$\Upsilon \rightarrow \gamma A_1$ in the NMSSM at large $\tan \beta$

Robert N. Hodgkinson

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

ABSTRACT

We investigate the effects of the radiatively-generated $\tan \beta$-enhanced Higgs-singlet Yukawa couplings on the decay $\Upsilon \rightarrow \gamma A_1$ in the NMSSM, where $A_1$ is the lightest CP-odd scalar. This radiative coupling is found to dominate in the case of a highly singlet Higgs pseudoscalar. The branching ratio for the production of such a particle is shown to be within a few orders of magnitude of current experimental constraints across a significant region of parameter space. This represents a potentially observable signal for experiments at present $B$-factories.
The Next to Minimal Supersymmetric extension of the Standard Model (NMSSM) is a well-motivated model of electroweak symmetry breaking which resolves both the hierarchy problem of the Standard Model (SM) and $\mu$ problem of the Minimal Supersymmetric extension of the Standard Model (MSSM) in a natural way [1]. The $\mu$ parameter of the MSSM is replaced with an additional gauge singlet Higgs superfield $\hat{S}$ and an effective doublet mixing term $\mu_{\text{eff}}$ is generated when the singlet field acquires a vacuum expectation value (VEV). It has long been known that the NMSSM suffers from the formation of electroweak scale cosmic domain walls [2], although mechanisms to resolve this problem have been suggested, e.g. [3]. In the NMSSM, all parameters are naturally predicted to be of the order the SUSY-breaking scale $M_{\text{SUSY}}$.

The Higgs sector of the NMSSM may be derived from the superpotential of the model, given by

$$W_{\text{Higgs}} = \lambda \hat{S} \hat{H}_1 \hat{H}_2 + \kappa \hat{S}^3,$$

where $\hat{H}_1(\hat{H}_2)$ is the doublet Higgs superfield which gives masses to the down-type quarks and leptons (up-type quarks). The corresponding soft SUSY-breaking terms are given by

$$L_{\text{soft}}^{\text{Higgs}} = \lambda A_\lambda S \Phi_1 \Phi_2 + \kappa A_\kappa S^3,$$

where $\Phi_1, 2$ and $S$ are the scalar components of $\hat{H}_1, 2$ and $\hat{S}$ respectively. At tree level only two further parameters are required, the ratio of doublet VEVs $\tan \beta = \frac{v_2}{v_1}$ and the effective doublet mixing parameter $\mu_{\text{eff}} = \lambda v S \sqrt{2}$. Radiative corrections due to the quarks and scalar quarks of the third generation must also be included in order to raise the mass of the SM-like Higgs $H_1$ above the LEP bound of 114 GeV.\footnote{It is also possible to evade the LEP bound if $H_1$ decays into the lightest pseudoscalars, with branching ratio $B(H_1 \to A_1 A_1) > 0.7$ [5]. This requirement leads to a lower bound on the doublet component of $A_1$, $O_{A_1}^1 > 0.04$. Since the radiative corrections are subdominant in this region, we do not include these points in our results.}

The superpotential of the NMSSM exhibits a global $U(1)_R$ symmetry which is spontaneously broken when $S$, the scalar component of $\hat{S}$, acquires a VEV. In addition, it is explicitly broken by the soft trilinear couplings $A_\lambda, A_\kappa$ [4]. The CP-odd scalar component of $\hat{S}$ is therefore a pseudo-Goldstone boson of this symmetry, and is massless in the limit $A_\lambda, A_\kappa \to 0$. For small values of the trilinear couplings, the lightest pseudoscalar in the NMSSM spectrum can therefore naturally be very light and highly gauge singlet in nature. Typically this requires $A_\lambda \sim 200$ GeV, $A_\kappa \sim 5$ GeV. Such a scenario can arise within the context of gauge- or gaugino-mediated SUSY breaking, where both couplings are zero at tree level, with non-zero $A_\lambda$ being radiatively generated at one loop and non-zero $A_\kappa$ at two loops [5].

For a sufficiently light $A_1$ boson, observation in the decay $\Upsilon(1s) \to \gamma A_1$ becomes a
possibility\footnote{In principle, this decay is also possible within other Higgs singlet extensions of the MSSM, such as the Minimal Non-minimal Supersymmetric extension of the Standard Model (MNSSM) [6], and Eq. (3) is also valid in this case. However, the Higgs bosons of the MNSSM obey a tree level mass sum rule, so that any light pseudoscalar boson is accompanied by a quasi-degenerate scalar boson. Singlet-doublet mixing in the scalar sector typically excludes such a scenario except in the MSSM limit of the theory $\lambda \to 0$ with $\mu_{\text{eff}}$ fixed. The radiative coupling of the singlet pseudoscalar to fermions, whose effects we consider here, will also vanish in this limit.}. Such a signal has previously been considered in [7, 8], with the pseudoscalar coupling to $b$-quarks only through tree level singlet-doublet mixing. It has recently been shown that the singlet Higgs bosons also receive a direct coupling to fermions at one loop [9]. Although loop suppressed, this coupling is enhanced by the ratio of Higgs doublet VEVs $\tan \beta$ and can become competitive with tree-level effects when this parameter is large.

In this Letter we consider the effects of such a direct coupling on the decay $\Upsilon \to \gamma A_1$. Experimental searches [10] for a light Higgs boson in $\Upsilon$ decays place a 90% confidence level upper bound on the branching ratio $B(\Upsilon \to \gamma A_1) < 1 \times 10^{-4}$ for a light particle $m_{A_1} < 8 \text{ GeV}$ decaying visibly within the detector. The upper bound rises to $\sim 10^{-3}$ for heavier particles due to the softness of the recoil photon and cuts placed on energy deposits in the detector. The branching ratio for $\Upsilon$ decays through the Wilczek mechanism [8, 12] is given by

$$
B(\Upsilon \to \gamma A_1) = \frac{B(\Upsilon \to \mu^+\mu^-)}{B(\Upsilon \to \mu^+\mu^-)} = \frac{G_F m_\Upsilon^2}{4\sqrt{2} \pi \alpha} \left( g_{P_{b\bar{b}}}^A \right)^2 \left( 1 - \frac{m_{A_1}^2}{m_\Upsilon^2} \right) F .
$$

Here $F \sim 1/2$ includes QCD corrections [13] and $B(\Upsilon \to \mu^+\mu^-) = (2.48 \pm 0.06)\%$. The SM-normalised pseudoscalar coupling $g_{P_{b\bar{b}}}^A$ is given by [9]

$$
g_{P_{b\bar{b}}}^A = \left( 1 + \frac{\sqrt{2} \langle \Delta_b \rangle}{v_1} \right)^{-1} \left[ - (\tan \beta + \Delta_{a^2}^b) O_{1i}^A + \Delta_{a^S}^b \frac{O_{2i}^A}{\cos \beta} \right],
$$

with $O^A$ the $2 \times 2$ orthogonal pseudoscalar mixing matrix, such that

$$
A_1 = O_{1i}^A a + O_{2i}^A a_S,
$$

where $a$ is the would-be CP-odd scalar in the MSSM limit and $a_S$ is the CP-odd singlet Higgs boson. In addition, $\Delta_{a^2,a^S}^b$ are the one-loop non-holomorphic Yukawa couplings of the states $a_{2,S}$ to $b$ quarks. At zero external momentum, they may be calculated by

$$
\Delta_{a^2,a^S}^b = i \sqrt{2} \left\langle \frac{\partial \Delta_b[\Phi_1, \Phi_2, S]}{\partial a_{2,S}} \right\rangle ,
$$

where $\Delta_b[\Phi_1, \Phi_2, S]$ is a Coleman-Wienberg type functional [14] of the background Higgs fields which encodes radiative corrections to the $b$ quark self-energy. Here $\langle \ldots \rangle$ denotes...
taking the VEV of the enclosed expression. The dominant contributions to $\Delta_b^{a_2,s}$ are
due to gluino-sbottom quark and chargino-stop quark loops. In the single-Higgs-insertion
approximation, neglecting subdominant terms proportional to the weak gauge coupling $\alpha_w$,
they may be given by

$$\Delta_b^{a_2} = -\frac{2\alpha_S}{3\pi} \tilde{M}_3 \mu I(\tilde{M}_Q^2, \tilde{M}_b^2, \tilde{M}_3^2) - \frac{h_t^2}{16\pi^2} \mu A_t I(\tilde{M}_Q^2, \tilde{M}_Q^2, \mu^2),$$

(7)

$$\Delta_b^{a_s} = -\frac{2\alpha_S}{3\pi} \tilde{M}_3 \mu \frac{v_2}{v_S} I(\tilde{M}_Q^2, \tilde{M}_b^2, \tilde{M}_3^2) - \frac{h_t^2}{16\pi^2} \mu A_t \frac{v_2}{v_S} I(\tilde{M}_Q^2, \tilde{M}_Q^2, \mu^2),$$

(8)

where $\tilde{M}_{Q,t,b}$ are the soft squark masses, $A_t$ is the top-squark soft trilinear coupling and
$\tilde{M}_3$ is the gluino mass. The one-loop function $I(x, y, z)$ is given by

$$I(x, y, z) = \frac{xy \ln(x/y) + yz \ln(y/z) + xz \ln(z/x)}{(x-y)(y-z)(x-z)}.$$  

(9)

In Fig. 1 we present results from a scan over the parameters

$$0 < \lambda < 0.5, \quad 0 < A_\lambda < 300 \text{GeV},$$

$$-0.5 < \kappa < 0.5, \quad 0 < A_\kappa < 20 \text{GeV},$$

(10)

whilst fixing $\tan \beta = 50$ and $\mu_{\text{eff}} = 120 \text{GeV}$. We require a light Higgs pseudoscalar $m_{A_1} < 9 \text{GeV}$ along with a lightest Higgs scalar $m_{H_1} > 114 \text{GeV}$, in agreement with constraints from LEP II. The soft-SUSY breaking parameters which enter the calculation of $\Delta_b^{a_2,s}$ are
taken to be equal at $M_{\text{SUSY}} = 600 \text{GeV}$. The branching ratio $B(\Upsilon \rightarrow \gamma A_1)$ is plotted against
the non-singlet fraction of $A_1$, described by the mixing matrix element $O_{11}^A$.

The threshold corrections are independent of the tree-level coupling proportional to
the pseudoscalar mixing, and enter the expression for $g_{A_1 bb}^P$ with opposing sign. For a
relatively large non-singlet component above few $\%$, the threshold corrections represent a
small suppression to the branching ratio of up to $\sim 10\%$. In the case of a highly singlet $A_1$
 boson, the threshold corrections become the dominant effect, producing a branching ratio of
the order $\sim 1 \times 10^{-6}$ across a significant region of parameter space. This prediction
is found to be generic for electroweak-scale soft SUSY-breaking terms around a TeV. At
the intersection of these regimes, the contributions cancel giving a highly suppressed decay
rate.

Fig. 2 shows results from a scan over the parameter range of Eq. (10) for $\tan \beta = 10$,
keeping $\mu = 120 \text{GeV}$ and the common soft-SUSY breaking scale $M_{\text{SUSY}} = 600 \text{GeV}$. Both the
doublet-singlet mixing and threshold correction contributions to $g_{A_1 bb}^P$ are $\tan \beta$
enhanced, such that the branching ratio at low $\tan \beta$ is smaller by $1 \sim 2$ orders of magnitude
across the full parameter space. Due to their common enhancement, the relative importance
Figure 1: The branching ratio $B(\Upsilon \rightarrow \gamma A_1)$ vs. the non-singlet fraction $O^A_{11}$ at $\tan \beta = 50$. The points in green (light grey) include the one-loop threshold effects $\Delta_a$, points in red (dark grey) neglect these corrections. Here $\mu_{\text{eff}} = 120$ GeV and $\lambda, \kappa, A_\lambda, A_\kappa$ are scanned over the range given in Eq. (10). All other soft-SUSY breaking parameters are taken to equal $M_{\text{SUSY}} = 600$ GeV. Experimental bounds are shown in dark blue (black) for a stable or invisibly decaying pseudoscalar and in light blue (grey) for a visibly decaying particle, assuming here $m_{A_1} \sim 5$ GeV. The full limits are strongly dependent on the value of $m_{A_1}$ and are less restrictive by one to two orders of magnitude for a heavy $A_1$ boson ($m_{A_1} > 8$ GeV).

of the two terms in Eq. (4) varies only slowly with $\tan \beta$, so that for all values of $\tan \beta \gtrsim 5$, minimal branching ratios are observed for singlet-doublet mixing around few $\times 0.1\%$. The magnitude of the branching ratio is not found to vary strongly with $M_{\text{SUSY}}$ or $\mu$, although the available parameter space consistent with our requirements $m_{H_1} > 114$ GeV, $m_{A_1} < 9$ GeV decreases as $\mu$ increases, such that small values of the singlet-doublet mixing $O^A_{11}$ do not appear.

The inclusion of threshold corrections can clearly alter the phenomenology of highly singlet light pseudoscalars in a dramatic way, allowing for the possibility of detectable $\Upsilon \rightarrow A_1 \gamma$ decays in a new corner of parameter space. In the limit of vanishing singlet-doublet mixing the tree level coupling of the $A_1$ boson to $\tau$ leptons also vanishes, however an analogous threshold correction also contributes to the $A_1\tau^+\tau^-$ coupling $g_{A_1\tau^+\tau^-}$, through a wino-stau loop. The pseudoscalar is therefore predicted to decay into $\tau^+\tau^-$ pairs with branching ratio of order one, for $2m_\tau < m_{A_1} < m_\Upsilon$, independently of the singlet-doublet
mixing. An order-of-magnitude estimate suggests that current $B$-factories should be sensitive to branching ratios of the order $\mathcal{B}(\Upsilon \to \gamma A_1) \lesssim 10^{-6}$, for observing such a final state.

At masses above $\sim 9$ GeV, the $A_1$ boson can mix with the $\eta_b$ meson. This can lead to significant enhancement or suppression of $\mathcal{B}(\Upsilon \to \gamma A_1)$ [15]. In addition, there is a broadening of the $A_1$ width, and the resonance in the energy spectrum of the recoil photon is less sharply peaked. There has been a suggestion to search for such a light Higgs boson through precision tests of lepton universality in the decays of the $\Upsilon$ [16]. Such searches would also be sensitive to decays in the zero-mixing limit. If the $A_1$ boson is below the $\tau^+\tau^-$ threshold, the dominant decay channels are $s\bar{s}, gg$ (and hence light mesons) or photon pairs, since the coupling to $c\bar{c}$ is $\tan\beta$ suppressed. This remains a favourable situation for the clean environment of an $e^+e^-$ collider, where these final states can be reliably measured.

Unfortunately, despite the tremendous production rates for $b$-mesons at the LHC, a discovery of the $A_1$ boson through this mechanism appears difficult. The final state consists of low-energy $\tau$ jets and a photon, neither of which presents a clean signal above $m_{A_1} < 2m_c$ in the limit of vanishing singlet-doublet mixing.

3The $\tan\beta$ suppression would also exclude the possibility of an observable signal $J/\psi \to \gamma A_1$ for $m_{A_1} < 2m_c$ in the limit of vanishing singlet-doublet mixing.
background activity. An alternative production mechanism has been suggested in [17], which considers instead the process $pp \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^- A_1$ in the limit of vanishing doublet-singlet mixing. The possibility for observing such a signal is strongly dependent on the masses and decay channels of both the lightest chargino and the $A_1$ boson.

If both terms contributing to $g_{A_{1}b\bar{b}}^{P}$ are of similar magnitude, typically for around $\sim 0.5\%$ mixing, detection of the $A_1$ boson may be extremely challenging. In this case, the branching fraction of $\Upsilon \rightarrow A_1 \gamma$ becomes extremely suppressed. An alternative experimental strategy is to look for $A_1$ pair production from Higgs boson decays [18]. In the small singlet-doublet mixing scenario at large $\tan\beta$, the lightest CP-even Higgs boson $H_1$ is highly SM-like, and the branching fraction $\mathcal{B}(H_1 \rightarrow A_1 A_1)$ is conservatively bounded from above at around $\sim 10^{-3}$. Associated production of the $A_1$ boson with a chargino pair would remain a possibility.

In conclusion, we have shown that the branching ratio for production of a light Higgs pseudoscalar in $\Upsilon(1s)$ decays does not vanish in the absence of doublet-singlet mixing. We found that the decay $\Upsilon \rightarrow \gamma A_1$ is predicted to be observable at existing experimental facilities if supersymmetry is broken at the TeV scale with large $\tan\beta$. In the event of a cancellation between the threshold corrections and tree-level mixing contributions to the $A_1 b\bar{b}$ coupling the branching ratio may be highly suppressed even though the doublet-singlet mixing is still significant, and further phenomenological considerations would be needed. We hope to return to this issue in a future communication.

Acknowledgments

The author would like to thank Apostolos Pilaftsis and Roger Barlow for helpful discussions. This research was supported by the U.K. Science and Technology Facilities Council.

References

[1] 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 [arXiv:hep-ph/9505326];
   M. Bastero-Gil, C. Hugonie, S. F. King, D. P. Roy and S. Vempati, Phys. Lett. B 489, 359 (2000) [arXiv:hep-ph/0006198];
   D. J. Miller, R. Nevzorov and P. M. Zerwas, Nucl. Phys. B 681 (2004) 3
U. Ellwanger, J. F. Gunion and C. Hugonie, JHEP 0507 (2005) 041 [arXiv:hep-ph/0503203].

[2] S. A. Abel, S. Sarkar and P. L. White, Nucl. Phys. B 454 (1995) 663 [arXiv:hep-ph/9506359].

[3] C. Panagiotakopoulos and K. Tamvakis, Phys. Lett. B 446 (1999) 224 [arXiv:hep-ph/9809475];
C. Panagiotakopoulos and K. Tamvakis, Phys. Lett. B 469 (1999) 145 [arXiv:hep-ph/9908351].

[4] B. A. Dobrescu, G. L. Landsberg and K. T. Matchev, Phys. Rev. D 63 (2001) 075003 [arXiv:hep-ph/0005308];
B. A. Dobrescu and K. T. Matchev, JHEP 0009 (2000) 031 [arXiv:hep-ph/0008192].

[5] R. Dermisek and J. F. Gunion, Phys. Rev. D 76 (2007) 095006 [arXiv:0705.4387 [hep-ph]].

[6] C. Panagiotakopoulos and A. Pilaftsis, Phys. Rev. D 63 (2001) 055003 [arXiv:hep-ph/0008268];
A. Dedes, C. Hugonie, S. Moretti and K. Tamvakis, Phys. Rev. D 63 (2001) 055009 [arXiv:hep-ph/0009125];
C. Balazs, M. S. Carena, A. Freitas and C. E. M. Wagner, JHEP 0706 (2007) 066 [arXiv:0705.0431 [hep-ph]]; 
S. Hesselbach, D. J. Miller, G. Moortgat-Pick, R. Nevzorov and M. Trusov, arXiv:0712.2001 [hep-ph].

[7] R. Dermisek, J. F. Gunion and B. McElrath, Phys. Rev. D 76 (2007) 051105 [arXiv:hep-ph/0612031].

[8] G. Hiller, Phys. Rev. D 70 (2004) 034018 [arXiv:hep-ph/0404220].

[9] R. N. Hodgkinson and A. Pilaftsis, Phys. Rev. D 76 (2007) 015007 [arXiv:hep-ph/0612188].

[10] H. Albrecht et al. [ARGUS Collaboration], Phys. Lett. B 154 (1985) 452;
H. Albrecht et al. [ARGUS Collaboration], Z. Phys. C 29 (1985) 167;
P. Franzini et al., Phys. Rev. D 35 (1987) 2883.

[11] R. Balest et al. [CLEO Collaboration], Phys. Rev. D 51 (1995) 2053.

[12] F. Wilczek, Phys. Rev. Lett. 39 (1977) 1304;
H. E. Haber, A. S. Schwarz and A. E. Snyder, Nucl. Phys. B 294 (1987) 301.
[13] M. I. Vysotsky, Phys. Lett. B 97 (1980) 159. P. Nason, Phys. Lett. B 175 (1986) 223.

[14] S. R. Coleman and E. Weinberg, Phys. Rev. D 7 (1973) 1888.

[15] E. Fullana and M. A. Sanchis-Lozano, Phys. Lett. B 653 (2007) 67
    arXiv:hep-ph/0702190

[16] M. A. Sanchis-Lozano, J. Phys. Soc. Jap. 76 (2007) 044101
    arXiv:hep-ph/0610046.

[17] A. Arhrib, K. Cheung, T. J. Hou and K. W. Song, JHEP 0703 (2007) 073
    arXiv:hep-ph/0606114.

[18] U. Ellwanger, J. F. Gunion, C. Hugonie and S. Moretti, arXiv:hep-ph/0305109
    U. Ellwanger, J. F. Gunion and C. Hugonie, JHEP 0507 (2005) 041
    arXiv:hep-ph/0503203;
    S. Chang, P. J. Fox and N. Weiner, JHEP 0608 (2006) 068
    arXiv:hep-ph/0511250;
    K. Cheung, J. Song and Q. S. Yan, Phys. Rev. Lett. 99 (2007) 031801
    arXiv:hep-ph/0703149;
    M. Carena, T. Han, G. Y. Huang and C. E. M. Wagner, arXiv:0712.2466 [hep-ph];
    J. R. Forshaw, J. F. Gunion, L. Hodgkinson, A. Papaefstathiou and A. D. Pilkington,
    arXiv:0712.3510 [hep-ph].