Higgs Limit and $b \to s\gamma$ Constraints in Minimal Supersymmetry

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Abstract

New limits on the Higgs mass from LEP and new calculations on the radiative (penguin) decay of the $b$-quark into $s\gamma$ restrict the parameter space of the Constrained Minimal Supersymmetric Standard Model (CMSSM).

We find that for the low tan $\beta$ scenario only one sign of the Higgs mixing parameter is allowed, while the high tan $\beta$ scenario is practically excluded, if one requires all sparticles to be below 1 TeV and imposes radiative electroweak symmetry breaking as well as gauge and Yukawa coupling unification. For squarks between 1 and 2 TeV high tan $\beta$ scenarios are allowed. We consider especially a new high tan $\beta = 64$ scenario with triple unification of all Yukawa couplings of the third generation, which shows an infrared fixed point behaviour.

The upper limit on the mass of the lightest Higgs in the low (high) tan $\beta$ scenarios is $97\pm6$ ($120\pm2$) GeV, where the largest error originates predominantly from the uncertainty in the top mass.
1 Introduction

Recently several new results have appeared, which have a considerable impact on the allowed parameter space of the Constrained Minimal Supersymmetric Model (CMSSM). As constraints we consider gauge unification and radiative electroweak symmetry breaking assuming the supergravity inspired scenario of the soft breaking of the SUSY masses with common scalar and gaugino masses at the GUT scale [1]. The additional reduction of the parameter space by relic density, $b - \tau$ Yukawa coupling unification and the new Higgs limits from LEP as well as the present $b \rightarrow s\gamma$ rate is calculated.

The most significant new results are:

- the SM Higgs mass limit is now 89.3 GeV at 95\% C.L. for the SM Higgs [2], which corresponds to the CMSSM case too, since all the other Higgses are heavy and decouple in the CMSSM limit.

- Higher order corrections to the $b \rightarrow s\gamma$ calculation became available [3, 4, 5]. As experimental value we use the combined value from the CLEO Coll. [6] and ALEPH Coll. [7]

$$\text{BR}(b \rightarrow s\gamma) = (2.54 \pm 0.56) \cdot 10^{-4},$$

which is consistent with the SM expectation: $\text{BR}(b \rightarrow s\gamma) = (3.6 \pm 0.3) \cdot 10^{-4}$. This SM value was calculated for $\alpha_s = 0.122$, which is the best value from the $Z^0$ data of the four LEP experiments [8]. Since the calculations are usually quoted as the ratio of the $b \rightarrow s\gamma$ rate and the semileptonic decay rate of the $b$-quark, one needs for absolute predictions $\text{BR}(b \rightarrow X_c e \nu)$, for which we use $10.5 \pm 0.4\%$ [9].

In a recent paper the theoretical uncertainties on the photon spectrum are studied in next-to-leading order as well [5]. They could increase the CLEO value by as much as 14\%.

Such a correction improves the agreement between data and the SM prediction. Given the present uncertainty of 22\% in the published data, this correction has not been taken into account, so we stick to the published values and added in the analysis the theoretical uncertainty from the SM prediction, the semileptonic branching ratio uncertainty and the experimental errors in quadrature.

Our statistical analysis, in which the probability for each point of parameter space is determined by a $\chi^2$ fit, has been described previously [10]. The masses of the third generation and the couplings constants are taken from LEP results presented at the Jerusalem Conf. [8], except for the top mass from the combined data from D0 and CDF, which was updated: $m_t = 173.9 \pm 5.2$ GeV at Moriond [11].

Since the common GUT masses $m_0$ and $m_{1/2}$ for the spin 0 and spin 1/2 particles are strongly correlated, we perform the fit for all combinations of $m_0$ and $m_{1/2}$ between 100 GeV and 1 TeV in steps of 100 GeV.

2 Results

2.1 Constraints from the top and bottom mass

The requirement of bottom-tau Yukawa coupling unification strongly restricts the possible solutions in the $m_t$ versus $\tan \beta$ plane, as discussed in our previous paper [10].

Only three values of $\tan \beta$ give an acceptable $\chi^2$ fit for $m_t = 174$ GeV, as shown in the bottom part of fig. [1]. Clearly, the value $\tan \beta = 64$ gives a worse $\chi^2$ than the values at 35 and 1.65, which is mainly due to the fact that this solution cannot fit simultaneously the top and
bottom mass with a single $\tan \beta$ if $b - \tau$ unification is required (either $m_t$ or $m_b$ too large by about $3\sigma$).

Therefore it was not considered in our previous paper[10]. We consider it here, since the solution with $\tan \beta \approx 35$ does not give a better $\chi^2$ anymore in the combined fit of all constraints because of the smaller errors from the gauge couplings, $m_t$ and $b \to s \gamma$.

Furthermore, this solution has the unique feature of a possible triple Yukawa unification: all 3 Yukawa couplings are driven to an approximate fixed point, which yield approximately the correct masses of the leptons and quarks of the third generation. Since the mass deviation are at the $3\sigma$ level, we consider it here, especially since the solution with $\tan \beta \approx 35$ does not give a better $\chi^2$ with the smaller errors from the gauge couplings, $m_t$ and $b \to s \gamma$. The fixed point is reached for a large range of initial values of the Yukawa couplings at the GUT scale, as shown in fig. 2.

The difference between the two solutions at high $\tan \beta$ corresponding to opposite signs of $\mu$ stems from finite loop corrections to the bottom quark mass involving squark-gluino and stop-chargino loops. These corrections are small for low $\tan \beta$ solutions, but can become as large as 10-20% for the high $\tan \beta$ values[10], since the dominant corrections are proportional to $\mu \tan \beta$. Consequently they change sign for different signs of $\mu$. As shown in the middle part of fig. 2 the Yukawa couplings become of the same order of magnitude at large $\tan \beta$. For $\tan \beta = 64$ $Y_b$ varies so rapidly that it is easy to find a solution where all Yukawa couplings are equal, but large at the GUT scale. The low energy value is then approximately given by a fixed point solution for each of the 3 Yukawa couplings, i.e. the solutions of the RGE are insensitive to the starting values at the GUT scale.

Radiative electroweak symmetry breaking is only possible in that case, if one assumes non-universal values for the mass parameters in the Higgs potential.

As will be discussed below, the $b \to s \gamma$ rate is also sensitive to the sign of $\mu$ and it turns out to be difficult to have solutions, where both $b - \tau$ unification and the $b \to s \gamma$ rate are correctly described.

2.2 $b \to s \gamma$ decay rate

Recently the Next-to-Leading Order (NLO) QCD contributions to the $b \to s \gamma$ decay rate have been calculated in the SM[3]. They increase the SM value by about 10%. The NLO electroweak (EW) contributions decrease the result by a few %[5, 4]. In NLO the large theoretical uncertainty from the scale dependence is now reduced to the 10% level (see e.g. ref. [5] for a recent summary on the discussion of the uncertainties). The NLO corrections can be simulated in the lowest order calculation by choosing a renormalization scale $\mu \approx 0.6m_b$. Since the NLO have not yet been calculated for the MSSM, we will stick to the first order calculations, but choose $\mu = 0.6m_b$. Since the NLO corrections are small compared with the experimental uncertainty of 22%, we expect that with the present large experimental uncertainties the NLO corrections in the MSSM will not change the conclusions significantly.

The $b \to s \gamma$ branching ratio is very sensitive to the presence of SUSY particles. Their effect has been calculated by Bertolini et al [12] where a complete but cumbersome matrix notation for the amplitudes was used. We use the more explicit expressions by Barbieri and Giudice [13]. However, in the latter paper the intergenerational mixing in squark mass matrices is omitted, which is not justified in all regions of parameter space: we find it can change the charged current amplitude with charginos in the loop by $\approx 10\%$ or more. So the usual assumption that intergenerational mixing predominantly introduces only the usually negligible flavour changing
Figure 1: The top quark mass as function of tan $\beta$ (top) for values of $m_0, m_{1/2} \approx 1$ TeV. The curve is hardly changed for lower SUSY masses. The middle part shows the corresponding values of the Yukawa couplings at the GUT scale and the lower part the $\chi^2$ values. If the top constraint ($m_t = 174 \pm 5$, horizontal band) is not applied, all values of tan $\beta$ between 1.2 and 70 are allowed, but if the top mass is constrained to the experimental value, only the regions tan $\beta = 1.65 \pm 0.3$, tan $\beta \sim 35$, and tan $\beta \sim 64$ are allowed.

neutral current with gluinos and down-type quarks in the loop, is not valid. Therefore we extended the explicit formulae of Barbieri and Giudice with the intergenerational mixing, thus obtaining an efficient code, which is important for numerical $\chi^2$ analysis.

At low tan $\beta$ the $b \to s\gamma$ rate is close to its SM value for most of the plane (see fig. 3). The charginos and/or the charged Higgses are only light enough at small values of $m_0$ and $m_{1/2}$ to contribute significantly.

However, for large tan $\beta$ the $b \to s\gamma$ rate strongly depends on the particle spectrum and can be larger than the SM value by factors 3 or more, as shown on the right hand side of fig. 3. The chargino contribution, which is proportional to $\mu \tan \beta$, changes sign for different signs of $\mu$. Through interference with the large operator mixing term it can lower or increase the total $b \to s\gamma$ rate. The gluino contribution is noticeable only if $m_{1/2}$ is small. However, in this region the charginos are light too, so the chargino contribution is always larger by an order of magnitude.

For negative $\mu$ values the $b \to s\gamma$ rate is usually too large, except if all sparticles are very heavy and don’t contribute (see fig. 3, right hand side). For positive $\mu$ one can reach the values
Figure 2: The running of the Yukawa couplings in case $Y_t = Y_b = Y_\tau$ at the GUT scale ($SO(10)$ type solution). One can clearly see the approach to the three different fixed points, i.e. the low energy value is largely independent of the GUT scale value. Consequently the GUT scale values can be chosen to be equal (triple unification). The fixed point values at low energy yield correct masses for bottom and tau for $\tan \beta \approx 64$; the fixed point of the top mass yields $m_t^2 = (4\pi)^2 Y_t v^2 \sin^2 \beta = 184$ GeV, which is about $2\sigma$ above the experimental value. The values of the Yukawa couplings at the GUT scale are shown in the maximum allowed range: higher values of $Y_t$ tend to produce negative Higgs masses, while lower values yield a too low top mass.

The $b \to s\gamma$ rate close to the experimental value in most of the parameter region, as is apparent from the $\chi^2$ distribution in the left hand bottom side of fig. 4. However it is difficult to fit simultaneously $m_b$, if the Yukawa coupling $Y_b$ is determined from $m_\tau$ via $b - \tau$ unification (see left hand top part of fig. 4). Note that for this sign of $\mu$ the $m_b$ mass is not just slightly wrong, for which threshold corrections could be blamed, but for our case off by more than $7\sigma$, where we used the conservative Particle Data Book errors for the running mass of $m_b = 4.2 \pm 0.15$ GeV. Recent determinations from Ypsilon mass spectra quote a 50% lower error[14].

On the right hand side of fig. 4 the sign of $\mu$ is reversed. In this case $m_b$ can be fitted very well, but $b \to s\gamma$ cannot be fitted for $m_{1/2} < 600$ GeV. The reason is that in this case the negative chargino loop contribution $\propto \mu \tan \beta$ yields a strong positive contribution through the interference with the negative operator mixing term, so $BR(b \to s\gamma)$ comes out too large, if the charginos are light. On the contrary, the large $m_b$ contributions from $\tilde{g} - \tilde{q}$ and $\tilde{\chi}^\pm - \tilde{t}$ loops proportional to $\mu \tan \beta$ now lower $m_b$ sufficiently.

Our conclusions may seem to contradict those of refs. [15] and [16]. In ref. [15] the discrepancy between the left bottom and left top of fig. 4 is solved by minimizing the corrections to $m_b$ proportional to $\alpha_s \mu \tan \beta$, which can be done by a small value of $\mu$ and a small value of $\alpha_s$. 
Giving up exact gauge unification by using larger errors on the gauge couplings and allowing for an offset of a few % for the strong coupling $\alpha_s$ at the GUT scale, they find acceptable fits.

Our $b \to s\gamma$ calculations are in reasonable agreement with the results from ref. [16] for both signs of $\mu$. However, since they do not impose $b - \tau$ unification, their preferred solution is $\mu > 0$, which corresponds to the left bottom corner of fig. 3 and ignoring the left top corner.

### 2.3 Excluded parameter space and the mass spectrum

In fig. 4 the total $\chi^2$ distribution is shown as a function of $m_0$ and $m_{1/2}$ for the three values of $\tan \beta$ determined above from $b - \tau$ unification. The areas at low $m_0$ and high $m_{1/2}$ are excluded by the LSP constraint, since in that case the lightest $\tilde{\tau}$ can become the LSP. If R-parity is conserved, a charged LSP is not allowed, since the vacuum would be filled with charged relics from the Big Bang.

The relic density constraint excludes $m_0 > 350$ GeV in case of small $\tan \beta$, as discussed previously [10]. For large $\tan \beta$ the Higgsino mixture of the LSP allows a fast enough decay via s-channel $Z^0$ exchange, so in that case the requirement $\Omega h^2 \leq 1$ is easily fulfilled.

As discussed in the previous section the combined requirements of the correct $b \to s\gamma$ rate and $b - \tau$ unification exclude quite a region of parameter space for large $\tan \beta$, as shown by the contours in the lower part.

One observes $\chi^2$ minima at $m_0, m_{1/2}$ around (200,500), (1000,900), and (500,500) for the different $\tan \beta$ values, respectively, as indicated by the stars.

Note that the squarks and gluinos are typically above 1 TeV for the high $\tan \beta$ solutions. Furthermore, the minimal $\chi^2$ values are not excellent for high $\tan \beta$: for $\tan \beta = 64$, $\chi^2_{\text{min}} = 6.1$ from the fitted top mass ($m_t = 189$ GeV), while for $\tan \beta = 35$, $\chi^2_{\text{min}} = 4.3$ from $b \to s\gamma$. All other $\chi^2$ contributions are negligible. For $\tan \beta = 1.65$, $\chi^2_{\text{min}} = 1.7$, basically from the $BR(b \to s\gamma)$ constraint alone.
Figure 4: The $\chi^2$ contributions from the bottom mass and $b \to s\gamma$ for different signs of $\mu$ in the high $\tan \beta$ solution. Note that $b \to s\gamma$ is perfect for $\mu < 0$, while the bottom mass is perfect for $\mu > 0$. Requiring both yields a low $\chi^2$ only for heavy sparticles ($m_{1/2} > 650$ GeV).

Apart from the heavy spectra for large $\tan \beta$, one has the problem that the Born term Higgs masses are strongly negative, as can be anticipated from the fast running of the soft mass terms of the two Higgs doublets, $m_1^2$ and $m_2^2$, which receive negative radiative corrections proportional to the Yukawa couplings. For large $\tan \beta$ both $Y_t$ and $Y_b$ are large (see fig. 4). Then the pseudoscalar Higgs mass, $(m_1^2 + m_2^2)$ at the Born level, and the lightest Higgs get negative masses at the Born level, which can only become positive by large radiative corrections. In this case it is extremely important to include the full 1-loop radiative corrections, i.e. one has to include both the sparticles and particles in the loops. For the full 1-loop corrections we used the formulae from ref. [17]. For the lightest Higgs the dominant second order corrections from ref. [18] have been included as well.

For low $\tan \beta$ the present Higgs limits severely constrains the parameter space, as can be seen from fig. 6, which shows the excluded regions in the $(m_0, m_{1/2})$ plane for different signs of $\mu$. As mentioned in the introduction the SM Higgs limit of 89.3 GeV is valid for the low $\tan \beta$
scenario (tan β < 4) of the MSSM too. As is apparent from fig. 8 this limit clearly rules out the μ < 0 solution, in agreement with other studies[21]. However, this figure assumes mt = 175 GeV. The top dependence on the Higgs mass is slightly steeper than linear in this range (see fig. 6). Adding about one σ to the top mass, i.e. mt = 180 GeV, implies that for the contours in fig. 8 one should add 6 GeV to the numbers shown. Even in this case the μ < 0 solution is excluded for a large region of parameter space. Only the small allowed region with ml/2 > 700 GeV is not yet excluded for mt = 180 GeV. Note that in this region the squarks are well above 1 TeV, so in this case the cancellation of the quadratic divergencies in the Higgs masses, which is only perfect if sparticles and particles have the same masses, starts to become worrying again.

In summary, if all squarks are required to be below 1 TeV the μ < 0 solutions are excluded for the low tan β scenario, even if the uncertainty from the top mass is taken into account. For μ > 0 practically the whole plane is allowed, except for the left bottom corner shown on the top right hand side of fig. 8.

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**Figure 5**: The $\chi^2$-distribution for the low and high tan β solutions. The different shades in the projections indicate steps of $\Delta \chi^2 = 4$, so basically only the light shaded region is allowed. The stars indicate the optimum solution. Contours enclose domains excluded by the particular constraints used in the analysis.
\[
\begin{align*}
\tan \beta &= 1.65, \, \mu < 0 & \tan \beta &= 1.65, \, \mu > 0
\end{align*}
\]

Figure 6: Contours of the Higgs mass (solid lines) in the \(m_0, m_{1/2}\) plane (above) and the Higgs masses (below) for both signs of \(\mu\) for the low tan \(\beta\) solution \(\tan \beta = 1.65\) for \(M_t = 175\) GeV.

The upper limit for the mass of the lightest Higgs is reached for heavy squarks, but it saturates quickly, as is apparent from the bottom row in fig. 6. For \(m_0 = 1000, m_{1/2} = 1000\), which corresponds to squarks masses of about 2 TeV\(^1\), one finds for the upper limit on the Higgs mass in the low tan \(\beta\) scenario:

\[
m_h^{\text{max}} = 97 \pm 6\text{ GeV},
\]

where the error is dominated by the uncertainty from the top mass. If one requires the squarks to be below 1 TeV, these upper limits are reduced by 4 GeV.

It agrees well with the value from Casas et al.: \(m_h = 97 \pm 2\) GeV\(^2\). In both analysis the Renormalization Group Equations are used to determine the trilinear coupling at low energies and \(\mu\) from EWSB, so the mixing in the stop sector is fixed, once the sign of \(\mu\) is choosen. Furthermore in both papers solutions close to the infrared fixed point are considered, which are required in our case from the EWSB. Our error on the upper limit is larger than the one from Casas et al., since they did not consider the error on the top mass.

For high tan \(\beta\) the upper limit on the Higgs mass in the CMSSM is:

\[
m_h^{\text{max}} = 120 \pm 2\text{ GeV}.
\]

The error from the top mass is small, since the high tan \(\beta\) fits prefer anyway top masses around 190 GeV.

\(^1\)Explicit analytical expressions for the sparticle masses as function of the SUSY parameters can be found in ref. \([19]\).
Figure 7: The mass of the lightest CP-even Higgs as function of the top mass at Born level (dotted lines), including complete one-loop contributions of all particles (dashed lines). Two-loop contributions [18] reduce the one-loop corrections [17] significantly as shown by the shaded area (the upper boundary corresponds to $\mu > 0$, the lower one to $\mu < 0$). The solid line just below the dashed line is the one-loop prediction from the third generation only, which apparently gives the main contribution. The upper scale indicates the value of $\tan \beta$, as calculated from the top mass by the requirement of $b - \tau$-unification.

3 Summary

It is shown that the CMSSM can fit simultaneously the constraints from

- gauge couplings unification;
- $b - \tau$ Yukawa couplings unification;
- radiative EWSB;
- life time of the Universe.

The present Higgs limit of 89.3 GeV largely excludes the $\mu < 0$ solution for the low $\tan \beta$ scenario. The $\mu > 0$ solution is just within reach of the coming LEP II runs above 189 GeV, since the upper limit on the Higgs mass is $97 \pm 6$ GeV. Here the central value of 97 GeV was calculated for $m_t = 175$ GeV and $m_h = 130$ GeV is reached for $m_t = 180$ GeV. If one requires the squarks to be below 1 TeV, these upper limits are reduced by 4 GeV.

The maximum value of the lightest Higgs mass in the high $\tan \beta$ scenarios is $120 \pm 2$ GeV, which is beyond the reach of LEP II. However, the high $\tan \beta$ solutions have serious finetuning problems, because of the fact that all Yukawa couplings are large, which cause the Higgs masses at tree level to be strongly negative (typically -1 TeV), so the radiative corrections have to be positive and very large to offset this large negative "starting" value. Furthermore, the solutions
with minimal $\chi^2$ require squarks above 1 TeV, which have an additional finetuning problem because of the non-cancellation of the quadratic divergencies to the Higgs masses. These quadratic corrections only cancel perfectly, if sparticles and particles have the same mass.

**Note added in proof.**

After completing this work, first results on the NLO corrections to the $b \to s \gamma$ rate in the MSSM became available\[22\]. They change the amplitudes by a similar amount as the SM amplitudes, but due to the large cancellations in particular regions of parameter space, the effect can be large sometimes. This will be studied in a future paper. We thank G.F. Giudice for a discussion on this point.

**References**

[1] For references see e.g. the review papers:
   H.-P. Nilles, Phys. Rep. 110 (1984) 1;
   H.E. Haber, G.L. Kane, Phys. Rep. 117 (1985) 75;
   A.B. Lahanas and D.V. Nanopoulos, Phys. Rep. 145 (1987) 1;
   R. Barbieri, Riv. Nuovo Cim. 11 (1988) 1;
   W. de Boer, Progr. in Nucl. and Particle Phys., 33 (1994) 201;
   D. Kazakov, Surveys in High Energy Phys., 10 (1997) 153.

[2] S. de Jong, Higgs Search at LEP, talk at the XXXIIIrd Rencontres de Moriond, (Electroweak Interactions and Unified Theories) Les Arcs, France, March 1998
   V. Kleider-Ruhlmann, Higgs search at LEP, talk at the XXXIIIrd Rencontres de Moriond (QCD), Les Arcs, France, March 1998
   W. de Boer, Frühjahrstagung German Physical Society, Freiburg, March 1998
   U. Schwickerath, Second Latin American Symposium on High Energy Physics, San Juan, April 1998.

[3] K. Adel and Y. Yao, Phys. Rev. D49 (1994) 4945;
   K. Chetyrkin, M. Misiak, M. Munz, Phys. Lett. B400 (1997) 206;
   A. Buras, A. Kuvatkowski, N. Pott, hep-ph/9707482;
   C. Greub and T. Hurth, hep-ph/9708214;
   M. Ciuchini et al, hep-ph/9710333.

[4] A. Czarnecki and W.J. Marciano, BNL-HET-98/11, hep-ph/9804252.

[5] A.L. Kagan and M. Neubert, hep-ph/9805303.

[6] M.S. Alam et al. (CLEO Coll.), Phys. Rev. Lett 74 (1995) 2885.

[7] R. Barate et al. (ALEPH Coll.), Preprint CERN-EP/98-044.

[8] W. de Boer et al., hep-ph/9712376, Proc. of International Europhysics Conference on High-Energy Physics (HEP 97), Jerusalem, Israel, 19-26 Aug 1997.

[9] P. S. Drell (Cornell U., LNS). CLNS-97-1521, hep-ex/9711020, Invited Talk at the 18th Int. Symposium on Lepton - Photon Interactions (LP 97), Hamburg, Germany, 28 Jul - 1 Aug 1997;
   M. Feindt, IEKP-KA-98-4, hep-ph/9802380, Invited talk at the Int. Europhysics Conference on High-Energy Physics (HEP 97), Jerusalem, Israel, 19-26 Aug 1997;
M. Neubert, CERN-TH/98-2, hep-ph/9801269. Invited talk at the Int. Europhysics Conference on High-Energy Physics (HEP 97), Jerusalem, Israel, 19-26 Aug 1997.

[10] W. de Boer et al., Z. Phys. C71 (1996) 415.

[11] M. Jones, for the CDF and D0 Coll., talk at the XXXIIIrd Rencontres de Moriond, Les Arcs, France, March 1998.

[12] S. Bertolini, F. Borzumati, A. Masiero, G. Ridolfi, Nucl. Phys. B353 (1991) 591.

[13] R. Barbieri, G.F. Giudice, Phys. Lett. B309 (1993) 86.

[14] A.H. Hoang, UCSD/PTH 98-02, hep-ph/9803454; K. Melnikov and A. Yelkhovsky, hep-ph/9805270.

[15] T. Blazek and S. Raby, hep-ph/9712254.

[16] H. Baer et al., hep-ph/9712305.

[17] A.V. Gladyshev et al. Nucl. Phys. B498 (1997) 3, and references therein.

[18] M. Carena et al., Higgs Physics at LEP 2, hep-ph/9602250.

[19] W. de Boer et al., Z. Phys. C67 (1995) 647.

[20] J.A. Casas, J.R. Espinosa and H. Haber, hep-ph/9801365.

[21] S.A. Abel and B.C. Allanach, hep-ph/9803476.

[22] M. Ciuchini et al. hep-ph/9806308.