CERN-TH.4496/86

GAUGINO MASSES FROM RADIATIVE CORRECTIONS IN
SUPERSTRING MODELS

John Ellis, D.V. Nanopoulos, M. Quirós^{*)} and F. Zwirner⁺⁾
CERN -- Geneva

ABSTRACT

We discuss contributions to the observable gaugino masses due to one-loop radiative corrections in the effective theory below the compactification and condensation scales of superstring models. We assume that some unspecified dynamics determines the compactification radius, and that gaugino condensation in the hidden sector generates a gravitino mass smaller than the Planck mass. The largest contribution to the gaugino mass arises from radiative corrections to the effective scalar potential, which generate $m_{1/2} \sim m_{3/2}/m_P^{2/3}$. Scalar masses appear in higher orders of perturbation theory. We comment on the implications for phenomenological model calculations.

^{*)} On leave from Instituto de Estructura de la Materia, Madrid, Spain.

⁺⁾ Also at International School for Advanced Studies, Trieste, and Istituto Nazionale di Fisica Nucleare, Padova, Italy.

One of the key problems in superstring phenomenology¹⁾ is to understand how supersymmetry is broken to give a gravitino mass, and how this breaking of supersymmetry is subsequently communicated to the observable sector. Another key problem is to understand how the expectation values of gauge-singlet "dilaton" and "axion" fields (S, T) ²⁾ are determined. A promising suggestion for solving parts of these problems is that hidden sector gauginos χ condense^{3), 4)} after compactification, thereby determining dynamically the expectation value of S ⁴⁾ and generating a non-zero gravitino mass. However, this mechanism does not determine the value of the scalar field T which fixes the size of the compactified dimensions: $R_c \sim T_R^{\frac{1}{2}}$, and also the mechanism for generating non-zero gaugino masses $m_{\frac{1}{2}}$ and scalar masses m_0 in the observable sector remains obscure. In certain models for compactification⁵⁾, non-perturbative world-sheet instanton effects can generate a non-trivial effective potential for the T field⁶⁾, and might drive it towards a preferred value. However, these models have theoretical and phenomenological defects⁷⁾ and we will not discuss them further here. We will just assume that some unspecified dynamics fixes $\langle 0|T|0\rangle = O(1)$ ⁸⁾, so that the compactification scale is comparable to the string excitation scale. This paper is concerned with supersymmetry breaking and its communication to the observable sector.

It has been argued previously⁹⁾ that hidden sector gauginos can only condense: $\langle 0|\bar{\chi}\chi|0\rangle \neq 0$ if the hidden sector gauge group E_8 is broken by the Hosotani mechanism to some $SU(N)$ subgroup. In this case, the scale Λ_c at which the hidden sector gauge interactions become strong is necessarily somewhat smaller than the compactification scale, and the gravitino mass $m_{3/2} \sim \langle 0|\bar{\chi}\chi|0\rangle/m_P^2 \sim \Lambda_c^3/m_P^2 \ll m_P$. If there is no inflation, or if this non-perturbative effect occurs after inflation, the Universe contains $SU(N)$ flux tubes which form cosmic strings that cause density perturbations $\delta\rho/\rho \sim 10^2(\Lambda_c^2/m_P^2)$ and could seed galaxy formation^{10), 11)}. The resulting perturbations are acceptably small¹¹⁾ if $N \leq 4$ so that $\Lambda_c \lesssim 10^{16}$ GeV and $m_{3/2} \lesssim 10^{12}$ GeV. This is further motivation for our assumption in the rest of this paper that $m_{3/2} \ll m_P$.

Supersymmetry breaking in the observable sector is absent at the tree level, but is presumably generated from $m_{3/2}$ by radiative corrections, and $m_{\frac{1}{2}}, m_0$ could be much smaller than $m_{3/2}$. One type of one-loop contribution¹²⁾ to $m_{\frac{1}{2}}$ would be $O(\langle 0|N|0\rangle \times m_{3/2}/m_P)$, where N is some generic gauge non-singlet field. This is negligible if $\langle 0|N|0\rangle$ is not much larger than m_W , as we believe⁷⁾, and $m_{3/2} \ll m_P$ as argued above. Subsequently, a possible two-loop contribution to m_0 has been

identified¹³⁾. We argue below that it yields $m_0 \sim m_{3/2}^2/m_P$ and is probably too small to be the dominant source of supersymmetry breaking in the observable sector. More recently, possible contributions to $m_{1/2}$ which are of order $m_{3/2}^3/m_P^2$ have been calculated¹⁴⁾. We will comment later on these contributions, which are also unlikely to be dominant^{*}).

The main purpose of this paper is to present another contribution to $m_{1/2}$ which is of order $m_{3/2}^{5/3}/m_P^{2/3}$, and hence likely to be the dominant source of supersymmetry breaking in the observable sector. It arises because one-loop radiative corrections to the effective scalar potential shift the non-perturbatively determined value of $\langle 0|S|0\rangle$, thereby enabling the S field to contribute to $m_{1/2}$, which it did not do at the tree level. Previous phenomenological model calculations¹⁷⁾ of low-energy spectra, which were made assuming that observable sector supersymmetry breaking was seeded by $m_{1/2}$, are not substantially modified if $m_{1/2}$ is in fact generated by the mechanism described here. This point is discussed at the end of our paper, together with comments on the relations of our calculations to previous ones, some cosmological remarks and some reminders of open problems.

At energies below the condensation scale Λ_c the tree-level potential for the minimal no-scale¹⁸⁾ supergravity theory obtained from the superstring is given by

$$V_0 = e^{\mathcal{G}} \left(\mathcal{G}^i (\mathcal{G}^{i-1})^j \mathcal{G}_j - 3 \right) + \frac{1}{2} D^\alpha D_\alpha \quad (1)$$

where $\mathcal{G}$ is the Kähler potential²⁾:

$$\mathcal{G} = -\ln(S+S^\dagger) - 3\ln(T+T^\dagger - 2\phi^l \phi_l^\dagger) + \ln|f(\phi) + W(S)|^2 \quad (2)$$

with $f(\phi) \equiv f_{\lambda mn} \phi^\lambda \phi^m \phi^n$ the superpotential for matter fields and⁴⁾

$$W(S) = c + h \exp(-\lambda S) : \lambda = \frac{8\pi^2}{C_2(G)} \quad (3)$$

is the effective superpotential induced by non-perturbative effects in the hidden sector. The second term in $W(S)$ (3) results from G' gaugino condensation, where

^{*}) We have nothing to add to previous comments¹⁵⁾ on the interpretation of Wess-Zumino terms, which some authors¹⁶⁾ had thought might also contribute to $m_{1/2}$ and m_0 .

G' is the biggest simple component of the subgroup into which E_8' is broken during compactification. The constant term in (3) may result from a vacuum expectation value of the antisymmetric tensor field $H_{\mu\nu\rho}$.

The vacuum expectation values of the matter fields are determined at the tree level by

$$g^a (g^{\prime\prime-1})^b_g g_b = 3 : a, b = T, \phi^L \quad (4)$$

which leads to $\langle 0 | \phi^L | 0 \rangle = 0$ but leaves $\langle 0 | T | 0 \rangle$ undetermined. As mentioned earlier, we assume that T_R is fixed by some other mechanism as yet unspecified. Using (4), the tree-level potential can be written as

$$V_0 = e^g e^s (g^{\prime\prime-1})^s_s g_s \quad (5)$$

which is positive semi-definite. Its vacuum state is therefore given by $g_s = 0$, which occurs when

$$\cos \lambda S_I = -1 \quad (6a)$$

$$(1 + 2z_0) e^{-z_0} = c/h \quad (6b)$$

where $z \equiv \lambda S_R$ and we have denoted $S = S_R + iS_I$. The vanishing of g_s translates into a vanishing gaugino mass, since

$$m_{Y_2} = \frac{1}{2S_R} e^{g/2} g_s (g^{\prime\prime-1})^s_s \quad (7)$$

For this reason, although local supersymmetry is spontaneously broken in the hidden sector, at the tree level this breaking does not communicate to the observable sector. Moreover, other sources of supersymmetry breaking including scalar masses and trilinear couplings are known to vanish at the tree¹⁸⁾ and one-loop levels¹⁹⁾.

However, when radiative corrections to the effective scalar potential are considered, the tree level vacuum expectation value (6) will be corrected by a very small amount $\Delta z = O(\hbar) \ll 1$. This smallness is in agreement with the consistency of the perturbation expansion, which leads us to expect small shifts in the tree-level vacuum state unless the tree-level potential is flat. Although small, the one-loop correction $\Delta z = \lambda \Delta S_R$ has dramatic consequences. The shifted value of g_s no longer vanishes, and is a source of supersymmetry breaking in the

observable sector at the one-loop level^{*)}.

The one-loop correction to the effective scalar potential, which arises from supersymmetry breaking below the condensation scale Λ_c , is in general given by^{13),20)}

$$V_1 = \frac{1}{64\pi^2} \left\{ \Lambda_c^2 \text{Str } M^2 + \text{Str} (M^4 \ln M^2 / \Lambda_c^2) + \Lambda_c^4 \text{Str} \ln (1 + M^2 / \Lambda_c^2) - \text{Str} (M^4 \ln (1 + M^2 / \Lambda_c^2)) \right\} \quad (8)$$

where the condensation scale is

$$\Lambda_c = \left(\frac{\lambda}{T_R} \frac{1}{z} \right)^{1/2} \exp(-1/3 z) \quad (9)$$

All masses in (8) are to be taken at the tree level, and so are proportional to the gravitino mass:

$$m_{3/2} = e^{g/2} = \left(\frac{\lambda}{16 T_R^3 z} \right)^{1/2} (c - h e^{-z}) \quad (10)$$

where we put S_I equal to its vacuum expectation value (6a). Therefore, if the tree-level vacuum expectation value (6b) is at $z_0 \gg 1$ (i.e., $m_{3/2} \ll 1$), then Λ_c (9) is also $\ll 1$, while $(m_{3/2} / \Lambda_c)^2 \ll 1$ and so $M^2 / \Lambda_c^2 \ll 1$ in Eq. (8). In this limit we can approximate (8) by²⁰⁾

$$V_1 = \frac{1}{32\pi^2} \Lambda_c^2 \text{Str } M^2 \quad (11)$$

The general expression for the supertrace in (11) using the vacuum (4) for the matter fields, is^{19),21)}

^{*)} This effect was considered in Ref. 12), but discarded because there the dynamical hypothesis was made that radiative corrections to the effective potential were $O(m_W^4)$.

$$St M^2 = 2m_{3/2}^2 (-1 - k \mathcal{G}^s (\mathcal{G}^{"-1})^s \mathcal{G}_s) \quad (12)$$

where k is a constant which in principle depends on the number of chiral fields and on the number of gauginos. Our final result does not depend on the value of k . Thus the one-loop effective scalar potential for z can be cast as

$$V_{eff}(z) = V_0(z) + V_1(z) \quad (13)$$

where

$$V_0(z) = \frac{\lambda}{16T_R^3} \frac{1}{z} (c - h(1+2z)e^{-z})^2 \quad (14)$$

is the tree-level potential, and

$$V_1(z) = -\frac{1}{16\pi^2} \Lambda_c^2(z) [m_{3/2}^2(z) + kV_0(z)] \quad (15)$$

is the one-loop correction, with the functions $\Lambda_c^2(z)$ and $m_{3/2}^2(z)$ defined in Eqs. (9) and (10) respectively.

The minimum of $V_{eff}(z)$ (13) can easily be obtained by expanding $z = z_0 + \Delta z$ around the tree-level vacuum. A straightforward calculation gives in leading order

$$\Delta z = -\frac{\lambda}{24\pi^2 T_R} \frac{z_0+3}{(z_0-1)^2} e^{-\frac{2}{3}z_0} \quad (16)$$

from which we see that indeed $\Delta z \ll z_0$ as required by our approximation. The function $\mathcal{G}_s$ corresponding to the vacuum (16) does not vanish, but takes the value

$$\mathcal{G}_s = \frac{\lambda^2}{24\pi^2 T_R} \frac{z_0+3}{z_0-1} \frac{1}{4z_0^2} e^{-\frac{2}{3}z_0} \quad (17)$$

This gives a common value before renormalization for all the gaugino masses (7), which is hierarchically smaller than $m_{3/2}$ and is given by

$$m_{1/2} = \frac{\alpha_{\text{GUT}}}{6\pi} \left(\frac{\lambda + 6\pi\alpha_{\text{GUT}}}{\lambda - 2\pi\alpha_{\text{GUT}}} \right) \left(\frac{2\pi\alpha_{\text{GUT}}}{h^2\lambda^2} \right)^{1/3} m_{3/2}^{5/3} / m_P^{2/3} \quad (18)$$

where we have used^{2),4)} $\alpha_{\text{GUT}} \equiv g_{\text{GUT}}^2/4\pi = (4\pi S_R^0)^{-1}$. We observe that the value (18) obtained for $m_{1/2}$ depends on T_R only implicitly through $m_{3/2}$. It does not have any explicit dependence on T_R .

Numerically, using h and $g_{\text{GUT}}^2 = 0(1)$, a value of $m_{1/2} \approx (10^2 \text{ to } 10^3) \text{ GeV}$ corresponds to a gravitino mass $m_{3/2} \approx (1.0 \text{ to } 4.1)\lambda^{2/5} \times 10^{10} \text{ GeV}$. This result combined with Eq. (10) leads to $\lambda \approx 20$ which yields

$$m_{3/2} \approx (0.34 \text{ to } 1.3) \times 10^{11} \text{ GeV} \quad (19a)$$

and hence

$$\Lambda_c \approx \left(\frac{32\pi^2\alpha_{\text{GUT}}^2}{h\lambda} \right)^{1/3} m_{3/2}^{1/3} m_P^{2/3} \approx (2.7 \text{ to } 4.2) \times 10^{15} \text{ GeV} \quad (19b)$$

This value of Λ_c is naturally obtained if G' is $SU(4)$. The implications of these results are discussed later, after some comments on other estimates of supersymmetry breaking effects in the observable sector.

Other possible one-loop sources of gaugino masses have recently been studied in Ref. 14). This paper analyzes: (1) gravitino and gravitino/gauge boson loops; (2) analogous diagrams with the gravitino replaced by an $\tilde{S}$ fermion; and (3) S boson and S boson/gaugino loops. The third class of diagrams is logarithmically divergent in four dimensions, with a physical cut-off being provided by the momentum $|p| \sim \Lambda_c$ above which the gaugino condensate dissolves and the theory becomes supersymmetric. In the limit of large z , the leading contribution of these diagrams is

$$m_{1/2}^{(3)} \approx \frac{m_{3/2}^3}{\pi m_P^2} \cdot 16 z_0^3 \ln \Lambda_c / m_{3/2} \approx \frac{32 z_0^4}{3\pi m_P^2} m_{3/2}^3 \quad (20)$$

which is much smaller than our contribution (18) to $m_{1/2}$. It is not possible to calculate reliably the non-logarithmic finite part to be added to (20) unless one knows how the condensate dissolves at $|p| \gtrsim \Lambda_c$. This means that one cannot

calculate reliably corrections to (20) which involve lower powers of z . Since $z = O(\alpha^{-1})$, such terms would be in any case of higher order in a perturbative expansion in α and should be compared to diagrams containing additional gauge loops that might not be universal. Next we note that all loop calculations in the low-energy effective theory should be done in a consistent regularization scheme. We believe that the physics of the superstring in general and the melting of the gaugino condensate in particular favour the momentum space regularization used above to evaluate (20). If this scheme is also used to evaluate the diagrams of classes (1) and (2) above, their net contributions to $m_{\frac{1}{2}}$ vanish^{*)}. We conclude that our contribution (18) to $m_{\frac{1}{2}}$ is larger than all the contributions studied in Ref. 14).

Chiral gauge non-singlet matter superfields ϕ^λ in the observable sector do not feel supersymmetry breaking at the one-loop level¹⁹⁾. One particular two-loop diagram which can contribute to the scalar masses has been mentioned in Ref. 14). It involves the superpotential $f(\phi) = f_{\lambda mn} \phi^\lambda \phi^m \phi^n$ and a scalar singlet S propagator with a supersymmetry breaking (mass)² insertion $\propto m_{3/2}^2$. However, the four-scalar $f_{\lambda mn} \phi^\lambda \phi^m \phi^n S$ vertex is $O(m_{3/2})$, and we estimate that this particular diagram induces a scalar mass $m_0 \sim m_{3/2}^2$, with a typical suppression factor $f/(4\pi^2)$ coming from the two loop integrations and the Yukawa coupling f . This source of supersymmetry breaking is also negligible compared to the one computed in this paper. We plan to return in the future to a more complete discussion²²⁾ of higher loop contributions to scalar masses.

We now comment on the phenomenological implications of our main result (18) for model calculations. In some recent papers¹⁷⁾, detailed phenomenological studies of radiative symmetry breaking and of the resulting particle spectrum were presented, in the framework of the minimal superstring model that could be obtained from E_6 after symmetry breaking by Wilson loops. These studies were made assuming that the gravitino mass $m_{3/2}$ and the condensation scale Λ_c were both of order of the Planck mass. The result (18) of the present paper suggests, however, that $m_{3/2}$ (19a) and Λ_c (19b) could be significantly smaller than m_P . Nevertheless, we believe that the dominant source of supersymmetry breaking in the observable sector is still a universal gaugino mass $m_{\frac{1}{2}}$, as was assumed in Ref. 17), with scalar masses and trilinear couplings generated by higher order radiative corrections. Apart from this, the only other ingredient used in the old approach¹⁷⁾ was a consistency condition requiring that the vacuum expectation

*) Even if dimensional regularization were used, their contributions would be of lower order in z , and hence higher order in α , than Eq. (20).

values of the Higgs doublets H and $\bar{H}$ and of a neutral singlet N [which transformed non-trivially under the additional $U(1)_E$] all be non-zero. However, this requirement must be imposed in any case, if one wants a phenomenologically acceptable spectrum of fermion and gauge boson masses. Thus the only practical difference between the old approach and that espoused here will be the scale at which the initial boundary conditions are imposed on the soft supersymmetry breaking parameters in the renormalization group equations. These only apply at renormalization scales below Λ_c , since above Λ_c the gaugino condensate dissolves and there is no supersymmetry breaking. In the old approach¹⁷⁾ we were taking $\Lambda_c \sim M_X \sim (10^{17} \text{ to } 10^{18}) \text{ GeV}$, whereas now we should use (19b). This difference has little effect on the previous discussion of radiative symmetry breaking triggered by the scalar squared mass m_N^2 being driven negative, and details of the resulting particle spectra will be discussed in Ref. 22). The slight reduction in the range of Q^2 available for evolution of the scalar masses can be compensated by an increase in the Yukawa couplings governing their evolution. Since the operative Yukawa couplings were those giving masses to the vectorlike charge $-1/3$ quarks, and to a lesser extent to the Higgsinos and top quark, these particles could be somewhat heavier.

Next we comment on astrophysical consequences of our results. It was already mentioned in the introduction of this paper that if E_8 is broken down to some $SU(N)$ subgroup which condenses after inflation, hidden sector flux tubes will create cosmic strings which cause density perturbations $\delta\rho/\rho \sim 10^2(\Lambda_c^2/m_P^2)^{10),11)}$. These perturbations are too large if $N > 4$, and are uninterestingly small if $N < 4$. However, they have just the right magnitude to seed galaxies if $N = 4$, which is the case favoured by our relation²¹⁾ between $m_{\frac{1}{2}}$ and $m_{3/2}$. If, on the other hand, inflation occurs after compactification during²³⁾ the formation of the gaugino condensate, it will generate scale-free perturbations of magnitude $\delta\rho/\rho \sim 20\lambda^{3/2}\Lambda_c^3/m_P^3$. These are of the right magnitude if $N = 4$ or 5 . Thus, galaxy formation and the magnitude of the gauge hierarchy may be intimately linked. We will return to cosmology in Ref. 22).

However, all is not a bed of roses in the scenario outlined in this paper, and we would like to close by recalling three important open problems. One is that we have no dynamical mechanism to fix T_R and hence the scale of the compactified dimensions. No way has been found to fix T_R using radiative corrections^{13),20)}, and here we have only discussed their impact on the determination of S_R . In certain circumstances, non-perturbative dynamics⁶⁾ can give a non-trivial potential for T_R , but the interpretation of this is unsatisfactory⁷⁾. A

second problem is that in order to fix $m_{3/2} \ll m_P$, and assuming $T_R = 0(1)$ as is usually argued⁸⁾, Eq. (10) tells us that we need $z_0 \gg 1$, and hence [according to Eq. (6b)] $c \ll h$ in the superpotential (3). Since $m_W \sim m_{\frac{1}{2}} \sim m_{3/2}^{5/3}/m_P^{2/3}$, it is actually this c/h hierarchy which guarantees the gauge hierarchy: $m_W \sim (c/h)^{5/3} m_P$. The superpotential term c is normally assumed⁴⁾ to be a relic of a vacuum expectation value of the antisymmetric tensor strength $H_{\mu\nu\rho}$ in the six compactified dimensions: $\langle 0|H_{IJK}|0\rangle = c\epsilon_{IJK}$, though any other effect which gave a non-trivial superpotential to beat against the $he^{-\lambda S}$ term (3) due to gaugino condensation could do as well. However, it is known⁶⁾ that non-perturbative world-sheet instanton effects do not give any interactions to the S field. Thus the origin of the hierarchy $c/h \ll 1$, and hence of $m_{3/2} \ll m_P$ and $m_W \ll m_P$, is still open. Finally, we note that the small value of the cosmological constant is also not yet understood. Radiative corrections of the form (8) or (15) make a contribution to the vacuum energy which is much larger than the m_W^4 which had been assumed previously¹²⁾, and which is in turn much larger than the experimental upper limit on the cosmological constant.

ACKNOWLEDGEMENTS

One of us (F.Z.) acknowledges partial financial support from the Fondazione Ing. A. Gini.

REFERENCES

- 1) For recent reviews, see:
 J. Ellis, CERN preprint TH.4439/86 (1986);
 H.-P. Nilles, CERN preprint TH.4444/86 (1986);
 L.E. Ibáñez, CERN preprint TH.4459/86 (1986).
- 2) E. Witten, Phys. Lett. 155B (1985) 151.
- 3) J.-P. Derendinger, L.E. Ibáñez and H.-P. Nilles, Phys. Lett. 155B (1985) 65.
- 4) M. Dine, R. Rohm, N. Seiberg and E. Witten, Phys. Lett. 156B (1985) 55.
- 5) E. Witten, Nucl. Phys. B268 (1986) 79;
 C.M. Hull, Phys. Lett. 167B (1986) 51;
 A. Strominger, Santa Barbara preprint "Superstrings with Torsion" (1985).
- 6) X.-G. Wen and E. Witten, Phys. Lett. 166B (1986) 397;
 J. Ellis, C. Gómez, D.V. Nanopoulos and M. Quirós, Phys. Lett. 173B (1986) 59;
 M. Dine, N. Seiberg, X.-G. Wen and E. Witten, Princeton preprint (1986).
- 7) J. Ellis, K. Enqvist, D.V. Nanopoulos, K. Olive, M. Quirós and F. Zwirner, CERN preprint TH.4449/86 (1986) and UMN-TH-563/86 (1986).
- 8) M. Dine and N. Seiberg, Phys. Rev. Lett. 55 (1985) 366 and Phys. Lett. 162B (1985) 299;
 V. Kaplunovsky, Phys. Rev. Lett. 55 (1985) 1036.
- 9) E. Cohen, J. Ellis, C. Gómez and D.V. Nanopoulos, Phys. Lett. 160B (1985) 62.
- 10) E. Witten, Phys. Lett. 153B (1985) 243.
- 11) B.A. Campbell, J. Ellis, D.V. Nanopoulos, K. Olive and S. Rudaz, CERN preprint TH.4479/86 (1986).
- 12) E. Cohen, J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys. Lett. 161B (1985) 85.
- 13) Y.J. Ahn and J.D. Breit, Nucl. Phys. B273 (1986) 75.
- 14) P. Binétruy, S. Dawson and I. Hinchliffe, preprint LBL-21606 and EFI-86-28 (1986).
- 15) J. Ellis, C. Gómez and D.V. Nanopoulos, Phys. Lett. 168B (1985) 215.
- 16) J.-P. Derendinger, L.E. Ibáñez and H.-P. Nilles, Nucl. Phys. B267 (1986) 365;
 L.E. Ibáñez and H.-P. Nilles, Phys. Lett. 169B (1986) 354.
- 17) E. Cohen, J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys. Lett. 165B (1985) 76;
 J. Ellis, K. Enqvist, D.V. Nanopoulos and F. Zwirner, Nucl. Phys. B276 (1986) 14 and Mod. Phys. Lett. A1 (1986) 57.
- 18) J. Ellis, C. Kounnas and D.V. Nanopoulos, Nucl. Phys. B247 (1984) 373.

- 19) J.D. Breit, B.A. Ovrut and G. Segrè, Phys. Lett. 162B (1985) 303;
P. Binétruy and M.K. Gaillard, Phys. Lett. 168B (1986) 347.
- 20) M. Quirós, Phys. Lett. 173B (1986) 265.
- 21) M.T. Grisaru, M. Roček and A. Karlhede, Phys. Lett. 120B (1983) 110.
- 22) J. Ellis, A.B. Lahanas, D.V. Nanopoulos, M. Quirós and F. Zwirner, CERN
preprint (in preparation).
- 23) J. Ellis, K. Enqvist, D.V. Nanopoulos and M. Quirós, Nucl. Phys. B277 (1986)
231.