The 2007 Europhysics Conference on High Energy Physics IOP Publishing
Journal of Physics: Conference Series **110** (2008) 072018 doi:10.1088/1742-6596/110/7/072018

Tan(beta) enhanced Yukawa couplings for supersymmetric Higgs singlets at one loop

Robert N. Hodgkinson[1]

School of Physics and Astronomy, University of Manchester, Manchester M13 9PL, United Kingdom

Abstract. Extensions of the MSSM generically feature gauge singlet Higgs bosons. These singlet Higgs bosons have $\tan\beta$-enhanced Yukawa couplings to down-type quarks and leptons at the one-loop level. We present an effective Lagrangian incorporating these Yukawa couplings and use it to study their effect on singlet Higgs boson phenomenology within both the mnSSM and NMSSM. It is found that the loop-induced couplings represent an appreciable effect for the singlet pseudoscalar in particular, and may dominate its decay modes in some regions of parameter space.

PACS. 12.60.Jv Supersymmetric models – 14.80.Cp Non-standard-model Higgs bosons

1 Introduction

The Minimal Supersymmetric Standard Model (MSSM) is a well-motivated extension of the Standard Model of particle physics (SM), which provides a technical solution to the gauge hierarchy problem. The model is minimal in the sense that it includes only those terms in the superpotential which are phenomenologically required, namely the Yukawa couplings $\mathbf{h}_u, \mathbf{h}_d, \mathbf{h}_e$ and a Higgs mass term μ.

One theoretical weakness of the MSSM is the so-called μ-problem [1,2]. In order to achieve a successful electroweak symmetry breaking scheme, the μ-parameter describing the mixing of the two Higgs superfields in the superpotential, i.e. $\mu\hat{H}_u\hat{H}_d$, must be of the order of the soft SUSY-breaking scale $M_{\rm SUSY} \sim 1$ TeV. Within the context of supergravity (SUGRA), the μ-parameter is not in general protected from gravity effects, and is expected to be of the order the Planck scale $M_{\rm Pl}$.

A natural solution to the μ-problem may be obtained by extending the MSSM to include a third Higgs superfield $\hat{S}$, which is a singlet under the SM gauge group, and replacing the μ-term in the superpotential by $\lambda\hat{S}\hat{H}_u\hat{H}_d$. When supersymmetry is softly broken, the scalar component S of $\hat{S}$ generically acquires a vacuum expectation value (VEV) of order $M_{\rm SUSY}$, giving rise to an effective μ-term of the required order.

The superpotential of such a singlet extension of the MSSM exhibits an unwanted global Peccei-Quinn (PQ) symmetry $U(1)_{\rm PQ}$, unless further additions or assumptions are made to the model. The PQ symmetry must be explicitly broken above the electroweak scale to avoid the appearance of visible axions after spontaneous symmetry breaking. Several models have been proposed in the literature based on different choices of discrete and gauged symmetries to break the PQ symmetry [2], including the Next-to-Minimal Supersymmetric SM (NMSSM) [3], the minimal nonminimal Supersymmetric SM (mnSSM) [4] and the $U(1)'$-extended Supersymmetric SM (UMSSM)[5].

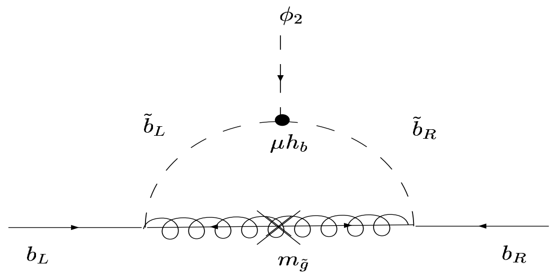

Fig. 1. *The dominant contribution to the inhomogenious coupling $\phi_2 b\bar{b}$ in the MSSM at large $\tan\beta$.*

A common feature of all these models is that the singlet Higgs boson has no tree level couplings to SM fermions or gauge bosons. It has long been known [6, 7] that within the MSSM, threshold corrections to the Yukawa couplings to b quarks and τ leptons can become significant in the limit of large $\tan\beta$, where $\tan\beta$ is the ratio of the two Higgs VEVs. This enhancement partially overcomes the loop suppression factor, and in regions where mixing between the Higgs particles is negligible, the one-loop correction can dominate the $H_1 \to b\bar{b}$ decay width [8]. The dominant contribution to the inhomogenous coupling $\phi_2 b\bar{b}$ is shown in Fig. 1.

An analogous $\tan\beta$ enhanced Yukawa coupling for the singlet Higgs boson is generated at one-loop through sfermion-gaugino loops in singlet extensions of the MSSM [9]. The dominant contribution to the $\phi_S b\bar{b}$ coupling is shown in Fig. 2. These effective couplings can be significant, e.g. of order the SM Yukawa couplings, and in the limit where the H_d doublet decouples from the

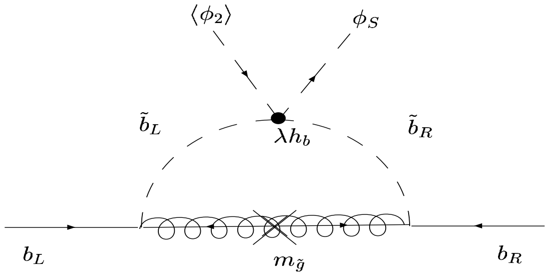

Fig. 2. *The dominant contribution to the coupling $\phi_S b\bar{b}$ in singlet extensions of the MSSM at large $\tan\beta$, analogous to the MSSM graph of Fig 1.*

low energy spectrum, they can provide the dominant decay mechanism for light singlets.

2 Effective Lagrangian framework

The general effective Lagrangian for the self-energy transition $f_L \to f_R$ in the nonvanishing Higgs background may be written as

$$-\mathcal{L}^f_{\rm self} = h_f \bar{f}_R \left(\Phi_1^{0*} + \Delta_f \left[\Phi_1^0, \Phi_2^0, S\right]\right) f_L + \text{H.c.} \quad (1)$$

where $\Phi^0_{1,2} = \frac{1}{\sqrt{2}}\left(v_{1,2} + \phi_{1,2} + i a_{1,2}\right)$ are the electrically neutral components of the two Higgs doublets $H_{d,u}$[1] and $S = \frac{1}{\sqrt{2}}\left(v_S + \phi_S + i a_S\right)$ is the singlet Higgs field. Here $\Delta\left[\Phi_1^0, \Phi_2^0, S\right]$ is a Coleman-Wienberg type functional which encodes the radiative corrections. The VEV of the effective Lagrangian $-\mathcal{L}^f_{\rm self}$ is equal to the fermion mass m_f, allowing us to substitute for the effective Yukawa coupling h_f.

We can use the self-energy effective Lagrangian $\mathcal{L}^f_{\rm self}$ to obtain the form of the effective Lagrangian for the Higgs boson couplings to the fermion f through a Higgs boson low energy theorem [10,11]. Written in terms of the physical Higgs eigenstates $H_{1,2,3}$ and $A_{1,2}$, the effective interaction Lagrangian is

$$-\mathcal{L}^{\rm eff}_{\phi\bar{f}f} = \frac{g_w m_f}{2M_W}\left[\sum_{i=1}^{3} g^S_{H_i ff} H_i \bar{f} f + \sum_{i=1}^{2} g^P_{A_i ff} A_i \left(\bar{f} i\gamma^5 f\right)\right], \quad (2)$$

where the effective couplings g^S and g^P are given by [9]

$$g^S_{H_i ff} = \left(1 + \frac{\sqrt{2}}{v_1}\langle\Delta_f\rangle\right)^{-1}\left[\frac{O^H_{1i}}{c_\beta} + \Delta_f^{\phi_2}\frac{O^H_{2i}}{c_\beta} + \Delta_f^{\phi_S}\frac{O^H_{3i}}{c_\beta}\right] \quad (3)$$

$$g^S_{A_i ff} = \left(1 + \frac{\sqrt{2}}{v_1}\langle\Delta_f\rangle\right)^{-1}\left[-\left(t_\beta + \Delta_f^{a_2}\right) O^A_{1i} + \Delta_f^{a_S}\frac{O^A_{2i}}{c_\beta}\right] \quad (4)$$

Here the orthogonal matrix $O^H(O^A)$ is related to the mixing of the CP-even (CP-odd) scalars and the loop corrections are given by the HLET

[1] Here we adopt the convention for the Higgs doublets: $H_u \equiv \Phi_2$, $H_d \equiv i\tau_2\Phi_1^*$, where τ_2 is the usual Pauli matrix.

$$\Delta_f^{\phi_{2,S}} = \sqrt{2}\left\langle\frac{\partial\Delta_f}{\partial\phi_{2,S}}\right\rangle, \quad \Delta_f^{a_{2,S}} = i\sqrt{2}\left\langle\frac{\partial\Delta_f}{\partial a_{2,S}}\right\rangle. \quad (5)$$

2.1 One-loop evaluation

As may be seen from the above discussion, the effective low-energy couplings of the Higgs bosons to fermions may be calculated from the fermion self-energies. The dominant contributions to the b quark self-energy at large $\tan\beta$ are due to squark-gluino and squark-higgsino loops, giving

$$\begin{aligned}\Delta_b = &-\frac{2\alpha_s}{3\pi} M_3 \left(A_b\Phi_1^{0*} - \lambda S^*\Phi_2^{0*}\right) I(m^2_{\tilde{b}_1}, m^2_{\tilde{b}_2}, M_3^2) \\ &+ \frac{h_t^2}{16\pi^2}\left(A_t\Phi_2^{0*} - \lambda S\Phi_1^{0*}\right) \\ &\times\Big[\, m_{\tilde{\chi}_1}\mathcal{V}^\dagger_{\{21\}}\mathcal{U}^*_{\{12\}}\, I(m^2_{\tilde{t}_1}, m^2_{\tilde{t}_2}, m^2_{\tilde{\chi}_1}) \\ &\quad + m_{\tilde{\chi}_2}\mathcal{V}^\dagger_{\{22\}}\mathcal{U}^*_{\{22\}}\, I(m^2_{\tilde{t}_1}, m^2_{\tilde{t}_2}, m^2_{\tilde{\chi}_2})\Big]. \end{aligned} \quad (6)$$

Here $I(a,b,c)$ is the usual 1-loop integral function

$$I(a,b,c) = \frac{ab\ln(a/b) + bc\ln(b/c) + ac\ln(c/a)}{(a-b)(b-c)(a-c)}. \quad (7)$$

Note that the chargino-mixing matrices $\mathcal{V}, U$ are functionals of $\Phi^0_{1,2}$ and S, as are the sbottom quark masses $m_{\tilde{b}_{1,2}}$, stop quark masses $m_{\tilde{b}_{1,2}}$ and chargino masses $m_{\tilde{\chi}_{1,2}}$.

Similarly, the dominant $\tan\beta$ enhanced contribution to the τ lepton self-energy is due to a stau-chargino loop and is easily derived. The effective Yukawa couplings $\Delta_f^{\phi,a}$ are obtained as the derivatives of these expressions. Note that the presence of the singlet in the model does not alter the form of the 1-loop $\tan\beta$ enhanced couplings of the doublet Higgs fields well known from the MSSM [13].

3 Phenomenology

As the one-loop couplings of the singlet Higgs boson to the b quark and the τ lepton become significant at large values of $\tan\beta$ and λ, we shall set $t_\beta = 50$ and $\mu = \frac{1}{\sqrt{2}}\lambda v_S = 110$ GeV throughout our discussion. The remaining default values of the SUSY parameters for our benchmark scenario, consistent with the constraints from LEP data, are

$$\begin{array}{lll} M_{\tilde{Q}} = 300\ \text{GeV}, & M_{\tilde{L}} = 90\ \text{GeV}, & \\ M_{\tilde{b}} = 110\ \text{GeV}, & M_{\tilde{t}} = 600\ \text{GeV}, & M_{\tilde{\tau}} = 200\ \text{GeV}, \\ A_\tau = 1\ \text{Tev}, & A_t = 1\ \text{TeV}, & A_b = 1\ \text{TeV}, \\ M_1 = 400\ \text{GeV}, & M_2 = 600\ \text{GeV}, & M_3 = 400\ \text{GeV}, \end{array}$$

The physical Higgs boson couplings to the b quark and τ lepton, i.e. $H_{1,2,3}\bar{f}f$ and $A_{1,2}\bar{f}f$, have contributions from both the proper vertex interaction, dominated by the tree-level ϕ_1 coupling, and also the mixing

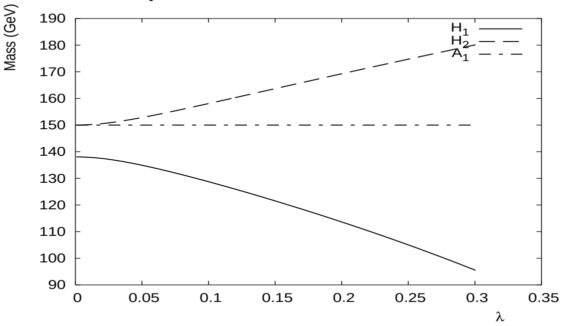

Fig. 3. *Masses of the H_1 (solid line), H_2 (dashed line) and A_1 (dot-dashed line) bosons in the mnSSM with $\mu = 110$ GeV, $\lambda t_S/\mu = (150\ \text{GeV})^2$ and $m_{12}^2 = -0.5\ \text{TeV}^2$. The values of other soft SUSY-breaking parameters are given in Section 3.*

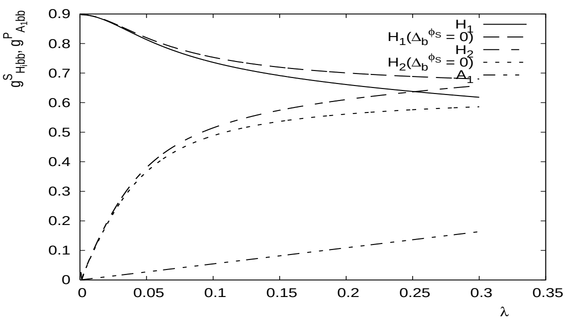

Fig. 4. *The SM-normalized couplings $H_1 b\bar{b}$, $H_2 b\bar{b}$ and $A_1 b\bar{b}$ in the mnSSM, as functions of λ, for the same model parameters as in Fig. 3. .*

of the fields $\phi_{2,S}$ with ϕ_1. This mixing is a tree level effect and is very significant for generic Higgs boson mass matrices. Since our interest is to assess the significance of the one loop singlet Higgs vertex effects, we focus on variants of the mnSSM and NMSSM where the mixing of $\phi_1(a_1)$ with the other scalars is suppressed.

Suppressing both the Higgs boson self-energy transitions $\phi_1 \to \phi_{2,S}$ simultaneously is difficult, except in the MSSM limit $\lambda \to 0$ with μ fixed, where the couplings $\Delta_f^{\phi_S(a_S)}$ also vanish. Instead we impose a constraint on the pseudoscalar mass matrix such that $\left(M_P^2\right)_{12} = 0$. Although this condition is arbitrarily applied here, it is robust against the dominant corrections to the pseudoscalar mass matrix, which can absorbed into the would-be MSSM pseudoscalar mass M_a, and can be generated naturally within certain SUSY-breaking scenarios, e.g. [12].

3.1 mnSSM results

The mnSSM is based on the renormalizable superpotential

$$\begin{aligned} \mathcal{W}_{\text{mnSSM}} = {} & h_l \widehat{H}_d^T i\tau_2 \widehat{L}\widehat{E} \;+\; h_d \widehat{H}_d^T i\tau_2 \widehat{Q}\widehat{D} \\ & + h_u \widehat{Q}^T i\tau_2 \widehat{H}_u \widehat{U} \\ & + \lambda \widehat{S}\widehat{H}_d^T i\tau_2 \widehat{H}_u + t_F \widehat{S}\ . \end{aligned} \tag{8}$$

The term linear in $\hat{S}$ is induced by supergravity quantum effects from Planck-suppressed non-renormalizable operators in the Kähler potential and superpotential [14].

In Fig. 3 we plot the masses of the two lightest CP-even Higgs bosons H_1 and H_2 and the lightest CP-odd Higgs A_1 in the mnSSM with $M_{H^\pm} = 5$ TeV and $\lambda t_S/\mu = (150\ \text{GeV})^2$. The remaining physical Higgs states $H_3 \sim \phi_1$ and $A_2 \sim a$ are heavy, of order $M_{H^\pm}$. For large values of $\lambda > 0.3$ the lightest Higgs boson mass M_{H_1} is well below the LEP limit from direct Higgs searches. Fig. 4 then shows the dependence of the b-quark Yukawa couplings $g^S_{H_{1,2}bb}$ and $g^P_{A_1 bb}$, for the above scenario. The CP-even Yukawa couplings $g^S_{H_{1,2}bb}$ receive appreciable contributions from the tree-level mixing of the state ϕ_1 with $\phi_{2,S}$, which are competitive with the loop-induced Yukawa couplings $\Delta_b^{\phi_{2,S}}$. The coupling $g^P_{A_1 b\bar{b}} \approx g^P_{a_S b\bar{b}}$ is completely dominated by the 1-loop contribution $\Delta_b^{a_S}$. For moderate values of $\lambda \sim 0.3$, we find that $g^P_{A_1 b\bar{b}} \sim 0.15$. Moreover, the decay $A_1 \to \bar{b}b$ is expected to be the dominant decay channel in this specific scenario of the mnSSM.

3.2 NMSSM results

We now turn our attention to the NMSSM. The superpotential of this model is given by

$$\begin{aligned} \mathcal{W}_{\text{NMSSM}} = {} & h_l \widehat{H}_d^T i\tau_2 \widehat{L}\widehat{E} \;+\; h_d \widehat{H}_d^T i\tau_2 \widehat{Q}\widehat{D} \\ & + h_u \widehat{Q}^T i\tau_2 \widehat{H}_u \widehat{U} \\ & + \lambda \widehat{S}\widehat{H}_d^T i\tau_2 \widehat{H}_u \;+\; \frac{\kappa}{3}\widehat{S}^3\ . \end{aligned} \tag{9}$$

The NMSSM spectrum contains a light singlet dominated pseudoscalar if the soft trilinear couplings are approximately $A_\lambda \sim 200$ GeV and $A_\kappa \sim 5$ GeV. This can be naturally arranged in gauge or gaugino mediated SUSY breaking scenarios, where these parameters are zero at tree level. The above scales are generated by quantum corrections if the gaugino masses are of the order 100 GeV.

In recent years there has been some interest in the phenomenology of light Higgs pseudoscalars in the NMSSM, which may provide an invisible decay channel for a light SM-like Higgs boson. If these CP-odd scalars have a large singlet component, it is possible for them to escape experimental bounds [15]. In Fig. 5 we plot the couplings of H_1 to $b\bar{b}$ and of A_1 to both $b\bar{b}$ and $\tau\bar{\tau}$ pairs for such a scenario with $M_{H^\pm} = 2$ TeV. Here M_{A_1} is in the range $6 \sim 9$ GeV and M_{H_1} the range $120 \sim 140$ GeV. The threshold corrections can clearly have a significant effect on the branching ratios of a light CP-odd singlet scalar for moderate to large values of λ. Previous studies have considered detection of these particles through decays to photon pairs as the dominant mode [16] in the limit of vanishing singlet-doublet pseudoscalar mixing. Our analysis shows that this need not be the case, and the impact of the hadronic decays of A_1 in so-called "invisible Higgs" scenarios should still be considered even in this limit.

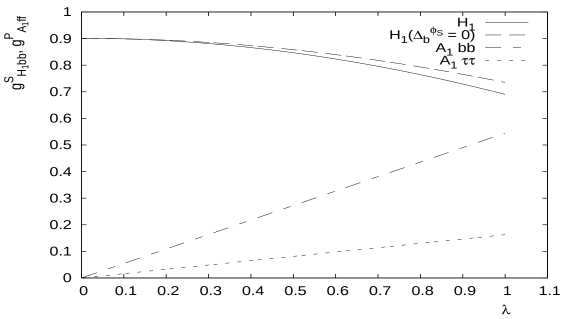

Fig. 5. *The SM-normalized couplings $H_1 b\bar{b}$ (solid line), $A_1 b\bar{b}$ (dashed line) and $A_1 \tau^+\tau^-$ (dot-dashed line) in the NMSSM, as functions of λ.*

4 Conclusions

Minimal extensions of the MSSM generically include singlet Higgs bosons. Although singlet Higgs bosons have no direct or proper couplings to the SM particles, their interaction with the observed matter can still be significant as a result of two contributions. The first one is their mixing with Higgs doublet states, which is often considered in the literature. The second contribution is novel and persists even if the Higgs doublet-singlet mixing is completely switched off. It results from $\tan\beta$ enhanced gluino, chargino and squark quantum effects at the 1-loop level.

In the absence of a Higgs doublet-singlet mixing, the 1-loop quantum effects we have been studying here will be the only means by which the CP-odd singlet may couple to quarks and leptons. For a sufficiently light CP-odd singlet scalar, with a mass below the squark threshold, the loop-induced Yukawa couplings will provide its dominant decay channel into b quarks. This has important phenomenological implications for studies of the NMSSM with light pseudoscalar.

References

1. J. E. Kim and H. P. Nilles, Phys. Lett. B **138** (1984) 150; L. J. Hall, J. D. Lykken and S. Weinberg, Phys. Rev. D **27** , (1983) 2359; G. F. Giudice and A. Masiero, Phys. Lett. B **206**, (1988) 480; E. J. Chun, J. E. Kim and H. P. Nilles, Nucl. Phys. B **370** , (1992) 105; I. Antoniadis, E. Gava, K. S. Narain and T. R. Taylor, Nucl. Phys. B **432**, (1994) 187.
2. E. Accomando *et al.*, arXiv:hep-ph/0608079.
3. P. Fayet, Nucl. Phys. B **90** (1975) 104; J. M. Frere, D. R. T. Jones and S. Raby, Nucl. Phys. B **222** (1983) 11; J. P. Derendinger and C. A. Savoy, Nucl. Phys. B **237** (1984) 307; J. R. Ellis, J. F. Gunion, H. E. Haber, L. Roszkowski and F. Zwirner, Phys. Rev. D **39** (1989) 844; S. F. King and P. L. White, Phys. Rev. D **52** (1995) 4183; M. Bastero-Gil, C. Hugonie, S. F. King, D. P. Roy and S. Vempati, Phys. Lett. B **489**, (2000) 359; U. Ellwanger, J. F. Gunion and C. Hugonie, JHEP **0507** (2005) 041.
4. C. Panagiotakopoulos and A. Pilaftsis, Phys. Rev. D **63** (2001) 055003; A. Dedes, C. Hugonie, S. Moretti and K. Tamvakis, Phys. Rev. D **63** (2001) 055009; A. Menon, D. E. Morrissey and C. E. M. Wagner, Phys. Rev. D **70** (2004) 035005; S. W. Ham, S. K. OH, C. M. Kim, E. J. Yoo and D. Son, Phys. Rev. D **70** (2004) 075001.
5. M. Cvetic, D. A. Demir, J. R. Espinosa, L. L. Everett and P. Langacker, Phys. Rev. D **56**, (1997) 2861 [Erratum-ibid. D **58**, (1998) 119905]; P. Langacker and J. Wang, Phys. Rev. D **58**, (1998) 115010; S. F. King, S. Moretti and R. Nevzorov, Phys. Rev. D **73** (2006) 035009.
6. T. Banks, Nucl. Phys. B **303** (1988) 172; E. Ma, Phys. Rev. D **39** (1989) 1922.
7. R. Hempfling, Phys. Rev. D **49** (1994) 6168; L. J. Hall, R. Rattazzi and U. Sarid, Phys. Rev. D **50** (1994) 7048; M. Carena, M. Olechowski, S. Pokorski and C. E. M. Wagner, Nucl. Phys. B **426** (1994) 269; D. M. Pierce, J. A. Bagger, K. T. Matchev and R. j. Zhang, Nucl. Phys. B **491** (1997) 3; F. Borzumati, G. R. Farrar, N. Polonsky and S. D. Thomas, Nucl. Phys. B **555** (1999) 53.
8. J. A. Coarasa, R. A. Jimenez and J. Sola, Phys. Lett. B **389** (1996) 312; R. A. Jimenez and J. Sola, Phys. Lett. B **389** (1996) 53; K. T. Matchev and D. M. Pierce, Phys. Lett. B **445** (1999) 331; P. H. Chankowski, J. R. Ellis, M. Olechowski and S. Pokorski, Nucl. Phys. B **544** (1999) 39; K. S. Babu and C. F. Kolda, Phys. Lett. B **451** (1999) 77.
9. R. N. Hodgkinson and A. Pilaftsis, Phys. Rev. D **76** (2007) 015007
10. J. R. Ellis, M. K. Gaillard and D. V. Nanopoulos, Nucl. Phys. B **106** (1976) 292; M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, Phys. Lett. B **78** (1978) 443; M. A. Shifman, A. I. Vainshtein, M. B. Voloshin and V. I. Zakharov, Sov. J. Nucl. Phys. **30** (1979) 711; A. I. Vainshtein, V. I. Zakharov and M. A. Shifman, Sov. Phys. Usp. **23** (1980) 429 [Usp. Fiz. Nauk **131** (1980) 537]; M. B. Voloshin, Sov. J. Nucl. Phys. **44** (1986) 478 [Yad. Fiz. **44** (1986) 738]; M. A. Shifman, Phys. Rept. **209** (1991) 341 [Sov. Phys. Usp. **32** (1989 UFNAA,157,561-598.1989) 289]; S. Dawson and H. E. Haber, Int. J. Mod. Phys. A **7** (1992) 107; B. A. Kniehl and M. Spira, Z. Phys. C **69** (1995) 77.
11. D. Binosi, J. Papavassiliou and A. Pilaftsis, Phys. Rev. D **71** (2005) 085007.
12. R. Dermisek and J. F. Gunion, Phys. Rev. Lett. **95** (2005) 041801; R. Dermisek and J. F. Gunion, Phys. Rev. D **73** (2006) 111701; S. Chang, P. J. Fox and N. Weiner, JHEP **0608** (2006) 068.
13. M. Carena, D. Garcia, U. Nierste and C. E. M. Wagner, Nucl. Phys. B **577** (2000) 88; M. E. Gomez, T. Ibrahim, P. Nath and S. Skadhauge, Phys. Rev. D **74** (2006) 015015; A. J. Buras, P. H. Chankowski, J. Rosiek and L. Slawianowska, Nucl. Phys. B **659** (2003) 3.
14. For a power-counting calculation of the size of the effective tadpole in the mnSSM, see the first reference in [4].
15. J. F. Gunion, D. Hooper and B. McElrath, Phys. Rev. D **73** (2006) 015011.
16. B. A. Dobrescu, G. Landsberg and K. T. Matchev, Phys. Rev. D **63** (2001) 075003; A. Arhrib, K. Cheung, T. J. Hou and K. W. Song, arXiv:hep-ph/0606114; arXiv:hep-ph/0611211.