CORNERING SUSY

Kari Enqvist and D.V. Nanopoulos

CERN - Geneva

and

A.B. Lahanas

University of Athens, Physics Department
Division of Elementary Particles
and Nuclear Physics
Panepistimiopolis, Athens 621
Greece

ABSTRACT

We confront two classes of supergravity models, the dimensional transmutation models and the no-scale models, with the UA1 monojet events and other experimental data. We perform a detailed search in the parameter space of these models with $m_{\tilde{q}} \approx 40-60$ GeV and $m_{\tilde{t}} = 30-50$ GeV and find large regions, with natural values of parameters, capable of providing a supersymmetric explanation of the monojets. The photino should have a mass between 5 and 8.5 GeV, and the right-handed slepton should not weigh more than 30 GeV.

1. Supersymmetry is a promising candidate for solving the hierarchy problem and other related matters. It predicts a plethora of new, exciting physics at the energy scales $O(10^2)$ GeV that may already be partly accessible to us. However, until recently experiments have yielded nothing that could be interpreted as an experimental manifestation of supersymmetry. The CERN SPS $p\bar{p}$ Collider has gathered very strong evidence for the Standard Model, and now some hints have appeared that supersymmetry may after all be lurking just round the corner. The reasons for such optimism are the so-called monojet (or ‘zen’) (i.e. single jet + missing energy) events observed by the UA1 Collaboration [1]. These very narrow jets are difficult to explain in the context of the Standard Model alone. For instance, if they were due to jet + ($Z^0 \rightarrow \nu\bar{\nu}$) one would expect to see also jet + ($Z^0 \rightarrow \ell^+\ell^-$) events as well as events with more jets. In supersymmetry such monojets can arise through an associated production of strongly interacting sparticles like squarks $\tilde{q}$ [2] or gluinos $\tilde{g}$ [2,3] with a mass in the range $O(40)$ GeV. Because $\tilde{q} \rightarrow q\tilde{\gamma}$ while $\tilde{g} \rightarrow q\tilde{q}\tilde{\gamma}$ the preferred interpretation [2] is the production of squarks rather than gluinos, as the former leads to a harder missing p_T distribution, to narrower monojets, and to fewer jets in the final state (typically two instead of four in the case of gluino pair production). This would imply that the gluino should not weigh much less than the squark. Otherwise they would be produced copiously in $p\bar{p}$ collisions, and the experimental signature would be radically altered.

2. In order to make contact with experiments one needs a supersymmetric model for calculating the spectrum of particles, and the most attractive ones are those based on spontaneously broken $N = 1$ supergravity, which lead to a highly restrictive phenomenologically acceptable breakdown of global supersymmetry [4]. Indeed, at the moment these are the only phenomenologically viable models, because both the spontaneous breaking of global supersymmetry and dynamical breakdown are very problematic [4]. The most interesting, and a rather restricted, class of $N = 1$ super-gravity models are those where the spontaneous breaking of $SU(2) \times U(1)$ is induced radiatively [5–7]. In the present paper we will confront such models with the possible supersymmetric explanation of the UA1 monojet events and other phenomenological constraints. First attempts in this direction have recently been made by Ellis and Sher [8], and we will comment on their results when appropriate.

The soft breaking terms at scales $O(M_p)$ can be written as [4]

$$\begin{aligned} \mathcal{L}_{\text{SOFT}} = & m^2 \sum_i |\phi_i|^2 + m \left[\sum_n (A - 3 + n) f_n(\phi_i) + \text{h.c.} \right] \\ & - \frac{1}{2} \sum_i [m_{\tilde{\nu}_i} \bar{\lambda}^i \lambda^i + \text{h.c.}] , \end{aligned} \quad (1)$$

where λ^i ($i = 1, 2, 3$) are gauginos and ϕ_i are the scalar fields (squarks, sleptons, and Higgses); $f_n(\phi_i)$ is the n^{th} order term of the superpotential $f(\phi_i) = \sum_n f_n(\phi_i) \equiv \sum_n c_n^i \phi_i^n$ with $n = 2, 3, \dots$ not in general terminating at $n = 3$ because of possible non-renormalizable terms in the effective Lagrangian (1). For our purposes we can neglect all light Yukawa couplings, which are inessential in the analysis of the renormalization group equations. Therefore we will consider only

$$f = h_t U_3^c Q_3 H_2 + m_4 H_1 H_2 , \quad (2)$$

where h_t is the top Yukawa coupling, U_3 is the superfield containing the top quark, and Q_3 is the corresponding weak doublet superfield. H_1 and H_2 are the two Higgs superfields necessary in supersymmetric theories, with the transformation properties $H_1 \sim H_2^\dagger$, and m_4 is a mass parameter naturally present in (2) which is also needed to avoid an unwanted Goldstone boson.

The soft SUSY breaking in the scalar sector is characterized by the mass parameter m , which may be related to the gravitino mass $m_{3/2}$ or to the common grand unified gaugino mass $m_{\tilde{V}}$. In the following we will parametrize the supersymmetry breaking due to gaugino masses by ξ_i , where

$$\xi_i = m_{\tilde{V}_i}/m \quad (i = 1, 2, 3) \quad (3)$$

The gaugino masses in supergravity are related to an arbitrary analytic function $f_{\alpha\beta}$ transforming as the symmetric product of two adjoint representations. In the simplest case $f_{\alpha\beta} \propto \delta_{\alpha\beta}$, grand unification would dictate that at M_X

$$\xi_1 = \xi_2 = \xi_3 \equiv \xi, \quad (4)$$

but in general this need not be true, and examples exist [9] where, despite grand unification, the gaugino masses are not equal at the superheavy mass scale because of the non-trivial form of the arbitrary function $f_{\alpha\beta}$. In the present paper we will concentrate mainly on the simplest case (4), but we will return to the more general situation in a forthcoming publication.

The soft SUSY breaking terms (1) contain also, using (2), trilinear and bilinear scalar interactions. We will denote in the following the strength of the bilinear piece by

$$B = A - 1. \quad (5)$$

The conditions (5) and (4) are expected to hold at the superheavy (grand unified) mass scale M_X . The general strategy [5 – 7] in radiative breaking of $SU(2) \times U(1)$ is to start with the set of the soft SUSY breaking terms given in (1) with certain boundary conditions on the parameters and to follow the evolution of the scalar masses down to low energies [10]. One then finds that the squark and the slepton masses get renormalized but stay positive, whereas the Higgs masses can become negative, thus signalling $SU(2) \times U(1)$ breakdown. In general, there are three different ways to induce this radiative breaking. One possibility is to have a ‘D-type’ breaking [5], which however leads to large sfermion and top-quark masses [$\gtrsim O(M_W)$]. In addition, because the breakdown is driven by D-terms, the ratio of the v.e.v.’s of the two Higgses goes to zero, $\langle H_1 \rangle / \langle H_2 \rangle = v_1/v_2 \rightarrow 0$. These properties are in conflict with the assumed SUSY explanation of the UA1 monojet events. This requires a spectrum of particles which is $O(40)$ GeV and experimentally it seems also that $v_1/v_2 \simeq 1$. Of course, the UA1 observation [11] of the top quark at mass $m_t = 40 \pm 10$ GeV is also in contradiction with the predictions of ‘D-type’ breaking [5]. Therefore we disregard this possibility and turn to the dimensional transmutation [6,7] and no-scale models [12].

In dimensional transmutation the scale of the $SU(2) \times U(1)$ breaking is determined dynamically through the renormalization group to be exponentially small with respect to the grand unified scale M_X . A detailed discussion on dimensional transmutation can be found in ref. 7. At $\mu = M_X$, one postulates the following boundary conditions [in addition to (4) and (5)] consistent with grand unification:

$$\begin{aligned} m_{\tilde{q}}^2 &= m_{\tilde{\ell}}^2 = m^2 \\ m_1^2 &= m_2^2 = m^2 + m_A^2, \end{aligned} \quad (6)$$

where $m_{\tilde{q}}$ and $m_{\tilde{\ell}}$ denote generic squark and slepton masses, respectively, and m_1 and m_2 are the masses of the Higgs doublets H_1 and H_2 . For a given value of the top mass m_t , $A(M_X)$, and ξ , the boundary conditions (4) to (6) determine the spectrum of sparticles at low energies through radiative corrections. In general one finds [6,7] that the sfermion masses are small ($\ll M_W$) as required, and that $\tan \theta \equiv v_2/v_1 \cong 1$, as indicated by experiments.

In the no-scale models [12] also the SUSY breaking scale m is determined dynamically to be $O(M_W)$. This takes place via the relation

$$m^2 = (M_W^2/\cos^2 \theta_W) f(\hat{m}_1^2, \hat{m}_2^2, \hat{m}_3^2) , \quad (7)$$

where f is a function of the dimensionless mass parameters $\hat{m}_k = m_k/m$ ($k = 1, 2, 3$) with $m_3^2 = m_{H_4}^2$ (for details see ref. 12a). This is an additional constraint with respect to the dimensional transmutation models, and therefore the no-scale models are even more restrictive than the dimensional transmutation models. Despite this, the no-scale models are somewhat more complicated, because now also the D-terms will contribute ($\theta \neq \pi/4$). Therefore, for instance, in the neutral fermion mass matrix which contains four states, the ‘axino’ does not decouple from the rest of the neutrals as it does in the dimensional transmutation models (see ref. 7). Because of this, some care must be taken with the definition of the photino.

In the simple case (4) the renormalization group equations can be solved to yield [7]

$$m_a^2 = m^2 (1 + C_a \xi^2) , \quad (8)$$

where $a = E, L, Q, U, D$ for all families of sleptons and for the 1st and 2nd family of squarks. The constant C_a is an integral, the value of which depends on the renormalization point which we have chosen to be M_W . We feel that this choice is appropriate for the $p\bar{p}$ Collider, where the mean c.m.s. energy of the colliding quarks is $O(M_W)$. In calculations we have taken three families and a minimal Higgs content. The stop masses also get sizeable contributions from the top quark, and among the three physical neutral Higgses one gets its mass radiatively. Therefore one can make the firm prediction that there should exist a Higgs with mass $m_H < O(10)$ GeV. In most cases, the mass of this Higgs is in the region of a few GeV and thus observable in $Y \rightarrow H\gamma$ decay. For details we refer to ref. 7*). We have tabulated the values of the constants C_a in Table 1.

The gaugino masses also get renormalized, and one finds for the gluino that

$$m_{\tilde{g}} = (\alpha_3/\alpha_{GUT}) \xi m = 3.0 \xi m , \quad (9)$$

where α_3 is the strong coupling constant and α_{GUT} is the grand unified coupling constant with the values $\alpha_3(M_W) = 0.122$ and $\alpha_{GUT} = 0.042$.

In the dimensional transmutation models the axino is decoupled from the rest of the neutrals because $\theta = \pi/4$, and for the photino mass one finds, in obvious notation,

$$m_{\tilde{\gamma}} = \frac{8}{3} \sin^2 \theta_W (\alpha_2/\alpha_{GUT}) m \xi + \dots \approx 0.46 \xi m , \quad (10)$$

*) Here we correct some misprints in table 2 of ref. 7b. In the expression for stop mass one should change $(m_Q^2 - m_A^2)^2$ to $(m_Q^2 - m_U^2)^2$, whereas in the expression for m_U^2 one should change $(C_Q - \frac{4}{3} C_L)\xi^2$ to $(C_Q - \frac{4}{3} C_U)\xi^2$ and finally in the expression for the radiative Higgs mass m_{rad}^2 , m_A^2 should read m_U^2 .

where the dots stand for additional contributions to the photino mass coming from the diagonalization of the 3×3 neutral sfermion mass matrix. In the ranges of interest these modifications contribute about $\frac{1}{2}$ GeV to the photino mass, thus accounting for roughly 10% of the total photino mass. The leading term is still that given by (10). In our subsequent calculations we have used the exact value for the photino mass. In fig. 1 we show it as a function of ξ for $m_{\tilde{q}} = 40, 50$, and 60 GeV. In the no-scale models there is some dependence on $A(M_X)$ in the photino mass, which we find however to be very slight; in all cases relevant to us we find this contribution to account for less than $O(\frac{1}{2})\%$ of the total photino mass.

3. Let us now turn to the experimental constraints on these models. The results from PETRA/PEP are telling us that there is no charged scalar particle with mass below about 18 GeV. The UA1 discovery [11] of the top quark at $m_t = 40 \pm 10$ GeV from the decay $t \rightarrow b\ell\nu$ indicates that the mode $t \rightarrow \tilde{t}\tilde{\gamma}$ is suppressed. This we will implement by requiring that $m_t \leq m_{\tilde{t}} + m_{\tilde{\gamma}}$. As mentioned before, we should forbid an abundant pair production of gluinos as compared to squarks or an excessive $\tilde{q} \rightarrow q\tilde{g}$ decay, and therefore we demand that $m_{\tilde{g}} \geq O(m_{\tilde{q}})$ and, more specifically, in the following we will make use of the bound $m_{\tilde{g}} \geq 0.85 m_{\tilde{q}}$. What is meant by abundant pair production of gluinos is, of course, sensitive to the experimental set-up and various cuts used in the data analysis. We feel that the limit $m_{\tilde{g}} \geq 0.85 m_{\tilde{q}}$ is illustrative, although our results do not depend critically on this limit, and should a more stringent limit be required, this can easily be translated into our results with the help of fig. 2. Also, it seems impossible to live [13] with stable particles carrying electric charge or colour, and thus we demand that $m_{\text{charged}} > m_{\text{neutral}}$.

Given the mass of the top, the free parameters of the dimensional transmutation model are $A_X \equiv A(M_X)$ and ξ , and the above experimental constraints can, in general, be satisfied in some domain of the (A_X, ξ) plane. However, apart from these experimental constraints there are other requirements that must also be satisfied. Naturally, the radiatively produced mass of one of the neutral higgses must be positive. Moreover, we must ascertain that the potential is bounded from below at the superheavy mass scale (this we call the stability condition). In fact, this stability of the vacuum at M_X is sufficient to guarantee a positive radiative mass. We must also ensure that there will be no other minima below the $SU(3)_c \times U(1)_{em}$ conserving one that break colour and/or electric charge or, if such minima exist, there must be a potential wall high enough so that the metastable $SU(3)_c \times U(1)_{em}$ minimum is very long lived (this condition we call the condition for a correct minimum). Finally, given the experimentally observed value of m_t , one can solve the value of the parameter $\hat{m}_4 = m_4/m$ as a function of ξ , $A(M_W)$, and $B(M_W)$, the latter being calculable from the renormalization group in terms of ξ and A_X because of the boundary condition (5). However, for $m_t = 40 \pm 10$ GeV there are regions in the (A_X, ξ) plane where no real solutions exist for m_4 , and accordingly these are disallowed.

We have performed a scan in A_X in the region $-5 \leq A_X \leq 5$ and in ξ in the region $-1 \leq \xi \leq 2$ to find the domain(s) where all the above constraints are satisfied in the dimensional transmutation models. In contrast to the analysis of Ellis and Sher [8] we have also considered negative values of A_X , where the situation is qualitatively completely different from that of the positive A_X as we will discuss shortly. Our result for the case $m_t = 30$ GeV, $m_{\tilde{q}} = 50$ GeV is shown in Fig. 3, where also the effects of the various relevant experimental and other constraints are illustrated. For positive A_X and ξ , we find an overall agreement with the

dimensional transmutation results of Ellis and Sher [8]. There are some small differences, which we attribute to the fact that they have chosen the renormalization point to be $\mu = M_W/2$ rather than $\mu = M_W$. We find also a large domain in the negative A_X , on which the constraints have a different origin from that of the case of positive A_X , as is indicated in Fig. 3. It is striking that the allowed domain with $A_X < 0$ satisfies all the UA1 experimental constraints automatically.

The reason why $A_X < 0$ gives rise to a completely different behaviour compared with the case $A_X > 0$ can be traced back to the form of the renormalization group equations. For instance, the boundary condition (5) on B will be numerically completely different if $A_X < 0$, which will affect, for example, the top mass, or conversely, the value of the parameter m_4 , which enters into the expressions for the masses of the radiative Higgs and the stop.

The allowed domain for positive A_X shrinks when the value of m_t is raised and has disappeared already when $m_t = 40$ GeV. This is in agreement with the corrected results of Ellis and Sher [8], who find $m_t \lesssim 37$ (34) GeV if $m_{\tilde{q}} \approx 50$ (45) GeV. In contrast, the domain for negative A_X is not very sensitive at all to the values of $m_{\tilde{q}}$ and m_t , as is shown in Fig. 4. There we display the allowed regions in the cases $m_t = 30$ to 50 GeV, $m_{\tilde{q}} = 40$ to 50 GeV. We have also checked that there is no considerable change in the shape of these regions if we choose $m_{\tilde{q}} = 60$ GeV. If the cuts in ξ due to the bounds $m_{\tilde{g}} \geq 0.85 m_{\tilde{q}}$ and $m_{\tilde{e}_R} \geq 18$ GeV must be made more stringent because of some future experimental information, this would shrink the allowed domains. The behaviour of $m_{\tilde{g}}/m_{\tilde{q}}$ and $m_{\tilde{e}_R}$ as a function of ξ has been shown in fig. 2, from where the corresponding bounds on ξ can be found and translated to figs. 3 and 4.

As can be seen from fig. 2, in practice the gluino cannot weigh much more than the squark but independently of the squark mass one has

$$m_{\tilde{g}}/m_{\tilde{q}} \lesssim 1.02 . \quad (11)$$

The parameter $\hat{m}_4 = m_4/m$ has different solution for $A_X > 0$ and $A_X < 0$. In the former case $\hat{m}_4$ varies roughly in the range (1.5, 2.5), whereas in the latter case $\hat{m}_4$ falls in the range $(-1.0, -0.1)$. It should be emphasized that these values are dictated by experiment, and therefore it is suggestive that they should turn out to be in a very natural region.

The behaviour of $\hat{m}_4$ as a function of A_X has been shown in fig. 5 for $m_t = 40$ GeV and is not very sensitive to the top mass. The ξ -dependence is approximately linear. As the supersymmetry-breaking parameter m is rather light in our scenario, m_4 will also be light, and thus the masses of the winos will be roughly equal and $O(M_W)$.

We should also like to point out that there exists an allowed region when $\xi < 0$. This is shown in fig. 6 for $m_t = 30$ GeV, $m_{\tilde{q}} = 40$ GeV and 50 GeV. For higher values of the top mass, the range of this region in A shrinks considerably but stays located at $A_X \approx -1$. As is evident from figs. 3, 4, and 6, the region $\xi < 0$ differs greatly from that where $\xi > 0$. This can be traced back to the fact that the renormalization group equations for A and B are linear in ξ and therefore sensitive to change of sign of ξ , and give rise to a completely different spectrum of sparticles. Theoretically there is no reason why ξ should be positive, and therefore we find it interesting that also $\xi < 0$ contains experimentally allowed regions.

If the simple assumption (4) does not hold but some or all ξ_i 's are different at the unification scale, the dimensional transmutation models have more freedom and some of the bounds are easier to satisfy. For instance, with $\xi_1 = \xi_2 \neq \xi_3$ it is possible to have a heavier photino than in

the case (4) and therefore, for example, the condition $m_t < m_{\tilde{\tau}} + m_{\tilde{\gamma}}$ is easier to fulfil. We will return to these questions in detail in a forthcoming publication.

Let us now turn to the no-scale models. We have performed a computer scan in the (A_X, ξ) plane to find the curve consistent with all relevant experimental and other constraints.

In the dimensional transmutation models [6,7] the whole mass spectrum, including the parameter m_4 , is easily calculated [7] up to the overall scale factor m for a given set of values of ξ , A_X , and m_t . The value of m can be fixed, again for a given ξ and A_X , from eq. (8) for a given value of $m_{\tilde{q}}$. The region allowed by the experiments dwells therefore in the four-dimensional space of the parameters ξ , A_X , $m_{\tilde{q}}$, and m_t .

The situation in the no-scale model is more complicated. For a given value of m_t , the function $r^2 = r^2(\hat{m}_4, A_W, B_W)$ as defined in eq. (52) of ref. 7, which in the dimensional transmutation models could be used to determine m_4 , is in the no-scale models related to m_t via the relation [12]

$$(m_t/M_W)^2 = (Nr^2/1 + 6r^2) 2 \sin^2 \theta , \quad (12)$$

which now involves also $\sin^2 \theta \neq 1/2$ unlike in the dimensional transmutation models. The value of the parameter N is given in table 1. Also, the function r^2 relates to the other parameters of the theory in a way different from the case of dimensional transmutation, because the critical scale which enters into the evaluation of the electroweak scale is defined through $m_1^2 m_2^2 = m_3^4$ rather than from $m_1^2 + m_2^2 = 2m_3^2$. All this amounts to a technical difficulty of extracting the value of m_4 (needed for the calculation of the full mass spectrum) that corresponds to a given value of m_t . We have resolved this problem by searching for points in the (ξ, A_X) plane that, for a given m_t , lead to a consistency with (12) and yield a value of the soft supersymmetry breaking parameter m from (7) that agrees with the one deduced from (8) for a given value of $m_{\tilde{q}}$. This method is not free from systematic errors, which we have however checked to be less than 5%. Now the allowed region is of course a three-dimensional subspace of the full four-dimensional parameter space of the dimensional transmutation models.

4. Our results are displayed in fig. 4, where we have superimposed a band corresponding to the no-scale models on the dimensional transmutation result. We present our results for the no-scale models in this form because, as stated before, our method of scanning contains some sources of systematic errors and moreover, there is also some uncertainty in the coupling constants entering into the renormalization group equations. Therefore we estimate that there is a possible overall error of about 10% in our results for the no-scale models.

A striking feature of the no-scale models is that we do not find any allowed curve in the region of positive A_X , and here we disagree with Ellis and Sher [8], who claim to have found a consistency with the no-scale models also for positive A_X . However, as can be seen from fig. 4, there is an allowed band of values in the (A_X, ξ) plane for negative values of A_X . In most cases we find $\cot \theta = 0.98-0.99$, very near to 1 but yet different from 1 so that there is some small contribution from D-terms also. Although in general the no-scale models cannot be said to be a special case of the dimensional transmutation models (there is no physically meaningful limit in which the former could be obtained from the latter); it is nevertheless encouraging that the no-scale curve appears to be in the same region as the allowed domain for the dimensional transmutation models.

In conclusion, we have investigated two supersymmetric models with a radiative breaking of $SU(2) \times U(1)$ (the dimensional transmutation and the no-scale models) taking into account several experimental and theoretical constraints with a special emphasis on the possible interpretation of the UA1 monojet events as production and subsequent decay of squarks^{*)}. We find that there indeed is an allowed region in the space of the parameters (A_X and ξ) in these models. However, although we find an allowed region for the dimensional transmutation models in positive A_X for light enough top mass (thus verifying the result of Ellis and Sher), we do not find any allowed region for the no-scale models if $A_X > 0$. In contrast, both models have allowed domains if $A_X < 0$, which region may therefore be considered to be as the more natural one. It is also encouraging that the experiments are not forcing us to consider unnatural values for the different parameters of these models, but that we find

$$\xi \simeq 0.5 - 1 \text{ [or } \xi \simeq -O(1)\text{]}, \quad |A(M_X)| \simeq O(3), \text{ and } |m_4/m| \simeq O(1) \text{ .}$$

As can be readily seen from figs. 1 and 2, these values indicate that the photino mass should fall in the range $5 \text{ GeV} \lesssim m_{\tilde{\gamma}} \lesssim 8.5 \text{ GeV}$ and that the right-handed selectron should be very light, $m_{\tilde{e}_R} \lesssim 30 \text{ GeV}$. As for Ellis and Sher [8], these values lie just outside the region forbidden by the MAC experiment, which may soon be able to give a decisive verdict for or against these models.

Acknowledgements

We wish to thank John Ellis and Marc Sher for helpful and illustrative discussions. One of us (K.E.) gratefully acknowledges the financial support of the Academy of Finland.

^{*)} We disregard here the possibility of light gluinos ($\sim 3\text{--}7 \text{ GeV}$) [14].

References

- [1] G. Arnison et al., (UA1 Collaboration), Phys. Lett. **139B** (1984) 115.
- [2] J. Ellis and H. Kowalski, Phys. Lett. **142B** (1984) 441; Nucl. Phys. **B246** (1984) 189; Preprint CERN-TH.4072 (1984).
- [3] E. Reya and D.P. Roy, Phys. Lett **141B** (1984) 442; Phys. Rev. Lett. **52** (1984) 881.
- [4] For recent reviews, see for example, H.P. Nilles, Phys. Rep. **C110** (1984) 2; D.V. Nanopoulos, Preprint-CERN TH-3995 (1984).
- [5] a) J. Ellis, D.V. Nanopoulos and K. Tamvakis, Phys. Lett. **121B** (1983) 123.
b) L. Alvarez-Gaumé, J. Polchinski and M.B. Wise, Nucl. Phys. **B221** (1983) 495.
c) L.E. Ibanez and C. Lopez, Phys. Lett. **126B** (1983) 54.
- [6] J. Ellis, J.S. Hagelin, D.V. Nanopoulos and K. Tamvakis, Phys. Lett. **125B** (1983) 275.
- [7] a) C. Kounnas, A.B. Lahanas, D.V. Nanopoulos and M. Quiros, Phys. Lett. **132B** (1983) 95.
b) C. Kounnas, A.B. Lahanas, D.V. Nanopoulos and M. Quiros, Nucl. Phys. **B236** (1984) 438.
- [8] J. Ellis and M. Sher, Phys. Lett. **148B** (1984) 309; Erratum to appear (1985).
- [9] J. Ellis, C. Kounnas and D.V. Nanopoulos, Nucl. Phys. **B247** (1984) 373.
- [10] K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, Prog. Theor. Phys. **68** (1982) 927.
- [11] G. Arnison et al. (UA1 Collaboration), Phys. Lett. **147B** (1984) 493.
- [12] a) J. Ellis, A.B. Lahanas, D.V. Nanopoulos and K. Tamvakis, Phys. Lett. **134B** (1984) 429.
b) J. Ellis, C. Kounnas and D.V. Nanopoulos, Nucl. Phys. **B241** (1984) 406 and Nucl. Phys. **B247** (1984) 373; Phys. Lett. **143B** (1984) 410.
c) J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys. Lett. **147B** (1984) 99; Preprint CERN-TH.4036 (1984).
- [13] J. Ellis, J.S. Hagelin, D.V. Nanopoulos, K. Olive and M. Srednicki, Nucl. Phys. Lett. **B238** (1984) 453.
- [14] M.J. Herrero et al., Phys. Lett. **132B** (1983) 199 and **145B** (1984) 430.
V. Barger et al., Phys. Rev. Lett. **53** (1984) 641 and Univ. Wisconsin preprint MAD/PH/197 (1984).

Table 1

The values of the integrals C_a entering into renormalization group equations of the sparticle masses, calculated at $\mu = M_W$. For a full elaboration of the renormalization group equations and these integrals, as well as the parameter N , see ref. 7.

C_Q	7.256
C_U	6.824
C_D	6.807
C_L	0.531
C_E	0.152
C_4	1.404
C_A	4.055
C_B	0.595
N	24.272

Figure captions

- Fig. 1 The photino mass $m_{\tilde{\gamma}}$ as a function of ξ for $m_{\tilde{q}} = 40, 50,$ and 60 GeV. $m_{\tilde{\gamma}}$ is not sensitive to the top mass in the range 30 – 50 GeV. The dotted line shows, for comparison, the naïve result $m_{\tilde{\gamma}} = \frac{8}{3} \sin^2 \theta_W (\alpha_2/\alpha_{\text{GUT}}) m_{\tilde{\nu}}$ [see eq. (10)].
- Fig. 2 $m_{\tilde{g}}/m_{\tilde{q}}$ as a function of ξ (dotted line) and the behaviour of $m_{\tilde{e}_R}$ as a function of ξ for $m_{\tilde{q}} = 40, 50,$ and 60 GeV.
- Fig. 3 The allowed regions in the case $m_t = 30$ GeV, $m_{\tilde{q}} = 50$ GeV. Also shown are the boundaries coming from various relevant constraints.
- Fig. 4 The allowed regions in the negative A_X for dimensional transmutation models and no-scale models (shaded bands), for a) $m_t = 30$ GeV; b) $m_t = 40$ GeV; c) $m_t = 50$ GeV. The regions are not very sensitive for $m_{\tilde{q}}$ in the range 40 – 60 GeV.
- Fig. 5 The $H_1 H_2$ mixing parameters $\hat{m}_4 = m_4/m$ as a function of A for $\xi = 0.4$ – 1.0 and $m_t = 40$ GeV. The dependence on ξ is approximately linear, and the result does not change much for top masses in the range 30 – 50 GeV.
- Fig. 6 The allowed regions in the negative ξ for a) $m_{\tilde{q}} = 40$ GeV and b) $m_{\tilde{q}} = 50$ GeV, both for $m_t = 30$ GeV.

Fig. 1

Fig. 2

Fig. 3

Fig. 4

(c)

Fig. 4

Fig. 5

Fig. 6