature of the IYIT model readily generalizes to a variety of other DSB models. This means that a number of D-terms (belonging to certain gauged flavor symmetries), which might appear very large at first sight, may actually turn out to be substantially suppressed, as long as one imposes the right flavor symmetry on the SUSY-breaking dynamics. While, in retrospective, this result may appear trivial, we emphasize the importance of having concrete examples at one's disposal that illustrate, within the context of specific models, how dynamically generated D-terms may indeed be suppressed by means of approximate flavor symmetries. For this reason, one of the main motivations behind the present paper is to provide just such an example.

This is, however, not the end of the story. So far, we have only considered the tree-level corrections to the masses of the MSSM sfermions. Besides that, we also have to take into account that the nonzero charges of the SUSY-breaking fields $Z_\pm$ result in a mass splitting within the $B{-}L$

vector multiplet. The $B-L$ gauge DOFs thus act as gauge messengers that induce gauge-mediated sfermion masses at the loop level. Here, the most important (one-loop) correction is given as [23]

$$\left(m_{\tilde{f}}^{1-\text{loop}}\right)^2 = -\frac{q_f^2 g^2}{32\pi^2}\, m_V^2 \,\ln\left[\frac{m_{\tilde{a}}^8}{m_V^6 m_\phi^2}\right] = -\frac{q_f^2 g^2}{32\pi^2}\, m_V^2 \,\ln\left[\frac{\left(m_V^2+m^2\right)^4}{m_V^6\left(m_V^2+2m^2\right)}\right]\,, \tag{55}$$

with m_V, $m_{\tilde{a}}$ and m_ϕ denoting the masses of the vector boson A_μ, gaugino $\tilde{a}$ and real scalar ϕ contained in the "massive $B-L$ vector multiplet" V_A, respectively, (see Eqs. (19), (24) and (25))[10]

$$m_V^2 = 2\,q^2 g^2 f_A^2\,,\quad m_{\tilde{a}}^2 = m_V^2 + m^2\,,\quad m_\phi^2 = m_V^2 + 2\,m^2\,,\quad m^2 = \rho^2\lambda^2\left(\frac{\Lambda}{\eta}\right)^2\,. \tag{56}$$

The effective one-loop correction in Eq. (55) contributes at $\mathcal{O}\left(\gamma^4\right)$ to the total sfermion mass,

$$\left(m_{\tilde{f}}^{1-\text{loop}}\right)^2 = \left(\frac{3}{2}\right)^{1/2}\frac{\gamma^4\lambda^3}{8\pi^2}\left(\frac{q_f}{q}\right)^2\frac{M_{\text{Pl}}}{m_{3/2}}\left[\ln 128 + 6\ln\gamma + \mathcal{O}\left(\gamma^2\right)\right] m_{3/2}^2\,, \tag{57}$$

which is always negative. That is, even when the tree-level, ξ-induced contribution to the MSSM sfermion masses is sufficiently suppressed, the gauge-mediated one-loop contribution in Eq. (55) may still render the MSSM sfermions tachyonic. To prevent this from happening, the gauge coupling constant g must remain small enough, so that $\left|m_{\tilde{f}}^{1-\text{loop}}\right|$ is always smaller than $m_{3/2}$.

This results in an absolute upper bound on the gauge coupling that cannot be lifted any further, even if we tune the suppression factor ϵ in Eq. (53) to an arbitrarily small value,

$$\left|m_{\tilde{f}}^{1-\text{loop}}\right| \leq m_{3/2} \quad\Rightarrow\quad g \leq g_{\max}^{\text{loop}} = \frac{\lambda}{2^{7/6}\,q}\exp\left(\frac{1}{4}\,W_{-1}\left[-\frac{2^{1/6}\,512\pi^2}{3^{3/2}\lambda^3}\left(\frac{q}{q_f}\right)^2\frac{m_{3/2}}{M_{\text{Pl}}}\right]\right)\,, \tag{58}$$

where W_{-1} denotes the lower branch of the Lambert W function or product logarithm (which can take values $-\infty \leq W_{-1} \leq -1$ and which satisfies $x = W_{-1}(x)e^{W_{-1}(x)}$, so that $W_{-1}\left(x\,e^x\right) = x$). For $m_{3/2} = 1000\,\text{TeV}$ and $\lambda = 1$, the bound $g_{\max}^{\text{loop}}$ evaluates to $g_{\max}^{\text{loop}} \sim 10^{-3}$, which means that a gauge coupling constant of $\mathcal{O}(1)$ is, in fact, unviable in our set-up. On the other hand, it is worth noting that, by imposing an approximate flavour symmetry, we were able to relax the naive bound on g resulting from the tree-level D-term scalar potential, $g \lesssim 10^{-6}$ (see our discussion below Eq. (43)), by three orders of magnitude, which is a remarkable improvement. Finally, we note that, depending on the value of the suppression factor ϵ, either the tree-level bound on g in Eq. (53) or the loop-level bound in Eq. (58) dominates. This is summarized in Fig. 1, where we show the maximally allowed value of g as a function of ϵ for different values of λ.

[10] More precisely, $\tilde{a}$ is the fermionic component of the $B-L$ goldstone multiplet A, which shares a Dirac mass term with the fermionic component of the linear combination Y (see Eq. (18)), and ϕ is the real part of the complex scalar contained in the goldstone multiplet. $\tilde{a}$ and ϕ therefore correspond to the fermionic and scalar partners of the goldstone phase a [22]. Similarly, if A was to be identified with the chiral axion multiplet in a supersymmetric implementation of the PQ mechanism, $\tilde{a}$ and ϕ would be referred to as the axino and the saxion, respectively [13].

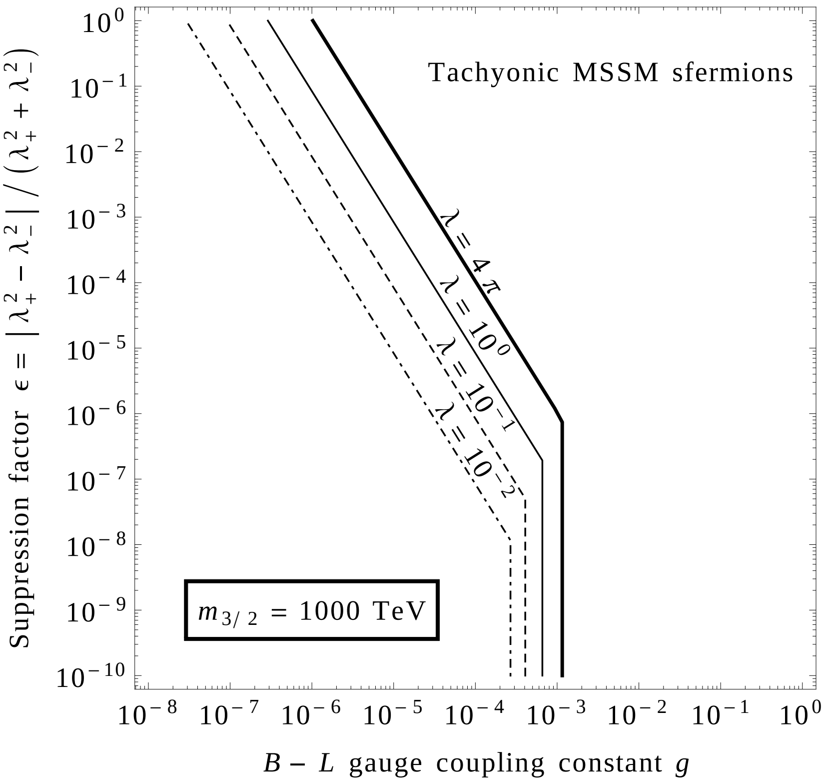

Figure 1: Bound on the $B-L$ gauge coupling constant g as a function of the suppression factor ϵ (see Eq. (50)) and the IYIT Yukawa coupling λ (see Eq. (15)) for $m_{3/2} = 1000\,\mathrm{TeV}$. For very small values of ϵ, the tree-level, ξ-induced contributions to the MSSM sfermion masses are negligible and g is constrained according to the loop-level bound in Eq. (58). For larger values of ϵ, the tree-level bound in Eq. (53) then becomes more stringent than the one in Eq. (58), so that g becomes even more strongly constrained. Note that λ should not be chosen much smaller than $\mathcal{O}\left(10^{-3}\right)$, since otherwise the VEV of the SUSY-breaking field X in SUGRA, $\langle X \rangle \sim 16\pi^2/\lambda^3\, m_{3/2}$ would begin to exceed the Planck scale [13, 22]. At the same time, unitarity restricts λ to take at most a value of $\lambda_{\max} \simeq \eta \simeq 4\pi$.

4 Conclusions and outlook

The IYIT model is an instructive and easy-to-handle toy model for examining how the dynamics of dynamical SUSY breaking might be related to other beyond-the-standard-model phenomena. In particular, the global $U(1)_A$ flavor symmetry present in the IYIT tree-level superpotential is well suited to be identified with other commonly studied local or global $U(1)$ symmetries: (i) In [22], e.g., this $U(1)_A$ symmetry has been promoted to a weakly gauged FI symmetry, $U(1)_A \to U(1)_{\mathrm{FI}}$, in order to demonstrate how dynamical SUSY breaking may entail the generation of a field-dependent FI-term in field theory. (ii) Meanwhile, in [13], the same $U(1)_A$ symmetry has been identified with the global PQ symmetry, $U(1)_A \to U(1)_{\mathrm{PQ}}$, in order to point out a possibility how the dynamical breaking of SUSY may also give rise to a QCD axion that is capable of solving the strong CP problem. (iii) And in the present paper, we have finally promoted the $U(1)_A$ symmetry to a weakly gauged $B-L$ symmetry, $U(1)_A \to U(1)_{B-L}$, in order to illustrate how the paradigm of pure gravity mediation (PGM) may be implemented into concrete models of dynamical SUSY breaking.

This has led us to a number of interesting conceptual and phenomenological observations. For one thing, we have described a mechanism by means of which one is able to sufficiently suppress the $B-L$ D-term, so that it no longer poses a threat to low-energy phenomenology: In the context of the IYIT model, we were able to derive an explicit expression for the $B-L$ FI parameter ξ in terms of the Yukawa couplings appearing in the IYIT superpotential. We then found that, by imposing an approximate flavor symmetry on the SUSY-breaking dynamics, the magnitude of the D-term in the $B-L$ gauge sector can be rendered arbitrarily small. We are confident that similar results also hold for D-terms associated with other gauged flavor symmetries in the context of other DSB models. For another thing, we have identified a direct relation between the heavy neutrino mass scale in the seesaw extension of the MSSM, Λ_N, and the gravitino mass $m_{3/2}$: If the spontaneous breakings of SUSY and $B-L$ should really be tied to each other similarly as in the set-up investigated in this paper, we expect that $\Lambda_N \sim m_{3/2}$. The heavy neutrino mass scale then ends up being much smaller than naively expected, i.e., much smaller than the GUT scale, $\Lambda_N \ll \Lambda_{\rm GUT} \sim 10^{16}\,\text{GeV}$, which has profound implications for cosmology. For heavy Majorana neutrinos as "light" as $m_{3/2} \sim 1000\,\text{TeV}$, we are, e.g., no longer able to rely on standard thermal leptogenesis to account for the origin of the baryon asymmetry in the universe. Instead, leptogenesis should proceed at a much lower energy scale, like in the case of resonant leptogenesis or nonthermal leptogenesis via inflaton decay [40]. Here, we note that such a scenario fits together particularly well with the notion of thermal wino dark matter in the framework of PGM. As has recently been shown, a *thermal* relic abundance of MSSM winos with a mass around 3 TeV allows to nicely reproduce the antiproton-to-proton ratio measured by the AMS-02 experiment in cosmic rays [41]. Therefore, in order to avoid overproduction of *nonthermal* winos in gravitino decays after reheating, the reheating temperature after inflation should not be too high. This favors some form of low-scale leptogenesis over standard thermal leptogenesis, which nicely agrees with the fact that our model predicts a low neutrino mass scale Λ_N. In addition to that, in the particular case of nonthermal leptogenesis via inflaton decay, the reheating temperature should also not be too low, $T_{\rm rh} \gtrsim 10^6\,\text{GeV}$, since otherwise leptogenesis fails to generate a sufficient baryon asymmetry. In this case, the heavy Majorana neutrinos must then have a mass of at least $\mathcal{O}\left(1000\right)$ TeV, which, in the context of our model, translates into a gravitino mass of at least $\mathcal{O}\left(1000\right)$ TeV. Under the specific assumption of nonthermal leptogenesis, the connection between Λ_N and $m_{3/2}$ discussed in the present paper therefore automatically entails a possible answer to the fundamental question as to why SUSY apparently needs to be broken at a scale that is much higher than naively expected according to electroweak naturalness (i.e., as to why we have not yet seen SUSY at colliders). This is an intriguing observation, which directly follows from the connection between the spontaneous breakings of SUSY and $B-L$ proposed in this paper (see also [39] for a similar argument).

Another prediction of our model is the fact that the $B-L$ gauge coupling constant can at most be as large as $\mathcal{O}\left(10^{-3}\right)$. For larger values of g, the SUSY-breaking mass splitting within the massive $B-L$ vector multiplet results in too large (negative) gauge-mediated contributions to the

MSSM sfermion masses. This upper bound on g justifies, *a posteriori*, that we have performed all of our calculations in the weakly gauged limit. From a theoretical point of view, it would, however, still be interesting to generalize our analysis to arbitrary values of the gauge coupling constant. We anticipate such a study to lead to conceptual insights, which may very well imply more general applications for dynamical SUSY breaking and/or gauge mediation than our study for the special case of a local $B-L$ symmetry. Moreover, such an analysis would allow to determine the global maximum of the ratio D_0/F_0 in the IYIT model (see Eq. (39)), which would also be of great theoretical interest. Last but not least, we point out that, if the bounds on ξ and g derived in this paper should only be *marginally* satisfied, we would expect a characteristic modulation of the MSSM sparticle spectrum compared to the "pure PGM" case which is determined by the $B-L$ charges of the MSSM sfermions. This could, in particular, result in a sizable mass gap between light sleptons and heavy squarks— an intriguing possibility, which deserves further study as well.

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