Boosting Higgs decays into gamma and a Z in the NMSSM

Geneviève Bélinger,1 Vincent Bizouard,1 and Guillaume Chalons2

1LAPTh, Université de Savoie, CNRS, 9 Chemin de Bellevue, B.P. 110, F-74941 Annecy le-Vieux, France
2LPSC, Université Grenoble-Alpes, CNRS/IN2P3, Grenoble INP, 53 rue des Martyrs, F-38026 Grenoble, France

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In this work we present the computation of the Higgs decay into a photon and a Z0 boson at one-loop level in the framework of the Next-to-Minimal Supersymmetric Standard Model (NMSSM). The numerical evaluation of this decay width was performed within the framework of the SloopS code, originally developed for the Minimal Supersymmetric Standard Model (MSSM) but which was recently extended to deal with the NMSSM. Thanks to the high level of automation of SloopS all contributions from the various sectors of the NMSSM are consistently taken into account, in particular the non-diagonal chargino and sfermion contributions. We then explored the NMSSM parameter space, using HiggsBounds and HiggsSignals, to investigate to which extent these signal can be enhanced.

I. INTRODUCTION

The discovery of a 125 GeV Higgs boson at the LHC on July 2012 [1, 2], is a milestone in the road leading to the elucidation the ElectroWeak Symmetry Breaking (EWSB) riddle. Since then, its couplings to electroweak gauge bosons, third generation fermions and the loop-induced couplings to photon and gluons have been measured with an already impressive accuracy by the ATLAS and CMS collaborations [3, 4] during the 7 and 8 TeV runs. This great achievement was made possible since for a 125 GeV Higgs bosons many different production and decay channels are detectable at the LHC. A spin and parity analysis of the Higgs boson in the decays $H \to \gamma\gamma$, $H \to ZZ^*$ and $H \to WW^*$ favor the $J^P = 0^+$ hypothesis [5, 6].

The couplings of the Higgs to photons $H\gamma\gamma$ and gluons $Hgg$ are induced at the quantum level, even in the Standard Model (SM), and thus are interesting probes of New Physics (NP) since both the SM and NP contributions enter at the same level. On the one hand, the updated CMS analysis [7] gives a signal strength for the diphoton which is now compatible with the SM, compared to the previous one. On the other hand, the updated ATLAS analysis [8] still observes a slight excess of events in this channel but recent measurements of the $H \to \gamma\gamma$ differential cross sections do not show significant disagreements with expectations from a SM Higgs [9].

The search for another important loop-induced Higgs decay channel, $H \to \gamma Z^0$, is also performed by the ATLAS and CMS experiments [10, 11]. Within the SM, the partial width...
for this channel $\Gamma_{\gamma Z}$ is about two third of that for the diphoton decay and its measurement can also provide insights about the properties of the boson, such as its mass, spin and parity \cite{12} thanks to a clean final state topology. No excess above SM predictions has been found in the 120-160 GeV mass range and first limits on the Higgs boson production cross section times the $H \rightarrow \gamma Z^0$ branching fraction have been derived \cite{10, 11}. The collaborations set an upper limit on the ratio $\Gamma_{\gamma Z}/\Gamma_{SM}^{\gamma Z} < 10$. A measurement of $\Gamma_{\gamma Z^0}$ can also provide insights about the underlying dynamics of the Higgs sector since new heavy charged particles can alter the SM prediction, just as for the $H \rightarrow \gamma \gamma$ channel, without affecting the gluon-gluon fusion Higgs production cross section \cite{13}. Moreover, the measurement of $H \rightarrow \gamma Z^0$ and its rate compared to $H \rightarrow \gamma \gamma$ is crucial for broadening our understanding of the EWSB pattern \cite{14, 15}. Testing the SM nature of this Higgs state and inspecting possible deviations in its coupling to SM particles will represent a major undertaking of modern particle physics and a probable probe of models going beyond the Standard Model (BSM).

Despite the fact that no significant deviation from the SM has been observed, there are many theoretical arguments and observations from astroparticle physics and cosmology supporting the fact that it cannot be the final answer for a complete description of Nature. If New Physics must enter the game, what the experimental results told us so far is that any BSM should exhibit decoupling properties. Among BSM, supersymmetry (SUSY), is probably the best motivated and most studied framework. Its minimal incarnation, the MSSM, although possessing such a decoupling regime, can reproduce laboriously a 125 GeV light Higgs boson since this value is close to the upper bound attainable when radiative corrections are taken into account \cite{17–21}. The introduction of an additional gauge-singlet superfield $S$ to the MSSM content, whose simplest version is dubbed as the Next-to-Minimal Supersymmetric Standard Model (NMSSM) \cite{22, 23}, relaxes such an upper bound and a 125 GeV Higgs mass appears more naturally within its parameter space than in the MSSM. The singlet term provides an extra tree level contribution to the Higgs mass matrix such that the MSSM limit can be exceeded, already at tree level \cite{24–26}. The neutral CP-even Higgs sector is then enlarged with three states $h^0_i$, where $i$ ranges from 1 to 3 and ordered in increasing mass. In this context the lightest CP-even Higgs state might well be dominantly singlet with reduced couplings to the SM and thus could remain essentially invisible at colliders: the SM-like Higgs state would then be the second lightest and a small mixing effect with the singlet would in turn shifts its mass towards slightly higher values. The NMSSM possesses another virtue, in addition to the ones of the MSSM: the so-called $\mu$ problem of the MSSM \cite{27} can be circumvented as it is dynamically generated once the singlet field gets a vacuum expectation value (vev) \cite{22–23}. All in all, the NMSSM appears now as more appealing than the MSSM and has received considerable attention \cite{28–49}.

In the present work, we investigate the $\gamma Z^0$ decay channel of the SM-like Higgs boson $h$ of the NMSSM and its correlation with $H \rightarrow \gamma \gamma$. We compute these two decay widths with the help of the automatic code SloopS \cite{50, 51}, initially designed to tackle one-loop calculations in the MSSM \cite{52–55}. This code has been recently further developed to deal with the NMSSM extended field content and applied to Dark Matter \cite{56–58} and Higgs phenomenology \cite{59}. Thanks to this implementation, all the relevant particles running in the loops are properly taken into account, in particular the non-diagonal contributions due to the non-diagonal couplings of the $Z$ boson to charginos and sfermions. Our results are consistent with those presented in \cite{60}. As compared to \cite{60}, we perform a more thorough exploration on the parameter space of the NMSSM, in particular considering also the small $\lambda$ region, and we compute the expected signal strengths for both the vector boson fusion...
production mode (VBF) and the gluon fusion mode ($gg$). In addition we impose the most recent collider constraints on the Higgs sector using HiggsBounds[61] and HiggsSignals[62] to require that one of the Higgses fits the properties of the particle observed at LHC thus illustrating that the most severe deviations from the SM expectations for $H \rightarrow \gamma \gamma (Z^0)$ are already constrained. We further explicitly distinguish the case where the 125 GeV Higgs is the lightest or second lightest CP-even Higgs in the NMSSM. 

This work is organised as follows. In the first part we quickly review the CP-even Higgs sector of NMSSM relevant for our work, in the second part we discuss the calculation of the $H \rightarrow \gamma \gamma (Z^0)$ partial widths and review the effects of SM and SUSY particles inside the loops. In the next section we present the implementation and the numerical evaluation of the partial width within SloopS and then we carry out a numerical investigation to explore to which extent the signal can be enhanced in the NMSSM after applying various experimental constraints. In the last section we draw our conclusions.

II. CP-EVEN HIGGS SECTOR OF THE NMSSM

In the NMSSM superpotential the $\mu$-term involving the two Higgs doublet superfields $\hat{H}_u$ and $\hat{H}_d$ is absent and a gauge singlet superfield $\hat{S}$ is added instead [22, 23],

$$W_{\text{NMSSM}} = W_{\text{MSSM}}^\mu + \lambda S \hat{H}_u \cdot \hat{H}_d$$

(1)

The MSSM $\mu$ bilinear term has now been replaced by the trilinear coupling of the singlet with the two doublets and any dimensionful terms are forbidden by requiring the superpotential to be $Z_3$ symmetric. Once the singlet acquire a vev $s = \langle S \rangle$, an effective $\mu$ term is generated with respect to the MSSM, which is then naturally of order the EW scale,

$$\mu_{\text{eff}} = \lambda s$$

(2)

The soft-SUSY breaking Lagrangian is also modified according to

$$- \mathcal{L}_{\text{soft}} = m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 + m_S^2 |S|^2$$

$$\quad \quad + (\lambda A_\lambda H_u \cdot H_d S + \frac{1}{3} \kappa A_\kappa S^3 + h.c)$$

(3)

where $H_u \cdot H_d$ stand for the usual $SU(2)$ product: $X \cdot Y = X^T \theta Y$ with $\theta = -i \sigma_2$. Given $M_Z$ and using conditions coming from the minimisation of the Higgs potential, one can choose six independent parameters for the Higgs sector

$$\lambda, \kappa, A_\lambda, A_\kappa, \mu_{\text{eff}}, t_\beta$$

(4)

where $t_\beta = \tan \beta = v_u/v_d$, the ratio of the two Higgs doublet VEV’s : $\langle H_u^0 \rangle = v_u, \langle H_d^0 \rangle = v_d$. From the superpotential in eq. (1) one derives the tree-level Higgs potential containing the $D$, $F$- and soft-SUSY breaking terms (we stick to real parameters),

$$V_0 = \left( m_{H_u}^2 + \lambda^2 |S|^2 \right) |H_u|^2 + \left( m_{H_d}^2 + \lambda^2 |S|^2 \right) |H_d|^2 + |\lambda H_u \cdot H_d + \kappa S|^2$$

$$+ \frac{g^2}{8} \left( |H_u|^2 - |H_d|^2 \right)^2 + \frac{g^2}{8} \left[ (|H_u|^2 + |H_d|^2)^2 - 4 |H_u \cdot H_d|^2 \right]$$

$$+ m_S^2 |S|^2 + \left[ \lambda A_\lambda S H_u \cdot H_d + \frac{1}{3} \kappa A_\kappa S^3 + h.c \right]$$

(5)
with $g'$ and $g$ being the $U(1)_Y$ and $SU(2)_L$ couplings respectively. The parameters $m_{H_u}^2, m_{H_d}^2$ and $m_S^2$ can be traded for the vev’s $v_u, v_d, s$ through the minimisation conditions of $V_0$. The neutral physical fields are obtained by expanding the full scalar potential eq. (5) around the the vev's as

$$H_d = \left( v_d + (h_d^0 + i a_d^0)/\sqrt{2} \right), \quad H_u = \left( v_u + (h_u^0 + i a_u^0)/\sqrt{2} \right), \quad S = s + (h_s^0 + i a_s^0)/\sqrt{2} \quad (6)$$

The $3 \times 3$ symmetric CP-even Higgs mass matrix is derived by collecting the real parts and reads, in the basis $(h_d^0, h_u^0, h_s^0)$,

$$M_S^2 = \begin{pmatrix}
M_Z^2 s_{\beta}^2 + \mu_{\text{eff}} B_{\text{eff}} c_{\beta} & \lambda v (2 \mu_{\text{eff}} s_{\beta} - (B_{\text{eff}} + \nu) c_{\beta}) & \lambda v (2 \mu_{\text{eff}} c_{\beta} - (B_{\text{eff}} + \nu) s_{\beta}) \\
\mu_{\text{eff}} B_{\text{eff}} c_{\beta} & M_Z^2 c_{\beta}^2 & \lambda v (2 \mu_{\text{eff}} c_{\beta} - (B_{\text{eff}} + \nu) s_{\beta}) \\
\lambda v (2 \mu_{\text{eff}} s_{\beta} - (B_{\text{eff}} + \nu) c_{\beta}) & \lambda v (2 \mu_{\text{eff}} c_{\beta} - (B_{\text{eff}} + \nu) s_{\beta}) & \lambda^2 v^2 A_\lambda s_{\beta} - 2 \mu_{\text{eff}} + \kappa s (B_{\text{eff}} + 3 \nu)
\end{pmatrix} \quad (7)$$

where we have traded the $SU(2)_L \times U(1)_Y$ gauge couplings for the gauge boson masses through $M_Z^2 = (g^2 + g_s^2)/2$. $M_W^2 = g^2 v^2/2$. We used the short-handed notations $\nu = \kappa s, B_{\text{eff}} = A_\lambda + \nu$ as well as $c_{\beta} = c_\beta, s_{\beta} = s_\beta$ and so forth for trigonometric functions. The physical eigenstates $h_i^0$ (with $i = 1 \to 3$) are obtained by diagonalising eq. (7) with an orhogonal matrix $S$, such that $\text{diag}(m_{h_1}^2, m_{h_2}^2, m_{h_3}^2) = S \cdot M_S^2 \cdot S^{-1}$. Although it is possible, as said in the Introduction, to reproduce a 125 GeV SM-like Higgs mass in the NMSSM already at the tree-level, radiative corrections (mostly from the stop sector) to the Higgs sector are significant [22, 63, 64] and should be taken into account. For this purpose we computed the radiatively corrected Higgs masses using the publicly available code NMSSMTools [65, 67].

### III. THE $H \rightarrow \gamma\gamma$ AND $H \rightarrow \gamma Z^0$ DECAY WIDTHS

The partial width $\Gamma_{\gamma\gamma}$ and $\Gamma_{\gamma Z}$ in the case of the CP-even $h_i^0$ Higgs bosons can be written in SUSY generically as,

$$\Gamma_{\gamma\gamma}(h_i^0) = \frac{\alpha^2 G_F m_{h_i}^3}{128 \sqrt{2} \pi^3} \left| \sum_j A_j^{\gamma\gamma} + A_W^{\gamma\gamma} + A_{H^\pm}^{\gamma\gamma} + \sum_j A_j^{\gamma\gamma} + A_{\chi^\pm}^{\gamma\gamma} \right|^2 \quad (8)$$

$$\Gamma_{\gamma Z}(h_i^0) = \frac{\alpha^2 G_F m_{W}^2 m_{Z}^2}{64 \pi^4} \left( 1 - \frac{m_Z^2}{m_{h_i}^2} \right) \left| \sum_j A_j^{\gamma Z} + A_W^{\gamma Z} + A_{H^\pm}^{\gamma Z} + \sum_j A_j^{\gamma Z} + A_{\chi^\pm}^{\gamma Z} \right|^2 \quad (9)$$

The analytic expression for each amplitude $A_j$ ($j = W, f, H^\pm, \tilde{f}, \chi^\pm$) can be found in [14, 17, 60]. For the sake of completeness let us first discuss the SM-like contributions $A_W$ and $A_f$. As the Higgs boson couples to SM particles proportionally to their mass, its couplings to neutral gauge bosons are dominantly mediated by the heaviest charged particles of the SM: the $W^\pm$ boson and third generation quarks ($f = t, b$). The growing of the couplings with mass counterbalance the decrease of the triangle amplitudes with increasing loop mass, thereby not decoupling the contribution of heavy particles. The remaining fermions contributions are much smaller, due to smaller masses. These two partial width are therefore interesting probes of the number of heavy charged particles which can couple to the Higgs boson. In both decay channel the $W$ loops are by far the dominant ones and about 4.5 times
larger than the top quark amplitude for a 125 GeV Higgs bosons for the $\gamma\gamma$ channel and about one order of magnitude in the $\gamma Z^0$ case \cite{14}. However the total width is significantly reduced by the destructive interference between the two contributions. The full two-loop corrections (EW+QCD) for the SM-like Higgs decay into $\gamma\gamma$ and is under 2\% \cite{68} below the $W^+W^-$ theshold. The complete QCD corrections at the three loop level are also known and were presented in \cite{69}. For the $\gamma Z^0$ decay width the two loop QCD corrections to top quark loops were computed in \cite{70} and the relative magnitude of the QCD correction to the partial width for a 125 GeV SM-like Higgs boson are below 0.3 \%.

In SUSY theories, the additional superpartners of the SM particles do not couple to the Higgs boson proportionally to their masses, as the masses are generated through the soft-SUSY breaking Lagrangian and not the Higgs mechanism. Hence these contributions are suppressed by the heavy masses running in the loops. However if some of the superpartners masses are not too large, most notably the lightest chargino and third generation squarks when the squark mixing angle is large, the decay channels can be affected and their contribution can enable a discrimination between the lightest SUSY and standard Higgs boson even in the decoupling regime \cite{17}. Although the $h^0_i \rightarrow \gamma Z^0$ partial width is generically suppressed with respect to $h^0_i \rightarrow \gamma\gamma$, this channel is of interest as the non-diagonal couplings of the Higgs and the $Z^0$ gauge boson to sfermions $\tilde{f}_1\tilde{f}_2$ and $\tilde{\chi}^\pm\tilde{\chi}_2^\mp$ pairs are active. They are absent in the two-photon case due to the $U(1)_{\text{QED}}$ gauge invariance. The complete analytical LO SUSY amplitudes can be found in \cite{14,17,60}. The common lore is that these non-diagonal contributions are ignored since they are in general small for most purposes \cite{17}. In the package NMSSMTools \cite{65,67} the diagonal chargino loop contributions are included only and all sfermions contributions are missing. It was pointed out in \cite{60} that sometimes they may play a role and that some numerical factors were missing in \cite{14}. In the present work we computed these partial widths with the numerical code Sloops\cite{50,51} which automatically generates the one-loop amplitudes and shuns hand calculation errors thanks to internal checks like ultraviolet (UV) finiteness and gauge invariance. We now turn to the description of the numerical computation in the next section.

\section*{IV. NUMERICAL EVALUATION OF $\Gamma_{\gamma\gamma}$ AND $\Gamma_{\gamma\gamma}$ WITH SLOOPS}

In Sloops, the complete spectrum and set of vertices are generated at tree level through the LanHEP package \cite{71,73}. The complete set of Feynman rules is then derived automatically and passed to the bundle FeynArts/FormCalc/LoopTools \cite{74,76}(that we denote FFL for short from now on). A powerful feature of Sloops is the ability to check not only the UV finiteness check as provided by FFL, but also the gauge independence of the result through a generalised gauge fixing Lagrangian, which was adapted to the NMSSM \cite{56}. The gauge-fixing Lagrangian can be written in a general form

\begin{equation}
\mathcal{L}_{GF} = -\frac{1}{\xi_W} F^+ F^- - \frac{1}{2\xi_Z} |F^Z|^2 - \frac{1}{2\xi_A} |F^A|^2
\end{equation}
For the $\Gamma_{\gamma\gamma}$ and $\Gamma_{\gamma Z}$ decay width only the nonlinear form of $F^{\pm}$ is of relevance and is given by\(^1\)

$$F^{+} = \left( \partial_\mu - ie\tilde{\alpha}A_\mu - igc_W\tilde{\beta}Z_\mu \right) W^{\mu +} + i\xi_W g \frac{2}{2} \left( v + \sum_{i=1}^{3} \tilde{\delta}_i \tilde{\kappa}_i^0 \right) G^+ \tag{11}$$

The parameters $\tilde{\alpha}$, $\tilde{\beta}$ and $\tilde{\delta}_i$ are dubbed as nonlinear gauge (NLG) parameters and $G^{\pm}$ are the charged Goldstones fields. We recover the usual 't Hooft-Feynman gauge by setting these parameters to vanishing values. The ghost Lagrangian $L_{gh}$ is established by requiring that the full Lagrangian is invariant under BRST transformations\(^2\). The gauge dependence is in turn transferred from the vector boson propagators to a modification of the ghost-goldstones-vector boson vertices (see for example [77]). Numerically the gauge invariance check is performed by varying the parameters $\tilde{\alpha}$, $\tilde{\beta}$ and $\tilde{\delta}_i$.

Similarly to the MSSM, radiative corrections to the Higgs masses in the NMSSM can be relatively large (see e.g [63] and references therein) and thus affect significantly the kinematics of the decay. As said previously we used NMSSMTools to compute the Higgs spectrum and rotation matrices. Since the Higgs spectrum and parameters entering the Higgs potential in NMSSMTools of the decay. As said previously we used NMSSMTools to compute the Higgs spectrum and rotation matrices. Since the Higgs spectrum and parameters entering the Higgs potential in NMSSMTools of the decay. As said previously we used NMSSMTools to compute the Higgs spectrum and rotation matrices. Since the Higgs spectrum and parameters entering the Higgs potential

$$V_{\text{eff}} = V_0 + V_{\text{rad}} \tag{12}$$

where $V_0$ is the same as eq. [5] and

$$V_{\text{rad}} = \frac{\lambda_1}{2} |H_d|^4 + \frac{\lambda_2}{2} |H_u|^4 + \lambda_3 |H_u|^2 |H_d|^2 + \lambda_4 |H_u| H_d|^2 + \tilde{\kappa}^2 |S|^4 + \frac{1}{3} (\bar{A}_S S^3 + h.c.)$$

$$+ \lambda_5^u |S|^2 |H_u|^2 + \lambda_4^d |S|^2 |H_d|^2 + [A_{ud} S H_u \cdot H_d + \lambda_5^M S^3 H_u \cdot H_d + h.c.] \tag{13}$$

Once the spectrum and mixing matrices are known from NMSSMTools we solve for the $\lambda$'s\(^3\) using equations derived from the effective Higgs mass matrices extracted from eq. [12]. We refer to [59] for further details concerning the extracting procedure of the $\lambda$ parameters. This procedure then ascertain the gauge independence of the computation. This was explicitly demonstrated analytically and numerically in [59] for the partial width $h_0^0 \to \gamma\gamma$. The nonlinear gauge fixing Lagrangian possesses another virtue: setting particular values to the NLG parameters can cancel specific vertices. For the case at hand, with the peculiar choice $\tilde{\alpha} = -1$ the coupling $\gamma W^{\pm} G^\mp$ is absent due to an underlying $U(1)_{\text{QED}}$ symmetry conserving gauge-fixing function $F^{\pm}$ and less diagrams have to be considered. This property is therefore a welcomed feature and was employed to simplify the analytic calculation of $H \to \gamma\gamma$ in [78] [79]. Although in such a gauge the calculation of the SM-like amplitude of $\Gamma_{\gamma Z}$ is not easily translated from $\Gamma_{\gamma \gamma}$, since the vanishing of the $\gamma W^{\pm} G^\mp$ vertex introduces an asymmetric treatment of the photon and the Z boson [80], the evaluation of the $Z - \gamma$ transition diagram $\Pi_{\gamma Z}(0)$ is not needed. Indeed, for $\tilde{\alpha} = 1$ this mixing self-energy vanishes

\(^1\) The other gauge-fixing functions $F^{Z}$ and $F^{A}$ can be kept in the usual $R_f$ form. The complete expressions can be found in [50]. For practical purposes we set $\xi_{W,Z,A} = 1$.

\(^2\) The BRST transformations for the gauge fields can be found in [77] and for the scalar fields in [59].

\(^3\) We denote generically as “$\lambda$” any parameter entering eq. [13].
at \( q^2 = 0 \) thanks to, once more, the fact that \( F^\pm \) preserves the \( U(1)_{\text{QED}} \) gauge symmetry [77]. This is of importance for our numerical evaluation of \( \Gamma_{\gamma Z} \) since we do not generate diagrams with external wave function correction and the introduction of the field-renormalisation constant \( \delta Z_{\gamma Z}^{1/2} = -\Pi_{\gamma Z}^0(0)/M_Z^2 \) would not be needed in such a gauge. In a general nonlinear gauge this transition is crucial to maintain the UV finiteness and the gauge invariance of the result. We checked thoroughly this feature numerically by varying the NLG parameters \( \tilde{\alpha}, \tilde{\beta} \) and \( \tilde{\delta} \).

V. NUMERICAL INVESTIGATION

Although the properties of the Higgs boson observed at the LHC are compatible with the SM predictions [1, 81], the precise measurements of all its decay channels can give some crucial information on new physics. In this analysis we examine the expectations for the decay \( h^0 \to \gamma Z \) in the framework of the NMSSM after taking into account the constraints on the Higgs observed at the LHC. In particular we quantify the importance of the sfermions and non-diagonal charginos contributions discussed in the previous section. For this we explore the parameter space of the NMSSM with emphasis on the regions which could lead potentially to large corrections, those with light charginos and/or light stop.

The chargino mass matrix is given by,

\[
\begin{pmatrix}
M_2 & gv_u \\
gv_d & \mu_{\text{eff}}
\end{pmatrix}
\]

while the stop mass matrix reads

\[
\begin{pmatrix}
m_{\tilde{U}_3}^2 & h_t^2 v_u^2 - (v_u^2 - v_d^2) \frac{g_2^2}{3} \\
h_t(A_t v_u - \mu_{\text{eff}} v_d) & m_{\tilde{Q}_3}^2 + h_t^2 v_u^2 + (v_u^2 - v_d^2) \left( \frac{g_1^2}{12} + \frac{g_2^2}{4} \right)
\end{pmatrix}
\]

where \( M_2 \) is the SU(2) gaugino mass, \( m_{\tilde{Q}_3}, m_{\tilde{U}_3} \), the soft masses for the stops and \( A_t \) the stop trilinear mixing. Since \( h_t \) is of order 1, the mixing between both stops can be important. The large mixing can lead to large radiative corrections to the SM-like Higgs mass and to one of the stop being quite light. Thus the main squark contribution to the Higgs loop-induced decays (\( \gamma\gamma, \gamma Z^0 \) and \( gg \)) is coming from the stop sector.

To restrict the number of free parameters of the phenomenological NMSSM, we perform a scan over only the parameters most relevant for the Higgs mass (eq. [7]) and couplings, specifically those of the chargino, squark and Higgs sectors which we take in the following range

\[
\begin{align*}
100 \text{ GeV} & < \mu < 500 \text{ GeV} \\
100 \text{ GeV} & < M_2 < 1000 \text{ GeV} \\
0 \text{ GeV} & < t_\beta < 20 \\
0 & < \lambda, \kappa < 0.7 \\
100 \text{ GeV} & < A_\lambda < 1000 \text{ GeV} \\
-1000 \text{ GeV} & < A_\kappa < -100 \text{ GeV} \\
-3000 \text{ GeV} & < A_t < 3000 \text{ GeV} \\
400 \text{ GeV} & < m_{\tilde{Q}_3}, m_{\tilde{U}_3} < 2000 \text{ GeV}
\end{align*}
\]
We assume that all squarks of the first and second generations are heavy ($m_{\tilde{Q}_i} = m_{\tilde{u}_i} = m_{\tilde{d}_i} = 2\text{TeV}$) as well as the right-handed sbottom mass, $m_{\tilde{d}_R} = 2\text{TeV}$, since they do not play an important role in Higgs physics. We also assume that all sleptons are heavy $m_{\tilde{l}_i} = m_{\tilde{\ell}_i} = m_{\tilde{\nu}_i} = 2\text{TeV}$  as well as the gluino, $M_3 = 1.5\text{TeV}$. The most important LHC constraints on supersymmetric particles are then automatically satisfied. Finally we set $M_1 = 150\text{GeV}$, the exact value of the neutralino LSP is not very important for our analysis, provided the neutralino is too heavy for the Higgs to decay invisibly. However, the value of $M_1$ could be adjusted to ensure that the limit on the stop mass from the LHC is satisfied (the limit on the lightest stop can easily be avoided when $m_{\tilde{t}_1} - m_{\tilde{\chi}^0_1} < m_{\tilde{t}_1}$) [82, 83]. Note that we concentrate on small values of tan $\beta$ since it is the region where the Higgs sector can differ significantly from that of the MSSM. In this region one can find large deviations in the $h^0 \rightarrow \gamma \gamma$ decay [29] due in particular to the singlet component of the Higgs. Thus large deviations are also expected for $h^0 \rightarrow \gamma Z^0$.

In the NMSSM, either of the light scalar, $h^0_1$ or $h^0_2$ could be the one observed at the LHC with a mass of 125 GeV, we consider both possibilities. We use NMSSMTools to compute the supersymmetric spectrum and to impose constraints on the parameter space, specifically: the LEP constraints on Higgs and chargino masses as well as the constraint that there be no unphysical global minimum on the Higgs potential. We select only the points with one Higgs in the mass range $122 - 128 \text{GeV}$. Finally, we impose the collider constraints on the Higgs sector from HiggsBounds [61] and require that one Higgs fits the properties of the particle observed at the LHC using HiggsSignals [62]. For these two codes, we chose a theoretical uncertainty of 2GeV for the Higgs masses. The allowed points are those for which $\chi^2 < \chi^2_{\text{best fit}} + 18.3$ corresponding to 95% confidence level for ten free parameters. Note that the best fit point is slightly better than the SM.

First we checked whether sfermions and non-diagonal charginos contributions have a significant impact on the the process $h^0_i \rightarrow \gamma Z^0$, for this we calculate the following ratio:

$$R = \frac{\Gamma(h^0_i \rightarrow \gamma Z^0)_{\text{total}}}{\Gamma(h^0_i \rightarrow \gamma Z^0)_{\text{restricted}}}, \quad (16)$$

where $\Gamma(h^0_i \rightarrow \gamma Z^0)_{\text{total}}$ is the decay width calculated with all the possible particles in the loops, and $\Gamma(h^0_i \rightarrow \gamma Z^0)_{\text{restricted}}$ is the one calculated by omitting sfermions and non-diagonal charginos contributions in the loop.

The results for the ratio are shown in figure[1]. In both cases, the ratio is plotted as function of the mass of the corresponding Higgs boson. We notice that for most of the points, the effect is less than $\sim 10\%$. In fact, the main effect is coming from the chargino contributions, we have checked that with only the sfermionic contribution the effect is less than 5%. However, for a few points, the variation can be as high as 70% in the case of $h^0_2$. These large deviations are found for points for which $h^0_2$ is almost singlet : in this case the partial decay width is suppressed since the W contribution becomes negligible and the chargino (higgsino) contribution can become relatively more important. Note that for the singlet case the total decay width is also suppressed leading to a branching ratio that can be either

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4 We have made additional scans to check the impact of varying the parameters of the stau sector. Corrections to $h^0 \rightarrow \gamma Z^0$ lie below a few percent except for very light staus (below the LEP limit). We expect that for values of tan $\beta$ much larger than those considered here, one can get large enhancement to the $\gamma Z^0$ branching ratio as was shown previously for the $h \rightarrow \gamma \gamma$ in the MSSM [13].
enhanced or suppressed relative to the SM. However most of these points are excluded by LHC constraints on the Higgs sector.

![Graph showing ratio of decay widths of h_i^0 → γZ^0 and h^0_i → γZ^0](image)

**FIG. 1:** Ratio of the decay width of h^0_i → γZ^0 (left panel) and h^0_i → γZ^0 (right panel) with all possible particles in the loop by the one without sfermions nor non diagonal Z-charginos contributions as function of the mass of h^0_i. Blue points are allowed and green triangles are excluded by LHC constraints on the Higgs.

We next consider the expectations for the h^0_i → γZ^0, γγ branching ratios in the NMSSM as compared to the SM when including all contributions. The case where h^0_1 is near 125 GeV leads to only mild variations of the γZ^0 branching ratio - typically ≈ 10%. In a few cases however deviations as large as 25% can be found, typically they are found when h^0_1 has a significant singlet component. In all cases we found a strong correlation to the h^0_1 → γγ branching ratio (within 10%). The two-photon mode should therefore provide a better probe of new physics effects in the NMSSM considering it can be measured with a much better precision.

The results for h^0_2 → γZ^0, γγ are more interesting. Figure 2 shows the branching ratio h^0_2 → γZ^0 as compared to the SM value as function of the mass of the lightest Higgs boson (recall that for these points 122 GeV < m_h^0 < 128 GeV). Although the bulk of the points are centered around the SM value, large deviations can occur, in particular when the lightest Higgs is near 120 GeV. In this case h^0_2 has a large singlet component and the Br(h^0_2 → γZ^0) can be up to 4 times larger than the SM or suppressed by more than two orders of magnitude. Note that most of these points are excluded by the LHC constraints as implemented in HiggsSignals either because the production of the singlet Higgs deviates significantly from the SM and/or the corresponding γγ channel which is correlated with γZ^0 is incompatible with current measurements. Nevertheless, we found few points that satisfy the HiggsSignals constraints even though the h^0_2 is almost pure singlet and thus have non SM couplings, see empty circles in Fig. 2 which have S_{h_23} > 0.9. The reason why such points avoid the LHC constraints is that they correspond to cases where h^0_1 and h^0_2 are almost degenerate and the superposition of the signal of both Higgs bosons is what is observed at the LHC (h^0_1 is mainly doublet and SM-like).

The branching ratio h^0_2 → γZ^0 can also be strongly suppressed when m_h^0 < 60 GeV and a new decay channel h^0_3 → h^0_1h^0_1 opens up thus increasing significantly the total width, see Fig. 2. Points with a large suppression are however incompatible with the LEP and LHC constraints.

Finally we comment on the correlation between the h^0_2 → γZ^0 and h^0_2 → γγ channels displayed in Fig. 2 right panel. As for the case of h^0_1 both channels are strongly correlated. The
correlation breaks down when $h_2^0$ has a strong singlet component and mainly for points with suppressed $\gamma Z^0/\gamma \gamma$ branching ratios that are to a large extent excluded by LHC constraints.

The LHC collaborations will not measure directly the branching ratio but the signal strengths ($\mu$) for gluon fusion ($gg$) and vector boson fusion (VBF) production modes in the $\gamma Z^0$ channel. The predictions for the signal strength in the gluon fusion mode are shown in figure 3 for $h_2^0$ as function of the scalar mixing $S_{h_23}$. Three regions present an enhancement as compared to the SM, they correspond to a) $S_{h_23} < -0.5$, b) $S_{h_23} > 0.7$, and c) $S_{h_23} \approx 0.4$. In all three cases, $h_2^0$ has significant singlet and doublet components and its couplings to u-type quarks and gauge bosons are somewhat suppressed while those to b-type quarks and leptons are strongly suppressed. As a result the total width of $h_2^0$ is much reduced and the branching ratio into $WW, ZZ, \gamma \gamma, gg$ and $\gamma Z^0$ are all larger than in the SM. The signal strength in the VBF mode is mostly correlated with that in the gluon fusion mode, although for some points the VBF is suppressed by more than a factor 2 as compared to the gluon fusion, see figure 3 right panel. In particular note that the enhancement in the gluon fusion mode ($\mu_{gg} > 1$) is more important than in the VBF mode. We had also mentioned that for $S_{h_23} > 0.9$ which corresponds to a $h_2^0$ dominantly singlet, the $h_2^0 \rightarrow \gamma Z^0$ branching ratio could still be enhanced, however the singlet coupling to gauge bosons becomes very small so that the VBF production mode is suppressed and to a lesser extent also the gluon fusion production mode leaving making it difficult to probe this dominantly singlet Higgs at the LHC.

VI. CONCLUSION

We have performed a complete computation of the branching ratio into $\gamma Z^0$ of either of the two lightest CP -even Higgs in the NMSSM using $S\text{loopS}$. We found that the previously neglected contributions of charginos and stops (as well as staus) generally lie below 10%. Exploring the parameter space of the NMSSM we found that $h_1^0 \rightarrow \gamma Z^0$ did not vary much from the SM expectations while regions with large enhancement (suppression) of $h_2^0 \rightarrow \gamma Z^0$ were possible, especially when $h_2^0$ had a significant singlet component. However most of these scenarios are constrained by measurements of the 125 GeV Higgs boson at the LHC.
FIG. 3: Left panel: Signal strength for $h_2^0 \rightarrow \gamma Z^0$ for gluon gluon fusion production in function of $Sh_{23}$. Right panel: Correlation between the gg and VBF signal strengths. Same color code as Fig. 2

since the singlet component significantly changes the coupling of $h_2^0$ to gauge bosons and fermions (in particular b quarks). We conclude that given the correlation between the $\gamma\gamma$ and $\gamma Z^0$ branching ratios expected in the NMSSM, a better measurement of the former at the next LHC run - together with a higher precision on other standard decay channels - will further constrain the range of values expected for $h \rightarrow \gamma Z^0$. Nevertheless an independent measurement of $h \rightarrow \gamma Z^0$ is useful in probing BSM physics, for example a large deviation from the SM expectations not correlated with a similar deviation in the $\gamma\gamma$ mode would put strong constraints on the NMSSM (and other MSSM-like models). Furthermore a suppressed signal strength in the VBF mode relative to the gluon fusion mode is characteristic of the partially singlet Higgs in the NMSSM.

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