Abstract

We explore the one-loop electroweak radiative corrections in the context of the traditional minimal $SU(5)$ and the string-inspired $SU(5) \times U(1)$ supergravity models by calculating explicitly vacuum-polarization and vertex-correction contributions to the $\epsilon_1$ and $\epsilon_b$ parameters. We also include in this analysis the constraint from $b \rightarrow s\gamma$ whose inclusive branching ratio $B(b \rightarrow s\gamma)$ has been actually measured very recently by CLEO. We find that by combining these three most important indirect experimental signatures and using the most recent experimental values for them, $m_t \gtrsim 170$ GeV is excluded for $\mu > 0$ in both the minimal $SU(5)$ supergravity and the no-scale $SU(5) \times U(1)$ supergravity. We also find that $m_t \gtrsim 175$ (185) GeV is excluded for any sign of $\mu$ in the minimal ($SU(5) \times U(1)$) supergravity model.

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I. INTRODUCTION

With the increasing accuracy of the LEP measurements, it has become extremely important performing the precision test of the standard model (SM) and its extensions. A standard model fit to the latest LEP data yields the top mass, $m_t = 178 \pm 11^{+18}_{-10} \text{ GeV}$ \cite{1}, which is in perfect agreement with the measured top mass from CDF \cite{2}, $m_t = 174 \pm 10^{+13}_{-12} \text{ GeV}$. Therefore, it would be desirable for one to narrow down the top mass in the vicinity of the above central value in studying the phenomenology of specific models of interest. Although the SM is in remarkable agreement with all the known experiments, there are a few experimental results that can be interpreted as a possible manifestation of new physics beyond the SM. First, $R_b \equiv \frac{\Gamma(Z\rightarrow b\bar{b})}{\Gamma(Z\rightarrow \text{hadrons})}$ measured at LEP is in disagreement at $2\sigma$ level with the SM predictions. Secondly, the flavor-changing radiative decay $b \rightarrow s\gamma$ \cite{3–6}, whose inclusive branching ratio has been actually measured by CLEO to be at 95% C. L. \cite{7},

\[
1 \times 10^{-4} < B(b \rightarrow s\gamma) < 4 \times 10^{-4},
\]

still leaves room for new physics. Large experimental value for $R_b$ would put rather perilously small upper bound on $m_t$ in the SM \cite{1}. One could certainly interpret this as a possible manifestation of new physics beyond the SM, where at one loop the negative standard top quark contributions are cancelled to a certain extent by the contributions from the new particles, thereby allowing considerably larger $m_t$ than in the SM. In fact, the minimal supersymmetric standard model (MSSM) realizes this possibility \cite{8,9}. Similarly in $b \rightarrow s\gamma$, the values of the branching ratio near the current lower bound can be accommodated for reasonable values of $m_t$ in the MSSM because suppression can occur due to the additional contributions in the model.

In supergravity (SUGRA) models, radiative electroweak symmetry breaking mechanism \cite{10} can be described by at most 5 parameters: the top-quark mass ($m_t$), the ratio of Higgs vacuum expectation values ($\tan \beta$), and three universal soft-supersymmetry-breaking
parameters \( (m_{1/2}, m_0, A) \). Since the entire range of sparticle mass spectrum is quite broad for the models we consider, the large hadron collider (LHC) and the next linear collider (NLC) are needed in order to explore all the regions of the parameter space of our interest. However, the present collider facilities have been successful in probing a good part of the allowed parameter space through indirect experimental signatures. In particular, we will concentrate here on the precision measurements at LEP and the flavor-changing radiative decay \( b \to s\gamma \) observed by CLEO. We adopt the \( \epsilon \)-scheme \cite{11,12} for a global analysis of the precision data at LEP. Among four \( \epsilon \)-parameters, \( \epsilon_i \) \((i = 1, 2, 3, b)\) in this scheme, \( \epsilon_b \) has been studied very recently in the context of the minimal \( SU(5) \) and the no-scale \( SU(5) \times U(1) \) supergravity models \cite{8}. In this work we expand the analysis to include the additional constraints from \( \epsilon_1 \) and \( b \to s\gamma \) in the minimal \( SU(5) \) and a larger class of \( SU(5) \times U(1) \) supergravity models. We will show that by combining above three most important constraints from indirect processes, \( m_t > 170 \) GeV is excluded for \( \mu > 0 \) in both the minimal \( SU(5) \) and the no-scale \( SU(5) \times U(1) \) SUGRA models.

\section*{II. THE MINIMAL SU(5) AND SU(5) \times U(1) SUGRA MODELS}

We consider the minimal \( SU(5) \) SUGRA model \cite{14} and \( SU(5) \times U(1) \) SUGRA model \cite{15} which can be regarded as traditional versus string-inspired unified models. These two models both contain, at low energy, the SM gauge symmetry as well as the particle content of the MSSM, that is, the SM particles with two Higgs doublets and their superpartners. There are, however, a few crucial differences between the two models which are:

(i) The unification groups are different, \( SU(5) \) versus \( SU(5) \times U(1) \).

(ii) The gauge coupling unification occurs at \( \sim 10^{16} \) GeV in the minimal \( SU(5) \) model whereas in \( SU(5) \times U(1) \) model it occurs at the string scale \( \sim 10^{18} \) GeV \cite{16}. In \( SU(5) \times U(1) \) SUGRA, the gauge unification is delayed because of the effects of an additional vector-like

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1See, however, Ref. \cite{13} for non-universal soft-supersymmetry breaking parameters
quark doublet with a mass $\sim 10^{12}$ GeV and one additional vector-like quark singlet of charge $-1/3$ with a mass $\sim 10^6$ GeV. The different heavy field content at the unification scale leads to different constraints from proton decay.

(iii) The minimal $SU(5)$ SUGRA is highly constrained by proton decay while $SU(5) \times U(1)$ SUGRA is not.

The above SUGRA models can be completely described by only five parameters under a few simplifying assumptions on the values of the soft-supersymmetric-breaking parameters at the unification scale. That is, all three gauginos are assumed to have a common mass $m_{1/2}$, and all squarks, sleptons, and two Higgs scalar doublets to have a common mass $m_0$, and three trilinear scalar couplings are taken to be identical to $A$. The Higgs mixing parameter $\mu$ and its associated bilinear coupling $B$ are in fact determined by imposing the radiative EW breaking condition. All these boil down to only five parameters, $m_{1/2}, m_0, A, \tan \beta$, and $m_t$.

One can also restrict further the above 5-dimensional parameter spaces as follows [17]. First, upon sampling a specific choice of $(m_{1/2}, m_0, A)$ at the unification scale and $(m_t, \tan \beta)$ at the electroweak scale, the renormalization group equations (RGE) are run from the unification scale to the electroweak scale, where the radiative electroweak breaking condition is imposed by minimizing the effective 1-loop Higgs potential to determine $\mu$ up to its sign and $B$. Here the sign of $\mu$ is given as usual [18], and differs from that of Ref. [38]; i.e., we define $\mu$ by $W_\mu = \mu H_1 H_2$. We also impose consistency constraints such as perturbative unification and the naturalness bound of $m_\tilde g \lesssim 1$ TeV. Finally, all the known experimental bounds on the sparticle masses are imposed [2]. This procedure yields the restricted parameter spaces for the two models.

Further reduction in the number of input parameters in $SU(5) \times U(1)$ SUGRA is made

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$^2$We use the following experimental lower bounds on the sparticle masses in GeV in the order of gluino, squarks, lighter stop, sleptons, and lighter chargino: $m_\tilde g \gtrsim 150$, $m_\tilde q \gtrsim 100$, $m_{\tilde t_1} \gtrsim 45$, $m_{\tilde t} \gtrsim 43$, $m_{\chi^\pm_1} \gtrsim 45$. 

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possible because in specific string-inspired models for \( (m_{1/2}, m_0, A) \) at the unification scale these three parameters are computed in terms of just one of them \(^{19}\). One obtains \( m_0 = A = 0 \) in the no-scale model and \( m_0 = \frac{1}{\sqrt{3}} m_{1/2}, A = -m_{1/2} \) in the dilaton model \(^{3}\).

The low energy predictions for the sparticle mass spectra are quite different in the two SUGRA models mainly due to the different pattern of supersymmetry radiative breaking. In the minimal \( SU(5) \) SUGRA model, all the squarks except the lighter stop and all the Higgs except the lighter neutral Higgs are quite heavy (\( \gtrsim \) a few hundred GeV) whereas in the \( SU(5) \times U(1) \) SUGRA model they can be quite light. This difference leads to strikingly different phenomenology in the two models, for example in the flavor changing radiative decay \( b \to s \gamma \). \(^{5}\)

### III. CONSTRAINTS FROM THE EW RADIATIVE CORRECTIONS AND THE FLAVOR CHANGING RADIATIVE DECAY

Parametrizing the electroweak vacuum polarization corrections with three parameters can be understood as follows. It can be shown, by expanding the vacuum polarization tensors to order \( q^2 \), that one obtains three independent physical parameters. Alternatively, one can show that upon symmetry breaking three additional terms appear in the effective lagrangian \(^{21}\). Among several schemes to parametrize the corrections \(^{21}\), in the \( (S, T, U) \) scheme \(^{23}\), the deviations of the model predictions from the SM predictions (with fixed SM values for \( m_t, m_{H} \)) are considered as the effects from “new physics”. This scheme is valid only up to the lowest order in \( q^2 \), and is therefore not applicable to a theory with light new particles comparable to \( M_Z \). In the \( \epsilon \)-scheme \(^{11,12}\), on the other hand, the model predictions are absolute and also valid up to higher orders in \( q^2 \), and therefore this scheme is more applicable to the electroweak precision tests of the MSSM \(^{25}\) and a class of supergravity models \(^{26}\).

There are two different \( \epsilon \)-schemes. The original scheme \(^{11}\) was considered in one of

\(^3\)Note, however, that one loop correction changes this relation significantly \(^{20}\).
author’s previous analyses [26,27], where $\epsilon_{1,2,3}$ are defined by a basic set of observables $\Gamma_{l}, A_{FB}^l$ and $M_{W}/M_{Z}$. Due to the large $m_t$-dependent vertex corrections to $\Gamma_b$, the $\epsilon_{1,2,3}$ parameters and $\Gamma_b$ can be correlated only for a fixed value of $m_t$. Therefore, $\Gamma_{tot}$, $\Gamma_{hadron}$ and $\Gamma_b$ were not included in Ref. [11]. However, in the new $\epsilon$-scheme, introduced recently in Ref. [12], the above difficulties are overcome by introducing a new parameter $\epsilon_b$ to encode the $Z \rightarrow b\bar{b}$ vertex corrections. The four $\epsilon$’s are now defined by an enlarged set of $\Gamma_l$, $\Gamma_b$, $A_{FB}^l$ and $M_{W}/M_{Z}$ without even specifying $m_t$. This new scheme was adopted in a previous analysis by one of us (G.P.) in the context of the $SU(5) \times U(1)$ SUGRA models [28]. In this work we use this new $\epsilon$-scheme. As is well known, the SM contribution to $\epsilon_1$ depends quadratically on $m_t$ but only logarithmically on the SM Higgs boson mass ($m_H$). Therefore upper bounds on $m_t$ have a non-negligible $m_H$ dependence and become around 20 GeV lower when going from $m_H = 1$ TeV to $m_H = 100$ GeV. It is also known in the MSSM that the largest supersymmetric contributions to $\epsilon_1$ are expected to arise from the $\tilde{t}\tilde{b}$ sector, and in the limiting case of a very light stop the contribution is comparable to that of the $t-b$ sector. The remaining squark, slepton, chargino, neutralino, and Higgs sectors all typically contribute considerably less. For increasing sparticle masses, the heavy sector of the theory decouples, and only SM effects with a light Higgs boson survive. However, for a light chargino ($m_{\chi^\pm} \rightarrow \frac{1}{2} M_Z$), a $Z$-wavefunction renormalization threshold effect coming from $Z$-vacuum polarization diagram with the lighter chargino in the loop can introduce a substantial $q^2$-dependence in the calculation [23]. This results in a weaker upper bound on $m_t$ than in the SM. The complete vacuum polarization contributions from the Higgs sector, the supersymmetric chargino-neutralino and sfermion sectors, and also the corresponding contributions in the SM have been included in our calculations [20]. However, the supersymmetric contributions to the non-oblique corrections except in $\epsilon_b$ have been neglected.

Following Ref. [12], $\epsilon_b$ is defined from $\Gamma_b$, the inclusive partial width for $Z \rightarrow b\bar{b}$, as

$$\epsilon_b = \frac{g_A^b}{g_A^l} - 1 \quad (1)$$

where $g_A^b$ ($g_A^l$) is the axial-vector coupling of $Z$ to $b$ ($l$). In the SM, the diagrams for $\epsilon_b$
involve top quarks and $W^\pm$ bosons \cite{29}, and the contribution to $\epsilon_b$ depends quadratically on $m_t$ ($\epsilon_b = -G_Fm_t^2/4\sqrt{2}\pi^2 + \cdots$). In supersymmetric models there are additional diagrams involving Higgs bosons and supersymmetric particles. The charged Higgs contributions have been calculated in Refs. \cite{30,31} in the context of a non-supersymmetric two Higgs doublet model, and the contributions involving supersymmetric particles in Refs. \cite{32,33}. The main features of the additional supersymmetric contributions are: (i) a negative contribution from charged Higgs–top exchange which grows as $m_t^2/\tan^2 \beta$ for $\tan \beta \ll m_t/m_b$; (ii) a positive contribution from chargino-stop exchange which in this case grows as $m_t^2/\sin^2 \beta$; and (iii) a contribution from neutralino(neutral Higgs)–bottom exchange which grows as $m_b^2\tan^2 \beta$ and is negligible except for large values of $\tan \beta$ (i.e., $\tan \beta \gtrsim m_t/m_b$) (the contribution (iii) has been neglected in our analysis).

In the MSSM, $b \to s\gamma$ decay receives significant contributions from penguin diagrams with $W^\pm - t$ loop, $H^\pm - t$ loop \cite{31} and the $\chi^\pm_{1,2} - \tilde{t}_{1,2}$ loop \cite{35}. The expression used for $B(b \to s\gamma)$ in the leading logarithmic (LL) calculations is given by \cite{36}

$$B(b \to s\gamma) \sim B(b \to ce\bar{\nu}) = 6\alpha \left(\frac{16}{23} A_\gamma + \frac{8}{3}(\eta^{14/23} - \eta^{16/23})A_g + C\right)^2,$$

where $\eta = \alpha_s(M_W)/\alpha_s(m_b)$, $I$ is the phase-space factor $I(x) = 1 - 8x^2 + 8x^6 - x^8 - 24x^4 \ln x$, and $f(m_c/m_b) = 2.41$ the QCD correction factor for the semileptonic decay. $C$ represents the leading-order QCD corrections to the $b \to s\gamma$ amplitude when evaluated at the $\mu = m_b$ scale \cite{36}. We use the 3-loop expressions for $\alpha_s$ and choose $\Lambda_{QCD}$ to obtain $\alpha_s(M_Z)$ consistent with the recent measurements at LEP. In our computations we have used: $\alpha_s(M_Z) = 0.118$, $B(b \to ce\bar{\nu}) = 10.7\%$, $m_b = 4.8$ GeV, and $m_c/m_b = 0.3$. The $A_\gamma, A_g$ are the coefficients of the effective $bs\gamma$ and $bsg$ penguin operators evaluated at the scale $M_W$. Their simplified expressions are given in Ref. \cite{4} in the justifiable limit of negligible gluino and neutralino contributions \cite{27} and degenerate squarks, except for the $\tilde{t}_{1,2}$ which are significantly split by $m_t$. Regarding large uncertainties in the LL QCD corrections, which is mainly due to the choice of renormalization scale $\mu$ and is estimated to be $\approx 25\%$, it has been recently demonstrated by Buras et al. in Ref. \cite{37} that the significant $\mu$ dependence in the LL
result can in fact be reduced considerably by including next-to-leading logarithmic (NLL) corrections, which however, involves very complicated calculations of three-loop mixings between certain effective operators and therefore have not been completed yet.

IV. RESULTS AND DISCUSSION

In Figure 1 we present our numerical results for $\epsilon_1$ versus $\epsilon_b$ in the two $SU(5) \times U(1)$ SUGRA models. Similar analysis, within the context of the infrared fixed point solution of the top quark mass in the MSSM, was recently performed in Ref. [38]. $\alpha(S(M_Z)) = 0.118$ and $m_b = 4.8$ GeV are used throughout the numerical calculations. We use in the figure the following experimental values for $\epsilon_1$ and $\epsilon_b$,

$$\epsilon_1^{exp} = (3.5 \pm 1.8) \times 10^{-3}, \quad \epsilon_b^{exp} = (0.9 \pm 4.2) \times 10^{-3},$$

determined from the latest $\epsilon$- analysis using the LEP and SLC data in Ref. [39]. In the figure points between the two horizontal lines are allowed by the $\epsilon_1$ constraint at the 90% C. L. while the arrow points into the region allowed by the $\epsilon_b$ constraint at the 90% C. L. The values of $m_t$ are as indicated. The combined constraint is stronger for $\mu > 0$, excluding $m_t \gtrsim 180$ GeV at the 90% C. L. The significant drop in $\epsilon_1$ comes from the threshold effect of Z-wavefunction renormalization as discussed in the previous section. The current experimental values for $\epsilon_{1,b}$ prefer light but not too light chargino: for $m_t = 170$ GeV, in the no-scale (dilaton) model,

$$50 \text{ GeV} \lesssim m_{\chi_1^\pm} \lesssim 70 \ (60) \text{ GeV}, \quad \mu > 0$$

$$50 \text{ GeV} \lesssim m_{\chi_1^\pm} \lesssim 150 \ (140) \text{ GeV}, \quad \mu < 0$$

where $m_{\chi_1^\pm} \lesssim 50$ GeV is disfavored by the $\epsilon_1$ constraint $^4$. 

$^4$Our calculation of $\epsilon_1$ near the threshold ($m_{\chi_1^\pm} \rightarrow \frac{1}{2} M_Z$) becomes unstable and loses its credibility.
Since $\epsilon_{1,b}$ constrains the models the most for $m_t = 170$ GeV, we show in Figure 2 the model predictions for $B(b \to s\gamma)$ versus $\epsilon_b$ for $m_t = 170$ GeV to see if $b \to s\gamma$ provides an additional constraint. It is very interesting to see that the combined constraint of $b \to s\gamma$ and $\epsilon_b$ can in fact exclude $m_t = 170$ GeV for $\mu > 0$ in the no-scale model. Therefore, combining $\epsilon_{1,b}$ and $b \to s\gamma$ allows one to exclude $m_t \gtrsim 170$ GeV for $\mu > 0$ altogether in the no-scale model. Similarly, in the dilaton model, $m_t \gtrsim 170$ GeV for $\mu > 0$ is almost excluded. The large suppression in $B(b \to s\gamma)$ for $\mu < 0$ in these models is worth further explanation. As first noticed in Ref. [5], what happens is that in Eq. (2), the $A_{\gamma}$ term nearly cancels against the QCD correction factor $C$; the $A_g$ contribution is small. The $A_{\gamma}$ amplitude receives three contributions: from the $W^{\pm}-t$ loop, from the $H^{\pm}-t$ loop, and from the $\chi^{\pm}_{1,2}-\tilde{t}_{1,2}$ loop. The first two contributions are always negative [36], whereas the last one can have either sign, making it possible having cancellations among three contributions.

In Figure 3-4, we also present the model predictions in the minimal $SU(5)$ SUGRA model for $\epsilon_1$ versus $\epsilon_b$ (top row) and for $B(b \to s\gamma)$ versus $\epsilon_b$ (bottom row) for $m_t = 160, 175, 190$ GeV (Fig. 3) and for $m_t = 170$ GeV (Fig. 4). As can be seen in the figure, the additional constraint from $B(b \to s\gamma)$ here is rather mild as compared to the one in the $SU(5) \times U(1)$ models. However, $\epsilon_b$ constraint turns out to be the strongest of all, excluding $m_t \gtrsim 175$ GeV at the 90% C. L. For $m_t = 170$ GeV, unlike in the $SU(5) \times U(1)$ SUGRA, $\mu > 0$ is excluded by the $\epsilon_1-\epsilon_b$ constraint alone, thereby excluding $m_t \gtrsim 170$ GeV with $\mu > 0$ in the minimal $SU(5)$ SUGRA model. For $\mu > 0$, $\epsilon_b$ prefers light chargino, $m_{\chi^\pm_1} \lesssim 110$ GeV.

In comparison with the results of $\epsilon_{1,b}$ in the $SU(5) \times U(1)$ SUGRA models (Fig. 1), especially the rise in $\epsilon_b$ in the minimal $SU(5)$ SUGRA model is less pronounced. This is mainly due to the fact that the stop mass, which is responsible for the rise in $\epsilon_b$ when it gets light, in fact scales with the chargino mass in the $SU(5) \times U(1)$ SUGRA model whereas it does not in the minimal $SU(5)$ SUGRA model. Therefore, the stop mass and the chargino mass become lighter simultaneously only in the $SU(5) \times U(1)$ SUGRA models, making the light chargino effect in $\epsilon_b$ more manifest than in the minimal $SU(5)$ SUGRA model. This difference of course leads to different $\epsilon_b$-deduced $m_t$ bounds in these models [3].
We would like to comment on the possible direct experimental signatures at the present and future colliders in the models considered here. At the Tevatron, with the estimated sensitivity range by the end of the on-going Run IB the trilepton signals from chargino-neutralino production can probe $m_{\chi_{1}^{\pm}} \lesssim 80 - 90$ GeV only in the dilaton model but not in the no-scale model. The preferred signal at LEPII for the chargino pair production is so-called “mixed-mode” ($1\ell + 2j$), with which one may be able to probe $m_{\chi_{1}^{\pm}} \lesssim 96$ GeV only in the dilaton model [40].

V. CONCLUSIONS

We have computed the one-loop electroweak radiative corrections in terms of $\epsilon_1$ and $\epsilon_b$ in the context of the traditional minimal $SU(5)$ and the string-inspired $SU(5) \times U(1)$ supergravity models. We have also considered the strongest constraint from the radiative flavor-changing decay $b \to s\gamma$ as prompted by the recent measurement of CLEO on the inclusive branching ratio $B(b \to s\gamma)$. We use the latest experimental values for $\epsilon_1$, $\epsilon_b$ and $B(b \to s\gamma)$ in order to constrain the models. We find that combining constraints from $\epsilon_{1,b}$ and $b \to s\gamma$ allows one to exclude $m_t \gtrsim 170$ GeV for $\mu > 0$ altogether in both the minimal $SU(5)$ SUGRA and the no-scale $SU(5) \times U(1)$ SUGRA. We also find that $m_t \gtrsim 175$ $(185)$ GeV is excluded for any sign of $\mu$ in the minimal $(SU(5) \times U(1))$ SUGRA model.

Our results on $B(b \to s\gamma)$ is subject to change due to a large uncertainty in QCD corrections. One could perhaps draw more precise conclusions when QCD corrections become better known.

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REFERENCES

[1] D. Schaile, talk given at 27th International Conference on High Energy Physics, Glasgow, July 1994.

[2] CDF Collaboration, Phys. Rev. Lett. 73, 225 (1994).

[3] V. Barger, M. Berger, and R. J. N. Phillips, Phys. Rev. Lett. 70, 1368 (1993); J. Hewett, Phys. Rev. Lett. 70, 1045 (1993); V. Barger, M. Berger, P. Ohmann, and R. J. N. Phillips, University of Wisconsin preprint, MAD-PH-842 (hep-ph/9407273); Pran Nath and R. Arnowitt, CERN prerint, CERN-TH-7214-94 (hep-ph/9406389).

[4] R. Barbieri and G. Giudice, Phys. Lett. B 309, 86 (1993).

[5] J. Lopez, D. Nanopoulos, and G. T. Park, Phys. Rev. D 48, R974 (1993).

[6] Y. Okada, Phys. Lett. B 315, 119 (1993); R. Garisto and J. N. Ng, Phys. Lett. B 315, 372 (1993); M. Diaz, Phys. Lett. B 322, 207 (1994); F. Borzumati, DESY 93-090, (hep-ph/9310212).

[7] E. Thorndike (CLEO Collaboration), talk given at XXVII International Conference on High Energy Physics (ICHEP), July 1994, Glasgow, Scotland.

[8] J. E. Kim and G. T. Park, Seoul National University preprint SNUTP-94-66 (hep-ph/9408218, to appear in Phys. Rev. D).

[9] J. D. Wells, C. Kolda and G. L. Kane, Michigan preprint UM-TH-94-23 (hep-ph/9408228).

[10] K. Inoue et al., Prog. Theor. Phys. 68, 927 (1982); L. Ibáñez and G. Ross, Phys. Lett. B 110, 215 (1982); L. Ibáñez, Nucl. Phys. B 218, 514 (1983) and Phys. Lett. B 118, 73 (1982); H. P. Nilles, Nucl. Phys. B 217, 366 (1983); J. Ellis, D. V. Nanopoulos, and K. Tamvakis, Phys. Lett. B 121, 123 (1983); J. Ellis, J. Hagelin, D. V. Nanopoulos, and K. Tamvakis, Phys. Lett. B 125, 275 (1983); L. Alvarez-Gaumé, J. Polchinski, and M.
Wise, Nucl. Phys. B 221, 495 (1983); L. Ibañez and C. López, Phys. Lett. B 126, 54 (1983) and Nucl. Phys. B 233, 545 (1984); C. Kounnas, A. Lahanas, D. V. Nanopoulos, and M. Quirós, Phys. Lett. B 132, 95 (1983) and C. Kounnas, A. Lahanas, D. V. Nanopoulos, and M. Quirós, Nucl. Phys. B 236, 438 (1984).

[11] G. Altarelli, R. Barbieri, and S. Jadach, Nucl. Phys. B 369, 3 (1992).

[12] G. Altarelli, R. Barbieri, and F. Caravaglions, Nucl. Phys. B 405, 3 (1993).

[13] D. Matalliotakis and H. P. Nilles, TUM-HEP-201/94.

[14] For reviews see R. Arnowitt and P. Nath, Applied N=1 Supergravity (World Scientific, Singapore 1983); H. P. Nilles, Phys. Rep. 110, 1 (1984).

[15] For a recent review see J. Lopez, D. Nanopoulos, and A. Zichichi, CERN-TH.6926/93 (unpublished).

[16] J. Lopez, D. V. Nanopoulos, and A. Zichichi, Phys. Rev. D 49, 343 (1994) and references therein.

[17] S. Kelley, J. Lopez, D. V. Nanopoulos, H. Pois, and K. Yuan, Nucl. Phys. B 398, 3 (1993).

[18] See, for example, V. Barger, M. Berger, and P. Ohmann, Phys. Rev. D 49, 4908 (1994).

[19] See e.g., L. Ibáñez and D. Lüst, Nucl. Phys. B 382, 305 (1992); V. Kaplunovsky and J. Louis, Phys. Lett. B 306, 269 (1993); A. Brignole, L. Ibáñez, and C. Muñoz, FTUAM-26/93.

[20] K. Choi, J. E. Kim and H. P. Nilles, Phys. Rev. Lett. 73, 1758 (1994).

[21] B. Holdom and J. Terning, Phys. Lett. B 247, 88 (1990); M. Golden and L. Randall, Nucl. Phys. B 361, 3 (1991); A. Dobado, D. Espriu, and M. Herrero, Phys. Lett. B 255, 405 (1991).
[22] D. Kennedy and B. Lynn, Nucl. Phys. B 322, 1 (1989); D. Kennedy, B. Lynn, C. Im, and R. Stuart, Nucl. Phys. B 321, 83 (1989).

[23] M. Peskin and T. Takeuchi, Phys. Rev. Lett. 65, 964 (1990); W. Marciano and J. Rosner, Phys. Rev. Lett. 65, 2963 (1990); D. Kennedy and P. Langacker, Phys. Rev. Lett. 65, 2967 (1990).

[24] G. Altarelli and R. Barbieri, Phys. Lett. B 253, 161 (1990).

[25] R. Barbieri, M. Frigeni, and F. Caravaglios, Phys. Lett. B 279, 169 (1992).

[26] J. Lopez, D. V. Nanopoulos, G. T. Park, H. Pois, and K. Yuan, Phys. Rev. D 48, 3297 (1993).

[27] J. Lopez, D. V. Nanopoulos, G. T. Park, and A. Zichichi, Phys. Rev. D 49, 355 (1994).

[28] J. Lopez, D. V. Nanopoulos, G. T. Park, and A. Zichichi, Phys. Rev. D 49, 4835 (1994).

[29] J. Bernabeu, A. Pich, and A. Santamaria, Phys. Lett. B 200, 569 (1988); W. Beenaker and W. Hollik, Z. Phys. C 40, 141 (1988); A. Akhundov, D. Bardin, and T. Riemann, Nucl. Phys. B 276, 1 (1986); F. Boudjema, A. Djouadi, and C. Verzegnassi, Phys. Lett. B 238, 423 (1990).

[30] A. Denner, R. Guth, W. Hollik, and J. Kühn, Z. Phys. C 51, 695 (1991).

[31] G. T. Park, Mod. Phys. Lett. A 9, 321 (1994).

[32] M. Boulware and D. Finnell, Phys. Rev. D 44, 2054 (1991).

[33] A. Djouadi, G. Girardi, C. Verzegnassi, W. Hollik, and F. Renard, Nucl. Phys. B 349, 48 (1991).

[34] B. Grinstein and M. Wise, Phys. Lett. B 201, 274 (1988); B. Grinstein et. al., Nucl. Phys. B 399, 269 (1990); T. Rizzo, Phys. Rev. D 38, 820 (1988); C. Geng and J. Ng, Phys. Rev. D 38, 2858 (1988); W. Hou and R. Willey, Phys. Lett. B 202, 591 (1988).
[35] S. Bertolini, F. Borzumati, A. Masiero, and G. Ridolfi, Nucl. Phys. B 353, 591 (1991).

[36] B. Grinstein, R. Springer, and M. Wise, Phys. Lett. B 202, 138 (1988); R. Grigjanis et. al., Phys. Lett. B 213, 355 (1988); Phys. Lett. B 224, 209 (1989); M. Misiak, Phys. Lett. B 269, 161 (1991); M. Misiak, Nucl. Phys. B 393, 23 (1993).

[37] A. Buras, M. Misiak, M. Münz and S. Pokorski, Nucl. Phys. B 424, 374 (1994).

[38] M. Carena and C.E.M. Wagner, CERN preprint CERN-TH/7393/94 (hep-ph/9408253).

[39] G. Altarelli, talk given at the 1st International Conference on Phenomenology of Unification from Present to Future, Rome, CENR-TH.7319/94 (June 1994).

[40] J. Lopez, D. Nanopoulos, G. T. Park, X. Wang and A. Zichichi, Phys. Rev. D 50, 2164 (1994); J. Lopez, D. Nanopoulos and A. Zichichi, CTP-TAMU-27/94.
FIGURES

FIG. 1. The correlated predictions for $\epsilon_1$ and $\epsilon_b$ in $10^{-3}$ in the no-scale (top row) and the dilaton (bottom row) $SU(5) \times U(1)$ supergravity model. Points between the two horizontal lines are allowed by the $\epsilon_1$ constraint at the 90% C. L. while the arrow points into the region allowed by the $\epsilon_b$ constraint at the 90% C. L. The values of $m_t$ are as indicated.

FIG. 2. The correlated predictions for $B(b \to s\gamma)$ and $\epsilon_b$ in the no-scale (top row) and the dilaton (bottom row) $SU(5) \times U(1)$ supergravity model for $m_t = 170$ GeV. Points between the two horizontal lines are allowed by the $b \to s\gamma$ constraint at the 95% C. L. while the arrow points into the region allowed by the $\epsilon_b$ constraint at the 90% C. L.

FIG. 3. The correlated predictions in the minimal $SU(5)$ supergravity model for $\epsilon_1$ versus $\epsilon_b$ in $10^{-3}$ (top row) and for $B(b \to s\gamma)$ versus $\epsilon_b$ (bottom row). Points between the two horizontal lines are allowed by the $\epsilon_1$ or $b \to s\gamma$ constraint while the arrow points into the region allowed by the $\epsilon_b$ constraint at the 90% C. L. The values of $m_t$ are as indicated.

FIG. 4. Same as in the Fig. 3 except that $m_t = 170$ GeV is used.