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

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ABSTRACT

We investigate the effects of the radiatively-generated $\tan\beta$-enhanced Higgs-singlet Yukawa couplings on the decay $\Upsilon \to \gamma A_1$ 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 μ problem of the Minimal Supersymmetric extension of the Standard Model (MSSM) in a natural way [1]. The μ parameter of the MSSM is replaced with an additional gauge singlet Higgs superfield $\hat{S}$ and an effective doublet mixing term μ_{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_{SUSY}.

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

$$\mathcal{W}_{\text{Higgs}} = \lambda \hat{S}\hat{H}_1\hat{H}_2 + \kappa \hat{S}^3, \tag{1}$$

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

$$\mathcal{L}^{\text{soft}}_{\text{Higgs}} = \lambda A_\lambda S \Phi_1 \Phi_2 + \kappa A_\kappa S^3, \tag{2}$$

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 $\mu_{\text{eff}} = \frac{\lambda v_S}{\sqrt{2}}$. 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.[1]

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_λ, A_κ [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_λ being radiatively generated at one loop and non-zero A_κ at two loops [5].

For a sufficiently light A_1 boson, observation in the decay $\Upsilon(1s) \to \gamma A_1$ becomes a possibility.[2] 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.

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[1] It is also possible to evade the LEP bound if H_1 decays into the lightest pseudoscalars, with branching ratio $\mathcal{B}(H_1 \to A_1 A_1) > 0.7$ [5]. This requirement leads to a lower bound on the doublet component of A_1, $\mathcal{O}^A_{11} > 0.04$. Since the radiative corrections are subdominant in this region, we do not include these points in our results.

[2] 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 μ_{eff} fixed. The radiative coupling of the singlet pseudoscalar to fermions, whose effects we consider here, will also vanish in this limit.

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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 Υ decays place a 90% confidence level upper bound on the branching ratio $\mathcal{B}(\Upsilon \to \gamma A_1) \lesssim 1 \times 10^{-4}$ for a light particle $m_{A_1} < 8$ 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 tighten the constraints to $\sim 10^{-5}$ for a stable or invisibly decaying A_1 boson [11].

The branching ratio for Υ decays through the Wilczek mechanism [8,12] is given by

$$\frac{\mathcal{B}(\Upsilon \to \gamma A_1)}{\mathcal{B}(\Upsilon \to \mu^+\mu^-)} = \frac{G_F m_\Upsilon^2}{4\sqrt{2}\pi\alpha}\left(g^P_{A_1 bb}\right)^2\left(1 - \frac{m_{A_1}^2}{m_\Upsilon^2}\right)F. \tag{3}$$

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

$$g^P_{A_i bb} = \left(1 + \frac{\sqrt{2}\langle\Delta_b\rangle}{v_1}\right)^{-1}\left[-\left(\tan\beta + \Delta_b^{a_2}\right)\mathcal{O}^A_{1i} + \Delta_b^{a_S}\frac{\mathcal{O}^A_{2i}}{\cos\beta}\right], \tag{4}$$

with $\mathcal{O}^A$ the 2×2 orthogonal pseudoscalar mixing matrix, such that

$$A_1 = \mathcal{O}^A_{11} a + \mathcal{O}^A_{21} a_S, \tag{5}$$

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, $\Delta_b^{a_{2,S}}$ 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

$$\Delta_b^{a_{2,S}} = i\sqrt{2}\left\langle\frac{\partial\Delta_b[\Phi_1, \Phi_2, S]}{\partial a_{2,S}}\right\rangle, \tag{6}$$

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 $\langle\cdots\rangle$ 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 α_w, they may be given by

$$\begin{aligned}\Delta_b^{a_2} &= -\frac{2\alpha_S}{3\pi}\tilde{M}_3\mu I\left(\tilde{M}_Q^2, \tilde{M}_b^2, \tilde{M}_3^2\right) \\ &\quad + \frac{h_t^2}{16\pi^2}\mu A_t I\left(\tilde{M}_Q^2, \tilde{M}_t^2, \mu^2\right), \end{aligned}\tag{7}$$

$$\begin{aligned}\Delta_b^{a_S} &= -\frac{2\alpha_S}{3\pi}\tilde{M}_3\mu\frac{v_2}{v_S} I\left(\tilde{M}_Q^2, \tilde{M}_b^2, \tilde{M}_3^2\right) \\ &\quad - \frac{h_t^2}{16\pi^2}\mu A_t\frac{v_2}{v_S} I\left(\tilde{M}_Q^2, \tilde{M}_t^2, \mu^2\right), \end{aligned}\tag{8}$$

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

$$I(x,y,z) = \frac{xy\ln(x/y) + yz\ln(y/z) + xz\ln(z/x)}{(x-y)(y-z)(x-z)}. \tag{9}$$

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

$$\begin{aligned}&0 < \lambda < 0.5, \quad 0 < A_\lambda < 300\ \text{GeV}, \\ &-0.5 < \kappa < 0.5, \quad 0 < A_\kappa < 20\ \text{GeV}, \end{aligned}\tag{10}$$

whilst fixing $\tan\beta = 50$ and $\mu_{\text{eff}} = 120$ GeV. We require a light Higgs pseudoscalar $m_{A_1} < 9$ GeV along with a lightest Higgs scalar $m_{H_1} > 114$ 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 $M_{\text{SUSY}} = 600$ GeV. The branching ratio $\mathcal{B}(\Upsilon \to \gamma A_1)$ is plotted against the non-singlet fraction of A_1, described by the mixing matrix element $\mathcal{O}^A_{11}$.

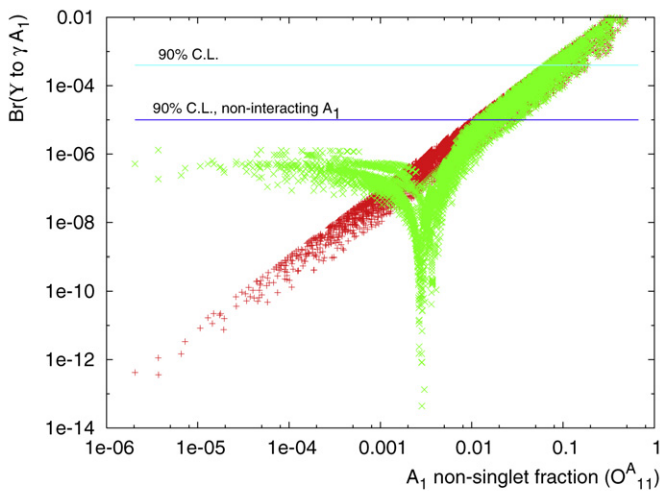

Fig. 1. The branching ratio $\mathcal{B}(\Upsilon \to \gamma A_1)$ vs. the non-singlet fraction O^A_{11} at $\tan\beta = 50$. The points in green (light grey) include the one-loop threshold effects $\Delta_b^{a_S}$, points in red (dark grey) neglect these corrections. Here $\mu_{\text{eff}} = 120$ GeV 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). (For interpretation of the references to colour in this figure legend, the reader is referred to the web version of this Letter.)

The threshold corrections are independent of the tree-level coupling proportional to the pseudoscalar mixing, and enter the expression for $g^P_{A_1 bb}$ 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 $\mu = 120$ GeV and the common soft-SUSY breaking scale $M_{\text{SUSY}} = 600$ GeV. Both the doublet-singlet mixing and threshold correction contributions to $g^P_{A_1 bb}$ 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 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_{SUSY} or μ, although the available parameter space consistent with out requirements $m_{H_1} > 114$ GeV, $m_{A_1} < 9$ GeV decreases as μ 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 $\Upsilon \to A_1\gamma$ 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 τ leptons also vanishes, however an analogous threshold correction also contributes to the $A_1\tau^+\tau^-$ coupling $g^P_{A_1\tau^+\tau^-}$, through a wino-stau loop. The pseudoscalar is therefore predicted to decay into $\tau^+\tau^-$

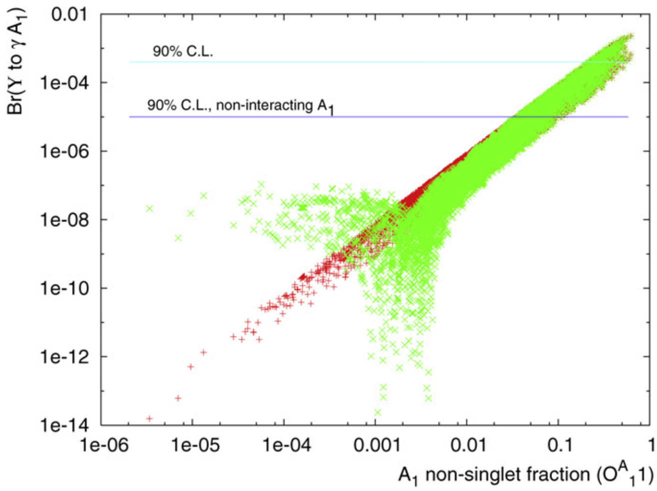

Fig. 2. The branching ratio $\mathcal{B}(\Upsilon \to \gamma A_1)$ vs. the non-singlet fraction O^A_{11} at $\tan\beta = 10$. Here $\mu_{\text{eff}} = 120$ GeV and $M_{\text{SUSY}} = 600$ GeV, with $\lambda, \kappa, A_\lambda, A_\kappa$ scanned over the range given in Eq. (10). Points in green (light grey) include the one-loop threshold effects $\Delta_b^{\alpha_s}$, points in red (dark grey) neglect these corrections. Experimental bounds are shown as for Fig. 1. (For interpretation of the references to colour in this figure legend, the reader is referred to the web version of this Letter.)

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 ~ 9 GeV, the A_1 boson can mix with the η_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 Υ [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.[3] 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 τ jets and a photon, neither of which presents a clean signal above background activity. An alternative production mechanism has been suggested in [17], which considers instead the process $pp \to \tilde{\chi}_1^+\tilde{\chi}_1^- A_1$ 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 $g^P_{A_1 b\bar{b}}$ 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 $\Upsilon \to A_1\gamma$ 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 $\mathcal{B}(H_1 \to A_1A_1)$ 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.

[3] The $\tan\beta$ 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.

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 $\Upsilon \to \gamma A_1$ 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.

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