CERN-TH.4128/85

PROTON DECAY CONFRONTS LIGHT SPARTICLE SPECTRUM

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ABSTRACT

We confront proton decay due to supersymmetric dimension-five operators with some recently proposed supergravity models with light spectrum of sparticles. We argue that none of these models in their simplest forms is consistent with limits on proton lifetime, and that the most economical solution to the problem is to rely on non-minimal kinetic terms of the vector supermultiplet in supergravity.

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Proton decay¹⁾ is usually considered as a probe of the possible underlying grand unified theory only. For ordinary non-supersymmetric GUTs, where the $d = 6$ $\Delta B \neq 0$ operator depends only on the masses of some superheavy Higgses or gauge bosons and on factors appearing because of scaling from $\mu = M_X$ down to $\mu \approx 0(1 \text{ GeV})$, this certainly is true. In supersymmetric GUTs, there exist also $d = 5$, $\Delta B \neq 0$ operators²⁾, which can be dressed by gaugino exchange to a $d = 6$ operator mediating nucleon decay (see Fig. 1). In a wide class of models, the leading contribution is expected to come from wino exchange^{3),4)}. As the dressing of the supersymmetric $d = 5$ operator involves supersymmetry breaking mass insertion on the wino line, the nucleon decay amplitude measures therefore⁵⁾ also some aspects of the low energy sparticle spectrum. Of course, because a loop of squarks and sleptons is also involved, proton decay due to supersymmetric $d = 5$, $\Delta B \neq 0$ operators will depend on the masses of these particles⁵⁾. This is of particular interest to the attempts⁶⁾⁻⁸⁾ to present a supersymmetric explanation of the UA1 monojets⁹⁾. In these phenomenological explanations, the spectrum of sparticles is more or less fixed. This is even more true in the models¹⁰⁾⁻¹²⁾ based on supergravity that have been used to obtain the phenomenologically desired spectra. These models employ the radiative breaking of $SU(2)_L \times U(1)$, and the mass spectrum of sparticles can be calculated in terms of few parameters restricted to a certain, usually narrow, range. Typically one finds¹⁰⁾⁻¹²⁾ that all sparticle masses are light ($O(M_W)$) and that the squark masses $m_{\tilde{q}} \lesssim 120 \text{ GeV}$.

However, at first glance any connection of the low energy sparticle spectra with proton decay may seem completely spurious, because the $d = 5$, $\Delta B \neq 0$ operators will naturally depend on the unknown coloured Higgs triplet mass m_{H_3} . In the present paper, we wish to point out that in fact there exists a connection which may already rule out some of the proposed models in their simplest forms.

It has been known for long³⁾⁻⁵⁾ that the $d = 5$ operators in simple SUSY GUTs like $SU(5)$ are potentially dangerous. The effective Lagrangian in minimal SUSY $SU(5)$ can be written^{5),13)}, assuming that the VEVs of the two 5-plets H and $\bar{H}$ are $\langle H \rangle \approx \langle \bar{H} \rangle$ as is true for dimensional transmutation models¹⁴⁾ and needed¹⁰⁾⁻¹²⁾ for the explanation of the monojet events, as

$$\mathcal{L} = - \frac{g_2^2}{16\pi} \frac{A}{m_{H_3}} K O_6 \quad (1)$$

where A is an enhancement factor, O_6 is a four-fermion operator, the precise form of which is not important here, and K is a kinematic function arising essentially from the loop integration in Fig. 1. We follow here closely the notation of Campbell et al.⁵⁾. K is given by:

$$K = m_+ M_W \cos^2 \theta \left[f(m_{\tilde{u}}, m_{\tilde{d}}, m_+) + f(m_{\tilde{u}}, m_{\tilde{e}}, m_+) \right] \\ + m_- M_W \sin^2 \theta \left[f(m_{\tilde{u}}, m_{\tilde{d}}, m_-) + f(m_{\tilde{u}}, m_{\tilde{e}}, m_-) \right] \quad (2)$$

where $m_{\tilde{u}}$ and $m_{\tilde{d}}$ are the left-handed squark masses for the first two generations, and $m_{\tilde{e}}$ is the mass of the left-handed slepton. [We assume that O_6 is a LLLL operator and does not contain³⁾ RRRR pieces^{*)}; we also disregard the effects of the third generation, as these are probably small.] The parameters m_+ and m_- are the charged wino mass eigenvalues, obtained by diagonalizing the mass matrix

$$m = \begin{pmatrix} M_2 & M_W \\ M_W & -\epsilon \end{pmatrix} \quad (3)$$

and θ is the corresponding mixing angle with $\sin\theta \cos\theta = M_W/(M_2 + \epsilon)$. In (3) M_2 is the direct gaugino mass and ϵ is the $H\bar{H}$ -mixing term. Only these two terms break SUSY and therefore proton decay amplitude will be proportional to them. The function f is given by

$$f(m_1, m_2, m_{\pm}) = \frac{1}{m_{\pm}^2} \frac{1}{\frac{x_{\pm}}{r} - 1} \left[\frac{1}{r-1} \ln r - \frac{1}{x_{\pm}-1} \ln x_{\pm} \right] \quad (4)$$

where

$$r \equiv m_1^2 / m_2^2 \quad ; \quad x_{\pm} = m_1^2 / m_{\pm}^2 \quad (5)$$

The dominant nucleon decay channel due to the $d = 5$, $\Delta B \neq 0$ operators is³⁾⁻⁵⁾ $p \rightarrow \bar{\nu} K^+$ and $n \rightarrow \bar{\nu} K^0$. It has been argued, however, that in some cases also $\bar{\nu} \pi$ modes might be the dominant ones. For reasons that we will discuss later, we do

^{*)}This is justified in dimensional transmutation models¹⁴⁾ where the L-R splitting of squarks and sleptons is small.

not think that this is going to happen in dimensional transmutation models¹⁴⁾. The nucleon lifetime for the νK -channel has been estimated in Ref. 5), where it was found that

$$\tau(n \rightarrow \nu K) = 3.6 \times 10^{26} \text{ yr } R^2 / K^2 \quad (6)$$

where the R-factor includes the Higgs triplet mass and the relevant matrix element $\beta = \langle 0 | (du)_L u_L | p \rangle$ and is given by

$$R = \frac{m_{H_3}}{10^{16} \text{ GeV}} \frac{0.03 \text{ GeV}^3}{\beta} \quad (7)$$

There are some uncertainties in the evaluation¹⁵⁾ of the three-quark overlap β . We take a reasonable range to be $0.003 \text{ GeV}^3 < \beta < 0.03 \text{ GeV}^3$. Then, with the recent IMB bound¹⁶⁾ $\tau(n \rightarrow \nu K) \gtrsim 1 \times 10^{31}$ years, and adopting an intermediate value of $\beta = 0.01$, one finds that

$$K \lesssim 2 \times 10^{-4} \frac{m_{H_3}}{10^{16} \text{ GeV}} \quad (8)$$

In Fig. 2 we display the bound (8) for $m_{H_3} = 10^{16} \text{ GeV}$ and 10^{17} GeV in the (M_2, ε) plane for $\varepsilon > 0$ in two cases: (i) $m_{\tilde{q}} \approx m_{\tilde{\chi}} = 100 \text{ GeV}$; this is the case relevant for the light gluino explanation⁸⁾ as well as the two-gluino explanation⁷⁾ of the UA1 monojets; (ii) $m_{\tilde{q}} \approx 50 \text{ GeV}$, $m_{\tilde{\chi}} = 30 \text{ GeV}$; this is the case relevant for the two-squark explanation⁶⁾ of the monojets. Figure 3 shows the same bounds for $\varepsilon < 0$.

Before we consider the implications of (8) for existing supersymmetric models, let us discuss the value of m_{H_3} .

In simple SU(5) there are two possible mechanisms for achieving the doublet-triplet splitting^{*)}. The first¹⁸⁾ uses artificial fine-tuning in the superpotential

*) We disregard the sliding singlet mechanism as it is not stable under radiative corrections¹⁷⁾.

$$\mathcal{P}(\Sigma, H, \bar{H}) = \lambda \bar{H}^\alpha (\Sigma_\alpha^\beta + \frac{3}{2} m \delta_\alpha^\beta) H_\beta \quad (9)$$

$[\alpha, \beta = 1, \dots, 5 \text{ are } SU(5) \text{ indices}]$, by choosing $m = \langle \Sigma \rangle + O(M_W)$. The mass of the Higgs triplets is then $m_{H_3} = 5/2 \lambda \langle \Sigma \rangle$, whereas the mass of the superheavy gauge bosons is $M_X = 5g_{GUT} \langle \Sigma \rangle / \sqrt{8}$. Therefore

$$\frac{m_{H_3}}{M_X} = \sqrt{2} \frac{\lambda}{g_{GUT}} \quad (10)$$

We demand that the four-point coupling $\lambda^2 < 1$ at scales relevant to proton decay in order that the theory remains perturbative at scales $\mu \approx M_X$ (this makes sense especially if $M_X \ll M_P \approx 1.2 \times 10^{19} \text{ GeV}$). In the minimal SUSY $SU(5)$ $\alpha_{GUT} \approx 1/25$; adding more coloured light particles will result in slower evolution of $\alpha_3(\mu)$, and adding more $SU(2)_L \times U(1)$ non-singlet light particles will result, because the one-loop beta-function coefficients b_2 and b_1 are already negative, in faster descent of $\alpha_2(\mu)$ or $\alpha_1(\mu)$. In both cases, the value of α_{GUT} will be larger than in the minimal case. Therefore, we set $\alpha_{GUT} \gtrsim 1/25$, and from (10) we read

$$m_{H_3} \lesssim 2 M_X \quad (11)$$

The missing-partner mechanism¹⁹⁾ splits doublets from triplets without fine-tuning by introducing an extra pair $\bar{\theta} + \theta \sim \bar{5}_0 + 5_0$ with the following superpotential

$$\mathcal{P}(H, \bar{H}, \theta, \bar{\theta}, \Sigma) = \frac{a_1}{M_P} H \theta \Sigma^2 + \frac{a_2}{M_P} \bar{H} \bar{\theta} \Sigma^2 + b M_P \theta \bar{\theta} \quad (12)$$

The fermionic Higgs triplet mass matrix is then

$$m_3 = \begin{pmatrix} 0 & a_1 \langle \Sigma \rangle^2 / M_P \\ a_2 \langle \Sigma \rangle^2 / M_P & b M_P \end{pmatrix} \quad (13)$$

Let us focus on the case $\langle \Sigma \rangle \sim M_X \ll M_P$. If $b \approx O(1)$ and $a_{1,2} \lesssim 1$, there will be one light eigenstate $m \approx a_1 a_2 \Sigma^4 / M_P^3 \ll \langle \Sigma \rangle$. With $\langle \Sigma \rangle \sim 10^{16} \text{ GeV}$, one would find

$m \lesssim 10^7$ GeV. If $b \ll a_1 a_2 \langle \Sigma \rangle^4 / M_P^4$ one finds $m \sim \sqrt{a_1 a_2} \langle \Sigma \rangle^2 / M_P$. If $\langle \Sigma \rangle \approx 10^{16}$ GeV, $m \gtrsim \langle \Sigma \rangle$ if $\sqrt{a_1 a_2} \approx 10^3$, a highly unlikely situation. Therefore, we conclude that in the case of missing partner mechanism, the coloured Higgs triplets are always lighter than the unification mass M_X .

In calculations of low energy sparticle spectra, one usually assumes unification which then implies equal soft SUSY breaking parameters for squarks and sleptons at M_X . Furthermore, to evaluate the sparticle spectrum one needs also to fix the light particle content of the theory. This then fixes uniquely the unification mass. For instance, in the minimal SUSY standard model with three generations and two Higgs doublets, the unification takes place at $M_X = 6 \times 10^{16} \times \frac{\Lambda_{\overline{MS}}}{10^{16} \text{ GeV}}^{3),20)$. Therefore, in these simplest models there is also an upper bound on the Higgs triplet mass, and as models explaining the UA1 monojets⁹⁾ also fix¹⁰⁾⁻¹²⁾ the parameters M_2 and ϵ , we are now ready to confront these models with proton decay constraint, as presented in Figs. 2 and 3.

Let us first consider the monojet explanation proposed by Ellis and Kowalski⁶⁾ with $m_{\tilde{q}} \approx 50$ GeV, $m_{\tilde{g}} \gtrsim m_{\tilde{q}}$ and $m_{\tilde{\chi}} \approx 30$ GeV. A typical value of the SUSY parameters A_X and ξ for $\epsilon > 0$ [in the notation of Ref. 11), $-m_4$] would be $A_X \approx -2.5$, $\xi \approx 0.7$ for which values $\epsilon \approx 7.7$ GeV (we take $m_t \approx 40$ GeV). The value of M_2 is fixed by the gluino mass to be $M_2 = (\alpha_2/\alpha_3)m_{\tilde{g}} \approx 15$ GeV. If we adopt, as seems natural to us, the bound corresponding to $M_{H_3} = 10^{16}$ GeV, we find from Fig. 3a that $M_2 \lesssim 1.5$ GeV and therefore this case is excluded. As can be seen from Fig. 3b, if we were able to raise unification mass by one order of magnitude, we would be able to accommodate this case.

If $\epsilon < 0$, this implies that $A_X > 0$, and for $m_{\tilde{q}} = 50$ GeV, $m_t = 40$ GeV there is no region compatible with the UA1 monojets. However, for $m_{\tilde{q}} = 60$ GeV such a region exists with $\xi \approx 1$ and $-\epsilon \approx 44$ GeV. From Fig. 2a, we find that for $\epsilon < 0$ and $M_2 \approx 18$ GeV, proton decay would imply $18 \text{ GeV} \lesssim -\epsilon \lesssim 25 \text{ GeV}$ if we demand that $m_{H_3} < 10^{16}$ GeV. Again the discrepancy is such that a raise in m_{H_3} of one order of magnitude would resolve it.

Let us now turn to models⁸⁾ which have a light gluino, $m_{\tilde{g}} \approx 5$ GeV, and heavy squarks and sleptons with $m_{\tilde{q}} \approx m_{\tilde{\chi}} \approx 100$ GeV. The model of Reya and Roy⁷⁾ with $m_{\tilde{g}} \approx 40$ GeV has also $O(100)$ GeV squarks and sleptons. If $m_{\tilde{g}} \approx 5$ GeV, one finds

$M_2 \approx 1.5$ GeV. In the examples of Table 2 of Ref. 12) for these models, one can find that if $\epsilon > 0$ [in the notation of Ref. 12), $-\mu/q$] and if $m_t > m_t + m_{\tilde{\gamma}}$ (so that the top does not decay into stop) $\epsilon \gtrsim 15$ GeV. In models with heavy squarks there is some propagator suppression of proton decay amplitudes so that the limits are not so stringent in this case. However, from Fig. 2a, we find that despite this, this case is also excluded for $m_{H_3} \lesssim 10^{16}$ GeV which would imply $\epsilon \lesssim 3.7$ GeV. It is also easy to see that for $\epsilon > 0$ too the case of $m_{\tilde{g}} \approx 40$ GeV is excluded. Therefore, we conclude that none of the present models as stated in the literature are consistent with proton decay limits. Let us again emphasize that these results were obtained by taking a generous value on the three-quark overlap β . Relaxing this limit would lead to even more strict constraints.

It has been pointed out by Arnowitt, Chamseddine and Nath²¹⁾ that in some cases it is possible to have destructive interference from zino or gluino dressing and left-right mixing. However, to enhance the L-R mixing one would need large gravitino masses typically of the order of $m_{3/2} \approx Am_t$. For zino contribution to be of interest, one would need large zino mass splitting. None of the above possibilities is easily obtained in radiative models. Finally, sizeable gluino dressing would imply large $\tilde{u}\tilde{d}$ -splitting which does not happen in radiative models. Although these possibilities are not totally excluded, we see no reason to believe that they can affect our result significantly.

In conclusion, we have shown that even taking the most generous value for β , $\beta = 0.01$ GeV³, the constraint on proton stability is not satisfied in the simplest form of the low energy supergravity models which have been proposed so far. How can one then reconcile the simple and predictive supergravity models with light sparticle spectrum with the proton decay constraint? One should be especially careful with solutions that introduce new light ($< M_X$) matter into the model because this will affect the renormalization group. For instance, one may introduce symmetries that forbid $d = 5$ operators but then cosmological baryogenesis necessitates²²⁾ also the introduction of new fields, some of which may remain light. Similarly, if one increases the light field content (e.g., new Higgs doublets) to increase M_X , this will have drastic effect on spectrum because not only renormalization group equations but also boundary conditions will change. In all these cases the spectrum of sparticles should be recalculated.

It seems that the most economical improvement, not involving new fields, is to consider supergravity models with non-minimal kinetic terms for the vector supermultiplet ($f_{\alpha\beta} \neq \delta_{\alpha\beta}$)²³⁾. Such theories have recently been analyzed in

Ref. 24) where it was found that the unification mass in the case of the minimal SUSY standard model is in the range $2.4 \times 10^{16} \text{ GeV} \lesssim M_X \leq 3.5 \times 10^{17} \text{ GeV}$ and therefore is not in contradiction with proton decay data which we found to imply $M_X \gtrsim 10^{17} \text{ GeV}$. It should be stressed, however, that in theories with non-minimal kinetic terms the gaugino mass ratios do not obey the simple formula $m_3/\alpha_3 : m_2/\alpha_2 : m_1/\alpha_1 = 1:1:1$. In the simplest SU(5) model, one finds²⁴⁾ rather the ratio 5:5:1. The physical photino would be only half as heavy as in the naive case. Although not problematic for monojet explanations with heavy gluinos ($m_{\tilde{g}} \gtrsim 40 \text{ GeV}$), this would spell trouble for the light gluino ($m_{\tilde{g}} \approx 5 \text{ GeV}$) explanation⁸⁾, where one would then find that $m_{\tilde{\gamma}} \approx 300 \text{ MeV}$. Such a very light photino is excluded because of cosmological considerations²⁵⁾.

It is remarkable that proton longevity¹⁶⁾ and the UAl monojet events⁹⁾, two seemingly completely unrelated phenomena, when combined together put such tight constraints on light s-spectra.

ACKNOWLEDGEMENTS

We thank R. Arnowitt and J. Ellis for useful discussions. One of us (K.E.) thanks the Academy of Finland for financial support.

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FIGURE CAPTIONS

Fig. 1: Supersymmetric $d = 5$, $\Delta B \neq 0$ operator dressed by gaugino exchange to a $d = 6$ operator capable of mediating nucleon decay.

Fig. 2: The domain in (M_2, ϵ) -plane [see (3)] not contradicting proton decay data for $\epsilon > 0$, for $m_{\tilde{q}} = m_{\tilde{\chi}} = 100$ GeV (below the unbroken line) and for $m_{\tilde{q}} = 50$ GeV, $m_{\tilde{\chi}} = 30$ GeV (below the dashed line); (a) $m_{H_3} < 10^{16}$ GeV, (b) $m_{H_3} < 10^{17}$ GeV.

Fig. 3: The allowed domain for $\epsilon < 0$ and for $m_{H_3} < 10^{16}$ GeV (between the dashed lines) and $m_{H_3} < 10^{17}$ GeV (between the unbroken lines); (a) $m_{\tilde{q}} = 50$ GeV, $m_{\tilde{\chi}} = 30$ GeV; (b) $m_{\tilde{q}} = m_{\tilde{\chi}} = 100$ GeV.

Fig. 1

Fig. 2a

Fig. 2b

Fig. 3a

Fig. 3b