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BROKEN SUPERSYMMETRIES IN HIGH ENERGY PHYSICS *

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ABSTRACT

The renormalisation group analysis of the running coupling constants in the hierarchies of N-extended supersymmetric simple unification schemes is presented. For proton lifetimes of order 10^{30} years the scale(s) of supersymmetry breaking are of order 10^{12} GeV. In local realisations of such supersymmetries, such high mass-scales lead to gravitinos with masses in the 10^{5} GeV range. Gravitinos this massive decay long before the time of helium synthesis to be of relevance in the early universe.

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The hierarchy problem[1] of grand unified theories that requires the unnatural fine tuning of the parameters in the scalar potential to effect the desired hierarchy of gauge symmetries has been the main motivation of the considerable recent activity in highlighting the relevance of global supersymmetry for elementary particle physics. Other than this and the mere aesthetics of combining fermions and bosons in oné multiplet there are indications from other directions of the important role that supersymmetry can play in formulating theories of elementary particles. Thus the extended version of supersymmetry, $N = 4$ in particular, offers the exciting possibility of a finite field theory and supersymmetry realised locally[2] has the effect of softening the ultraviolet divergences encountered in coupling gravity with matter.

If Nature does make use of either supersymmetry or extended supersymmetry then it must be broken as is evident from the known spectrum of particles. An important parameter is the scale at which supersymmetry breaking occurs. If it is in the few TeV range then presumably it solves the previously mentioned "hierarchy" problem. But this need not be the case. The scale of supersymmetry breaking could be in the (10^{12} - 10^{19}) GeV range[3]. In this case the solution to the hierarchy problem and the naturalness of fine tuning require a different approach. If the scalars making up the potential are elementary then they would have to be strongly interacting in order to compensate for their enormous mass in order to keep fixed the scale of weak interactions. Another possibility is that at low energies the strongly interacting scalar of very high mass actually form scalar-isoscalar low mass bound states[4] that serve the purpose of symmetry breaking of electroweak interactions. Thus phenomenologically, at least, strongly interacting scalar fields are as interesting and probably as useful as the weakly interacting ones. However due to the difficulty of making specific calculations with them at present we will not address this problem any further. In what follows we study the evolution of the gauge couplings in the hierarchical descents of extended supersymmetries to SU(3) x SU(2) x U(1) symmetry. From these and the constraints on phenomenologically accessible parameters the mass-scales of supersymmetry breaking are deduced. These are then discussed in relation to the ones deduced from cosmological arguments.

Let G_U^* represent the grand unifying gauge group. The asterisk denotes that supersymmetry is not broken. The strong and electroweak symmetry of particle interactions in the supersymmetric and ordinary states are denoted by $G_{SEW}^* = SU(3)_C^* \times SU(2)_L^* \times U(1)^*$ and $G_{SEW} = SU(3)_C \times SU(2)_L \times U(1)$. Consider N = 1 supersymmetry in which G_U^* descends to $SU(3)_C \times U(1)_{em}$ as follows

$$G_U^*(N=1) \xrightarrow{M_U^*} G_{SEW}^*(N=1) \xrightarrow{M_1^*} G_{SEW} \xrightarrow{M_W} SU(3)_C \times U(1)_{em} \quad (1)$$

Depending on the choice of G_U^*, clearly there are other gauge sub-symmetries between G_U^* and G_{SEW}^* but for simplicity we will assume that G_U^* is supersymmetrised SU(5). The running coupling constants of the hierarchy in (1) at the renormalisation point μ equal to the scale of electroweak interactions are given by

$$g_3^{-2}(M_W) = g_U^{-2} + (6f^* + 9n_A - 27)K \ln\frac{M_U^*}{M_1^*} + (4f-33)K \ln\frac{M_1^*}{M_W} , \quad (2)$$

$$g_2^{-2}(M_W) = g_U^{-2} + (6f^* + 6n_A - 18)K \ln\frac{M_U^*}{M_1^*} + (4f-22+\frac{n_D}{2})K \ln\frac{M_1^*}{M_W} , \quad (3)$$

$$g_1^{-2}(M_W) = g_U^{-2} + 6f^*K \ln\frac{M_U^*}{M_1^*} + (4f+\frac{3n_D}{10})K \ln\frac{M_1^*}{M_W} . \quad (4)$$

Here f^* and n_A are the number of matter multiplets in the fundamental and adjoint representations before supersymmetry is broken at the mass-scale M_1^*, f and n_D are the number of ordinary fermion families and light scalar doublets after supersymmetry is broken, $K = 1/24\pi^2$. It is assumed that no light scalar triplets survive. The intermediate mass-scales (M_U^*, M_1^*, M_W) and the weak angle $\sin^2\theta_W$ satisfy the following two relations:

$$(9-3n_A)\ln\frac{M_U^*}{M_W} + (2+3n_A+\frac{n_D}{6})\ln\frac{M_1^*}{M_W} = \frac{\pi}{\alpha_{em}}\left(1-\frac{8\alpha_{em}}{3\alpha_s}\right) , \quad (5)$$

$$\sin^2\theta_W(M_W) = \frac{1}{6} + \frac{5\alpha_{em}(M_W)}{9\alpha_s(M_W)} + \frac{n_D}{18}\frac{\alpha_{em}(M_W)}{\pi}\ln\frac{M_1^*}{M_W} . \quad (6)$$

In deriving the above relations the definitions of the electric charge e and the weak angle $\sin^2\theta_W$ used are

$$e^{-2}(M_W) = g_2^{-2}(M_W) + (5/3)\, g_1^{-2}(M_W) , \quad (7)$$

$$\sin^2\theta_W(M_W) = e^2(M_W)/g_2^2(M_W) . \quad (8)$$

Notice that complete supersymmetric matter multiplets do not contribute to the weak angle $\sin^2\theta_W$ (Eqn. 6), only the light doublets contribute. From equation (5), unless the number of light scalar doublets are quite large ($\gtrsim$ 20), the masses $\{M_U^*, M_1^*\}$ are constrained only by the number n_A of supersymmetric adjoint matter multiplets. For proton lifetimes[F1] in the range of 10^{30} years to 10^{46} years which correspond to M_U^* in the range 10^{15} GeV to 10^{19} GeV, the predictions of M_1^* are given in table I. The values of the strong and electromagnetic couplings $\{\alpha_s = g_3^2/4\pi,\ \alpha_{em} = e^2/4\pi\}$ taken are {0.12, 1/128}. Note also that for an arbitrary large number of scalars in the adjoint contributing, the unification mass-scale M_U^* and the supersymmetry breaking mass-scale M_1^* are comparable (Equation 5). In this limit as M_U^* approaches the Planck scale ($\sim 10^{19}$ GeV) so does M_1^*. This situation is contrary to the one encountered in the absence of adjoints contributing where the supersymmetry breaking mass-scale M_1^* approached the electroweak scale ($\sim$100 GeV) as M_U^* approached Planck energies. In equations 2 - 4, the SU(2) and the SU(3) gauge couplings are unrenormalised by the unifying mass-scale M_U^* if $n_A = 3$. Also if $n_A > 3$ then in the energy range characterised by M_U^* and M_1^* the theory is 'strongly' interacting. Perhaps supersymmetric theories are after all strongly interacting before supersymmetry is broken. The strongly interacting particles help to form bound states of light scalars which break supersymmetry and at the same time serve to break the remaining symmetry. However this scenario although tempting, is highly speculative and it is unclear what dynamics to propose to make it viable. (F2) From table I, the scale of supersymmetry breaking M_1^*/M_W is clustered around the value 10^{13}, as long as the scalar contribution in the adjoint representation is not too large. For these value of M_1^*/M_W, the values of $\sin^2\theta_W$ as expected from equation (6) are given in table II. These value lie comfortably in the range deduced experimentally[5] i.e. $\sin^2\theta_W|_{exp} = 0.23 \pm 0.015$ without radiative corrections to neutral current data and $\sin^2\theta_W|_{exp} = 0.215 \pm 0.012$ with radiative corrections to neutral current data. These experimental numbers impose the following constraint on $\eta_1^* = \log_{10} M_1^*/M_W$ and on n_D.

$$40 < \eta_1^* n_D < 140 \qquad \text{without radiative corrections} \qquad (9)$$

$$0 < \eta_1^* n_D < 80 \qquad \text{with radiative corrections} . \qquad (10)$$

As many as four light doublets are easily tolerated even if the supersymmetry breaking scale is as high as the Planck mass. We conclude the discussion of N = 1 supersymmetry by noting that if the scale of supersymmetry breaking is of the order of the weak interaction scale as is entertained in theories that address solutions to the hierarchy problem, then the presence of scalar adjoints lead to the supersymmetric unification scale to be much higher than that encountered in the simple SU(5) scheme. This can be seen from equation (5). The scale M_U^* and M_U of the simple SU(5) are related by

$$M_U^* = (M_U)^{\frac{11 + n_D/6}{9 - 3n_A}} . \qquad (11)$$

Thus a single adjoint contribution leads to $M_U^* > 10^{30}$ GeV for the standard value of $M_U \simeq 10^{15}$ GeV which is in the macroscopic rather than the microscopic regime. We now turn to the discussion on extended supersymmetry.

The general features of N = 2, 3 and 4 extended global supersymmetry in its applications to supergrand unification have already been considered[6]. The matter and gauge vector multiplets of N = 2 and 4 have as physical particles the 'on shell' following content:

N	Gauge	Matter
2	$1, \frac{1}{2}^2, 0^2$	$\frac{1}{2}^2, 0^2$
4	$1, \frac{1}{2}^4, 0^6$	-

where j^n denotes an n-plet of spin j particles and all particles are described by real (or Majorana) wave functions. At present spontaneous supersymmetry breaking in unification schemes with extended supersymmetry breaking seems difficult. In the case of N = 2 we failed to find any mechanism for spontaneous supersymmetry breaking at the tree level[7]. While either explicit (soft) or non-perturbative N = 2 supersymmetry breaking seems the only method of relevance at present, we are of the opinion that the technology is not sufficiently developed to achieve N = 2, or for that matter any N > 2, spontaneous supersymmetry breaking. In this spirit we entertain the following hierarchy in N = 2 grand unification with an arbitrary internal gauge symmetry G_U^*,

$$G_U^*(N=2) \xrightarrow{M_U^*} G_{SEW}^*(N=2) \xrightarrow{M_2^*} G_{SEW}^*(N=1) \xrightarrow{M_1^*} G_{SEW} \xrightarrow{M_W} SU(3)_c \times U(1)_{em} \qquad (12)$$

The above hierarchy is plausible in view of the fact that N = 2 can be completely expressed in terms of N = 1 multiplets. Since no 'mirror' or 'gauge' matter has been observed at present energies, it is not unreasonable to suppose that these decouple themselves at some high energy M_2^*. The renormalised gauge couplings at the renormalisation point μ equal to the mass-scale M_W of electroweak interactions are given by

$$g_3^{-2}(M_W) = g^{-2}(M_U^*) + K(12f^{(2)} + 6n_A^{(2)} - 18)\ln\frac{M_U^*}{M_2^*} + K(6f^{(1)} + 9n_A^{(1)} - 27)\ln\frac{M_2^*}{M_1^*} + K(4f - 33)\ln\frac{M_1^*}{M_W}, \qquad (13)$$

$$g_2^{-2}(M_W) = g^{-2}(M_U^*) + K(12f^{(2)} + 4n_A^{(2)} - 12)\ln\frac{M_U^*}{M_2^*} + K(6f^{(1)} + 6n_A^{(1)} - 18)\ln\frac{M_2^*}{M_1^*} + K(4f + \frac{n_p}{2} - 22)\ln\frac{M_1^*}{M_W}, \qquad (14)$$

$$g_1^{-2}(M_W) = g^{-2}(M_U^*) + 12f^{(2)}K\ln\frac{M_U^*}{M_2^*} + 6f^{(1)}K\ln\frac{M_2^*}{M_1^*} + (4f + \frac{3n_p}{10})K\ln\frac{M_1^*}{M_W}. \qquad (15)$$

These equations, together with the definitions in equations (7) and (8) lead to the predictions of the weak angle $\sin^2\theta_W$ as in equation (6) while the masses $\{M_U^*, M_2^*, M_1^*, M_W\}$ satisfy the following relation

$$(6 - 2n_A^{(2)})\ln\frac{M_U^*}{M_2^*} + (9 - 3n_A^{(1)})\ln\frac{M_2^*}{M_1^*} + (11 + \frac{n_p}{6})\ln\frac{M_1^*}{M_W} = \frac{\pi}{\alpha_{em}}\left(1 - \frac{8\alpha_{em}}{3\alpha_s}\right). \qquad (16)$$

In equations (13) - (15) $n_A^{(2)}(n_A^{(1)})$ is the number of adjoint matter multiplets and $f^{(2)}$ $(f^{(1)})$ is the number of matter families for N = 2 (N = 1) supersymmetry. With $\alpha_{em} \sim 1/128$, $\alpha_s \simeq 0.012$ Equation (16) satisfies the constraint $(\eta_{U,1,2}^* = \log_{10}\frac{M_{U,1,2}^*}{M_W})$

$$(6 - 2n_A^{(2)})\eta_U^* + (3 - 3n_A^{(1)} + 2n_A^{(2)})\eta_2^* + (2 + 3n_A^{(1)} + \frac{n_p}{6})\eta_1^* = 144.4 . \qquad (17)$$

An interesting case is that for which no adjoint matter survives when N = 2 $SU(3)^* \times SU(2)_L^* \times U(1)^*$ breaks to N = 1 $SU(3)^* \times SU(2)_L^* \times U(1)$ through M_2^* i.e $(n_A^{(1)} = 0)$. Table III gives the values of $M_{1,2}^*$ allowed by the hierarchy in Equation (12). For $n_A^{(1)} = n_A^{(2)} = 0$ as M_U^* approaches Planck energies, M_1^* lies between (10^5 to 10^7) GeV. For $m_A^{(1)} = 0$ and $1 \leqslant n_A^{(2)} \leqslant 3$, M_1^* is found to lie between (10^2 and 10^6) GeV. The precise value of M_1^* will depend on the number of light scalar doublets that will emerge in the model of one's taste. Table III also gives values of $M_{1,2}^*$ for $n_A \neq 0$. For simplicity we have taken $n_A^{(1)} = n_A^{(2)}$. In this case the mass-scales M_1^*, M_2^* of supersymmetry breaking are greater than 10^{10} GeV for $10^{13.5} \leqslant M_U^*/M_W \leqslant 10^{17}$. Thus a realistic N = 2 supersymmetric grand unification can lead to satisfactory predictions of the weak angle $\sin^2\theta_W$ and mass-scales of supersymmetry breaking that are greater than 10^{10} GeV or the N = 2 breaking scale is high $(10^{10} \leq M_2^* \leq 10^{19})$ and the N = 1 breaking scale is low (10^2 GeV $< M_1^* \lesssim 10^7$ GeV). We now turn to discuss the N = 4 supersymmetry for unification.

The N = 4 super Yang-Mills theory has been widely discussed in the literature because of its novel feature of giving finite results when loop effects are taken into account. For instance it has been explicitly demonstrated that the gauge couplings are unrenormalised up to three loops and the general concensus is that N = 4 Yang-Mills supersymmetry will ultimately lead to a field theory that is free from the divergences that 'plague' the field theories of everyday life. As the name implies, all matter is in the adjoint representation of the internal symmetry group of one's taste. It was remarked in a previous publication that this is unsatisfactory from the point of view of neutral current phenomenology. Also if Nature is indeed N = 4 supersymmetric then it must be broken and the scale of supersymmetry will be high. This can be seen from the following argument: suppose we take the conventional assignment of electric charge for the fermions, then the bare value of the weak angle $\sin^2\theta_W$ is three-eighths (3/8) and this will not be renormalised as long as the theory is N = 4 supersymmetric. Phenomenologically $\sin^2\theta_W \simeq 0.21$ and this implies N = 4 supersymmetry must be broken to reproduce this value. For simplicity it will be assumed that N = 4 breaks to N = 0 as is implied by the supersymmetry algebra satisfied by the generators. The hierarchy, in this case is the same as in equation (1) and the $SU(3)_C$, $SU(2)_L$, U(1) gauge couplings are given by

$$g_3^{-2}(M_W) = g^{-2}(M_U^*) + (4f-33)K \ln M_1^*/M_W , \qquad (18)$$

$$g_2^{-2}(M_W) = g^{-2}(M_U^*) + (4f-22+\tfrac{n_D}{2})K \ln M_1^*/M_W , \qquad (19)$$

$$g_1^{-2}(M_W) = g^{-2}(M_U^*) + (4f+3n_D/10) \ln M_1^*/M_W . \qquad (20)$$

These are exactly the equations one deduces in theories without supersymmetry[8]. Due to the restrictive nature of the N = 4 theories regarding representations, there is the ambiguity in equations (17) - (19) of whether fermion families f and light doublets (n_D) will ever emerge. However, barring this problem, equations (17) - (19) give $\sin^2\theta_W \simeq 0.20$ and the mass-scale of N = 4 supersymmetry breaking $M_1^* = (10^{15} \sim 10^{16})$ GeV. As remarked the M_4^* of N = 4 supersymmetry has been interchanged with the unification mass-scale of the standard SU(5) grand unified theory. For the hierarchy to make sense, $M_U^* \geq M_4^*$. The alternative pattern of hierarchy

$$G^*(N=4) \xrightarrow{M_U^*} G^*_{SEW}(N=4) \xrightarrow{M_4^*} G^*_{SEW}(N=3) \xrightarrow{M_3^*} G^*_{SEW}(N=2) \xrightarrow{M_2^*} G^*_{SEW}(N=1) \xrightarrow{M_1^*} G_{SEW} \xrightarrow{M_W} SU(3)_c \times U(1)_{em} \qquad (21)$$

predicts the magnitudes of $\{M_3^*, M_2^*, M_1^*\}$ as in the N = 2 breaking (Equations (13) - (15) with $M_3^* \equiv M_U^*$). This is because N = 3 and N = 4 supersymmetries have identical gauge multiplets and these, as remarked earlier, do not renormalise the gauge couplings. Also $\{M_U^*, M_4^*\}$ are at least equal or greater than $\{M_3^*, M_2^*, M_1^*\}$.

Finally, although the scale(s) of supersymmetry breaking can lie between 10^2 GeV and 10^{19} GeV depending on the kind and amount of matter fields [tables I, II, III], cosmological ideas provide constraints on the allowed range of these scales as has been emphasised by Pagels and Primak[9] and more recently by Weinberg[10]. The point is that realising supersymmetry locally leads to the gravitino(s) to acquire mass(es) given by[11]

$$M_{g_{3/2}} = \frac{1}{2} \varkappa F^2 \tag{22}$$

and at the same time eliminates the undesirable vacuum energy through the cosmological constant. Here $= (16\pi G_N/3)^{\frac{1}{2}}$, G_N = Newton's constant and F denotes the scale of supersymmetry breaking which in our case is M_i^* (i = 1,,4). The upper bound in the lower end[F3] of the $(10^2 - 10^{19})$ GeV span foll[illegible] the gravitinos being light enough to be the sole contributors of the cosmic mass density. This gives $M_{g_{3/2}} < 1$ keV assuming that the number of degrees of freedom of thermally interacting particles at the time of gravitinos decoupling is of the order of 1000. From equation (22) it follows that $F = M_i^* < 10^6$ GeV. It is assumed that "R-symmetries" exist to guarantee a stable enough gravitino to survive to the present. The lower bound in the upper end of the $(10^2 - 10^{19})$ GeV span follows from considering the increase in the entropy density of the universe when gravitinos decay. This is amplified by $(G_N M_{g_{3/2}})^{-\frac{1}{2}}$, leading to helium formation far in excess of what is observed at present (∿25%). To remedy this the decay energy of the gravitinos should exceed the binding energy of helium (∿0.4 MeV) so as to disassociate it into neutrons and protons. The decay energy of the gravitinos is of order $(M^3_{g_{3/2}} G_N)^{\frac{1}{2}}$. From this and equation (22) the scale of supersymmetry breaking M_i^* satisfies $(1/8 \varkappa M_i^{*6} G_N)^{\frac{1}{2}} > 4 \times 10^{-4}$ leading to $M_i^* > 10^{11}$ GeV. This bound depends strongly on the assumption that helium nuclei are present at the time. Notice that if the gravitino mass is greater than one GeV then the bound on the scale of supersymmetry is two orders lower,

i.e. $> 2 \times 10^9$ GeV. Gravitinos this heavy would decouple before quarks synthesize to hadrons and hence fewer of them are present to contribute significantly to helium abundance. Only the very massive gravitini with mass $M_{g_{3/2}} \geq 10^4$ GeV ($F(= M_i^*) \geq 2.10^{11}$ GeV) "topple" the entropy balance. It is interesting to note that if the proton life time is of order 10^{30} years ($M_U^* \sim 10^{15}$ GeV) then the supersymmetry breaking mass-scales $\{M_1^*, M_2^*\}$ are high ($> 10^{10}$ GeV) favouring massive unstable gravitinos. In the N = 2 case there exists the possibility of a light gravitino ($F(\equiv M_1^*) < 10^6$ GeV) and a very massive gravitino ($F(\equiv M_2^*) > 10^{11}$ GeV) for proton lifetimes of order 10^{36} years.

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Footnotes

F1) Here it is assumed that the dominant mode of proton decay is through dimension six operators. The dimension four and five operators are suppressed by 'discrete' R invariance and heavy scalar triplets.

F2) At present all hopes of realizing this through dynamical supersymmetry breaking have faded away.

F3) An improved lower bound of (>50) GeV in the lower end has recently been deduced by M. Fukugita and N. Sakai, KEK preprint TH-42.

n_a \ M_U^*/M_W	$10^{13.5}$		10^{15}		10^{17}	
	$n_D = 0$ M_1^*/M_W	$n_D = 6$ M_1^*/M_W	$n_D = 0$ M_1^*/M_W	$n_D = 6$ M_1^*/M_W	$n_D = 0$ M_1^*/M_W	$n_D = 6$ M_1^*/M_W
$n_A = 1$	$10^{12.7}$	$10^{10.6}$	$10^{10.9}$	$10^{9.1}$	$10^{8.5}$	$10^{7.1}$
$n_A = 2$	10^{13}	$10^{11.5}$	$10^{12.4}$	$10^{11.0}$	$10^{11.7}$	$10^{10.4}$
$n_A = 3$	10^{13}	10^{12}	10^{13}	10^{12}	10^{13}	10^{12}
$n_A = 4$	$10^{13.3}$	10^{12}	$10^{13.2}$	$10^{12.3}$	10^{14}	10^{13}

TABLE I Variation of the mass-scale of supersymmetry breaking M_1^* with the unifying mass-scale M_U^* and the numbers (n_A, n_D) of scalar representations.

n_D \ M_1^*/M_W	10^7	10^9	10^{11}	10^{13}	10^{15}	10^{17}
0	0.203	0.203	0.203	0.203	0.203	0.203
2	0.207	0.208	0.210	0.212	0.202	0.213
4	0.211	0.213	0.217	0.219	0.221	0.223
6	0.215	0.218	0.224	0.226	0.230	0.234

TABLE II Weak angle predictions as the number of doublets and the supersymmetry breaking mass-scale M_1^* are varied.
[$\alpha_s = 0.12$, $\alpha_{em} = 1/128$]

M_U^*/M_W	$10^{13.5}$			10^{15}			10^{17}		
$n_A^{(2)}$ $(n_A^{(1)}=0)$	M_2^*/M_W	M_1^*/M_W $n_D=0$	M_1^*/M_W $n_D=6$	M_2^*/M_W	M_1^*/M_W $n_D=0$	M_1^*/M_W $n_D=6$	M_2^*/M_W	M_1^*/M_W $n_D=0$	M_1^*/M_W $n_D=6$
0	$10^{13.5}$	$10^{11.5}$	10^8	10^{15}	10^5	10^3	10^{11}	10^5	10^3
1	$10^{13.5}$	$10^{11.5}$	$10^{8.7}$	10^{15}	$10^{5.5}$	$10^{3.7}$	10^{14}	$10^{4.2}$	$10^{2.8}$
2	$10^{13.5}$	$10^{11.5}$	$10^{8.7}$	10^{15}	$10^{4.8}$	10^3	10^{15}	$10^{2.5}$	10^1
3	$10^{13.5}$	$10^{11.5}$	$10^{8.7}$	10^{15}	$10^{4.8}$	10^3	10^{15}	$10^{2.5}$	10^1
$n_A^{(2)}=n_A^{(1)}$									
1	$10^{13.5}$	$10^{12.7}$	$10^{10.6}$	$10^{13.5}$	$10^{11.5}$	$10^{9.3}$	10^{15}	$10^{9.3}$	$10^{7.7}$
2	$10^{13.5}$	10^{13}	$10^{11.6}$	10^{15}	$10^{12.5}$	$10^{11.1}$	10^{17}	$10^{11.7}$	$10^{10.4}$
3	-	$10^{13.1}$	10^{12}	-	$10^{13.1}$	10^{12}	-	$10^{13.1}$	10^{12}

TABLE III Predictions of the supersymmetry breaking mass-scales M_1^*, M_2^* for various values of the unifying mass-scale M_U^* and number of scalar adjoints ($\alpha_{em} = 1/128$, $\alpha_s = 0.12$, $0 \le n_D \le 6$).

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