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A NEW GAUGE MEDIATION THEORY

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Abstract

We propose a class of models with gauge mediation of supersymmetry breaking, inspired by simple brane constructions, where R -symmetry is very weakly broken. The gauge sector has an extended $N = 2$ supersymmetry and the two electroweak Higgses form an $N = 2$ hypermultiplet, while quarks and leptons remain in $N = 1$ chiral multiplets. Supersymmetry is broken via the D -term expectation value of a *secluded* $U(1)$ and it is transmitted to the Standard Model via gauge interactions of messengers in $N = 2$ hypermultiplets: gauginos thus receive Dirac masses. The model has several distinct experimental signatures with respect to ordinary models of gauge or gravity mediation realizations of the Minimal Supersymmetric Standard Model (MSSM). First, it predicts extra states as a *third* chargino that can be observed at collider experiments. Second, the absence of a D -flat direction in the Higgs sector implies a lightest Higgs behaving exactly as the Standard Model one and thus a reduction of the ‘little’ fine-tuning in the low $\tan\beta$ region. This breaking of supersymmetry can be easily implemented in string theory models.

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1 INTRODUCTION

The supersymmetric flavor problem is one of the guidelines for constructing realistic models at the electroweak scale with deep implications at LHC. Supersymmetry is generically broken in a hidden sector and transmitted to the observable sector either by gravity or by gauge interactions giving rise to the so-called gravity [1] or gauge-mediated [2] models. In gravity-mediated scenarios the soft-breaking masses are generated at the Planck scale and there is no symmetry reason why they should be flavor-invariant or why they should not mediate large contributions to flavor-changing neutral current processes. In fact this idea motivated the introduction of gauge-mediated theories where the generated soft terms are flavor-blind and therefore they feel flavor breaking only through Yukawa interactions.

In gauge mediated (GMSB) theories supersymmetry is broken in a *secluded* sector such that the Goldstino field belongs to a chiral superfield X which acquires a vacuum expectation value (VEV) along its auxiliary F -component as $\langle X \rangle = M + \theta^2 F$, where it is usually assumed that $F \ll M^2$ so that supersymmetry breaking can be treated as a small perturbation. The theory also contains a vector-like *messenger* sector (Φ, Φ^c) coupled to X by a superpotential coupling $\int d^2\theta \Phi^c X \Phi$ and with Standard Model (SM) quantum numbers providing one-loop Majorana masses to gauginos $\sim \alpha/4\pi F/M$ and two-loop positive squared masses to all sfermions as $\sim (\alpha/4\pi)^2 F^2/M^2$, where $\alpha = g^2/4\pi$ and g is the corresponding gauge coupling. The corresponding effective operators giving rise to these masses are $(1/M) \int d^2\theta X \text{Tr} W^\alpha W_\alpha$ and $(1/M^2) \int d^4\theta X^\dagger X \mathcal{Q}^\dagger \mathcal{Q}$, where W^α are the chiral field strengths of the SM vector superfields and \mathcal{Q} the SM chiral superfields. Fixing all masses in the TeV range one obtains a relation between F and M . In particular if gravity mediated contributions are required not to reintroduce the flavor problem, the flavor changing Planck suppressed contributions should remain smaller than the flavor conserving gauge-mediated ones, i.e. $M \lesssim (\alpha/4\pi) M_{Pl} \sim 10^{16}$ GeV, implying $\sqrt{F} \lesssim 10^{11}$ GeV. In gauge mediation the R -symmetry is broken by the same mechanism that breaks supersymmetry which constrains the corresponding ultraviolet (UV) completion of the theory. The main problem of any fundamental theory, as e. g. string theory, is then to provide the required mechanism of supersymmetry and R -symmetry breaking.

In the simplest compactifications of a class of string theories, leading to intersecting branes at angles, the gauge group sector is often in multiplets of extended supersymmetry while matter states come in $N = 1$ multiplets [3]. A simple way of breaking supersymmetry is by deforming the intersection angles from their special values corresponding to the su-

persymmetric configuration. A small deformation ϵ of these angles breaks supersymmetry via a D -term vacuum expectation value, associated in the T -dual picture to a magnetized $U(1)$ in the internal compactification space, $\epsilon = D\ell_s^2$ with ℓ_s the string length. In this case, the *secluded* sector is given by an (extended) vector multiplet containing an $N = 1$ vector superfield V^0 . The Goldstino is then identified with the gaugino of V^0 which acquires a VEV: $\langle V^0 \rangle = \frac{1}{2}\theta^2\bar{\theta}^2 D$ or $\langle W_\alpha^0 \rangle = \theta_\alpha D$. Of course, this supersymmetry breaking preserves the R -symmetry and can only give Dirac masses to the gauginos of the extended gauge sector.

A prototype model is based on type II string compactifications on a factorizable six-torus $T^6 = \otimes_{i=1}^3 T_i^2$ with appropriate orbifold and orientifold planes and two sets of brane stacks: the observable set \mathcal{O} and the hidden set \mathcal{H} [4, 5]. The SM gauge sector corresponds to open strings that propagate with both ends on the same stack of branes that belong to \mathcal{O} : it has therefore an extended $N = 2$ or $N = 4$ supersymmetry. Similarly, the secluded gauge sector corresponds to strings with both ends on the hidden brane \mathcal{H} . The SM quarks and leptons come from open strings stretched between different stacks of branes in \mathcal{O} that intersect at fixed points of the three torii T_i^2 and have therefore $N = 1$ supersymmetry. The Higgs sector however corresponds to strings stretched between different stacks of branes in \mathcal{O} that intersect at fixed points of two torii and that are parallel along the third one: it has therefore $N = 2$ supersymmetry. Finally the messenger sector contains strings stretched between stacks of branes in \mathcal{O} and the hidden branes \mathcal{H} , that intersect at fixed points of two torii and are parallel along a third torus. It has therefore also $N = 2$ supersymmetry. Moreover, the two stacks of branes along the third torus are separated by a distance $1/M$, which introduces a supersymmetric mass M to the hypermultiplet messengers. The latter are also charged under the supersymmetry breaking $U(1)$ and they are given corresponding supersymmetry breaking squared-masses $\pm D$.

In this paper (inspired by the previous brane constructions) we propose a new gauge mediated theory (NGMSB), alternative to the usual GMSB, where R -symmetry remains unbroken by the mechanism of supersymmetry breaking ². The observable gauge sector is a set of $N = 2$ gauge multiplets corresponding to the SM gauge group. In general an $N = 2$ gauge multiplet contains an $N = 1$ vector multiplet $V = (A_\mu, \lambda_1, D)$ and a chiral multiplet $\chi = (\Sigma, \lambda_2, F_\chi)$. It can be described by the $N = 2$ chiral vector superfield [6]

$$\mathbb{A} = \chi + \tilde{\theta}W + \tilde{\theta}^2 DD\chi, \quad (1.1)$$

²As we will see later, gravitational interactions produce R -symmetry breaking at a subleading level.

where $\tilde{\theta}$ is the second $N = 2$ Grassmannian coordinate and \mathcal{D} the $N = 1$ super-covariant derivative. The secluded sector is identified with the $N = 2$ chiral vector superfield \mathbb{A}^0 (whose $N = 1$ vector superfield is V^0) where supersymmetry is broken by a hidden D -term as ³

$$\langle \mathbb{A}^0 \rangle = \tilde{\theta} \langle W^0 \rangle = \tilde{\theta} \theta D \quad (1.2)$$

For the messenger sector, we choose a (set of) $N = 2$ hypermultiplet(s) denoted as (Φ^c, Φ) , with field contents $\Phi = (\phi, \psi, F_\Phi)$ and $\Phi^c = (\phi^c, \psi^c, F_{\Phi^c})$, charged under the secluded $U(1)$ (with charges ± 1) and with a supersymmetric mass M .

The content of the paper is as follows. In section 2, we present the structure of our model. In particular the transmission of supersymmetry breaking from the hidden to the observable sector is calculated by loop diagrams involving messenger fields. Its main feature is that Dirac masses for gauginos of the $N = 2$ gauge sector are obtained (R -symmetry is preserved by the D -breaking) as well as soft masses for $N = 1$ sfermions. In section 3, we study the generation of the soft-breaking terms for the $N = 2$ Higgs sector, as well as the electroweak symmetry breaking. The main departure of the latter from MSSM is the absence of D -flat direction along which the potential becomes unstable. As a consequence, after electroweak symmetry breaking, the SM-like Higgs has a $\tan \beta$ -independent mass and couplings to SM fermions while the rest of the Higgs sector has $\tan \beta$ -independent masses (at the tree-level) and $\tan \beta$ -dependent couplings to ordinary matter. In section 4, we present a few comments about the gravitino mass and the generation of a tiny F -breaking by gravitational interactions, possible dark matter candidates in our scenario, its experimental signatures at the next high energy colliders (LHC and ILC) and the possibility of unification at a GUT scale. In section 5, our conclusions are drawn. Finally in appendix A we give a description of the supersymmetry breaking potential, based on a Fayet-Iliopoulos (FI) term, for the scalar messengers and the scalar field in the chiral multiplet in the secluded gauge sector. We show that the supersymmetry D -breaking minimum is a local (metastable) one. However we also show that on cosmological times it is absolutely stable.

³We are assuming here that some dynamical mechanism in the hidden sector generates the D -breaking of supersymmetry given by the VEV in Eq. (1.2). A particular realization is given by the angle deformation of intersecting branes described above, leading to an $N = 2$ Fayet-Iliopoulos term, while the corresponding effective potential is analyzed in appendix A.

2 THE MODEL STRUCTURE

The interaction Lagrangian of the messengers with the different gauge sectors is written as

$$\int d^4\theta \left\{ \Phi^\dagger e^V \Phi + \Phi^c e^{-V} \Phi^{c\dagger} \right\} + \left\{ \int d^2\theta \Phi^c \left[M - \sqrt{2} \chi \right] \Phi + h.c. \right\} \quad (2.1)$$

where $V = \sum_A g_A T^A V^A$ contains all gauge fields (in the hidden and observable sectors) and similarly for $\chi = \sum_A g_A T^A \chi^A$ ⁴.

From (2.1) and after replacing the V^0 VEV from Eq. (1.2), we find that the Dirac spinor $\Psi = (\psi, \bar{\psi}^c)^T$ acquires a Dirac mass M while the scalar components have a squared mass matrix given by

$$(\phi^\dagger, \phi^c) \mathcal{M}^2 \begin{pmatrix} \phi \\ \phi^{c\dagger} \end{pmatrix}, \quad \mathcal{M}^2 = \begin{pmatrix} M^2 + D & 0 \\ 0 & M^2 - D \end{pmatrix} \quad (2.2)$$

Notice that in the absence of the supersymmetric mass M , the origin in the (ϕ, ϕ^c) field space is a saddle point. However since we are assuming that $M^2 > D$ the origin becomes a minimum along the (ϕ, ϕ^c) directions. Assuming that supersymmetry is broken by a FI mechanism the Σ^0 -scalar in the extended gauge sector is a flat direction of the potential. However in the presence of supersymmetry breaking it will acquire soft radiative masses, see Eq. (2.12), and the origin $\langle \Sigma^0 \rangle = \langle \phi \rangle = \langle \phi^c \rangle = 0$ becomes a local minimum. Of course there is a global minimum, provided that the SM gauge interactions are ignored, where supersymmetry is restored and gauge symmetry is broken, at $\langle \Sigma^0 \rangle = M/\sqrt{2}$ and $|\langle \phi^c \rangle|^2 = D$. The barrier height separating these minima is very small as compared to the distance between them if $M^2 \gg D$. In that case the tunneling probability per unit space-time volume from the local minimum to the global one is

$$\mathcal{P} \sim e^{-\kappa \frac{M^4}{D^2}} \quad (2.3)$$

where $\kappa > 1$ is a dimensionless constant. For $M^2 \gg D$ the probability (2.3) is so small that the false vacuum is essentially stable on cosmological times. More details are given in appendix A.

⁴We are normalizing the generators such that $\text{Tr} T^A T^B = 1/2 \delta^{AB}$ in the fundamental representation.

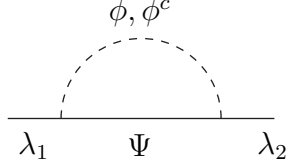


Figure 1: *Feynman diagrams contributing to the Dirac mass of the gaugino λ^a*

Using now the Lagrangian (2.1), the coupling of the SM gauginos with the messenger fields is written as

$$-\sqrt{2} (\phi^c \bar{\lambda}_1 \Psi - \phi^\dagger \bar{\lambda}_2 \Psi) + h.c. \quad (2.4)$$

where we have defined the four-component symplectic-Majorana spinors $\lambda_i = (\lambda_i, \epsilon_{ij} \bar{\lambda}_j)^T$. Notice that if we define a Dirac spinor $\lambda = (\lambda_1, \bar{\lambda}_2)^T$, it satisfies the following identity:

$$\bar{\lambda} \lambda = \frac{1}{2} (\bar{\lambda}_1 \lambda_1 - \bar{\lambda}_2 \lambda_2) . \quad (2.5)$$

A Dirac mass m_a^D for the Dirac gaugino λ_a is radiatively generated from the diagram of Fig. 1 which gives a finite value

$$m_a^D = k_a \frac{\alpha_a}{4\pi} \mathcal{N} \frac{D}{M} \left[1 + \mathcal{O} \left(\frac{D^2}{M^4} \right) \right], \quad a = 1, 2, 3 \quad (2.6)$$

with $k_1 = 5/3$, $k_3 = k_2 = 1$, $\alpha_a = g_a^2/4\pi$, and \mathcal{N} the number of messengers. This value for the Dirac gaugino mass can be equivalently understood from the effective operator [7, 4]

$$\sim \frac{1}{M} \int d^2\theta W^0 \text{Tr}(W\chi) + h.c. \quad (2.7)$$

This operator is actually consistent with $N = 2$ supersymmetry since it is generated by a manifest $N = 2$ supersymmetric Lagrangian (2.1). The operator (2.7) can be rewritten in an explicit $N = 2$ supersymmetric way:

$$\sim \frac{1}{M} \int d^2\theta d^2\tilde{\theta} \mathbb{A}^0 \text{Tr}(\mathbb{A})^2 + h.c. \quad (2.8)$$

where \mathbb{A}^0 is the *secluded* $U(1)$ $N = 2$ vector superfield.

In Ref. [5] such an operator was computed in string theory for the same physical setup of brane configurations discussed in section 1. The result was found to be topological, in the

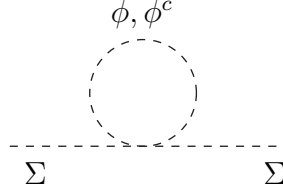


Figure 2: *Feynman diagrams contributing to squared masses for the adjoint scalars Σ^a*

sense that it is independent of the massive string oscillator modes. It receives contributions only from the field theory Kaluza-Klein (KK) part of the torus along which the messengers brane intersection of the observable stack \mathcal{O} with the hidden stack \mathcal{H} is extended. The separation ℓ of the two stacks along a direction within this torus determines the messengers mass $M = \ell$ in string units. The gaugino mass (in the limit $D\ell_s^2 \ll 1$) can then be written as an integral over the real modulus parameter t of the worldsheet annulus having as boundaries the two brane stacks ⁵:

$$m_{1/2}^D \simeq \frac{\alpha}{2} \mathcal{N} D \int_0^\infty \sum_{n=-\infty}^{+\infty} (nR + \ell) e^{-2\pi t(nR + \ell)^2}, \quad (2.9)$$

where for simplicity we have chosen the brane separation to be along one of the two dimensions of an orthogonal torus of radius R . The integration can be performed explicitly with the result:

$$m_{1/2}^D \simeq \frac{\alpha}{4\pi} \mathcal{N} D \left\{ \frac{1}{\ell} - \frac{2}{R^2} \sum_{n \geq 1} \frac{\ell^2/R^2}{n^2 - \ell^2/R^2} \right\}. \quad (2.10)$$

The first term in the right-hand side reproduces precisely the field theory expression (2.6), while the second term represents the (sub-leading) contribution of the messengers KK excitations.

The second superpotential term in Eq. (2.1) gives rise to the F -term potential

$$2 \left\{ \text{Tr} (\phi^\dagger \Sigma^\dagger \Sigma \phi) + \text{Tr} (\phi^c \Sigma^\dagger \Sigma \phi^{c\dagger}) \right\}. \quad (2.11)$$

By using now the one-loop diagrams of Fig. 2 we obtain the mass term $m_{\Sigma^a}^2 |\Sigma^a|^2$ as⁶

$$m_{\Sigma^a}^2 = k_a \frac{\alpha_a}{4\pi} \frac{D^2}{M^2} \left[1 + \mathcal{O} \left(\frac{D^2}{M^4} \right) \right] \quad (2.12)$$

⁵Note a factor of t misprint in Eq. (3.22) of Ref. [5]. Here, we also restored the numerical prefactor.

⁶In the following, we set $\mathcal{N} = 1$ for simplicity.

This value can be understood from the effective operator

$$\sim \frac{1}{M^2} \int d^2\theta (W^0)^2 \mathcal{P}_C \text{Tr} [\chi^\dagger \chi] + h.c. \quad (2.13)$$

where \mathcal{P}_C is the non-local operator [8, 9]

$$\mathcal{P}_C = \square^{-1} \bar{D} \bar{D} D D \quad (2.14)$$

that, acting on a real superfield, produces a chiral one such that its lowest component contains the lowest component of the real superfield. In our case the lowest component of $\mathcal{P}_C \text{Tr} \chi^\dagger \chi$ contains $\text{Tr} \Sigma^\dagger \Sigma$ and (2.12) follows. As in the case of gaugino masses discussed above, one can also show that the operator (2.13) is actually consistent with $N = 2$ supersymmetry.

The sector of quarks and leptons is made of $N = 1$ chiral multiplets, that we generically denote as $\mathcal{Q} = (Q, q_L)$. Its interactions with the gauge sector are given by the Lagrangian

$$\int d^4\theta \mathcal{Q}^\dagger e^V \mathcal{Q} \quad (2.15)$$

In principle, since the messenger sector has SM quantum numbers, there are quartic interactions (from integration of the SM D -terms) as

$$\frac{1}{2} g_a^2 \left(\phi_i^\dagger \phi_j - \phi_i^c \phi_j^{c\dagger} \right) Q_k^\dagger Q_l \sum_A (T_\phi^A)_{ij} (T_Q^A)_{kl} \quad (2.16)$$

where T_ϕ^A are the generators of the gauge group in the representation of ϕ . From Eq. (2.16) and by performing a one-loop contraction of the messenger fields ϕ and ϕ^c , one could in principle provide the SM sfermions Q with a phenomenologically unacceptable large squared mass $\sim D$. For non-abelian group factors this contribution vanishes since it is proportional to $\text{Tr}(T_\phi^A) = 0$. However, this cancellation does not automatically take place for the case of $U(1)$ factors, as the hypercharge. Moreover, the sign of the different sfermion squared masses depends on their hypercharge and tachyonic masses can thus be generated. A solution to this problem appears if ϕ is in a complete representation of $SU(5)$ in which case it is guaranteed that $\text{Tr}(Y_\phi) = 0$ ⁷. This mechanism is similar to that proposed in

⁷Notice that this solution uses the fact that within the states of the representation ϕ the mass matrix is proportional to the unity and hence does not involve the vector-like character of the hypermultiplet.

Ref. [10] in conventional gauge mediation. In fact, by making a $\pi/4$ rotation on the fields $(\Phi, \Phi^{c\dagger})$, the mass matrix (2.2) rotates to

$$\begin{pmatrix} M^2 & D \\ D & M^2 \end{pmatrix} \quad (2.17)$$

which coincides with the supersymmetry breaking mass matrix in usual gauge mediation via an F -term breaking when the messenger sector is invariant under a ‘‘messenger parity’’ by just making the identification $F = F^\dagger = D$.

In view of the previous identification and given that the adjoint superfields Σ do not have direct interactions with $N = 1$ matter, the two-loop diagrams that contribute to the sfermions masses are computed in the same way as in usual gauge mediation models (see the diagrams shown for instance in Ref. [11]) with the result:

$$m_Q^2 = 2 \sum_{j=1}^3 k_j C_j(Q) \left(\frac{\alpha}{4\pi}\right)^2 \frac{D^2}{M^2} \left[1 + \mathcal{O}\left(\frac{D^2}{M^4}\right)\right], \quad (2.18)$$

where $C_j(N) = (N^2 - 1)/2N$ for the fundamental representation of $SU(N)$. This value of the sfermion masses can be understood from the effective operator

$$\sim \frac{1}{M^2} \int d^2\theta (W^0)^2 \mathcal{P}_C[\mathcal{Q}^\dagger \mathcal{Q}] \quad (2.19)$$

where \mathcal{P}_C is the chiral projection operator defined in Eq. (2.14).

Finally associated with the superpotential term $\int d^2\theta h_t \mathcal{H}^c \mathcal{Q} \mathcal{U}$ there exists the relevant supersymmetry breaking parameter A_t that appears in the (R -symmetry breaking) Lagrangian $-A_t h_t H^c \mathcal{Q} \mathcal{U}$. This term generates a mixing between the left and right-handed stops and plays a crucial role in the radiative corrections contributing to the lightest Higgs mass. In our model we have $A_t = 0$ at the scale M and furthermore, since the D supersymmetry breaking does not break the R -symmetry, this null value is not modified by the renormalization group equations to any order in perturbation theory, unlike the case of usual gauge mediation where the gaugino Majorana masses contribute to the one-loop renormalization of A_t . However since, as we shall see, R -symmetry is broken by tiny effects required for electroweak symmetry breaking (and in particular also by gravitational interactions) there will be a small trilinear parameter A_t but irrelevant for phenomenological purposes.

3 ELECTROWEAK BREAKING

Concerning the Higgs sector, it belongs to an $N = 2$ hypermultiplet $\mathbb{H} = (\mathcal{H}^c, \mathcal{H})$ and its interactions with the gauge sector are given by the Lagrangian

$$\int d^4\theta \{ \mathcal{H}^\dagger e^V \mathcal{H} + \mathcal{H}^c e^{-V} \mathcal{H}^{c\dagger} \} - \left\{ \sqrt{2} \int d^2\theta \mathcal{H}^c \chi \mathcal{H} + h.c. \right\} \quad (3.1)$$

The Higgs scalars acquire two-loop masses from Feynman diagrams leading to doublet slepton masses (2.18) and those of Fig. 3 that come from the $N = 2$ superpotential gauge interactions in Eq. (3.1).

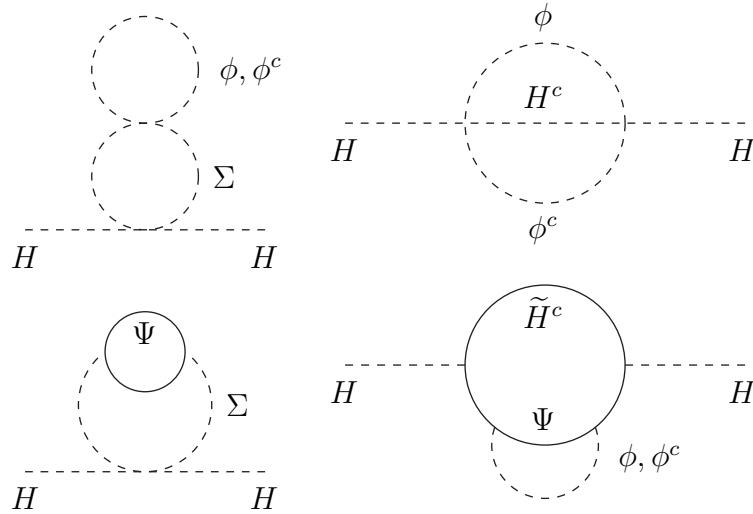


Figure 3: *Feynman diagrams from $N = 2$ superpotential contributing to the H squared mass. The same diagrams contribute also to the H^c squared mass (make the change $H \leftrightarrow H^c$ in all graphs)*

It turns out that the contribution from the superpotential interactions equals that from the gauge couplings and so the result is easily written as

$$m_H^2 = 4 \sum_{j=1}^2 k_j C_j(H) \left(\frac{\alpha}{4\pi} \right)^2 \frac{D^2}{M^2} \left[1 + \mathcal{O} \left(\frac{D^2}{M^4} \right) \right] \quad (3.2)$$

that can be understood as coming from the effective operator

$$\sim \frac{1}{M^2} \int d^2\theta \mathcal{W}^2 \mathcal{P}_C [\mathcal{H}^\dagger \mathcal{H} + \mathcal{H}^c \mathcal{H}^{c\dagger}]. \quad (3.3)$$

Note that the Higgs potential is corrected with some additional contributions coming from Eq. (2.7) [7]. The resulting Lagrangian for the Σ^A scalars, including the radiative mass terms in Eq. (2.12) can then be written as

$$-m_2^D D^a (\Sigma^a + \Sigma^{a\dagger}) - m_1^D D_Y (\Sigma_Y + \Sigma_Y^\dagger) - m_{\Sigma^a}^2 |\Sigma^a|^2 - m_{\Sigma_Y}^2 |\Sigma_Y|^2 \quad (3.4)$$

where Σ^a and D^a are the adjoint scalar and auxiliary field of the $SU(2)$ vector superfield, Σ_Y and D_Y are the neutral scalar and auxiliary field of the $U(1)$ -hypercharge vector superfield, and m_2^D and m_1^D represent the corresponding gaugino Dirac masses. We can now integrate out the adjoint fields Σ^a and Σ_Y . In the absence of the mass terms (2.12) this integration would yield $D^a = D_Y = 0$ which would be a phenomenological drawback for this kind of models. However, in the presence of the mass terms (2.12) the integration yields:

$$\begin{aligned} \Sigma^a &= -\frac{\alpha_2}{4\pi} \frac{D^a}{m_2^D} \\ \Sigma_Y &= -\frac{k_1 \alpha_1}{4\pi} \frac{D_Y}{m_1^D} \end{aligned} \quad (3.5)$$

where we used the relation $(m_a^D)^2 \simeq (\alpha_a/4\pi)m_{\Sigma^a}^2$, following from the expressions (2.6) and (2.12) for $k_2 = 1$.

Replacing now (3.5) in the Lagrangian generates an extra term (quartic in the Higgs fields)

$$\frac{\alpha_2}{4\pi} \vec{D}^2 + \frac{k_1 \alpha_1}{4\pi} D_Y^2 \quad (3.6)$$

which is a small correction to the tree-level Lagrangian $\frac{1}{2}(\vec{D}^2 + D_Y^2)$ and can thus be neglected in the calculation of the Higgs scalar potential. On the other hand, from Eq. (3.5), the VEV of the neutral component of the $SU(2)$ triplet Σ^3 gives rise to a small violation of the $SU(2)$ custodial symmetry when the neutral Higgses acquire a VEV. It would contribute to the ρ -parameter as

$$\rho_0 - 1 = 4 \frac{\langle \Sigma^3 \rangle^2}{v^2} = g^2 \frac{v^2 M^2 \cos^2 2\beta}{D^2} \simeq 3 \times 10^{-6} \left(\frac{v}{m_2^D} \right)^2 \cos^2 2\beta \quad (3.7)$$

where $v = 174$ GeV, $m_2^D = (\alpha_2/4\pi)D/M$ is the chargino mass, and as usual $\tan \beta$ is the ratio of the two Higgs VEVs, $\tan \beta = v_2/v_1$. Here we also used $\langle D^3 \rangle = g^2 v^2 \cos^2 2\beta$, following

from the usual minimization of the Higgs potential. Present experimental bounds [12], $\rho_0 - 1 < 6 \times 10^{-4}$ are then always satisfied. Similarly, there is an induced VEV for the singlet Σ_Y that contributes to the μ parameter but without any phenomenological impact.

Apart from the Higgs soft masses, there are two other key ingredients of the Higgs sector in order to successfully break $SU(2) \times U(1)$: namely the μ and m_3^2 terms. The former represents a supersymmetric mass for Higgses and Higgsinos that can be written as $\mu \int d^2\theta \mathcal{H}^c \mathcal{H} + h.c.$ Notice that this superpotential mass term is perfectly consistent with the $N = 2$ supersymmetry of the Higgs sector. However, its origin and correct size is one of the main problems of supersymmetric model building ⁸. Moreover the m_3^2 term is a soft mass for the scalar components of \mathcal{H} and \mathcal{H}^c , as $m_3^2 H^c H + h.c.$, which is required to explicitly break the Peccei-Quinn symmetry. The simplest possibility to generate a μ -term would be the existence in the $N = 1$ observable sector of an extra singlet chiral superfield S , coupled to the Higgs sector in the superpotential $\int d^2\theta S \mathcal{H}^c \mathcal{H}$ and that acquires a VEV $\langle S \rangle \sim \text{TeV}$ when supersymmetry is broken, together with an F -component VEV generated (by some O’Raifeartaigh mechanism or even by gravitational interactions as we will see in the next section) at the scale $\sqrt{F_S} \sim \text{TeV}$. This breaking would be subleading with respect to the D -term breaking and thus will not alter any of the general conclusions previously obtained. Finally notice that, as we will comment in the next section, D -breaking implies, in the presence of gravitational interactions and through the cancellation of the cosmological constant, an F -breaking that could be used to solve the μ -problem by some of the existing solutions in the literature either using gravitational [13] or gauge [14, 15] interactions.

The soft masses in the Higgs sector of Eq. (3.2) are ‘lowest order’ masses ⁹, that correspond to given boundary conditions at the scale M , while the renormalization group equations (RGE) evolution should be considered from M to the weak scale m_Z . In particular, the Higgs coupled to the top quark will get a large negative squared mass proportional to the squared top-Yukawa coupling, as in the MSSM ¹⁰. We can now write the Higgs potential in the usual notation where $H_{1,2}$ are the lowest components of the chiral superfields

⁸Actually, as mentioned above, in our model such a term is generated through the $N = 2$ superpotential in Eq. (3.1), upon the VEV of the scalar component Σ_Y of the superfield χ_Y in Eq. (3.5). However its magnitude is too low for phenomenological purposes and an additional source of μ is needed.

⁹Both masses are equal because gauge interactions respect the $N = 2$ structure of the Higgs hypermultiplet.

¹⁰Since the Yukawa couplings only respect $N = 1$ supersymmetry we expect the renormalization of the two Higgs masses to be different.

\mathcal{H} and \mathcal{H}^c respectively. For the neutral components of the Higgs doublets the potential is

$$\begin{aligned}
V &= m_1^2 |H_1|^2 + m_2^2 |H_2|^2 - m_3^2 (H_1 H_2 + h.c.) \\
&+ \frac{1}{8} (g^2 + g'^2) (|H_1|^2 - |H_2|^2)^2 + \frac{1}{2} (g^2 + g'^2) |H_1 H_2|^2
\end{aligned} \tag{3.8}$$

where $m_i^2 = m_{H_i}^2 + \mu^2$ are the mass parameters at low energy including the soft- and the μ -terms. In the last line, the first quartic term is the usual D -term of the MSSM, whereas the second is a genuine $N = 2$ effect similar to the one in Eq. (3.1).

Having an extra quartic term on the potential has interesting consequences for its minimization. Indeed, the origin in the Higgs field space can either be a maximum or a saddle-point, whereas in the MSSM it can only be a saddle-point due to the existence of a flat direction of the D -term ($|H_1| = |H_2|$). The conditions to have a vacuum that breaks electroweak symmetry at the correct value are:

$$\frac{m_Z^2}{2} = -\mu^2 + \frac{1}{\tan^2 \beta - 1} (m_{H_1}^2 - m_{H_2}^2 \tan^2 \beta) \tag{3.9}$$

$$m_A^2 = m_{H_1}^2 + m_{H_2}^2 + 2\mu^2 + m_Z^2 \tag{3.10}$$

where $m_A^2 = 2m_3^2 / \sin 2\beta$ is the squared mass of the pseudo-scalar. Notice the important fact that the MSSM stability condition $2|m_3^2| < m_{H_1}^2 + m_{H_2}^2 + 2\mu^2$ is not required in the case of an $N = 2$ Higgs sector. Actually the MSSM minimization condition $m_A^2 = m_{H_1}^2 + m_{H_2}^2 + 2\mu^2$ is replaced by Eq. (3.10) so that for the same input mass parameters the value of m_A is larger than that of the MSSM. Another feature of the potential is that the ‘‘little’’ electroweak fine-tuning of the MSSM is reduced for values of $\tan \beta$ close to one, as we will see below.

We can now calculate the spectrum of the Higgs sector. As in the MSSM, it is controlled by m_A . However it has no dependence on $\tan \beta$:

$$\begin{aligned}
m_h &= m_Z \\
m_H &= m_A \\
m_{H^\pm}^2 &= m_A^2 + 2m_W^2
\end{aligned} \tag{3.11}$$

Moreover, the rotation matrix R_α from H_1, H_2 to h, H is trivial leading to $\alpha = \beta - \pi/2$ which means that the coupling g_{zhh} is the SM one, while $g_{ZHH} = 0$; therefore h behaves always like a SM Higgs and H plays no role in electroweak symmetry breaking. This leads to a different phenomenology from that of the MSSM as we will describe in the next section.

The spectrum from Eq. (3.11) is modified by radiative corrections. At leading order, the mass matrix for the neutral states can be written as:

$$\mathcal{M} = \begin{pmatrix} m_Z^2 \cos^2 \beta + m_A^2 \sin^2 \beta & (m_A^2 - m_Z^2) \cos \beta \sin \beta \\ (m_A^2 - m_Z^2) \cos \beta \sin \beta & m_Z^2 \sin^2 \beta + m_A^2 \cos^2 \beta + \epsilon \end{pmatrix} \quad (3.12)$$

where ϵ is the leading one loop correction which can be written as:

$$\epsilon = \frac{3m_t^4}{4\pi^2 v^2} \left(\log \frac{m_{\tilde{t}_1}^2 + m_{\tilde{t}_2}^2}{m_t^2} + \frac{X_t^2}{2M_S^2} \left(1 - \frac{X_t^2}{6M_S^2} \right) \right) \quad (3.13)$$

where $m_{\tilde{t}_{1,2}}^2$ are the two stop masses, $X_t = A_t - \mu/\tan\beta \simeq -\mu/\tan\beta$ is the stop mixing mass parameter, and $M_S \simeq \sqrt{(m_{\tilde{t}_1}^2 + m_{\tilde{t}_2}^2)/2}$ represents the common soft supersymmetry breaking scale. It should be noted that for large m_A and *any value of* $\tan\beta$ this corresponds to the limit of the MSSM *at large* $\tan\beta$ so the Higgs mass in this model is always maximal in the decoupling limit for the given value of X_t . In particular for values of $\tan\beta \sim \mathcal{O}(1)$ we get a mixing $X_t^2 \sim \mu^2$. On the other hand, for values of m_A close to m_Z the value for m_h should be below the present bound on the Higgs mass although such small values should be normally excluded by the electroweak symmetry breaking condition (3.10).

Finally it is easy to see that the “little” fine-tuning in this model is greatly reduced with respect to that of the MSSM in the low $\tan\beta$ region. The origin of the “little” fine-tuning in the MSSM is that the Higgs mass only increases logarithmically with the stop mass m_Q while it appears quadratically in the determination of m_Z^2 . In fact since the tree-level mass of the Higgs in the MSSM (in the large m_A limit) goes to zero in the limit $\tan\beta \rightarrow 1$, in order to cope with the LEP limit on the Higgs mass, one should go to the region of very large values of m_Q and thus to a very severe “little” fine-tuning. This fine-tuning is softened in the region of large $\tan\beta$. Since in our model the tree-level Higgs mass does not depend on $\tan\beta$ (it coincides with the MSSM one in the $\tan\beta \rightarrow \infty$ limit) the “little” fine-tuning for any value of $\tan\beta$ coincides with that of the MSSM in the large $\tan\beta$ limit.

4 SOME PHENOMENOLOGICAL ASPECTS

In this section we discuss different phenomenological aspects of the model presented above.

THE GRAVITINO MASS AND GRAVITATIONAL EFFECTS

At the supergravity level, enforcing the cancellation of the cosmological constant leads to the presence of an extra source of supersymmetry breaking through an F -term for some chiral field. The goldstino is then a combination of the fermionic partners of this field and the $U(1)$ gauge boson with non-vanishing D -term. Through the super-Higgs mechanism it is absorbed by the gravitino which gets a mass of order [16]

$$m_{3/2} \simeq \frac{D}{M_{Pl}} \quad (4.1)$$

where $M_{Pl} = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. We will assume here that F is of the same order as D :

$$F \simeq m_{3/2} M_{Pl} \simeq D \quad (4.2)$$

In fact, the relative sizes of D and F are very model dependent. While normally $F \gtrsim D$, models with $F \ll D$ can also be engineered [17].

The additional F -breaking source can be arranged to reside in the hidden sector. It will be communicated only through gravitational interactions to the observable sector. The associated effects qualitatively differ from those from D -terms by the fact that they could break R -symmetry, generate Majorana masses and are not compelled to be flavor independent. This implies in particular that they must be subleading with respect to the gauge mediated mechanism presented in this paper in order to not re-introduce a flavor problem. As the typical size of soft-mass terms induced by gravitational interactions is $\sim F/M_{Pl}$, if Eq. (4.2) holds the condition for gravity mediated contributions to be subleading is that $M \lesssim (\alpha/4\pi) M_{Pl} \sim 10^{16}$ GeV which in turn implies the bound $m_{3/2} \lesssim 1$ TeV. However if $F \ll D$ then it may be possible to suppress gravitational interactions even with $m_{3/2} \gtrsim 1$ TeV.

Note also that anomaly-mediated supersymmetry breaking (AMSB) [18] is subleading with respect to the soft-breaking induced by the gauge mediation mechanism, provided that $M \ll M_{Pl}$. For instance, the gaugino Majorana masses M_a induced by anomaly mediation are

$$M_a \sim \frac{\alpha_a}{4\pi} m_{3/2} \sim \left[\frac{M}{M_{Pl}} \right] m_a^D \quad (4.3)$$

and similarly for the squark and Higgs masses.

Fixing $m_a^D \sim 1$ TeV in Eq. (2.6) one can write D as a function of M as

$$D \sim M \times 10^5 \text{ GeV} \quad (4.4)$$

and plugging it into (4.1) one obtains

$$m_{3/2} \sim \left[\frac{M}{10^9 \text{ GeV}} \right] \text{ keV} \quad (4.5)$$

Preferred ranges for values of the scales M and \sqrt{D} can then be obtained if one imposes certain cosmological constraints on the gravitino mass [19]. For instance as a warm dark matter component a light gravitino mass should lie below ~ 16 eV [20] in order not to have unwanted cosmological consequences, which translates into $M \lesssim 10^7$ GeV. Furthermore as a cold dark matter component the mass of the gravitino should be above few keV, which translates into $M \gtrsim 10^9$ GeV and the reheating temperature has to be such that the density of gravitinos does not overclose the universe [20].

COLLIDER PHENOMENOLOGY

There are three different signatures of these models. The first one, and common to the usual gauge mediation scenarios, arises when the gravitino is very light (the LSP) and therefore the NLSP could have a decay suppressed by the SUSY breaking scale. In this situation and if the NLSP is charged (a very common case) it could yield a track in the detector before decaying.

The second feature in our models is the $N = 2$ structure of the Higgs sector. As we noticed in section 3, the lightest Higgs h behaves as a SM Higgs with no $\tan\beta$ dependence on its couplings to fermions; therefore its production and decay rates are those typical of the SM rather than the MSSM. On the other hand, the couplings to matter of the heaviest Higgs H , the pseudo-scalar A and the charged Higgses H^\pm , as well as all self-couplings within the Higgs sector do depend on $\tan\beta$: in principle one could then distinguish these models from others (such as the MSSM or the non-supersymmetric two-Higgs doublet models) by measuring these couplings, although this may be the realm for ILC rather than LHC.

Finally the last signature comes from the $N = 2$ structure of the gauge sector and consequently from the fact that gauginos have their main mass contribution forming a Dirac particle. This translates into having different decay channels of gluinos, which are

difficult to measure but, more importantly, into the existence of *three* charginos and *six* neutralinos, the latter being paired into three pseudo-Dirac neutral fermions. Discovery of three charginos should constitute the smoking gun of this scenario.

UNIFICATION

Let us finish this phenomenological section with some comments on gauge coupling unification. The model as it is, with $N = 2$ gauge sector and a single Higgs hypermultiplet, together with the usual $N = 1$ chiral matter, has the following beta-function coefficients: $\beta_1 = \frac{33}{5}$, $\beta_2 = 3$, $\beta_3 = 0$. Evolution of the three gauge couplings with the previous beta-functions does not lead to unification. However, this can be achieved by adding appropriate extra states [21, 7]. For instance, if one adds to the above spectrum one hypermultiplet with the quantum numbers of the Higgs plus two hypermultiplets with the quantum numbers of a right-handed lepton (*unifons*) the new beta-function coefficients read:

$$\beta_1 = \frac{48}{5}, \quad \beta_2 = 4, \quad \beta_3 = 0 \quad (4.6)$$

and one recovers the one-loop unification at the MSSM GUT scale $M_{\text{GUT}} \sim 2 \times 10^{16}$ GeV. It should be noted that the unification scale is not affected by the messengers since they form complete $SU(5)$ representations. Moreover, the *unifons* come in vector-like representations and can be given the desired (supersymmetric) mass.

5 CONCLUSIONS

In this paper we have proposed a new model of gauge mediation where supersymmetry is broken by a D -term expectation value in a *secluded* local $U(1)$ sector, thus conserving R -symmetry in the global limit. The NGMSB model, an alternative to the usual GMSB where supersymmetry is broken by the F -term of a *secluded* chiral sector, is easily realized in intersecting brane models of type I string theory [5]. Its main feature is the presence of an extended $N = 2$ supersymmetry in the gauge, as well as in the Higgs and messengers sectors. As a result, the transmission of supersymmetry breaking by the hypermultiplet messengers generates Dirac masses for gauginos at the one-loop level. The observable matter sector is contained in chiral $N = 1$ multiplets (localized in brane intersections) that get radiative masses at the two-loop order. The NGMSB model shares some features with the usual GMSB models:

- It solves the supersymmetric flavor problem.
- It provides a common supersymmetry breaking scale, the masses of all supersymmetric particles being proportional to gauge couplings.
- In sensible models, the gravitino is very light and thus a candidate to describe the Dark Matter of the Universe.

However, NGMSB departs from usual GMSB in a number of features:

- The gaugino sector contains twice as many degrees of freedom as that of the MSSM assembled into (quasi)-Dirac fermions with Dirac masses. Finding three (instead of two) charginos, and six (instead of four) neutralinos should then be the smoking gun of this class of models at LHC.
- The two Higgs superfields of the MSSM are contained in one hypermultiplet. Thus, the Higgs phenomenology for low $\tan\beta$ at LHC is completely different from that of the MSSM. This alleviates the fine-tuning problem of low $\tan\beta$ and opens up the corresponding window in the supersymmetric parameter space.

There is a number of issues whose detailed analysis was outside the scope of the present paper but that should be worth of further study. On the phenomenological side, one should translate LEP data into bounds on the Higgs masses ¹¹ and work out more in detail the LHC phenomenology of both the scalar Higgs and the chargino and neutralino sectors. Finally at the string theory level, the construction of realistic intersecting brane models realizing the ideas contained in this paper.

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¹¹Exclusion plots should look very different from those in the MSSM.

A TUNNELING PROBABILITY DENSITY

In this appendix we will first analyze the potential structure of the $N = 2$ supersymmetric $U(1)$ secluded sector \mathbb{A}^0 with gauge coupling g_0 in the presence of the messenger hypermultiplet (Φ, Φ^c) with charges normalized to ± 1 . We will assume that $N = 2$ supersymmetry is broken by the FI parameter ξ . The interaction of the messenger hypermultiplet (Φ, Φ^c) with the secluded $U(1)$ is described by the Lagrangian

$$\int d^4\theta \left\{ \Phi^\dagger e^{g_0 V^0} \Phi + \Phi^c e^{-g_0 V^0} \Phi^{c\dagger} + \xi V^0 \right\} + \left\{ \int d^2\theta \Phi^c \left[M - \sqrt{2} g_0 \chi^0 \right] \Phi + h.c. \right\} \quad (\text{A.1})$$

that gives rise to the scalar potential

$$V = \frac{D^2}{2g^2} + \left\{ |M - \sqrt{2} \Sigma|^2 + D \right\} |\phi|^2 + \left\{ |M - \sqrt{2} \Sigma|^2 - D \right\} |\phi^c|^2 + \frac{g^2}{2} (|\phi|^2 + |\phi^c|^2)^2 \quad (\text{A.2})$$

where we have used in (A.2), to simplify the notation, $D = g_0^2 \xi$, $\Sigma = g_0 \Sigma^0$ and $g = g_0$.

The potential (A.2) has two minima:

- A local non-supersymmetric minimum at the origin, $\phi = \phi^c = 0$, which is a flat direction along the Σ -field. This minimum has a vacuum energy $\langle V \rangle = D^2/2g^2$. The flat direction is lifted by quantum corrections that provide a radiative mass to Σ as in Eq. (2.12)

$$V_{\text{rad}}(\phi, \phi^c, \Sigma)|_{\phi=\phi^c=0} = m^2 |\Sigma|^2, \quad m^2 \simeq \frac{1}{16\pi^2} \frac{D^2}{M^2} \quad (\text{A.3})$$

Using Eq. (A.3) the local supersymmetry breaking minimum is then at the origin

$$\langle \phi \rangle = \langle \phi^c \rangle = \langle \Sigma \rangle = 0 \quad (\text{A.4})$$

- A global supersymmetric minimum at

$$\langle \phi \rangle = 0, \quad |\langle \phi^c \rangle|^2 = D/g^2, \quad \langle \text{Re}(\Sigma) \rangle = M/\sqrt{2}, \quad \langle \text{Im}(\Sigma) \rangle = 0 \quad (\text{A.5})$$

with zero vacuum energy ¹².

¹²The general features of our analysis should remain true after introducing the interaction of the messengers with the observable gauge sector. In particular, when the gauge sector \mathbb{A}^A is introduced, this minimum is uplifted and the $\langle |\phi^c|^2 \rangle$ VEV is shifted by the observable D -term $g_A^2/2 (\phi^c T^A \phi^c)^\dagger$, while the Σ direction should have components along the Σ^A -fields of the observable sector. However radiative corrections (2.12) provide masses proportional to $\text{tr} T_\phi^A T_\phi^B + \text{tr} T_{\phi^c}^A T_{\phi^c}^B$. Given that $T_{\phi, \phi^c}^{U(1)} \propto 1$ and $\text{tr} T_{\phi, \phi^c}^Y = 0$ there is no mixing between the *secluded* $U(1)$ and $U(1)_Y$ and all flat directions are lifted. Since we intend to do only a semi-quantitative analysis of the tunneling probability from the local (A.4) to the global (A.5) minimum we do not include in this appendix gauge interactions from the observable sector.

For any value of Σ at zero mass, the potential has a minimum at the origin along the ϕ -direction and for values of Σ such that $|M - \sqrt{2}\Sigma|^2 > D$ the potential also has the minimum (A.5) along the ϕ^c direction. However when $|M - \sqrt{2}\Sigma|^2 = D$ there is an almost flat direction along ϕ^c at the minimum that becomes a maximum for $|M - \sqrt{2}\Sigma|^2 < D$. Since the flat direction Σ is lifted only by radiative corrections [see Eq.(A.3)] the path followed by the instanton which controls the tunneling from the local (A.4) to the global (A.5) minimum goes along the $\text{Re}(\Sigma)$ -direction, from the origin up to values of $\text{Re}(\Sigma)$, $\text{Re}(\Sigma) = M/\sqrt{2} + \mathcal{O}(D^{1/2})$ where the potential becomes unstable along the ϕ^c direction. Since the instanton has to jump over the path of least slope along $\text{Re}(\Sigma)$, up to values of $\text{Re}(\Sigma)$ where the potential becomes unstable, the tunneling problem can be very well approximated by a one-dimensional problem where the instanton field that satisfies the euclidean equation of motion [22]

$$\frac{d^2\varphi}{dr^2} + \frac{3}{r} \frac{d\varphi}{dr} = V'(\varphi) \quad (\text{A.6})$$

is $\varphi \equiv \text{Re}(\Sigma)/\sqrt{2}$. The variable $r = \sqrt{t^2 + \vec{x}^2}$ in (A.6) makes the $O(4)$ symmetry of the solution manifest with the boundary conditions $\varphi \rightarrow 0$ at $r \rightarrow \infty$ and $d\varphi/dr = 0$ at $r = 0$.

In order to make an estimate of the probability of bubble formation by tunneling from the local (A.4) to the global (A.5) we will follow Ref. [23]. First of all notice that the depth of the global minimum $\sim D^2/2g^2$ is much larger than the height of the barrier $\sim D^2/32\pi^2$ so that the bubble solution is outside the domain of validity of the thin wall approximation. Second, to compute the tunneling probability one can disregard the details of the behaviour of the potential $V(\varphi)$ at $\varphi \gg \varphi_1$ where $V(\varphi_1) \simeq V(0)$. In particular considering a potential $V_{\text{app}}(\varphi)$ that approximates $V(\varphi)$ for $\varphi \lesssim \varphi_1$ is usually a sufficiently good approximation [23]. In our case $\varphi_1 \simeq M/2$ and the simplest such potential is

$$V_{\text{app}}(\varphi) = \frac{1}{2}m^2\varphi^2 - \frac{\delta}{3}\varphi^3 \quad (\text{A.7})$$

where the mass term is given by Eq. (A.3) and δ is chosen such that $V(\varphi_1) \simeq 0$, i. e.

$$\delta \simeq \frac{3m^2}{2\varphi_1} \sim \frac{3}{16\pi^2} \frac{D^2}{M^3} \quad (\text{A.8})$$

The euclidean action S_4 corresponding to the solution φ of Eq. (A.6) is given by

$$S_4 \sim 2 \left(\frac{10M}{3m} \right)^2 \sim \left(\kappa_4 \frac{M^2}{D} \right)^2 \quad (\text{A.9})$$

where $\kappa_4 \simeq 59$.

At finite temperature one should replace $V(\varphi)$ by $V(\varphi, T)$ where finite temperature effects have been considered. In our case the dominant thermal effects in Eq. (A.7) can be encoded in the thermal mass

$$m^2(T) \simeq \frac{1}{16\pi^2} \left(\frac{D^2}{M^2} + aT^2 \right) \quad (\text{A.10})$$

where a is some order one coefficient. At finite temperature the $O(4)$ symmetric solution of (A.6) should be replaced by the $O(3)$ symmetric one that satisfies the equation [23]

$$\frac{d^2\varphi}{dr^2} + \frac{2}{r} \frac{d\varphi}{dr} = V'(\varphi, T) \quad (\text{A.11})$$

where now $r^2 = \bar{x}^2$. For the potential (A.7) it is found [23]

$$\frac{S_3}{T} \sim 5 \frac{M^3 m^3(T)}{m^4 T} \gtrsim \left(\kappa_3 \frac{M^2}{D} \right)^2 \quad (\text{A.12})$$

where $\kappa_3 \simeq 45$ and the last inequality holds for any temperature. The tunneling probability per unit time per unit volume is then

$$\mathcal{P} \sim e^{-\mathcal{B}} \quad (\text{A.13})$$

where $\mathcal{B} = \min [S_4, S_3(T)/T]$ at any temperature. This justifies our Eq. (2.3).

Of course we do expect expressions (A.9) and (A.12), based on the approximated potential (A.7) to be accurate only within factors of order a few. In fact if we approximate the potential by a different one with a negative quartic term, $V_{\text{app}} = 1/2 m^2 \varphi^2 - 1/4 \lambda \varphi^4$ as in Ref. [23], and fix λ such that $V_{\text{app}}(\varphi_1) \simeq 0$ the corresponding expressions (A.9) and (A.12) have coefficients $\kappa_4 \simeq 23$ and $\kappa_3 \simeq 19$ respectively, thus related to those in (A.9) and (A.12) by order two factors. Given that in all cases the euclidean actions are dominated by the large M^4/D^2 -factors, these results are very consistent to each other.

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