Scope of Higgs production in association with a bottom quark pair in probing the Higgs sector of the NMSSM at the LHC

M. M. Almarashi∗ and S. Moretti∗∗

∗Department of Physics, Faculty of Science,
Taibah University, P.O.Box 30002, Madinah, Saudi Arabia
∗∗School of Physics & Astronomy,
University of Southampton, Southampton, SO17 1BJ, UK

Abstract

We review the potential of the LHC to detect a very light CP-odd Higgs boson of the NMSSM, \( a_1 \), through its direct production in association with a bottom-quark pair at large \( \tan\beta \). We also review the LHC discovery potential of the two lightest CP-even Higgs states, decaying into two lighter Higgs states or into the lightest CP-odd Higgs state and the \( Z \) gauge boson.

1 Introduction

The Minimal Supersymmetric Standard Model (MSSM) \([1]\) is probably one of the most studied Beyond the SM (BSM) scenarios. However, this model suffers from two critical flaws: the \( \mu \)-problem \([2]\) and the little hierarchy problem. The former flaw results from the fact that the Superpotential has a dimensional parameter, \( \mu \) (the so-called ‘Higgs(ino) mass parameter’), whose natural value would be either 0 or \( m_{\text{Pl}} \) (the Planck mass). However, phenomenologically, in order to achieve Electro-Weak Symmetry Breaking (EWSB), \( \mu \) is required to take values of the order of the EW scale or possibly up to the TeV range. The latter flaw emerged first from LEP, which failed to detect a light CP-even Higgs boson, \( h \), thereby imposing severe constraints on \( m_h \). For this kind of Higgs state to pass the experimental constraints, large higher order corrections from both the SM and SUSY particle spectrum are required. The largest contributions come from the third generation, top quarks and squarks. However, these required large corrections seem quite unnatural. Recall in fact that at tree level the lightest CP-even Higgs boson mass of the MSSM is less than \( M_Z \). Even

\footnote{al_marashi@hotmail.com}

\footnote{s.moretti@soton.ac.uk}
recent LHC results, hinting at the possible existence of a SM-like Higgs state with mass of 124–126 GeV [3, 4], weaken the MSSM assumption, as such mass values are really extreme in such a SUSY realisation, towards the very end of the allowed mass range.

The simplest SUSY realisation beyond the MSSM that can solve these two problems at once is the NMSSM (for reviews see [5, 6]). This scenario includes a Higgs singlet Superfield in addition to the two MSSM-type Higgs doublets, giving rise to seven Higgs states: three CP-even Higgses $h_{1,2,3}$ ($m_{h_1} < m_{h_2} < m_{h_3}$), two CP-odd Higgses $a_{1,2}$ ($m_{a_1} < m_{a_2}$) and a pair of charged Higgses $h^\pm$. When the scalar component of the singlet Superfield acquires a Vacuum Expectation Value (VEV), an ‘effective’ $\mu$-term, $\mu_{\text{eff}}$, will be automatically generated and can rather naturally have values of order of the EW to TeV scale, as required [7]. In addition, in the NMSSM, the little hierarchy problem can be relieved [8, 9], since a SM-like scalar Higgs boson in the NMSSM context requires less (s)quark corrections than those in the MSSM or it can have mass less than the LEP bound due to unconventional decays over some regions of the NMSSM parameter space. In fact, currently, the NMSSM can also explain not only the LHC excess [10, 11] a possible LEP excess and is definitely preferred by EW global fits [12, 13, 14].

2 The NMSSM Superpotential

The Superpotential of the NMSSM is given by

$$W = h_u \hat{Q} \hat{H}_u \hat{U}^c - h_d \hat{Q} \hat{H}_d \hat{D}^c - h_u \hat{L} \hat{H}_u \hat{E}^c + \lambda \hat{S} \hat{H}_u \hat{H}_d + \frac{1}{3} \kappa \hat{S}^3,$$

(1)

where $h_u$, $h_d$, $h_e$, $\lambda$ and $\kappa$ are dimensionless couplings. The term $\lambda \hat{S} \hat{H}_u \hat{H}_d$ has been introduced to solve the $\mu$-problem of the MSSM Superpotential. However, the Superpotential in Eq. (1) without the term $\frac{1}{3} \kappa \hat{S}^3$ gives rise to an extra global $U(1)$ symmetry, the so-called Peccei-Quinn symmetry $U(1)_{\text{PQ}}$ [15, 16]. Once the Higgs bosons take on VEVs, this symmetry will break spontaneously and lead to the appearance of a CP-odd scalar, called a Peccei-Quinn axion. In fact, this axion has not been seen experimentally. In addition, there are severe astrophysical and cosmological constraints on $\lambda$, that is $10^{-7} < \lambda < 10^{-10}$ [17]. These constraints necessitate a very large value of $< S >$ in order to solve the $\mu$-problem. So, this is not a satisfactory way to solve the latter.

One elegant way to solve the $\mu$-problem is to break the $U(1)_{\text{PQ}}$ by introducing an additional term in the Superpotential. This is the last term in Eq. (1) and consequently the axion can be avoided. However, introducing this new term in the Superpotential enables one to break the PQ symmetry and the Superpotential still have a discrete $Z_3$ symmetry. This discrete symmetry is spontaneously broken when the additional complex scalar field acquires a VEV and that will lead to the domain wall problem. That is, during the EW phase transition of the early universe, this broken symmetry causes a dramatic change of the universe evolution and creates unobserved large anisotropies in the cosmic microwave background [18].

\footnote{We will give more explanations of this in Sec. 4.}
In order to solve the domain wall problem, one needs to break the \( \mathbb{Z}_3 \) symmetry by introducing higher order operators at the Plank scale. However, these operators generate quadratic tadpoles for the singlet. So, one also needs to impose a new discrete invariance, a \( \mathbb{Z}_2 \) symmetry, on these operators in order to get rid of the dangerous tadpole contributions, see [5] for more details.

3 The Higgs sector of the NMSSM

The NMSSM Higgs sector contains two Higgs doublets and one Higgs singlet:

\[
H_d = \begin{pmatrix} H_d^0 \\ H_d^+ \end{pmatrix}, \quad H_u = \begin{pmatrix} H_u^0 \\ H_u^+ \end{pmatrix}, \quad S.
\]

The scalar potential for the Higgs fields can be written as [19]:

\[
V_H = V_F + V_D + V_{\text{soft}},
\]

where

\[
V_F = |\lambda S|^2(|H_u|^2 + |H_d|^2) + |\lambda H_u H_d + \kappa S|^2,
\]

\[
V_D = \frac{g_1^2 + g_2^2}{8}(|H_d|^2 - |H_u|^2)^2 + \frac{1}{2}g_2^2|H_d^+H_d|^2,
\]

\[
V_{\text{soft}} = m_{H_u}^2 H_u H_u^+ + m_{H_d}^2 H_d H_d^+ + m_S^2 S S^+ + \left( \lambda A \lambda S H_u H_d + \frac{1}{3}\kappa A_S S^3 + \text{h.c.} \right).
\]

To generate EWSB, the Higgs fields should have VEVs. In fact, if one assumes that the VEVs are real and positive, they can be described by

\[
<H_d> = \frac{1}{\sqrt{2}} \begin{pmatrix} v_d \\ 0 \end{pmatrix}, \quad <H_u> = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_u \end{pmatrix}, \quad <S> = \frac{1}{\sqrt{2}} v_s.
\]

At the physical minimum of the scalar potential, \( V_H \), the soft mass parameters of the Higgs fields are related to the VEVs through the following relations [19]:

\[
m_{H_d}^2 = \frac{g_1^2}{8} (v_u^2 - v_d^2) - \frac{1}{2} \lambda^2 v_u^2 + \frac{1}{2} (\sqrt{2} A \lambda + \kappa v_s) \lambda u \frac{v_u}{v_d} - \frac{1}{2} \lambda^2 v_s^2,
\]

\[
m_{H_u}^2 = \frac{g_1^2}{8} (v_d^2 - v_u^2) - \frac{1}{2} \lambda^2 v_d^2 + \frac{1}{2} (\sqrt{2} A \lambda + \kappa v_s) \lambda v_s \frac{v_d}{v_u} - \frac{1}{2} \lambda^2 v_s^2,
\]

\[
m_S^2 = -\kappa^2 v_s^2 - \frac{1}{2} \lambda^2 v_d^2 + \kappa \lambda v_u v_d + \frac{1}{2} \lambda^2 A \lambda u \frac{v_u v_d}{v_s} - \frac{1}{2} \kappa^2 A_S v_s.
\]

The physical Higgs states arise after the Higgs fields acquire VEVs and rotate away the Goldstone modes. As a result, the potential can be written as

\[
V_H = m_{h^+}^2 h^+ h^- + \frac{1}{2} (P_1 P_2) M_P \left( \begin{array}{c} p_1 \\ p_2 \end{array} \right) + \frac{1}{2} (S_1 S_2 S_3) M_S \left( \begin{array}{c} s_1 \\ s_2 \\ s_3 \end{array} \right).
\]
The masses of charged Higgs fields, $h^\pm$, at tree level are
\[ m_{h^\pm}^2 = m_A^2 + M_W^2 - \frac{1}{2}(\lambda v)^2, \] (12)
where
\[ m_A^2 = \sqrt{2} \frac{\mu_{\text{eff}}}{\sin 2\beta} \left( A + \frac{\kappa \mu_{\text{eff}}}{\lambda} \right). \] (13)

Using the minimisation conditions, one can obtain the mass matrices in the scalar and pseudoscalar sectors. First, the mass matrix for CP-even Higgs states at tree level has the following entries [19]:
\[ M_{S11} = m_A^2 + \left( M_Z^2 - \frac{1}{2}(\lambda v)^2 \right) \sin^2 2\beta, \] (14)
\[ M_{S12} = -\frac{1}{2} \left( M_Z^2 - \frac{1}{2}(\lambda v)^2 \right) \sin 4\beta, \] (15)
\[ M_{S13} = -\frac{1}{2} \left( m_A^2 \sin 2\beta + 2 \frac{\kappa \mu_{\text{eff}}^2}{\lambda} \right) \left( \frac{\lambda v}{\sqrt{2} \mu_{\text{eff}}} \right) \cos 2\beta, \] (16)
\[ M_{S22} = M_Z^2 \cos^2 2\beta + \frac{1}{2}(\lambda v)^2 \sin^2 2\beta, \] (17)
\[ M_{S23} = \frac{1}{8} m_A^2 \sin^2 2\beta \left( \frac{\lambda^2 v^2}{\mu_{\text{eff}}^2} + \frac{4 \kappa \mu_{\text{eff}}^2}{\lambda^2} + \frac{\kappa A}{\lambda} \right) - \frac{1}{4} \lambda \kappa v^2 \sin 2\beta. \] (18)

Second, the mass matrix for CP-odd Higgs states at tree level has the following entries [19]:
\[ M_{P11} = m_A^2, \] (20)
\[ M_{P12} = \frac{1}{2} \left( m_A^2 \sin 2\beta - \frac{6 \kappa \mu_{\text{eff}}^2}{\lambda} \right) \left( \frac{\lambda v}{\sqrt{2} \mu_{\text{eff}}} \right), \] (21)
\[ M_{P22} = \frac{1}{8} \left( m_A^2 \sin 2\beta + 6 \frac{\kappa \mu_{\text{eff}}^2}{\lambda} \right) \left( \frac{\lambda^2 v^2}{\mu_{\text{eff}}^2} \sin 2\beta - 3 \frac{\kappa \mu_{\text{eff}} A}{\lambda} \right). \] (22)

To a good approximation, at large \(\tan\beta\) and large \(m_A\), the tree level neutral Higgs boson masses are given by the following expressions [19]:
\[ m_{a_1}^2 = \frac{-3 \kappa \mu_{\text{eff}} A}{\lambda}, \] (23)
\[ m_{a_2}^2 = m_A^2 \left( 1 + \frac{1}{8} \left( \frac{\lambda^2 v^2}{\mu_{\text{eff}}^2} \right) \sin^2 2\beta \right), \] (24)
\[ m_{h_{1/2}}^2 = \frac{1}{2} \left\{ M_Z^2 + \frac{\kappa \mu_{\text{eff}}}{\lambda} \left( \frac{4 \kappa \mu_{\text{eff}}}{\lambda} + A_\kappa \right) \right\} \]
\[ \pm \sqrt{ \left( M_Z^2 - \frac{\kappa \mu_{\text{eff}}}{\lambda} \left( \frac{4 \kappa \mu_{\text{eff}}}{\lambda} + A_\kappa \right) \right)^2 + \frac{\lambda^2 v^2}{2 \mu_{\text{eff}}^2} \left[ 4 \mu_{\text{eff}}^2 - m_A^2 \sin^2 2\beta \right]^2} \} \]  
(25)

\[ m_{h_3}^2 = m_A^2 \left( 1 + \frac{1}{8} \left( \frac{\lambda^2 v^2}{\mu_{\text{eff}}^2} \right) \sin^2 2\beta \right). \]  
(26)

4  LHC phenomenology of the NMSSM Higgs sector

Because of the existence of a singlet Superfield in the NMSSM, the latter is phenomenologically richer than the MSSM. In fact, the NMSSM has seven Higgs states and five neutralinos compared to only five Higgs states and four neutralinos in the MSSM. As a consequence, the search for Higgs bosons in the context of the NMSSM at present and future colliders is a big challenge and more complicated than in the MSSM.

It was mentioned before that the mass of the lightest CP-even Higgs boson in the MSSM, \( m_h \), at tree level should be less than \( M_Z \). So, large radiative corrections, mainly from top and stop loops, are required to pass the LEP lower limit on the Higgs mass. In fact, to achieve this we need large stop masses, which only contribute logarithmically in the loop corrections. This large discrepancy between top and stop masses causes essentially a fine tuning problem [5] (the aforementioned little hierarchy problem).

As for the NMSSM, the situation is quite different. Assuming CP-conservation in the Higgs sector, the upper mass bound for the lightest CP-even Higgs boson at tree level is given by

\[ m_{h_1}^2 \leq M_Z^2 \left( \cos^2 (2\beta) + \frac{2\lambda^2 \sin^2 (2\beta)}{g_1^2 + g_2^2} \right). \]  
(27)

The last term in this expression can lift \( m_{h_1} \) up to 10 GeV higher than the corresponding mass of the SM-like Higgs mass. So, smaller loop corrections are required to pass the lower bound on the SM-like Higgs mass. However, since the higher order corrections are similar to those in the MSSM, the upper mass bound reaches 135 – 140 GeV for maximal stop mixing and \( \tan \beta = 2 \) [20, 21], however, this configuration is already excluded in the MSSM by LEP data. Finally, notice that the corrections to the lightest CP-even Higgs boson mass are already calculated at complete one loop level [22, 23, 24] and also at the dominant two loop level [21].

Furthermore, the most interesting property of the NMSSM that can solve the little hierarchy problem of the MSSM comes from the fact that in large areas of the NMSSM parameter space Higgs-to-Higgs decays are kinematically open. For instance, the existence of the lightest CP-odd Higgs boson \( a_1 \) with mass less than \( \frac{1}{2} m_{h_1} \) is quite natural in the NMSSM, see, e.g., [25]. In fact, the Branching ratio (Br) for the decay \( h_1 \to a_1 a_1 \), \( \text{Br}(h_1 \to a_1 a_1) \), can be dominant in large regions of parameter space and as a result the \( \text{Br}(h_1 \to b\bar{b}) \) is suppressed.
This unconventional decay channel is so important as it could explain the $2.3\sigma$ excess observed at LEP for a Higgs mass, $m_H$, around 100 GeV as shown in figure 1. The reduced coupling in the figure is defined as follows:

$$\xi^2 = \left( \frac{g_{HZZ}}{g_{HZZ}^{SM}} \right)^2. \quad (28)$$

Here, $g_{HZZ}^{SM}$ denotes the SM $HZZ$ coupling while $g_{HZZ}$ the non-standard coupling. As it is clear from the plot the excess occurs when the $\text{Br}(H_{SM} \rightarrow b\bar{b})$ times $\xi^2$ gives about 20%. In the context of the NMSSM, one can explain this excess in two ways. Firstly, a SM-like Higgs boson, $h_{1,2}$, can decay dominantly into a pair of $a_1$'s and so the $\text{Br}(H \rightarrow b\bar{b})$ is suppressed [12, 13, 14]. This scenario can relieve the little hierarchy problem but requires that $m_{a_1} < 2m_b$. (Notice that this mass region is currently highly constrained by ALEPH [26] and BaBar [27] data.) In fact, there is also another possibility in the NMSSM that can explain the LEP excess due to the fact that the $\text{Br}(a_1 \rightarrow \gamma\gamma)$ can be dominant when the $a_1$ is highly singlet and again, as a result, the $\text{Br}(a_1 \rightarrow b\bar{b})$ is suppressed even with $m_{a_1} > 2m_b$ [25, 28]. Secondly, a CP-even Higgs boson, $h_1$, has a reduced coupling with $\xi \lesssim 0.4$ [29], due to the mixing between the Higgs singlet and doublets. Notice that neither the SM nor the MSSM can explain such modest excess, as they have a $\text{Br}(H/h \rightarrow b\bar{b})$ which is always dominant, hence yielding an excess much above the experimental limit. Besides, the NMSSM could also explain the recent excess observed at the LHC for a Higgs mass around 125 GeV [11, 30, 31, 32].

**Figure 1:** Upper limit on the ratio $\xi$ from LEP, where the SM $\text{Br}(H \rightarrow b\bar{b})$ and $\text{Br}(H \rightarrow \tau^+\tau^-)$ are assumed. Full line represents the observed limit and dashed line represents the expected limit. The green band and yellow band are within 68% and 95% probability, respectively [6].
The discovery of one or more Higgs boson at present or future colliders will open a new era in the realm of particle physics. In fact, many efforts have been made to detect such type of particles at colliders. In regard to the Higgs sector of the NMSSM, there has been some work devoted to explore the detectability of at least one Higgs boson at the LHC and the Tevatron. In particular, some efforts have been made to extend the ‘No-lose theorem’ of the MSSM (recall that this states that at least one Higgs boson of the MSSM will be found at the LHC via the usual SM-like production and decay channels throughout the entire MSSM parameter space [33, 34, 35]) to the case of the NMSSM [25, 28, 36, 37, 38, 39, 40, 41]. By assuming that Higgs-to-Higgs decays are not allowed, it was realised that at least one Higgs boson of the NMSSM will be discovered at the LHC. However, this theorem could be violated if Higgs-to-SUSY decays are kinematically allowed (e.g., into neutralino or chargino pairs, yielding invisible Higgs signals).

Because of the large number of input parameters of the NMSSM, it is practically very difficult to make a continuous scan over all the NMSSM parameter space. The alternative way to do the scan is by resorting to benchmark points in parameter space. (For example, for benchmark points in the NMSSM parameter space, see Refs. [36] and [42], for the unconstrained and constrained case, respectively.) Either way, one can distinguish between two scenarios in which Higgs-to-Higgs decays are either kinematically allowed or not.

So far, there is no conclusive evidence that the ‘No-lose theorem’ can be confirmed in the context of the NMSSM. In order to establish the theorem for the NMSSM, Higgs-to-Higgs decays should be taken into account, in particular the decay \( h_1 \to a_1a_1 \). Such a decay can in fact be dominant in large regions of the NMSSM parameter space, for instance, for small \( A_k \) [25], and may not give Higgs signals with sufficient significance at the LHC. However, a very light CP-odd Higgs boson, \( a_1 \), can be produced in association with chargino pairs [43] and in neutralino decays [44] at the LHC.

The importance of Higgs-to-Higgs decays in the context of the NMSSM has been emphasised over the years in much literature in all the above respects, see, e.g., Refs. [9, 15, 40, 47]. Eventually, it was realised that Vector Boson Fusion (VBF) \(^4\) could be a viable production channel to detect \( h_{1,2} \to a_1a_1 \) at the LHC, in which the Higgs pair decays into \( j j \tau^+\tau^- \) [36, 48]. Some scope could also be afforded by a \( 4\tau \) signature in both VBF and Higgs-strahlung (off gauge bosons) [49]. The gluon-fusion channel too could be a means of accessing \( h_1 \to a_1a_1 \) decays, so long that the two light CP-odd Higgs states decay into four muons [50] or into two muons and two taus [51]. Such results were all supported by simulations based on parton shower Monte Carlo (MC) programs and some level of detector response. For a recent survey of the ‘No-lose theorem’ in the NMSSM context, see Ref. [52].

Besides, there have also been some attempts to distinguish the NMSSM Higgs sector from the MSSM one, by affirming a ‘More-to-gain theorem’ [25, 28, 39, 40, 41, 53, 54] (that is, to recall, to assess whether there exist some areas of the NMSSM parameter space where more and/or different Higgs bosons can be discovered at the LHC compared with what is expected from the MSSM). Some comparisons between NMSSM and MSSM phenomenology, specifically in the Higgs sectors of the two SUSY realisations, can be found in [55].

In this paper, we review the LHC discovery potential for the NMSSM Higgs states assum-

\(^4\)Which is dominated by \( W^+W^-\)-fusion over \( ZZ\)-one.
ing as production mechanism of these states associated production with bottom-antibottom quark pairs. Generally, the heaviest CP-even Higgs, \( h_3 \), the heaviest CP-odd Higgs, \( a_2 \), and the \( h^\pm \) states have very large masses, above the TeV scale, in particular at large values of \( \tan\beta \), making their discovery at the LHC very difficult. So, we will focus on the LHC discovery potential through this production mode of the lightest CP-odd Higgs boson, \( a_1 \), and of the lightest two CP-even Higgs states, \( h_1 \) and \( h_2 \).

### 5 Parameter space scan

As intimated already, due to the large number of parameters in the NMSSM, it is practically not feasible to do a comprehensive scan over all of them. These parameters can however be reduced significantly by assuming certain conditions of unification. Here, since the mechanism of SUSY breaking is still unknown, to explore the NMSSM Higgs sector, we have performed a general scan in parameter space by fixing the soft SUSY breaking terms at high scale to reduce their contributions to the outputs of the parameter scans. Consequently, we are left with six independent inputs. Our parameter space is in particular defined through the Yukawa couplings \( \lambda \) and \( \kappa \), the soft trilinear terms \( A_\lambda \) and \( A_\kappa \) plus \( \tan\beta \) (the ratio of the VEVs of the two Higgs doublets) and \( \mu_{\text{eff}} = \lambda \langle S \rangle \) (where \( \langle S \rangle \), recall, is the VEV of the Higgs singlet). In our numerical analyses we have taken \( m_b(m_b) = 4.214 \text{ GeV} \), \( m_{\tau}^{\text{pole}} = 1.777 \text{ GeV} \), \( m_{\mu}^{\text{pole}} = 0.1057 \text{ GeV} \) and \( m_t^{\text{pole}} = 171.4 \text{ GeV} \) respectively for the running bottom-quark mass and the (pole) tau-lepton, muon-lepton and top-quark masses, respectively.

We have used here the fortran package NMSSMTools, developed in Refs. [56, 57]. This code computes the masses, couplings and decay widths of all the Higgs bosons of the NMSSM, including radiative corrections, in terms of its parameters at the EW scale. NMSSMTools also takes into account theoretical as well as experimental constraints from negative Higgs searches at LEP [58] and the Tevatron, including the unconventional channels relevant for the NMSSM. Notice that the NMSSMTOOLS version used, version 2.3.1, does not include the latest LHC constraints [29]. However, as we shall see below, since we keep the SUSY mass scales very high and, over the phenomenologically interesting region to this analysis, our \( h_1 \) state is not very SM-like, the parameter points tested here are safely beyond current LHC limits.

We have used the code to scan over the six tree level parameters of the NMSSM Higgs sector in the following intervals:

\[
\lambda : 0.0001 - 0.7, \quad \kappa : 0 - 0.65, \quad \tan\beta : 1.6 - 54, \\
\mu_{\text{eff}} : 100 - 1000 \text{ GeV}, \quad A_\lambda : -1000 - +1000 \text{ GeV}, \quad A_\kappa : -10 - 0 \text{ GeV}.
\]

(Notice that our aim is exploring the parameter space which has very low \( m_{a_1} \) and one way to do that is by choosing \( A_\kappa \) small, in which case its negative values are preferred [19]. Also, notice that small \( A_\kappa \) is preferred to have small fine-tuning [14].)

Remaining soft terms, contributing at higher order level, which are fixed in the scan include:

\footnote{We have used NMSSMTools_2.3.1.}
• $m_{\tilde{Q}} = m_{\tilde{u}_R} = m_{\tilde{b}_R} = m_{\tilde{L}} = m_{\tilde{e}_R} = 1$ TeV,
• $A_t = A_b = A_\tau = 1.2$ TeV,
• $m_{\tilde{q}} = m_{\tilde{u}_R} = m_{\tilde{d}_R} = m_{\tilde{t}} = m_{\tilde{e}_R} = 1$ TeV,
• $M_1 = M_2 = M_3 = 1.5$ TeV.

As intimated, we have fixed soft term parameters at the TeV scale to minimise their contributions to parameter space outputs but changing values of some of those parameters such as $A_t$ could decrease or increase the number of successful points emerging from the NMSSM-Tools scans but without a significant impact on the $m_{a_1}$ distribution. Also, notice that the sfermion mass parameters and the $SU(2)$ gaugino mass parameter, $M_2$, play crucial roles in constraining $\tan \beta$. Decreasing values of those parameters allow smaller values of $\tan \beta$ to pass experimental and theoretical constraints, however, this is a less interesting region of the NMSSM parameter space for our analysis, as our Higgs production mode is only relevant at large values of $\tan \beta$. The effect of heavy gaugino mass parameters on the outputs, in particular $m_{a_1}$, would be small except for $M_2$ through its effect on $\tan \beta$. In fact, when $\tan \beta$ is large, the sfermion masses should be large to avoid the constraints coming from the muon anomalous magnetic moment [59]. The dominant Supersymmetric contribution at large $\tan \beta$ is due to a chargino-sneutrino loop diagram [60]. Also, notice that the chargino masses depend strongly on $M_2$.

Guided by the assumptions made in the reference [36], the possible decay channels for neutral NMSSM CP-even Higgs boson $h$, where $h = h_{1,2,3}$, and neutral CP-odd Higgs boson $a$, where $a = a_{1,2}$, are:

\[
\begin{align*}
    h, a &\to gg, \quad h, a \to \mu^+ \mu^-, \quad h, a \to \tau^+ \tau^-, \quad h, a \to b\bar{b}, \quad h, a \to t\bar{t}, \\
    h, a &\to s\bar{s}, \quad h, a \to c\bar{c}, \quad h \to W^+ W^-, \quad h \to Z Z, \\
    h, a &\to \gamma \gamma, \quad h, a \to Z \gamma, \quad h, a \to \text{Higgses}, \quad h, a \to \text{sparticles}.
\end{align*}
\]

(Notice that the CP-odd Higgses are not allowed to decay into vector boson pairs due to CP-conservation.) Also, notice that here ‘Higgses’ denotes any possible final state involving two neutral or two charged Higgs bosons or one Higgs boson and one gauge boson.

We have performed a random scan over millions of points in the specified parameter space. The output of the scan, as mentioned above, contains masses, Br’s and couplings of the NMSSM Higgses for all the successful points which have passed the various experimental and theoretical constraints.

6 Inclusive event rates

For the successful data points, we used CalcHEP [61] to calculate the cross sections for NMSSM Higgs production. Some new modules have been implemented for this purpose.

We focus here on the process

\[ gg \to b\bar{b} a_1 \]  

\[ \text{(29)} \]

\footnote{\text{We adopt herein CTEQ6L [62] as parton distribution functions, with scale } Q = \sqrt{s}, \text{ the centre-of-mass energy at parton level, for all processes computed.}}
i.e., Higgs production in association with a $b$-quark pair. (The production mode $q\bar{q} \to b\bar{b} a_1$ is negligible at the LHC with $\sqrt{s} = 14$ TeV.) We chose the production mode $gg \to b\bar{b} a_1$ because it is the dominant one at large $\tan\beta$. The gluon fusion channel is instead burdened by huge SM backgrounds and $a_1$ does not couple to gauge bosons in Higgs-strahlung and Vector Boson Fusion (VBF) processes due to CP-conservation, see [63]. In addition, Higgs production in other modes has been studied before, see for example [38]. In fact, Higgs production in association with a $b\bar{b}$ pair has an extra advantage, whereby the associated $b\bar{b}$ pair can be tagged, allowing a useful handle for background rejection.

In the NMSSM, the $a_1$ state is a composition of the usual doublet component of the CP-odd MSSM Higgs boson, $a_{\text{MSSM}}$, and the new singlet component, $a_S$, coming from the singlet Superfield of the NMSSM. This can be written as [14]:

$$a_1 = a_{\text{MSSM}} \cos \theta_A + a_S \sin \theta_A. \quad (30)$$

For very small values of $A_k$, the lightest CP-odd Higgs, $a_1$, is mostly singlet-like with a tiny doublet component, i.e., the mixing angle $\cos \theta_A$ is small, see the top-pane of figure 2 which shows the relation between $m_{a_1}$ and $\cos \theta_A$. The bottom-pane of the figure shows that the $\text{Br}(a_1 \to \gamma\gamma)$ can be dominant in some regions of the NMSSM parameter space with the possibility of reaching unity when $\cos \theta_A \approx 0$.

To a good approximation, $m_{a_1}$ can be written in the NMSSM as [14]:

$$m_{a_1}^2 = -3\frac{\kappa A_k \mu_{\text{eff}}}{\lambda} \sin^2 \theta_A + \frac{9A_\lambda \mu_{\text{eff}}}{2\sin 2\beta} \cos^2 \theta_A. \quad (31)$$

The first term of this expression is dominant at large $\tan\beta$. Furthermore, it is clear that a combination of all the tree level Higgs sector parameters affects $m_{a_1}$ in general.

7 Photon and tauon signals of very light CP-odd Higgs states of the NMSSM at the LHC

In our attempt to test the two aforementioned theorems, we consider in this section the case of the $\gamma\gamma$ and $\tau^+\tau^-$ decay channels of a very light CP-odd Higgs boson. The first mode is the most important one to detect a CP-even Higgs boson below 130 GeV in the SM and MSSM despite the smallness of its branching ratio, of $O(0.001)$. In addition, this decay mode gives a clean signature and can be resolved efficiently at the LHC. The second one is used in the MSSM as a search channel of rather heavy CP-even and CP-odd states, in particular at large $\tan\beta$, and its exploitation has not been proved at very low masses, say, below $M_Z$.

In the NMSSM, because of the introduction of a complex singlet Superfield, the lightest CP-odd Higgs boson, $a_1$, can be a singlet-like state with a tiny doublet component in large regions of parameter space. In this section (and also in the next one) we are looking for direct production of the $a_1$ rather than looking for its traditional production through $h_{1,2}$ decay. We examine the discovery potential of the $a_1$ produced in association with a bottom-antibottom pair at the LHC through the $\gamma\gamma$ and $\tau^+\tau^-$ decay modes.
We will show that in the NMSSM there exist regions of its parameter space where one can potentially have a dominant di-photon branching ratio of $\mathcal{O}(1)$ for the lightest CP-odd Higgs boson with small mass. This possibility emerges in the NMSSM because of the fact that such a CP-odd Higgs state has a predominant singlet component and a very weak doublet one. As a consequence, all partial decay widths are heavily suppressed as they employ only the doublet component, except one: the $\gamma\gamma$ partial decay width. This comes from the fact that the $a_1\tilde{\chi}^+\tilde{\chi}^-$ coupling is not suppressed, as it is generated through the $\lambda H_1 H_2 S$ Lagrangian term and therefore implies no small mixing. Although the direct decay $a_1 \to \tilde{\chi}^+\tilde{\chi}^-$ is kinematically not allowed, the aforementioned coupling participates in the $a_1\gamma\gamma$ effective coupling [64].

Furthermore, we will show that the $\tau^+\tau^-$ decay mode can be a promising decay mode for detecting the $a_1$ state of the NMSSM with very low mass. The detection of such a very low mass Higgs state would then unmistakably signal the existence of a non-minimal SUSY Higgs sector.

Figure 3 shows the distribution of the event rates $\sigma(gg \to b\bar{b}a_1) \cdot Br(a_1 \to \gamma\gamma)$ and $\sigma(gg \to b\bar{b}a_1) \cdot Br(a_1 \to \tau^+\tau^-)$ as functions of $m_{a_1}$ and of Br’s of the corresponding channel. As expected, the inclusive cross section decreases with increasing $m_{a_1}$, see the top panes of the figure. It is worth mentioning that the $Br(a_1 \to \gamma\gamma)$ can be dominant over a sizable expanse of the NMSSM parameter space, which originates from tiny widths into all other channels due to the dominant singlet nature of $a_1$ as mentioned in Sec. 5.1. However, the dominance of $Br(a_1 \to \gamma\gamma)$ does not correspond to the region that maximises the yield of $\sigma(gg \to b\bar{b}a_1) \cdot Br(a_1 \to \gamma\gamma)$, as the maximum of the latter occurs for Br’s in the region of some $10^{-5}$ to $10^{-4}$, see the bottom-left pane of the figure. Therefore, one can not take full advantage of the phenomenon described in the introduction of this section with respect to the singlet nature of the $a_1$ state, at the LHC, which couples to $\gamma\gamma$ through charginos. Thus, if $a_1$ were highly singlet, it would be difficult for the LHC to discover this particle as the doublet component (necessary to enable a large $a_1 b\bar{b}$ coupling at production level) would be suppressed. The tension between the two components is such that the cross section times Br rates are less than 100 fb.

The outlook for the $\tau^+\tau^-$ decay mode is much brighter where the corresponding signal rates are at nb level for $Br(a_1 \to \tau^+\tau^-) \approx 0.1$ or even 10 nb for $Br(a_1 \to \tau^+\tau^-) \approx 1$, see the bottom-right pane of figure 3. Also, notice that such large rates naturally hold for different values of $m_{a_1}$, in the allowed interval, but they decrease with increasing $m_{a_1}$ (see the top-right pane of this figure).

In the NMSSM, there is a large area of parameter space where one Higgs state can decay into two, e.g., $h_1 \to a_1 a_1$; see figure 4. As it is clear from the top-pane of this figure, the majority of points generated here have $m_{h_1} > 110$ GeV and $m_{a_1} < 55$ GeV, thereby allowing the possibility of $h_1 \to a_1 a_1$ decays. Moreover, this decay can be dominant and can reach unity as shown in the bottom-pane of the figure. Despite this, such a decay may not give Higgs signals with sufficient statistical significance at the LHC (as discussed in previous literature). Therefore, we are well motivated to study, in the forthcoming sections, the scope of direct production of the $a_1$ state in single mode at the LHC, through $gg \to b\bar{b}a_1$, over
overlapping regions of NMSSM parameter space.

8 Muon and $b$-quark signals of very light CP-odd Higgs states of the NMSSM at the LHC

The di-muon decay mode has an advantage, that it has a clean signature with excellent mass resolution. However, the $\mu^+\mu^-$ branching ratio is small in most regions of parameter space but this decay mode is enhanced for large $\tan\beta$.

Figure 5 shows the correlations between the $a_1$ mass and the di-muon decay rate. One can see from this figure that the $\text{Br}(a_1 \to \mu^+\mu^-)$ can be of $\mathcal{O}(10\%)$, $\mathcal{O}(1\%)$ and $\mathcal{O}(0.1\%)$ or less for the mass intervals $2m_\mu < m_{a_1} < 2m_\tau$, $2m_\tau < m_{a_1} < 2m_t$ and $2m_b < m_{a_1}$, respectively. The first region of parameter space ($m_{a_1} < 2m_\tau$) is rather small, the second one ($2m_\tau < m_{a_1} < 2m_b$) more significant and the third one ($2m_b < m_{a_1}$) is by far the widest one.

Figure 6 illustrates the distribution of the inclusive event rates as a function of the $\text{Br}(a_1 \to \mu^+\mu^-)$ and of $m_{a_1}$. It is remarkable to notice that the inclusive event rates are sizable in all such mass regions. These event rates reach the $10^4$ $fb$ level in the two lower mass intervals and the $10^3$ $fb$ level in the higher mass range and clearly decrease by increasing $m_{a_1}$, as expected. Finally, notice that the mass region below the $\mu^+\mu^-$ threshold is very severely constrained [65].

As for $4b$-quark final states, at large $\tan\beta$ values, the cross-section of the $a_1$ produced in association with a bottom-antibottom pair followed by the decay $a_1 \to bb$ is strongly enhanced, in general. However, since the channel is a 4-quark final state, it is plagued by very large irreducible and reducible backgrounds. In this section, we examine whether or not the production mode $gg \to b\bar{b}a_1 \to b\bar{b}bb$ can be exploited to detect the $a_1$ at the LHC. In fact, the existence of $b$-jets in the final states offers the advantage of $b$-tagging, which can be exploited to trigger on the signal and enable us to require up to four displaced vertices in order to reject light jets. The ensuing $4b$ signature has already been exploited to detect neutral Higgs bosons of the MSSM at the LHC and proved useful, provided that $\tan\beta$ is large and the collider has good efficiency and purity in tagging $b$-quark jets, albeit for the case of rather heavy Higgs states (with masses beyond $M_Z$, typically) [65, 67].

Figure 7 illustrates the inclusive signal production cross section $\sigma(gg \to bba_1)$ multiplied by the branching fraction $\text{Br}(a_1 \to bb)$ as a function of the $\text{Br}(a_1 \to bb)$ and of $m_{a_1}$ and the plots in figure 8 display instead the correlations between the $a_1 \to bb$ decay rate and the $a_1$ mass (top-pane) and between the $a_1 \to bb$ decay rate and the $a_1 \to \gamma\gamma$ decay rate. From a close look at the bottom-left pane of figure 7 it is clear that the $\text{Br}(a_1 \to bb)$ is dominant for most points in the parameter space, about 90% and above. In addition, by looking at the the bottom-right pane of the figure, it is remarkable to notice that also these event rates are sizable in most regions of parameter space, topping the $10^7$ $fb$ level for small values of $m_{a_1}$ and decreasing rapidly with increasing $m_{a_1}$. One can also notice that there

---

7A partonic signal-to-background $(S/B)$ analysis for $\gamma\gamma$ and $\tau^+\tau^-$ final states has been done in [25], where extraction of the latter signature was proven for several benchmark scenarios.
are some points in the parameter space with \( m_{a_1} \) between 40 and 120 GeV, as shown in the top-pane of figure 8, in which the \( \text{Br}(a_1 \rightarrow b \bar{b}) \) is suppressed due to the enhancement of the \( \text{Br}(a_1 \rightarrow \gamma \gamma) \) (see the bottom-pane of the same figure), a phenomenon peculiar to the NMSSM that depends upon the amount of Higgs singlet-doublet mixing\(^8\), see [25]. (We discussed this in the previous section\(^9\).)

9 The ‘No-lose theorem’ for NMSSM Higgs discovery at the LHC in difficult scenarios

In this section, we continue to investigate whether or not the ‘No-lose theorem’ of the NMSSM at the LHC can be proven considering Higgs boson production in association with a \( b \)-quark pair. We do so based on the interesting results obtained in the previous two sections. In this section we will, however, no longer consider direct \( a_1 \) production, i.e., \( gg \rightarrow b \bar{b}a_1 \). Rather, we will initially produce either a \( h_1 \) or \( h_2 \), eventually decaying to one or more \( a_1 \)'s. In this case one may wonder though whether also Higgs boson production in association with a \( t \)-quark pair could play a role, owing to different couplings of the \( h_{1,2} \) state to fermions, with respect to the \( a_1 \) field. Production rates for \( gg \rightarrow t\bar{t}h_{1,2} \) were studied in [38], where they were found to be very subleading over the entire NMSSM parameter space.

We will be looking at inclusive event rates in presence of various Higgs-to-Higgs decays, namely, \( h_{1,2} \rightarrow a_1 a_1 \), \( h_1 \rightarrow h_1 h_1 \) and \( h_{1,2} \rightarrow Z a_1 \). We will also be studying the decay patterns of the lightest Higgs boson pairs, \( a_1 a_1 \) or \( h_1 h_1 \), and of the gauge boson and a light CP-odd Higgs boson, \( Z a_1 \), into different final states. Further details on material contained in this section can be found in [40] and [41]. We have again used NMSSMTools to perform a random scan over the usual parameter space, mentioned in section 5, and further required that \( m_{h_2} \leq 300 \text{ GeV} \), as corresponding production rates become negligible for heavier masses. We used CalcHEP [61] to determine the cross sections for NMSSM \( h_{1,2} \) production for the following two processes:

\[
gg \rightarrow b \bar{b} \, h_1 \quad \text{and} \quad gg \rightarrow b \bar{b} \, h_2,
\]

which were computed separately (i.e., without the interferences emerging whenever \( h_1 \) and \( h_2 \) have the same decay products).

The lightest two CP-even neutral Higgs boson masses are given by Eq.(25). Recall though that the equation is at tree level, mainly for guidance in interpreting the upcoming figures, while NMSSMTools includes radiative corrections as well.

Figure 9 shows the correlations between all three Higgs masses \( m_{a_1}, m_{h_1}, \) and \( m_{h_2} \). Since the successful points emerging from the scan have small values of \( \lambda, \kappa \) and also \( A_\kappa \), only rather small values of \( m_{a_1} \) are allowed. It is remarkable that the smaller \( m_{a_1} \), the smaller

\(^8\)Notice that constraints coming from Tevatron [68] do not affect our results since the singlet field plays a primary role in the NMSSM. But under severe conditions such as \( \lambda \rightarrow 0 \) and \( \kappa \rightarrow 0 \), the NMSSM and MSSM become similar and those constraints may be applied.

\(^9\)A partonic signal-to-background \((S/B)\) analysis for \( \mu^+\mu^- \) and \( b \bar{b} \) final states has been done in [28] [39], extracting both signals for several benchmark scenarios.
$m_{h_1}$ and $m_{h_2}$ too (two top-panes). In the bottom-pane of the same figure, for $m_{h_2}$ around 120 GeV, $m_{h_1}$ can have values from just above 0 up to slightly less than 120 GeV, showing the possibility that the two Higgs states can have the same mass, i.e., $m_{h_1} \sim m_{h_2}$. Notice also that the majority of points have $m_{h_1}$ between 115 and 120 GeV.

### 9.1 Production of $h_1$ and $h_2$ decaying into two lighter Higgs bosons

The production times decay rates of $h_1$ and $h_2$, in which $h_1$ decays into two lighter $a_1$’s and $h_2$ decays into either a pair of $a_1$’s or a pair of $h_1$’s, are shown in figure 10. This figure displays all the correlations between the three discussed production and decay processes. It is quite remarkable that the overall trend, despite an obvious spread also in the horizontal and vertical directions, is such that when one channel grows in event yield there is also another one which also does, hence opening up the possibility of the simultaneous discovery of several Higgs states of the NMSSM (three neutral Higgses at the same time: $h_1$, $h_2$ and $a_1$), an exciting prospect in order to distinguish the NMSSM Higgs sector from the MSSM one (in fact, a clear manifestation of a possible More-to-gain theorem being established).

### 9.2 Production of $h_1$ and $h_2$ decaying into a gauge boson and a light CP-odd Higgs

In this subsection, we examine the LHC discovery potential of the lightest two CP-even Higgs states $h_{1,2}$, followed by the decay $h_{1,2} \rightarrow Z a_1$. Figure 11 shows that the production rate for the $h_1$, produced in association with a $b\bar{b}$ pair, is small, topping the 0.001 fb level. Such a production rate is presumably not enough to discover the $h_1$ at the LHC. The top-pane of the figure shows that there is a linear relation between the $h_1$ production rate and the $\text{Br}(h_1 \rightarrow Z a_1)$ because the production rate $\sigma(gg \rightarrow b\bar{b} h_1)$ is nearly constant in our parameter space, which has large $\tan\beta$. The bottom-pane of the figure shows that the points passing the constraints have $m_{h_1} > 100$ GeV.

Figure 12 illustrates the inclusive $h_1$ production rates ending up with $Z a_1 \rightarrow \mu^+ \mu^- b\bar{b}$, $Z a_1 \rightarrow \mu^+ \mu^- \tau^+ \tau^-$ and $Z a_1 \rightarrow j j \tau^+ \tau^-$ (where $j =$ jet). It is clear that the production and decay rates are definitely too small, topping $10^{-4}$ fb for the first and last channels and $10^{-5}$ fb for the second one. Such rates are obviously not enough to discover the $h_1$ neither at the LHC nor at the SLHC with 1000 fb$^{-1}$ of luminosity.

In contrast, the situation for $h_2$ is promising as one can notice that $\sigma(gg \rightarrow b\bar{b} h_2)\text{Br}(h_2 \rightarrow Z a_1)$ is sizable, topping the 10000 fb level (figure 13). The highest values of the cross section are accompanied by an intriguingly large $\text{Br}(h_2 \rightarrow Z a_1)$, reaching up to 10%. It is clear from the top-pane of the figure that the distribution over the branching ratio for $h_2$ is not as uniform as that for the $h_1$ because the production rate $\sigma(gg \rightarrow b\bar{b} h_2)$ depends strongly on the tree level parameters unlike that for $h_1$. The bottom-pane of the figure shows that the highest cross section occurs for $m_{h_2} > 220$ GeV.

\footnote{A partonic signal-to-background ($S/B$) analysis for several $a_1 a_1$ and $h_1 h_1$ decays is currently being done in \cite{69}.}
In order to study the detectability of \( h_2 \) decaying into a gauge boson and a light CP-odd Higgs state at the LHC, we have calculated the inclusive production rates ending up with \( \mu^+\mu^-b\bar{b} \), \( \mu^+\mu^-\tau^+\tau^- \) and \( jj\tau^+\tau^- \) (figure [13]). The event rates for these processes are at the \( \mathcal{O}(100) \) fb level at the most. While clearly this number is not very large, signal events may still be detectable at planned LHC luminosities, especially if the background can be successfully reduced to manageable levels\(^{11}\). In short, there is a small but well defined region of the NMSSM parameter space where the \( h_2 \) and \( a_1 \) states, both with a mixed singlet and doublet nature, could potentially be detected at the LHC if \( 220 \text{ GeV} \lesssim m_{h_2} \lesssim 300 \text{ GeV} \) and \( 15 \text{ GeV} \lesssim m_{a_1} \lesssim 60 \text{ GeV} \), in the \( h_2 \to Z a_1 \to \mu^+\mu^-b\bar{b} \), \( h_2 \to Z a_1 \to \mu^+\mu^-\tau^+\tau^- \) and \( h_2 \to Z a_1 \to jj\tau^+\tau^- \) modes, when the CP-even Higgs state is produced in association with a \( b\bar{b} \) pair for rather large \( \tan\beta \).

### 10 Conclusions

The NMSSM has a singlet Superfield in addition to the usual Higgs doublets of the MSSM. This singlet gives rise to a more varied phenomenology in the case of the NMSSM, compared to that of the MSSM. For instance, this singlet Superfield mixes with the neutral components of the doublets, giving rise to one CP-even Higgs, one CP-odd Higgs and one extra neutralino in addition to the usual spectrum of the MSSM. Therefore, in the NMSSM, by assuming CP-conservation, there are seven Higgses: three CP-even, two CP-odd and a pair of charged Higgses. We have investigated whether or not at least one Higgs boson of the NMSSM can be discovered at the LHC (‘No-lose theorem’) and/or is possible to find some regions in the parameter space where more and/or different Higgs states of the NMSSM are detectable at the LHC, compared to those available within the MSSM (‘More-to-gain theorem’).

Because of the mixing between the Higgs singlet and doublets, Higgs-to-Higgs decays are kinematically possible for large regions of the NMSSM parameter space even for small masses of the Higgs states, which is impossible in the MSSM. For instance, a SM-like Higgs can decay into a pair of the lightest NMSSM CP-odd Higgses. This decay can be dominant in sizable areas of the NMSSM parameter space. Such a decay has a significant meaning if one notices that it can explain a \( 2.3\sigma \) excess occurred at LEP for the process \( e^+e^- \to Zb\bar{b} \) for \( M_{b\bar{b}} \sim 98 \text{ GeV} \) and the \( 2.6\sigma \) excess recently emerged at the LHC (primarily in the \( \gamma\gamma \) decay mode). Moreover, a SM-like Higgs with mass of order 100 GeV, which has no-fine tuning, can naturally occur in the NMSSM and this scenario is preferred by precision EW data. In addition, the NMSSM can solve both the \( \mu \)-problem and the little hierarchy problem of the MSSM.

In the context of the NMSSM, we have proven that a very light CP-odd Higgs state with mass \( m_{a_1} \lesssim M_Z \), which has a large singlet component and a small doublet one, can be discovered at the LHC via Higgs production in association with a bottom-antibottom pair. This mode is dominant at large \( \tan\beta \). After performing several analyses for signals and dominant backgrounds, not documented here yet referred to, we have proven that this

\(^{11}\)A partonic signal-to-background (\( S/B \)) analysis for \( jj\tau^+\tau^- \) final state has been done in [11], showing very promising results.
production mode is the ideal one to discover the $a_1$ through the following signatures: (i) $\tau^+\tau^-$ decay mode, in which $a_1$ can be discovered with mass up to $M_Z$; (ii) $\mu^+\mu^-$ decay mode, if $10 \lesssim m_{a_1} \lesssim 60$ GeV. Further, despite the fact that the $b\bar{b}$ decay mode is dominant in most regions of parameter space that have light $a_1$, this channel has huge QCD background and a smaller signal-to-background ratio. Finally, we also looked at the detectability of $a_1$ through the $\gamma\gamma$ decay mode but this proved unuseful despite the fact that this decay mode can be dominant in some areas of the NMSSM parameter space.

We believe that the results presented in sections 7 and 8 have a twofold relevance. Firstly, they support the ‘No-lose theorem’ by looking for direct $a_1$ production rather than looking for its production through the decays $h_{1,2} \rightarrow a_1 a_1$, which may not give a sufficient signal significance. Secondly, they corroborate the ‘More-to-gain theorem’ as such very light $a_1$’s (with $m_{a_1} \lesssim M_Z$) are not at all possible in the MSSM. Altogether, the existence of such a light neutral Higgs state is a direct evidence for the non-minimal nature of the SUSY Higgs sector.

Finally, we have mentioned in section 9 the importance of Higgs-to-Higgs decays in the NMSSM, here occurring after Higgs boson production in association with $b\bar{b}$ pairs (unlike in most previous literature), and have shown that such decays should be taken seriously before proving, or otherwise, the ‘No-lose theorem’. In fact, we also have shown that such decays are dominant in sizable regions of the NMSSM parameter space. We have studied the LHC discovery potential of a CP-even Higgs boson $h_1$ or $h_2$, decaying into a pair of light CP-odd Higgses $a_1$’s, and also $h_2$ decaying into a pair of $h_1$’s. We have found that these channels can give sizable signal rates, which could allow one to detect simultaneously two Higgs bosons: $h_1$ and $a_1$, $h_2$ and $a_1$ or $h_2$ and $h_1$. In addition, we have shown that the LHC has the potential to discover the three neutral Higgs bosons at the same time. Furthermore, we have studied the LHC discovery potential for $h_1$ and $h_2$ decaying into $Z a_1$ and have shown that, while the discovery of the $h_1$ through this channel is impossible, there is a small but well defined region of the NMSSM parameter space where the $h_2$ state could potentially be discovered.

Acknowledgments

M.M.A. gratefully acknowledges financial support from Taibah University in Saudi Arabia. S.M. is financially partially supported through the NExT Institute.

References

[1] A. Djouadi, Phys. Rept. 459 (2008) 1.
[2] J. E. Kim and H. P. Nilles, Phys. Lett. B 138 (1984) 150.
[3] The ATLAS Collaboration, ATLAS-CONF-2012-019.
[4] The CMS Collaboration, arXiv:1202.1488v1 [hep-ex].
[5] M. Maniatis, Int. J. Mod. Phys. A 25 (2010) 3505.
[6] U. Ellwanger, C. Hugonie and A. M. Teixeira, Phys. Rept. 496 (2010) 1.

[7] J. R. Ellis, J. F. Gunion, H. E. Haber, L. Roszkowski and F. Zwirner, Phys. Rev. D 39 (1989) 844.

[8] M. Bastero-Gil, C. Hugonie, S. F. King, D. P. Roy and S. Vempati, Phys. Lett. B 489 (2000) 359.

[9] R. Dermisek and J. F. Gunion, Phys. Rev. Lett. 95 (2005) 041801.

[10] U. Ellwanger and C. Hugonie, arXiv:1203.5048 [hep-ph].

[11] U. Ellwanger, JHEP 1203 (2012) 044

[12] R. Dermisek and J. F. Gunion, Phys. Rev. D 73 (2006) 111701.

[13] R. Dermisek and J. F. Gunion, Phys. Rev. D 75 (2007) 075019.

[14] R. Dermisek and J. F. Gunion, Phys. Rev. D 76 (2007) 095006.

[15] R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. 38 (1977) 1440.

[16] R. D. Peccei and H. R. Quinn, Phys. Rev. D 16 (1977) 1791.

[17] K. Hagiwara et al. [Particle Data Group], Phys. Rev. D 66 (2002) 010001.

[18] Y. B. Zeldovich, I. Y. Kobzarev and L. B. Okun, Zh. Eksp. Teor. Fiz. 67 (1974) 3.

[19] D. J. Miller, R. Nevzorov and P. M. Zerwas, Nucl. Phys. B 681 (2004) 3.

[20] M. Masip, R. Muñoz-Tapia and A. Pomarol, Phys. Rev. D 57 (1998) 5340.

[21] U. Ellwanger and C. Hugonie, Eur. Phys. J. C 25 (2002) 297.

[22] U. Ellwanger, Phys. Lett. B 303 (1993) 271.

[23] T. Elliott, S. F. King and P. L. White, Phys. Rev. D 49 (1994) 2435.

[24] P. N. Pandita, Phys. Lett. B 318 (1993) 338.

[25] M. M. Almarashi and S. Moretti, Eur. Phys. J. C 71 (2011) 1618.

[26] The ALEPH Collaboration, JHEP 1005 (2010) 049.

[27] The BABAR Collaboration, Phys. Rev. Lett. 103 (2009) 181801.

[28] M. M. Almarashi and S. Moretti, Phys. Rev. D 84 (2011) 015014.

[29] R. Dermisek and J. F. Gunion, Phys. Rev. D 77 (2008) 015013.

[30] J. F. Gunion, Y. Jiang and S. Kraml, Phys. Lett. B 710 (2012) 454.
[31] S. F. King, M. Muhlleitner and R. Nevzorov, Nucl. Phys. B 860 (2012) 207.

[32] D. A. Vasquez, G. Belanger, C. Boehm, J. Da Silva, P. Richardson and C. Wymant, arXiv:1203.3446 [hep-ph].

[33] The ATLAS Collaboration, arXiv:0901.0512 [hep-ex].

[34] The CMS Collaboration, J. Phys. G 34 (2007) 995.

[35] J. Dai, J. F. Gunion and R. Vega, Phys. Lett. B 315 (1993) 355 and Phys. Lett. B 345 (1995) 29; J. R. Espinosa and J. F. Gunion, Phys. Rev. Lett. 82 (1999) 1084.

[36] U. Ellwanger, J. F. Gunion and C. Hugonie, JHEP 0507 (2005) 041.

[37] U. Ellwanger, J. F. Gunion and C. Hugonie, hep-ph/0111179; C. Hugonie and S. Moretti, hep-ph/0110241; D. J. Miller and S. Moretti, hep-ph/0403137; U. Ellwanger, J. F. Gunion, C. Hugonie and S. Moretti, hep-ph/0305109 and hep-ph/0401228; A. Belyaev, S. Hesselbach, S. Lehti, S. Moretti, A. Nikitenko and C. H. Shepherd-Themistocleous, arXiv:0805.3505 [hep-ph]; J. R. Forshaw, J. F. Gunion, L. Hodgkinson, A. Papaefstathiou and A. D. Pilkington, JHEP 0804 (2008) 090; A. Belyaev, J. Pivarski, A. Safonov, S. Senkin and A. Tatarinov, Phys. Rev. D 81 (2010) 075021.

[38] S. Moretti, S. Munir and P. Poulose, Phys. Lett. B 644 (2007) 241.

[39] M. M. Almarashi and S. Moretti, Phys. Rev. D 83 (2011) 035023.

[40] M. M. Almarashi and S. Moretti, Phys. Rev. D 84 (2011) 035009.

[41] M. M. Almarashi and S. Moretti, Phys. Rev. D 85 (2012) 017701.

[42] A. Djouadi et al., JHEP 0807 (2008) 002.

[43] A. Arhrib, K. Cheung, T. J. Hou and K. W. Song, JHEP 0703 (2007) 073.

[44] K. Cheung and T. J. Hou, Phys. Lett. B 674 (2009) 54.

[45] J. F. Gunion, H. E. Haber and T. Moroi, In the Proceedings of 1996 DPF / DPB Summer Study on New Directions for High-Energy Physics (Snowmass 96), Snowmass, Colorado, 25 Jun - 12 Jul 1996, pp LTH095 arXiv:hep-ph/9610337.

[46] B. A. Dobrescu, G. L. Landsberg and K. T. Matchev, Phys. Rev. D 63 (2001) 075003.

[47] B. A. Dobrescu and K. T. Matchev, JHEP 0009 (2000) 031.

[48] U. Ellwanger, J. F. Gunion, C. Hugonie and S. Moretti, in Ref. [37].

[49] A. Belyaev, S. Hesselbach, S. Lehti, S. Moretti, A. Nikitenko and C. H. Shepherd-Themistocleous, in Ref. [37].

[50] A. Belyaev, J. Pivarski, A. Safonov, S. Senkin and A. Tatarinov, in Ref. [37].

18
[51] M. Lisanti and J. G. Wacker, Phys. Rev. D 79 (2009) 115006.

[52] U. Ellwanger, Eur. Phys. J. C 71 (2011) 1782.

[53] S. Moretti and S. Munir, Eur. Phys. J. C 47 (2006) 791.

[54] S. Munir, talk given at the ‘International School of Subnuclear Physics, 43rd Course’, Erice, Italy, Aug. 29 – Sept. 7, 2005, to be published in the proceedings, preprint SHEP-05-37, October 2005.

[55] F. Mahmoudi, J. Rathsman, O. Stal and L. Zeune, Eur. Phys. J. C 71 (2011) 1608.

[56] U. Ellwanger, J. F. Gunion and C. Hugonie, JHEP 0502 (2005) 066; U. Ellwanger and C. Hugonie, Comput. Phys. Commun. 175 (2006) 290.

[57] See the Web site “NMSSMTools: Tools for the Calculation of the Higgs and Sparticle Spectrum in the NMSSM: NMHDECAY, NMSPEC and NMGMSB”, http://www.th.u-psud.fr/NMHDECAY/nmssmtools.html.

[58] S. Schael et al., Eur. Phys. J. C 47 (2006) 547.

[59] F. Domingo and U. Ellwanger, JHEP 0807 (2008) 079.

[60] A. Czarnecki and W. J. Marciano, Phys. Rev. D 64 (2001) 013014.

[61] A. Pukhov, arXiv:hep-ph/0412191.

[62] See the Web site “CTEQ6 Parton Distribution Functions”, http://hep.pa.msu.edu/cteq/public/cteq6.html.

[63] M. M. Almarashi, talk given at ‘NExT meeting at RAL’, Didcot, United Kingdom, January 26, 2011, http://conference.ippp.dur.ac.uk/conferenceDisplay.py?confId=304.

[64] U. Ellwanger and C. Hugonie, private communication.

[65] S. Andreas, O. Lebedev, S. R. Sanchez and A. Ringwald, JHEP 1008 (2010) 003.

[66] J. Dai, J. F. Gunion and R. Vega, Phys. Lett. B 345 (1995) 29 in Ref. 35.

[67] J. Dai, J. F. Gunion and R. Vega, Phys. Lett. B 387 (1996) 801.

[68] The D0 Collaboration, Phys. Lett. B 707 (2012) 323.

[69] M. M. Almarashi and S. Moretti, in progress.
Figure 2: The lightest CP-odd Higgs mass $m_{a_1}$ and the $\text{Br}(a_1 \rightarrow \gamma\gamma)$ plotted against the mixing angle in the CP-odd Higgs sector $\cos \theta_A$. 

20
Figure 3: The rates for $\sigma(gg \to b\bar{b}a_1) \text{ Br}(a_1 \to \gamma\gamma)$ (left) and for $\sigma(gg \to b\bar{b}a_1) \text{ Br}(a_1 \to \tau^+\tau^-)$ (right) as functions of $m_{a_1}$ and of the Br of the corresponding channel.
Figure 4: The lightest CP-odd Higgs mass $m_{a_1}$ plotted against the lightest CP-even Higgs mass $m_{h_1}$ and against $\text{Br}(h_1 \rightarrow a_1a_1)$. 
Figure 5: The CP-odd Higgs mass $m_{a_1}$ as a function of the $\text{Br}(a_1 \rightarrow \mu^+\mu^-)$.

Figure 6: The rates for $\sigma(gg \rightarrow b\bar{b}a_1) \text{ Br}(a_1 \rightarrow \mu^+\mu^-)$ as a function of $\text{Br}(a_1 \rightarrow \mu^+\mu^-)$ and of $m_{a_1}$. 

23
Figure 7: The rates for $\sigma(gg \rightarrow b\bar{b}a_1) \ Br(a_1 \rightarrow b\bar{b})$ as a function of the $Br(a_1 \rightarrow b\bar{b})$ and of $m_{a_1}$.
Figure 8: The $\text{Br}(a_1 \to \bar{b}b)$ as a function of the CP-odd Higgs mass $m_{a_1}$ and of the $\text{Br}(a_1 \to \gamma\gamma)$.
Figure 9: The correlations between the lightest CP-odd Higgs mass $m_{a_1}$ and the lightest two CP-even Higgs masses $m_{h_1}$ and $m_{h_2}$ and between the latter two.
Figure 10: The rates for $\sigma(gg \rightarrow b\bar{b}h_2) \ Br(h_2 \rightarrow a_1a_1)$ versus $\sigma(gg \rightarrow b\bar{b}h_1) \ Br(h_1 \rightarrow a_1a_1)$, $\sigma(gg \rightarrow b\bar{b}h_1) \ Br(h_1 \rightarrow a_1a_1)$ versus $\sigma(gg \rightarrow b\bar{b}h_2) \ Br(h_2 \rightarrow h_1h_1)$ and for $\sigma(gg \rightarrow b\bar{b}h_2) \ Br(h_2 \rightarrow a_1a_1)$ versus $\sigma(gg \rightarrow b\bar{b}h_2) \ Br(h_2 \rightarrow h_1h_1)$. 
Figure 11: The signal rate for $\sigma(gg \to b\bar{b}h_1) \text{ Br}(h_1 \to Za_1)$ as a function of the $\text{Br}(h_1 \to Za_1)$ and of $m_{h_1}$. 
Figure 12: The signal rate for $\sigma(gg \rightarrow b\bar{b}h_1) \text{Br}(h_1 \rightarrow Z_{a1})$ times $\text{Br}(Z_{a1} \rightarrow \mu^+\mu^-b\bar{b})$, times $\text{Br}(Z_{a1} \rightarrow \mu^+\mu^-\tau^+\tau^-)$ and times $\text{Br}(Z_{a1} \rightarrow jj\tau^+\tau^-)$ as functions of $m_{h_1}$. 
Figure 13: The signal rate for $\sigma(gg \to b\bar{b}h_2) \, \text{Br}(h_2 \to Z a_1)$ as a function of the $\text{Br}(h_2 \to Z a_1)$ and of $m_{h_2}$. 
Figure 14: The signal rate for $\sigma(gg \rightarrow b\bar{b}h_2) \Br(h_2 \rightarrow Za_1)$ times $\Br(Za_1 \rightarrow \mu^+\mu^-b\bar{b})$, times $\Br(Za_1 \rightarrow \mu^+\mu^-\tau^+\tau^-)$ and times $\Br(Za_1 \rightarrow jj\tau^+\tau^-)$ as functions of $m_{h_2}$. 