Naughty or Nice? The Role of the ‘N’ in the Natural NMSSM for the LHC

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In this work, we present mass limits on gluinos and stops in a natural Next-to-Minimal Supersymmetric Standard Model (NMSSM) with a singlino as the lightest supersymmetric particle. Motivated by naturalness, we consider spectra with light higgsinos, sub-TeV third generation sparticles and gluinos well below the multi-TeV regime while the electroweak gauginos, the sleptons and the first and second generation squarks are decoupled. We check that our natural supersymmetry spectra satisfy all electroweak precision observables and flavour measurements as well as theoretical constraints. By reinterpreting the results from the 8 TeV ATLAS supersymmetry searches we present the 95% CL exclusion limits on the model. The results show that the presence of a singlino LSP can lengthen decay chains and soften the final state particle energies. Whilst this does reduce the strength of the bounds in some areas of parameter space, the LHC still displays good sensitivity to the model.

I. INTRODUCTION

The discovery of a Higgs Boson at the Large Hadron Collider in 2012 [1]–[3] was a triumph for experimental particle physics. Its measured mass of ≈ 125 GeV fits perfectly into the framework of the Standard Model (SM), which required — and hence predicted — a scalar particle with mass of the order ≤ 2 TeV from unitarity constraints [4]. Moreover, the observed Higgs mass falls into the narrow mass window $m_h = 96^{+31}_{-24}$ GeV predicted in global fits of the SM model to precision electroweak observables [5]. However, this theory suffers from a well known Hierarchy Problem due to the quadratic sensitivity of the Higgs mass to new physical scales. Supersymmetry (SUSY) [6]–[9] is able to ameliorate this problem and to stabilise the Higgs mass at the electroweak scale by cancelling the quadratic divergences. However, the minimal incorporation into the Standard Model, called the Minimal Supersymmetric Standard Model (MSSM), leads to the prediction that the CP-even Higgs should be lighter than the Z–Boson at tree-level. Under the assumption that the LHC has measured the lightest supersymmetric Higgs boson, we thus require significant radiative corrections in order to raise the mass to the experimentally measured value. These corrections are often provided by the supersymmetric partners of the top quark, called stop squarks, since their large Yukawa couplings dictate that they provide the leading one-loop correction.

In the MSSM, the Higgs is expected to have a mass between 113 and 135 GeV [10]. However, the problem is that these corrections only reproduce the correct Higgs mass when both stop masses have very large masses for negligible mixing in the stop sector, or the trilinear $A_t$ term is very large with at least one heavy stop (e.g. [11]). This is an issue since a large separation between the electroweak and SUSY breaking scale introduces the little hierarchy problem [12]. Thus, in this case the model is deemed to be ‘unnatural’ since fine-tuned cancellations are still required. To confront this issue in the MSSM, extended models that already predict a heavier Higgs mass at the tree level have become popular. The simplest example of such a model is known as the Next-to-Minimal Supersymmetric Standard Model (NMSSM) where an extra gauge singlet chiral superfield is added to the spectrum [13]. Since the singlet superfield couples both to the up- and the down-type Higgs superfield, the singlet scalar components contribute to the Higgs potential and can thus raise the tree-level Higgs mass, reducing the need for heavy stops.

In addition, the only dimensionful SUSY conserving parameter, $\mu$, of the MSSM can be dynamically generated in the NMSSM by a non-vanishing vacuum expectation value $s$ of the extra singlet scalar. To get a phenomenologically acceptable scenario of electroweak symmetry breaking, $|\mu|$ should lie within $M_Z$ and $M_{\text{SUSY}}$, the scale where Supersymmetry is broken. In the MSSM, the scale of $\mu$ is in principle arbitrary and no theoretical reasoning binds it to low scales, which leads to the so-called $\mu$-problem [14]. In the NMSSM, however, the effective $\mu$ parameter is determined by the scale of the vev $s$, which is automatically of the right order.

The large Higgs mass is not the only experimental evidence from the LHC that puts the idea of SUSY solving the hierarchy problem under strain. The fact that
no SUSY particles have yet been seen pushes the limits on SUSY gluons, called gluinos, and SUSY quarks, called squarks, to masses $\geq 1.5$ TeV (see e.g. [15]) that would already be deemed unnatural in constrained models like the CMSSM. These results have motivated a deeper study of exactly which pieces of the SUSY spectrum are required to be light for a theory to be considered natural [19]. Firstly, since the singlet itself now generates the $\mu$ parameter that sets the Higgsino masses, all of these particles, including the fermionic partner of the singlet (singlino), can be expected to have masses of the same order. Furthermore, the dominant one-loop corrections to the Higgs sector come from the stops and consequently these cannot be too heavy. Also, since the gluino yields a sizeable correction to the stop masses at one loop, we also have another, looser constraint on the mass of this particle for the same reason. Finally due to the weak isospin symmetry, the partners of the left handed bottom quarks (sbottoms), must have a mass similar to that of the left handed stops.

Consequently we are drawn to a SUSY spectrum with light singlinos and higgsinos, stops and sbottoms that may be a little more massive and a gluino that can be heavier still. Since none of the other SUSY partners are required by naturalness principles to be light enough to be seen at the LHC, we simply decouple these from our spectrum in this study.

In the context of the MSSM, naturalness is now used as a guiding principle for many LHC searches for gluinos, stops, sbottoms. These studies set bounds on the gluino of $m_{\tilde{g}} \geq 1150$ GeV in the case of a light ($m_{\tilde{\chi}_1^0} \leq 100$ GeV) LSP, but this can be reduced to, $m_{\tilde{g}} \geq 500$ GeV in the limit that the gluino becomes degenerate with the LSP [15, 17, 21]. For stops, the bounds can reach up to $m_{\tilde{t}_1} \geq 700$ GeV, if the dominate decay mode is $t \rightarrow t\tilde{\chi}_1^0$, and $m_{\tilde{t}_1} \geq 600$ GeV for $t \rightarrow b\tilde{\chi}_1^+$. Again, if the spectrum is compressed, the bounds weaken significantly and the limit is only $m_{\tilde{t}_1} \geq 255$ GeV for $m_{\tilde{t}_1} - m_{\tilde{\chi}_1^0} \approx m_t$ [22] In addition there are regions of parameter space ($m_{\tilde{t}_1} \sim m_t + m_{\tilde{\chi}_1^0}$) where no limit can be set at all since the kinematics very closely resemble the SM tt background but with a substantially smaller production cross-section [19, 23, 26]. Sbottom limits are similar to those of stops (up to $m_{\tilde{b}} \geq 650$ GeV for light $\chi_1^0$ and $m_{\tilde{b}} \geq 250$ GeV in compressed regions) but are more robust and do not contain holes as we move across the mass plane [23, 27, 29].

Since the LHC direct production constraints still allow for relatively light gluinos and have no model independent limit on the stop mass, the question of naturalness is driven by the Higgs mass in the MSSM. In the NMSSM however, the reduced need for heavy stops to contribute to the Higgs means that the direct production constraints become far more relevant. In addition, the limits can be expected to be different since a light singlino will be present in the spectrum. However, as the singlino does not couple directly to the squarks the state does not normally contribute to LHC phenomenology unless it is the lightest particle in the spectrum (LSP). The effect of a singlino LSP has now been examined in a number of studies and it has been claimed that it generally weakens the LHC limits since the longer decay chains softens the $p_T$ spectra and reduces the $E_T^{\text{miss}}$ [30, 31]. Other studies have also looked at purely Higgsino-singlino spectra [32, 33], direct stop [34, 35] or gluino [36] production and the possibility that the singlino may be light [37, 42]. A comprehensive list of the expected signatures of the NMSSM is given in [43] whilst [44] has explored possible methods to distinguish the NMSSM from the more commonly discussed MSSM.

In this study we wish to explore in detail the claim that a singlino LSP generally weakens the LHC bounds. As stated above this is expected and seen [30, 31] because the longer decay chain produce soft particles for similar LSP masses. However if we examine the particles produced in the extra NMSSM decay we see that this may not always be true. In particular the decays that may occur are $\tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0 X^0$ where $X^0$ is either a $Z^0$ or Higgs and $\tilde{\chi}_2^\pm \rightarrow \tilde{\chi}_1^0 X^\pm$ where $X^\pm$ is either a $W^\pm$ or a charged Higgs. In the case of $W^\pm$ or $Z^0$ production we can expect increased production of leptons over the MSSM that may improve the bounds but the branching ratio suppression makes it unlikely that this will result in a large change. However a bigger difference can be expected when a Higgs is produced that will decay to a $bb$ final state. If the mass splitting $m_{\tilde{\chi}_2^0} - m_{\tilde{\chi}_1^0} > m_h$, then this decay dominates in a large portion of the natural NMSSM parameter space. The reason that this final state can be so important for LHC phenomenology is that many SUSY searches use $b$-tags as a way to suppress the SM background (e.g. [45, 48]) and some even search for the presence of on-shell Higgs bosons e.g.([49, 53]). Both of these strategies give the possibility that the natural NMSSM may be even more constrained than the MSSM.

In order to fully test the effect of the additional $b$-quarks we require many LHC searches to be simultaneously checked. For this reason we use the Checkmate [54] tool which now contains over 40 analyses implemented via AnalysisManager [55]. In addition, we also test various theoretical and experimental constraints via NMSSTools [56, 60], HiggsSignals [61] and HiggsBounds [62].

We begin the paper by describing the Lagrangian of the natural NMSSM in Sec. [II] and the spectrum that we decide to investigate along with the LHC signatures this will lead to. In Sec. [III] we describe exactly how the model parameters are chosen and the experimental and theoretical constraints that are applied. Here we also introduce how we perform the LHC phenomenology in this paper. Sec. [IV] displays the results of our study, concentrating on the LHC bounds now present on the natural NMSSM. Finally in Sec. [V] we conclude.
II. A PHENOMENOLOGICALLY NATURAL NMSSM

A. Lagrangian, Masses and Parameters

In the following we present the Lagrangian formulation of our model, the resulting mass matrices and the features that motivate our spectra. Most definitions and relations are taken from [13] and we refer readers to check this source and references therein for more information.

The Next-to-Minimal Supersymmetric Standard Model (NMSSM) extends the well-known Minimal Supersymmetric Standard Model (MSSM) by an additional chiral superfield $\hat{S}$ which is uncharged under the Standard Model gauge groups. In this work we consider a simplified, natural version of a $Z_3$–invariant NMSSM. Here, only terms involving exactly three fields are allowed to appear in the superpotential, for reasons explained below. Furthermore, only the scalar partners of the gluon and the third generation quarks plus the fermionic components of the three Higgs superfields $H_u, H_d$ and $\hat{S}$ are assumed to be phenomenologically observable among all supersymmetric particles. This setup can be described by the following superpotential:

$$W = h_t(\hat{Q}_3 \cdot H_u)\hat{t}_R^c + h_b(\hat{Q}_3 \cdot \hat{H}_u)\hat{b}_R^c + \lambda(\hat{H}_u \cdot \hat{H}_d)\hat{S} + \kappa/3 \hat{S}^3,$$

(1)

where the $\cdot$ symbol denotes the usual SU(2) invariant antisymmetric product of the respective isospin doublets $\hat{H}_u \equiv (\hat{H}_u^+ \cdot \hat{H}_u^0), \hat{H}_d \equiv (\hat{H}_d^0 \cdot \hat{H}_d^-)$ and $\hat{Q}_3 \equiv (\hat{t}_L, \hat{b}_L)$. Here, $h_t$ and $h_b$ are the dimensionful Yukawa couplings, while $\lambda$ and $\kappa$ correspond to dimensionless Yukawas. Note that the assumed additional $Z_3$ symmetry prohibits the term $\mu(H_u \cdot H_d)$ which is present in the MSSM and hence provides a superpotential without any dimensionful parameters. A vacuum expectation value (vev) of the scalar singlet $\langle S \rangle \equiv s$ of electroweak scale order introduces this term after expanding the scalar field $S$ around its minimum and thus generates an effective $\mu$ term $\lambda s(H_u \cdot H_d) \equiv \mu_{\text{eff}}(H_u \cdot H_d)$ of naturally the correct scale, evading the known $\mu$–problem of the MSSM.

In addition to the terms derived from this superpotential, the following dimensionful ‘soft’ parameters have to be added to the Lagrangian of the theory:

$$-\mathcal{L}_{\text{soft}}^\text{mass} = m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 + m_S^2 |S|^2 + \frac{1}{2} M_3 \tilde{g} \tilde{g} + m_{\hat{Q}_3}^2 |\hat{Q}_3|^2 + m_{\hat{t}_R}^2 |\hat{t}_R|^2 + m_{\hat{b}_R}^2 |\hat{b}_R|^2,$$

(2)

Thus, the effective Lagrangian of the NMSSM reads

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{soft}}^\text{mass} + \mathcal{L}_{\text{scalar}} + \mathcal{L}_{\text{trilinear}} (3)$$

Here, $\tilde{g}$ denotes the gluino, i.e. the fermionic part of the vector superfield associated to the SU(3) gauge group. All other field names denote the scalar component of the respective chiral superfield in the superpotential. We assume all couplings to be real–valued to simplify the discussion.

The other SUSY particles, namely the squarks of the first two generations, the sleptons as well as the SU(2)×U(1) gauginos are assumed to be decoupled from the experimentally accessible spectrum as explained in Sec. IIIB. Consequently they are not listed here.

Adding $\mathcal{L}_{\text{soft}}$ to the supersymmetric F– and D–terms yields the full scalar potential from which three minimisation conditions for the non-vanishing singlet vevs and the doublet vevs $\langle H_u/d \rangle \equiv v_u/d$ can be derived:

$$m_{H_u}^2 + \mu_{\text{eff}}^2 + \lambda^2 v_u^2 + \frac{(\tilde{g}_1^2 + \tilde{g}_2^2)}{4(v_u^2 - v_d^2)} = \mu_{\text{eff}}B_{\text{eff}} \cot \beta,$$

(4)

$$m_{H_d}^2 + \mu_{\text{eff}}^2 + \lambda^2 v_d^2 + \frac{(\tilde{g}_1^2 + \tilde{g}_2^2)}{4(v_u^2 - v_d^2)} = \mu_{\text{eff}}B_{\text{eff}} \tan \beta,$$

(5)

$$m_S^2 + \kappa s (\lambda + 2\kappa s) + \lambda^2 v^2 - \kappa v_u v_d = \lambda v_u v_d s B_{\text{eff}}$$

(6)

with $B_{\text{eff}} \equiv (A_{\lambda} + \kappa s), v_u \equiv v_u^2 + v_d^2$. The first two of these can be reformulated as follows:

$$\frac{M_Z^2}{2} = \frac{2}{\tan \beta^2 - 1} \left( m_{H_d}^2 - \tan \beta^2 m_{H_u}^2 \right) - \mu_{\text{eff}}^2,$$

(7)

$$\frac{\sin 2\beta}{2} = \frac{\mu_{\text{eff}} B_{\text{eff}}}{m_{H_u}^2 + m_{H_d}^2 + 2\mu_{\text{eff}}^2 + \lambda^2 v^2}.$$

(8)

Here we have used the fact that $M_Z^2 = v^2(\tilde{g}_1^2 + \tilde{g}_2^2)/2$ is fixed by the known mass of the $Z$ boson. The above equations allow one to choose the parameters in the set $\{ \lambda, \kappa, A_\lambda, A_\kappa, \mu_{\text{eff}}, \tan \beta \equiv v_u/v_d \}$ to be independent, where $\mu_{\text{eff}}, \tan \beta$ and the known Standard Model parameter $M_Z$ replace the Lagrangian parameters $m_{H_u}^2, m_{H_d}^2$ and $m_S^2$. After expanding $H_u, H_d$ and $S$ around their minima, one gets the following symmetric mass matrix for their CP-even components $\{h_u, h_d, h_s\}$ at tree level:

$$\mathcal{M}_{\text{scalar}}^2 = \begin{pmatrix} g^2 v_d^2 + \mu_{\text{eff}} B_{\text{eff}} \tan \beta & (2\lambda^2 - \tilde{g}^2) v_u v_d - \mu_{\text{eff}} B_{\text{eff}} & \lambda(2\mu_{\text{eff}} v_d - (B_{\text{eff}} + \kappa s v_u) \\
\cdots & g^2 v_u^2 + \mu_{\text{eff}} B_{\text{eff}} \cot \beta & \lambda(2\mu_{\text{eff}} v_u - (B_{\text{eff}} + \kappa s v_d) \\
\cdots & \cdots & \lambda A_{\lambda} v_u v_d/s + \kappa s A_{\kappa} + 4\kappa^2 s^2 \end{pmatrix}.$$
with $g^2 = M_Z/(v_u^2 + v_d^2)$ given by the Standard Model gauge sector. We call the diagonalised mass eigenstates $h, H$ and $H_3$ which have increasing mass from left to right.

Similarly, the matrix of the respective CP–odd components $\{a_u, a_d, a_s\}$ reads

$$\mathcal{M}^2_{\text{pseudoscalar}} = \begin{pmatrix} \mu_{\text{eff}}B_{\text{eff}}\cot\beta & \mu_{\text{eff}}B_{\text{eff}} & \lambda v_u (A_\lambda - 2k_s) \\ \mu_{\text{eff}}B_{\text{eff}} & \mu_{\text{eff}}B_{\text{eff}}\cot\beta & \lambda v_d (A_\lambda - 2k_s) \\ \ldots & \ldots & \lambda (A_\lambda + 4ks)v_uf_s - 3k_A s \end{pmatrix},$$

(10)

using the matrix

$$\mathcal{M}_{\text{neutralinos}} = \begin{pmatrix} 0 & -\mu_{\text{eff}} & -\lambda v_u \\ -\mu_{\text{eff}} & -\lambda v_d & \ldots \\ \ldots & \ldots & 2k_s \end{pmatrix}. \quad (13)$$

The two charged higgsino components combine to a single Dirac chargino $\tilde{\chi}^\pm_1$ with mass term $\frac{1}{2}\mu_{\text{eff}}h^+_u h^+_d + h.c.$.

In the following, we will use the collective term ‘higgsino’ $(\tilde{h})$ for the two higgsino-like neutralinos and the chargino. Furthermore, for the sake of simplicity, we will use ‘electroweakino’ $(\tilde{\chi})$ collectively for all three neutralinos and the chargino, even though strictly speaking $\tilde{s}$ does not have any electroweak charge.

The stop and sbottom tree level mass matrices in the bases $(\tilde{t}_R, \tilde{t}_L)$ and $(\tilde{b}_R, \tilde{b}_L)$ read

$$\mathcal{M}^2_{\text{stops}} = \begin{pmatrix} m_{\tilde{t}_3}^2 + h_t^2 v_u^2 - (v_u^2 - v_d^2)g_d^2/3 \\ \ldots \\ m_{\tilde{D}_3}^2 + h_b^2 v_d^2 - (v_u^2 - v_d^2)g_u^2/6 \end{pmatrix},$$

$$\mathcal{M}^2_{\text{sbottoms}} = \begin{pmatrix} m_{\tilde{q}_3}^2 + h_t^2 v_u^2 + (v_u^2 - v_d^2)(g_d^2/12 - g_d^2/4) \\ \ldots \end{pmatrix},$$

(14)

(15)

with eigenstates $\tilde{t}_{1/2}, \tilde{b}_{1/2}$.

Even though we haven’t shown the full NLO corrections to these tree level masses, it can be understood that the whole model is fixed by Standard Model parameters plus the set $\{\lambda, k, A_\lambda, A_\kappa, \mu_{\text{eff}}, \tan\beta, m_{\tilde{Z}_3}^2, m_{\tilde{\ell}_3}^2, m_{\tilde{D}_3}^2, A_t, A_b, M_3\}$

**B. Natural Spectrum**

Naturalness comes into play in the context of Eq. (7). For a model to be natural all of the individual terms should be of order $M_Z^2$ and no fine-tuned cancellations should be present. In contrast to the $\mu$ parameter in the MSSM, which is a free parameter of the superpotential without any a priori relation to the electroweak scale, the $\mu_{\text{eff}}$ parameter in the NMSSM is itself induced by electroweak symmetry breaking and the vacuum expectation value of $S$. Thus, it is naturally of right order and determines the expected mass scale of the higgsinos which are mainly determined by $\mu_{\text{eff}}$, see Eq. (13) and below. In the limit of vanishing mixing, the tree level singlino mass reads $k_s = \mu_{\text{eff}}(k/\lambda) \lesssim \mu_{\text{eff}}$, where we have used that the stability of the $s \neq 0$ vacuum usually requires $k/\lambda < 1$.

We therefore expect a singlino that is lighter than the higgsinos in a natural setup.

This tree level relation is affected by loop corrections to the respective parameters. As an example, the large Yukawa coupling to the stops and their $O(h_s)$ correction from gluino loops induces a sizeable effect on $m_{\tilde{H}_u}$ in Eq. (7) while running from the SUSY breaking scale $\Lambda_s$ down to the TeV scale. In the leading log approximation [10], these corrections read

$$\Delta m_{\tilde{H}_u}^{2, l} \approx -\frac{3y_t^2}{8\pi^2} (m_{\tilde{q}_3}^2 + m_{\tilde{c}_3}^2 + |A_t|^2) \ln \left(\frac{\Lambda_s}{\text{TeV}}\right),$$

(16)

$$\Delta m_{\tilde{H}_u}^{2, g} \approx -\frac{2y_t^2}{\pi^2} \alpha_s M_3^2 \ln \left(\frac{\Lambda_s}{\text{TeV}}\right).$$

(17)
Naturalness requires these corrections to be moderately small which translates into mass bounds $m_t \lesssim m_{\tilde{b}} \lesssim O(1 \text{ TeV})$. Note that the naturalness bound on $m_{\tilde{Q}_3}^2$ also sets the scale of the $\tilde{b}_L$-like scalars, as they lie in the same SU(2)$_L$ doublet as the $\tilde{t}_L$ field. Though no equivalent naturalness constraint applies to the $\tilde{b}_R$ scalar, we assume that there is no a priori reason why the SUSY breaking mechanism should induce large splittings $m_{T_R}^2 - m_{T_L}^2$ or $A_t - A_b$ and thus we assume the soft breaking parameters to be degenerate (see Sec. III). The appearance of a second light sbottom, however, does not affect the collider results significantly.

Parameters related to the SU(2)$\times$U(1) gauginos and the squarks of the first two generations have negligible effect on the parameters in Eq. (7) and thus are not constrained by naturalness arguments. They can therefore safely be set to experimentally inaccessible scales while keeping the electroweak breaking scale small.

Note that the above consideration of naturalness is only performed on the qualitative level and solely serves as a motivation for the hierarchies and mass scales of our following collider study. More quantitative analyses in terms of so-called fine tuning are possible but require a more specific formulation of the decoupled supersymmetric sector to get a valid dependence of low-scale observables on independent high-scale parameters (see e.g. [63]).

While trying to keep the parameters natural, our model should still not violate experimental observation, i.e. a Standard Model like scalar boson with mass of order 125 GeV should emerge. In the limit where the lightest CP-odd Higgs boson decouples ($M_A \equiv 2 \mu_{\text{eff}}(A_\lambda + \kappa S)/\sin 2\beta \to \infty$), the lightest SM-like eigenvalue of Eq. (9) including leading order tree-level correction\(^1\) in $\lambda < 1$ and dominant radiative corrections of tops and stops reads

$$m_{h,\text{SM}}^2 \approx M_Z^2 \cos^2 2\beta + \frac{\lambda^2}{g^2} \sin^2 2\beta \frac{M_A^4}{1 + \tan^2 \beta} + \frac{3m_t^4}{4\pi^2 v^2} \left( \ln \left( \frac{M_{\tilde{q}_3}^2}{m_t^2} \right) + A_t^2 \left( 1 - \frac{A_t^2}{12M_{\tilde{q}_3}^2} \right) \right)$$

assuming degenerate soft breaking stop masses $M_{\tilde{q}_3}^2 = m_{\tilde{b}_3}^2 \gg m_t^2$. While the MSSM contribution shows an upper limit $M_Z^2$ in the decoupling limit and thus requires a significant contribution from heavy stops, the NMSSM contribution of $O(\lambda^2)$ can lead to a sizeable enhancement of the Higgs boson mass itself. Consequently no heavy stops are needed, making it easier to acquire a natural spectrum as explained above.

The benchmark spectra we are going to consider in the upcoming analysis are sketched in Fig. 1. We distinguish two main limits of the NMSSN, steered by the size of the dimensionless coupling parameter $\lambda$:

**large $\lambda \equiv \lambda_L$**: When the coupling $\lambda$ is large, Eq. (18) suggests that we can reach a large enough Higgs mass if $\sin 2\beta$ is large. In our analysis we choose $(\lambda = 0.7, \tan \beta = 2)$, with the value for $\lambda$ chosen at the maximum possible value which does not run into Landau poles at higher scales. In this setup, no large radiative corrections are required and as such it is expected that one can keep both stops (and the respective sbottoms) rather light while still being able to reach the correct Higgs mass. As a consequence all third generation scalars may be kinematically accessible at the LHC. The neutralinos can mix largely in this scenario and direct decays of coloured scalars into singlets and singlinos are possible [65].

**small $\lambda \equiv \lambda_S$**: In case of a very small $\lambda$, the Higgs mass is very MSSM like. To maximise the tree level value one needs a larger $\tan \beta$, which is why we define this point via $(\lambda = 0.01, \tan \beta = 15)$. Large radiative corrections are needed, which asks for at least one heavy stop. The sparticles of the MSSM sector decouple from the singlet states and experimentally, the only difference between the MSSM and the NMSSM would be sparticle decays into the singlino LSP. For very small $\lambda$, the scalar vev $s$ must be large in order to have a sufficiently large $\mu$ term. This generally translates into $s \gg v_u, v_d$. Contrarily to the previous case, this scenario will come along with a rather split sector of third generation squarks, mostly degenerate higgsinos and a mostly decoupled singlino and a singlet scalar sector.

In both scenarios, to avoid having an LSP-like chargino we always require the singlino to be lighter than the higgsinos. Furthermore, we always require the gluino to be heavier than the stops to avoid the consideration of loop-induced 2-body or off-shell 3-body decays. Whilst the relative hierarchy between stops/sbottoms and the higgsinos is not fixed by our setup, we will nevertheless find that it is often as depicted in Fig. 1.

### C. Signatures of Interest

The spectrum described in the previous section leads to interesting signatures for the LHC: due to the light $\tilde{t}$ and $b$ scalars we expect final states with many $t$ and $b$ quarks. The hadronised jets originating from $b$-quarks have a high probability of being correctly tagged as so-called $b$-jets and many analyses from both ATLAS and CMS have been designed to specifically tag final states with these objects, see Sec. [III.B]. In the following we only focus on final states with these objects and neglect other signatures:

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\(^1\) We have set $\kappa = 0$ in Eq. (18) due to the relation $\kappa < \lambda$ and the expansion in orders of $\lambda$. 
In this work, we consider third generation squark and gluino hadro-production via the strong interaction. In general, the gluon fusion diagrams will be the dominant production channel for not too heavy gluino and third generation scalar masses and the cross section is only determined by the respective mass and spin of the respective final state sparticle at leading order,

$$pp \rightarrow \tilde{g} \tilde{t}_i, \quad \tilde{b}_i \tilde{b}_i^* \quad i \in \{1, 2\}.$$  

Here, we have omitted the production of electroweakino pairs since the cross section is negligible compared to the production of coloured sparticles unless the higgsino and the singlino are the only kinematically accessible sparticles at the LHC — a scenario we are not going to assume. In addition, we have not considered compressed spectra where a hard initial state radiation jet has to be taken into account.

Let us now turn to the discussion of the decay modes. The decay chains can be very complicated in natural SUSY and typically depends on the details of the mass spectrum and the mixing angles. Since we define that the gluino is always the heaviest sparticle in our setup, the following strong two body decay modes are the dominant gluino decay channels

$$\tilde{g} \rightarrow \tilde{t}_i t_i, \quad \tilde{b}_i b_i \quad i \in \{1, 2\}.$$  

There is no tree level coupling between the squarks and the singlino in the NMSSM and thus the decays to neutralino states with a significant singlino component are suppressed. As the $\lambda_S$ scenario contains an almost pure singlino LSP, direct decays of the squarks to the LSP very rarely occur. Even in the $\lambda_L$ scenario, which can contain an LSP that is a higgsino-singlino mix, the large singlino component significantly suppresses the direct decay to this state. This situation is different to the natural MSSM where stop decays into the LSP are common and consequently we expect longer and more complicated decay chains in the natural NMSSM. A by-product of such longer decay chains is that the individual particles produced are necessarily softer and $E_T^{miss}$ can be expected to be reduced.

However in common with the natural MSSM, the particles will decay in final states with third generation SM particles which will give rise to a high $b$ jet multiplicity,

$$\tilde{t}_i \rightarrow \tilde{\chi}_{2/3}^0 t_i, \quad \tilde{\chi}_{1/1}^0 b_i,$$

$$\tilde{b}_i \rightarrow \tilde{\chi}_{2/3}^0 b_i, \quad \tilde{\chi}_{1/1}^0 t_i.$$  

In addition, the large expected squark mass splittings in the $\lambda_S$ scenario can lead to the following squark-to-squark decays with additional gauge bosons and Higgs scalars

$$\tilde{t}_2 \rightarrow \tilde{t}_1 X^0, \quad \tilde{b}_i \rightarrow \tilde{b}_1 X^0, \quad \tilde{t}_1 X^0,$$

with $X^0 \in \{Z^0, h, H, H_3, A_1, A_2\}$ and $X^\pm \in \{H^\pm, W^\pm\}$. As the production rate of the heavy squark in such a case is largely suppressed compared to $\tilde{t}_1/b_1$, this decay is however not expected to contribute significantly to the observed event rates.

The biggest difference between the natural MSSM and the natural NMSSM is that we can now have a singlino LSP. This leads to additional decays of the (now NLSP) higgsinos $\tilde{\chi}_{2/3}^0$, such as,

$$\tilde{\chi}_{2/3}^0 \rightarrow \tilde{\chi}_{1/1}^0 X^0$$

$$\tilde{\chi}_{1/1}^0 \rightarrow \tilde{\chi}_{1/1}^0 X^\pm.$$  

Generally, a light singlino is accompanied by relatively light singlet scalars. Depending on the mass difference between the NLSP and the LSP and the mass of the decay products $X$, differences between the MSSM and the NMSSM will arise, which may modify the decay patterns of the higgsino in a MSSM scenario.

Of course, for each of the above listed decays there exists a mode with all involved particles charge conjugated. Obviously the listed decay modes are only possible subject to kinematic constraints and all decay modes mentioned above can have a related three (or four) body decay mode if one (or more) of the final state particles are virtual.

We have not listed the tediously large list of possible decays for the neutral scalars $\{h, H, H_3, A_1, A_2\}$: They generally involve Standard Model like Higgs decays, decays of heavy into light scalars and decays of heavy scalars into pairs of lighter squarks or electroweakinos. However, in most cases the heavy scalars $H, H_3$ and $A_2$ do not appear in the observed decay chains and thus their decay modes are of no relevance in the following. It is mostly the Standard Model like Higgs and the singlet like scalars which are of importance and their decays are practically Standard Model like after having applied.
the experimental constraints as explained in upcoming Sec. III A.

III. MODEL TEST METHODOLOGY

As described at the end of Sec. II A, our model of interest can be described by 12 free parameters. To simplify the discussion, we assume a degeneracy of the soft parameters in the third generation, i.e.

\[ A_{q_3} = A_t = A_b, \]
\[ M_{q_3}^2 = m_{t_3}^2 = m_{U_3}^2 = m_{D_3}^2. \]  

This assumption always fixes the mass of the bottom squarks for given stop masses in a way as depicted in Fig. 1. In the following we explain how we fix the free parameters of our model

\[ \lambda, \kappa, A_\lambda, A_\kappa, \mu_{\text{eff}}, \tan \beta, A_{q_3}, M_{q_3}, M_3 \]  

with respect to the hierarchies of the models we want to consider. We follow with a discussion on how we test the respective parameter combination.

A. Scan Setup and Definitions

Each of the data points that we are going to analyse is solely defined by the following set of information:

1. The NMSSM scenario \( \lambda_S \) or \( \lambda_L \),
2. the mass \( m_{\tilde{g}} \) of the gluino,
3. the mass \( m_{\tilde{t}_1} \) of the lightest stop,
4. the higgsino mass parameter \( \mu_{\text{eff}} \) and
5. the singlino mass parameter \( m_S \equiv 2\kappa s \)

This fixes the following parameters in Eq. (21):

\[ \lambda = 0.7 \ (0.01) \quad \text{for } \lambda_L \ (\lambda_S), \]
\[ \tan \beta = 2 \ (15) \quad \text{for } \lambda_L \ (\lambda_S), \]
\[ \kappa = \sqrt{2}/m_S/\mu_{\text{eff}} \quad \text{due to } \mu_{\text{eff}} = \lambda_s. \]  

The remaining five parameters are found as follows: We require a natural, realistic particle content, that is we aim for a spectrum with as light as possible stops while having a Higgs boson at the correct mass. In addition we demand that the Higgs boson passes the most relevant theoretical and phenomenological constraints. Such a spectrum is found by using the public tool NMSSMTools [56–60]. This allows us to specify the above mentioned parameters at scale \( Q_t \) to get the corresponding physical particle masses, mixing matrices, branching ratios and test against a variety of observational tests (see below).

In order to find a parameter combination with a viable, natural spectrum, we perform the following chain of actions:

Loop over the heavy stop mass \( m_{\tilde{t}_1} \): We are interested in stops that are as light as possible, i.e. we aim to find the lightest spectrum that passes the most important phenomenological constraints. For that purpose, with \( m_{\tilde{t}_1} \) set above, we perform a loop over \( m_{\tilde{t}_1} \). Starting from \( m_{\tilde{t}_1} + 25 \) GeV and using a step-size of 5 GeV, we steadily increase the heavy stop mass and try to find a valid parameter point according to the steps described next. As soon as a valid point is found, that one is taken for the further collider study.

Fix the strong sector \( M_3, A_{q_3}, M_{q_3} \): The masses of the stops and the gluino are mostly determined by these three parameters. Given the target values \( m_{\tilde{t}_1}, m_{\tilde{g}} \) and the looped value for \( m_{\tilde{t}_1} \), we use NMSSMTools to scan over \( M_3, A_{q_3} \) and \( M_{q_3} \) and find the combination that reproduces the desired mass [70] best. For this scan, the values of \( A_\lambda \) and \( A_\kappa \) are barely of relevance as they have only a minor impact on the third generation stop masses. Consequently they are therefore fixed to the central values of the “scalar sector scan” described below. Note that at this stage we use NMSSMTools solely to find the correct mapping of physical masses to parameters. No phenomenological constraints are applied at this stage.

Explore the scalar sector \( A_\lambda, A_\kappa \): Having the strong sector fixed we start a new grid scan over the scalar trilinear parameters \( A_\lambda, A_\kappa \) in order to find a phenomenologically allowed scalar sector. We test \( A_\lambda \) uniformly in the range 0 to \( 2(\mu_{\text{eff}}/\sin 2\beta - m_{\tilde{g}}) \), which is chosen such that the central value minimises the higgsino-singlino mixing in Eq. (9) and hence maximises the SM-like Higgs boson mass [70]. \( A_\kappa \) is uniformly scanned in the range \([-550 \text{ GeV}, 450 \text{ GeV}] \).

For each point, NMSSMTools tests

- the absence of tachyonic masses and charge or colour breaking minima in the scalar potential,


| Ref. | identifier | identifier | Sensitive to which decay scenario(s) |
|------|------------|------------|---------------------------------------|
| 47   | atlas_conf_2013_0244 | stop/sbottom decay chains leading to purely hadronic final states |
| 50   | atlas_conf_2013_061 | \( gg \to t\bar{t}h^* \), \( b\bar{b}b\bar{b}^* \) and/or decays involving \( h \to b\bar{b} \). |
| 65   | atlas_conf_2013_062 | stop/sbottom decay chains with 1 isolated lepton from \( W/Z \) |
| 67   | atlas_1308_2631 | \( t \to b\chi^0_{1,2}, t\tilde{t} \) with a purely hadronic final state |
| 85   | atlas_1404_4853 | \( t \to b\chi^0_{1,2}, t\tilde{t} \) with an OS isolated lepton pair in the final state |
| 89   | atlas_1404_2500 | \( gg \) with decays into stop/sbottom producing 2 SS or 3 isolated leptons |
| 46   | atlas_1407_0583 | stop/sbottom decay chains with 1 isolated lepton from \( W/Z \) |

TABLE I. Summary of the expected most sensitive analyses within CheckMATE to the considered natural model, listed in alphabetical order. All analyses require a significant amount of missing transverse momentum in the final state and have at least one signal region which requires b-tagged jets. All other ATLAS analyses implemented in CheckMATE are tested in parallel, but are always found to be less sensitive than those listed.

- that there is a SM-like Higgs boson in the mass window 121 to 129 GeV;\(^5\)
- consistency with all other implemented collider constraints (mostly LEP limits on the Higgs sector, neutralinos and charginos)
- consistency with all other implemented low energy observables. (e.g. \( b \to s\gamma, B_s \to \mu^+\mu^- \), ...) apart from \( (g-2)\mu \) where our natural model will reproduce the SM expectation.

To consider more recent collider results from LEP, Tevatron and the LHC that constrain the scalar sector, we further use HiggsBounds 4.1.2\(^6\) and HiggsSignals 1.2.0\(^6\) to perform final tests on the scalar sector of the considered parameter points. For that purpose we fix the mass uncertainty for all Higgs bosons to be 4 GeV. HiggsBounds is used with the Landau setup. A parameter combination is discarded if HiggsBounds returns “excluded”. In HiggsSignals, the both setting is used that performs both a mass centred and a peak centred method using latestresults. A point is discarded if it produces a p-value smaller than 0.05.

Exit \( m_\ell_2 \) scan: If at the end of this stage no allowed \( A_\lambda, A_\kappa \) combination is left, the \( m_\ell_2 \) loop starts with the next iteration. If however a parameter combination of \( M_3, A_{\tilde{q}_3}, M_{\tilde{q}_3}, A_\lambda \) and \( A_\kappa \) passes all the aforementioned constraints, this parameter point is used for collider phenomenology part described next.

For completeness it should be noted that the 5 parameters mentioned at the beginning of this section are closely related to the physical electroweakino masses. Firstly due to the decoupled wino, the mass of the chargino, \( m_{\chi^+_1} \), is practically identical to the input parameter \( \mu_{\text{eff}} \) and we will therefore use both variables synonymously in the following. As depicted in Fig. 1 \( \mu_{\text{eff}} \) (or \( m_{\chi^+_1} \)) is also very close to the mass of the two neutral higgsinos within \( \lambda_S \). Likewise, the singlino mass parameter \( m_{\tilde{S}} \) sets the mass of the lightest neutralino, \( m_{\chi^0_1} \). Within \( \lambda_S \) however, large mixing in the neutralino sector will lead to deviations from these identities. In the following, instead of the input variable \( m_{\tilde{S}} \) we will only show the physical mass of the lightest neutralino, \( m_{\chi^0_1} \), which by construction is predominantly singlino like.

B. Collider Phenomenology

As explained in Sec. II C we assume that pair production of the light \( \tilde{g}, \tilde{t}_1 \) and \( \tilde{b}_1 \) dominates the expected signal. Production cross sections for these particles are calculated using NLLFast 2.1\(^7\) using CTEQ6.6L0 PDF \(^7\). Uncertainties due to scale variations, parton density functions and \( \alpha_s \) are provided and we take the quadratic sum of these to set the total theory error \( \Delta \sigma \). For each production mode, 50 000 signal events are generated using Herwig++ 2.7.0 \(^8\) \(^9\) with the NMSSM model setting. For practical reasons, decay tables of all relevant particles are calculated within Herwig++, which contains all tree level 2– and 3–body decays and effective implementations of the loop-induced decays \( h_i \to \gamma\gamma, gg \).

To test the model against a variety of LHC results, we use CheckMATE \(^5\) \(^5\) \(^5\). This tool applies an ATLAS tuned version of the Delphes 3 \(^8\) detector simulation which uses FastJet with the anti-\( k_T \) jet algorithm \(^8\) \(^5\). Reconstructed events are tested against various ATLAS analyses and the derived number of signal events is tested against observation and the Standard Model expectation. The compatibility of signal and observation is tested by comparing the predicted signal \( S \pm \Delta S \) to the model independent 95% CL limit S95, determined by using the CLS method \(^8\). Here, \( \Delta S \) considers both the MC error on our statistics as well as the theory error on the total cross sections. CheckMATE considers a

\(^5\) The window for \( m_h \) is motivated by theory uncertainties and the fact that the decoupled sector, most importantly the electroweakinos, can influence the Higgs mass by higher order corrections if they are of order \( O(\text{few TeV}) \), see e.g. \(^1\). The exact details of the heavy electroweakino sector would not affect our collider analysis at all and thus are incorporated by a looser constraining on the light Higgs boson mass.
large list of ATLAS analyses, however due to the signatures described in Sec. II C, it is expected that only a subset of these will be sensitive to the characteristics of our model. We list these analyses in Tab. I. They all require a significant amount of missing transverse momentum due to the expected undetected LSP in the final state and have signal regions that check for b–jets. They mainly differ by the final state jet multiplicities and the total amount and relative charge of final state isolated leptons (i.e. electrons and muons). The analyses also differ in the kinematics of the respective signal regions that are designed and tuned for particular final states. As we expect different final state signatures in our model, it is highly favourable to check all these possibilities in parallel and filter out the most sensitive one for each case. Fortunately, CheckMATE allows for an easy comparison of that kind.

IV. RESULTS

In the following we show exclusion lines in the parameter space of the model explained above. Since we still have $m_{\tilde{g}}$, $m_{\tilde{t}_1}$, $m_{\tilde{\chi}_1^\pm}$ and $m_{\tilde{\chi}_1^0}$ as continuous degrees of freedom, we choose to present results for specific chosen benchmark scenarios.

As one of our considered decay chains in Sec. II C starts with the production of gluinos and ends with the decay into the singlino LSP, we first choose to show exclusion lines in the plane spanned by the masses of these two particles. We do so for various choices of $m_{\tilde{t}_1}$, $m_{\tilde{\chi}_1^\pm}$ and always compare the results for $\lambda_L$ and $\lambda_S$. As it will turnout, light gluinos mostly lead to severely constrained models. Thus we will follow with a scenario in which the gluino is decoupled from the spectrum as well. We then show exclusion lines in the $m_{\tilde{t}_1}$–$m_{\tilde{\chi}_1^0}$ plane for different chargino masses, again putting the results for $\lambda_L$ and $\lambda_S$ side by side. For the specific case of a light LSP, we also present results in the $m_{\tilde{t}_1}$–$m_{\tilde{\chi}_1^0}$ plane to illustrate the dependence on the chargino mass for both $\lambda$ scenarios.

To keep the discussion compact, we only show 95% exclusion lines in different parameter planes within this section. An exhaustive list of plots showing distributions of masses, cross sections and branching ratios can be found in the appendix.

A. Gluino–LSP–Plane

In Figs. 2–6 we show the 95% exclusion region in the gluino-LSP mass plane, using fixed stop masses in the range $m_{\tilde{t}_1} = 400$ to 800 GeV. For each case, the $\lambda_L$ and $\lambda_S$ scenarios are compared in the left and right panel, respectively. Within each panel we compare the exclusion regions for different chargino mass values that obey $m_{\tilde{\chi}_1^+} < m_{\tilde{t}_1}$. Since the chargino must not be lighter than the LSP, each exclusion line has an individual upper limit on the $m_{\tilde{g}}$ axis, drawn by dashed horizontal lines. Chargino mass values that are listed in the legend but do not appear in the plot should be interpreted as being entirely excluded across the whole mass plane.

Generally, the exclusion lines split the parameter space into two regions of interest and we discuss these regions separately:

1. Light Gluinos

For $m_{\tilde{g}} \lesssim 1100$ GeV, Fig. 2–6 show that the limits are mostly independent of the chargino mass and apparently primarily driven by the gluino decay products in the decays $\tilde{g} \to b\bar{b}, t\bar{t}$.

As the bounds in that region do not seem to vary significantly as we change the mass of the electroweakinos (and only barely if we change the mass of the lightest stop), we conclude that the details of the decay chain of the third generation scalars into the LSP is almost irrelevant when setting limits on the model. The only exception is if very small mass splittings occur in the decay chain, for example between the gluino and the stop or the stop and the higgsinos. We can see the effect in the left parts of Fig. 6 and also can be observed for all scenarios with $m_{\tilde{\chi}_1^0} < m_{\tilde{t}_1}$ in Figs. 4–6.

When we compare the $\lambda_S$ and $\lambda_L$ limits we also see that the limits are stable between the two scenarios once gluino production is dominant. Consequently, we again conclude that the precise decay modes of the $t(b)$ and the various $\tilde{\chi}_1^\pm, \tilde{\chi}_1^0$ do not effect the LHC phenomenology in this region of parameter space.

The above conclusions may be different to the thoughts we had before commencing this study. In fact we may have guessed that the additional decay step present due to the singlino would have made the model more difficult to see at the LHC. The reason is that the extra decay can reduce the individual final state particle energies and also the total missing energy (e.g. [87]). We believe the reason that this does not occur here is the number of studies and therefore signal regions contained within the CheckMATE program. For example, in Fig. 7, we can compare the respective most constraining signal regions in the gluino dominated region for a specific benchmark scenario. We see that the signal regions used to constrain the models are different between the two scenarios. In particular the $\lambda_S$ scenario which generically contains longer decays is better constrained by signal regions that have a larger final state particle multiplicity. For instance, in the gluino...
FIG. 2. Observed 95% C.L. exclusion limits for $m_{\tilde{t}_1} = 400$ GeV. Left: $\lambda_L$. Right: $\lambda_S$

FIG. 3. Observed 95% C.L. exclusion limits for $m_{\tilde{t}_1} = 500$ GeV. Left: $\lambda_L$. Right: $\lambda_S$

FIG. 4. Observed 95% C.L. exclusion limits for $m_{\tilde{t}_1} = 600$ GeV. Left: $\lambda_L$. Right: $\lambda_S$

FIG. 5. Observed 95% C.L. exclusion limits for $m_{\tilde{t}_1} = 700$ GeV. Left: $\lambda_L$. Right: $\lambda_S$

FIG. 6. Observed 95% C.L. exclusion limits for $m_{\tilde{t}_1} = 800$ GeV. Left: $\lambda_L$. Right: $\lambda_S$
dominated region, the ATLAS search with at least 3 b-jets \[\text{atlas}\text{conf}\text{2013}\text{061}\] is the most powerful but whilst the $\lambda_L$ scenario is best constrained with the 4-jet signal region, the 6-jet + 1-lepton region dominates for $\lambda_S$. In addition, the multi b-jet ATLAS search demands moderate missing transverse momentum and hence the reduction of the total missing energy in the NMSSM does not significantly change the efficiency in the signal regions. The demands of the signal region therefore translates into the necessity of a sufficiently large gluino production cross section and a sizeable mass splitting of gluinos and squarks as well as squarks and electroweakinos. It is thus expected that limits should not depend significantly on the $\lambda$ scenario and only on the masses if they are close to threshold, as can be seen in our results.

2. Heavy Gluinos

For gluinos with mass above the production threshold of about 1.2 TeV, the exclusion sensitivity will be dominantly driven by the production of the third generation sparticles $\tilde{t}_1/2, \tilde{b}_1/2$ if they are sufficiently light. To illustrate this, we show the total production cross section for gluinos in Fig. 14 of the appendix and third generation squark production for fixed $m_{\tilde{t}_1}$ in Fig. 15. Comparing the cross section values in regions with $m_{\tilde{g}} > 1.2$ TeV, $m_{\tilde{t}_1} < 800$ GeV, one expects far more $\tilde{t}_1$ than gluinos to be produced. Depending on the $\lambda$ scenario large numbers of events with sbottoms and heavier stops are expected in addition. Therefore, beyond the gluino threshold we observe a gluino-independent upper limit on the mass of the lightest neutralino.

However, contrarily to the gluino-dominated region, one now finds significant dependencies of the limits on the chargino mass parameter and the $\lambda$ scenario in Fig. 2-6. In general, we observe that for a fixed mass of the lightest stop, limits on the LSP mass become weaker the lighter we chose the intermediate chargino. Also, throughout all cases we find consistently better limits in the $\lambda_L$ scenario than for $\lambda_S$.

To understand these differences, we first have to shed light on the analyses and signal regions which define our exclusion limits in this part of parameter space. In Fig. 7 we take the specific example of a light stop mass of 500 GeV and show the most sensitive signal regions for chargino masses of 400 and 300 GeV, comparing $\lambda_L$ on the left to $\lambda_S$ on the right. One finds two main classes of final states to be of importance here:

1. Signal regions from $\text{atlas}\text{conf}\text{2013}\text{024}$ and $\text{atlas}\text{1407}\text{0583}$ focus on final states that originate from direct $\tilde{t}_1 \rightarrow t \tilde{\chi}_0^0$ decays. That is, they require missing transverse momentum, b-jets and final state objects whose invariant mass lie close to the top-quark mass.

2. ‘BC-type’ regions in $\text{atlas}\text{1407}\text{0583}$ have been designed to tag events of type $t \rightarrow b \tilde{\chi}_\pm^0, \tilde{\chi}_\pm^0 \rightarrow W^\pm \tilde{\chi}_1^0$ by using kinematic variables that are sensitive to intermediate decay steps.
In the following, we will refer to these as ‘tN-like’ and ‘bC-like’ analyses and signal regions, respectively.

In our model setup, the choice of the Higgs mass parameter $\mu$ (which sets the $m_{\tilde{t}_1}^2$ and $m_{\tilde{t}_2}^2$) is crucial to determine how many events are expected to be counted for the above most sensitive signal regions. Its value sets the kinematically open channels from the full list in Sec. IIIC fixes the branching ratios and determines the energy distribution among the final state particles.

For $m_{\tilde{\chi}^+} > m_{\tilde{\chi}^0}$, the branching ratio for $\tilde{t}_1 \to t\tilde{\chi}^0$ is almost 100% — regardless of $\lambda$ — and thus the upper LSP mass limits in both scenarios are determined by results from tN-like signal regions. If the $\tilde{t}_1$ was the only squark kinematically available, the limits of $\lambda_L$ and $\lambda_S$ would be expected to coincide. Comparison of the corresponding $m_{\tilde{t}_1} = m_{\tilde{\chi}^+}$ lines in Figs. 2 [2] however shows that $\lambda_L$ yields stronger limits, with the difference being larger for lighter $m_{\tilde{t}_1}$. The reason here is that $\lambda_L$ can allow for additional lighter 3rd generation squarks while still being able to get the right SM Higgs mass, as in Eq. [18]. These lighter squarks have a larger production cross section and thus contribute more to the observable events, e.g. via decays $b_1 \to t\tilde{\chi}^+_1$ which can also pass the signal region cuts. If a light $\tilde{t}_1$ is present in a $\lambda_S$ scenario however, the additional 3rd generation squarks are required to be much heavier.

For lighter chargino masses, the decay $\tilde{t}_1 \to b\tilde{\chi}_1^\pm$ opens kinematically. Within the $\lambda_S$ scenario we have an almost purely singlino LSP which causes the branching ratio for $t\tilde{\chi}_1^0$ final states to become almost immediately disfavoured below the chargino threshold. Thus in this scenario almost all stops have to decay via intermediate electroweakinos. Interestingly, tN-like analyses are still most significant to set the limit if the charginos are not too light (see Fig. 7, top right). The reason is that events with intermediate charginos can lead to $bW^\pm\tilde{\chi}_1^0$ final states misidentified as top quarks within tN-like signal region selections if the neutralino is light enough (the top mass window is very large in this analysis, as wide as 130 < $m_t$ < 250 GeV). In addition one expects a significant contribution of sbottoms decaying into $t\tilde{\chi}_1^\pm$ final states which also look tN-like.

For even lighter charginos, the limit is however only set by bC-like analyses (see Fig. 7). Decreasing the chargino mass further leads to softer decay products in the decay $\tilde{\chi}_1^\pm \to \tilde{\chi}_2^0X^\pm$, which weakens the resulting upper limits on the $\chi_1^0$ mass. Finally, decays into $t\tilde{\chi}_2/3$ can reduce the branching ratio into the above mentioned decays once the chargino becomes light enough (see Fig. 16).

It should also be mentioned that the branching ratios of the stop into neutral and charged higgsinos are fixed by the stop mixing matrix and $\tan \beta$ [88, 89]. This results in a significant number of events displaying an ‘asymmetric’ topology in which each of the initially produced sparticles decays differently. However, the signal regions within the analyses that we use are mainly designed for symmetric decay scenarios, which leads to a reduction of the overall sensitivity.

Most of the explanations in the above discussion apply similarly to the $\lambda_L$ scenario. However, a distinctive feature is the strong mixing in the neutralino sector which allows for the LSP to have a large higgsino component and thus $t\tilde{\chi}_1^0$ decays still having a large branching fraction below the chargino threshold. For example one finds that for $m_{\tilde{t}_1} - m_{\tilde{\chi}_1^0} < 150$ GeV direct stop-to-top decays still happen with more than 20% probability (see Fig. 16). We therefore expect, and observe, that also within $\lambda_L$ the tN signal regions set the limit for charginos within that mass region (see Fig. 7, top left).

For lighter charginos, the limits become weaker due to the decreasing branching ratio of the ‘golden channel’ $\tilde{t}_1 \to t\tilde{\chi}_1^0$ and eventually the bC signal regions dominate and sets the limits thereafter (see Fig. 7, bottom left). The overall stronger exclusions within the $\lambda_L$ scenario can therefore be attributed to two different reasons. Firstly, the other 3rd generation squarks will again be lighter in the $\lambda_L$ scenario due to the additional singlet contributions to the Higgs mass. Secondly, the increased branching ratio of $\tilde{t}_1 \to t\tilde{\chi}_1^0$ which the LHC analyses are particularly sensitive to also helps.

Interestingly, in both $\lambda$ scenarios, $\mu_{\text{eff}}$ lighter than $m_{\tilde{t}_1} - m_{\tilde{\chi}_1^0}$ opens decay channels of the type $\tilde{t}_1 \to t\tilde{\chi}_2^0/3$. These could lead to NMSSM specific final states as discussed in Sec. IIIC. However, we do not observe any improvement on the LSP limits in these cases. Quite the contrary, the reduction of the branching ratio into $b\tilde{\chi}_1^\pm$ final states resulting from the new decay channel and asymmetric final states mentioned above weakens the limits even more as can be observed when comparing the limits in Figs. 2 [3] above or below this threshold. We investigate the impact of this more closely in the upcoming section.

### B. Stop-Electroweakino-Plane

As shown in the last set of results, below the gluino production threshold, the LHC limits only have a small dependence on the details of the natural spectrum. However, as we decouple the gluino, the masses and couplings of the electroweakino sector become more important. For that reason we also show results in the $m_{\tilde{b}_2} - m_{\tilde{\chi}_1^0}$-plane for a decoupled gluino of mass 2 TeV in Figs. 8 [11]. With one degree of freedom less, we are now able to show one exclusion limit per plot for specific values of $m_{\tilde{\chi}_1^0}$, again comparing $\lambda_L$ (left) to $\lambda_S$ (right). The parameter space that we investigate does not include the region where $m_{\tilde{t}_1}$ becomes close to $m_{\tilde{\chi}_1^0}$. This is shown by the diagonal dashed line within each plot which shows the kinematic range for which $m_{\tilde{t}_1} < m_t + m_{\tilde{W}} + m_{\tilde{\chi}_1^0}$ and only 4-body final states or flavour changing neutral current decays such as $\tilde{t}_1 \to c\tilde{\chi}_1^0$ are possible. Given the small mass difference, initial state radiation searches provide the most constraining limits in this region [90, 91]. These searches are relatively insensitive to the details of the decay chain...
FIG. 8. Observed 95% C.L. exclusion limit and most sensitive analysis per point for $m_{\tilde{\chi}^\pm_1} = 250$ GeV. Left: $\lambda_L$. Right: $\lambda_S$

FIG. 9. Observed 95% C.L. exclusion limit and most sensitive analysis per point for $m_{\tilde{\chi}^\pm_1} = 350$ GeV. Left: $\lambda_L$. Right: $\lambda_S$

FIG. 10. Observed 95% C.L. exclusion limit and most sensitive analysis per point for $m_{\tilde{\chi}^\pm_1} = 500$ GeV. Left: $\lambda_L$. Right: $\lambda_S$

FIG. 11. Observed 95% C.L. exclusion limit and most sensitive analysis per point for $m_{\tilde{\chi}^\pm_1} = 750$ GeV. Left: $\lambda_L$. Right: $\lambda_S$
in question and thus we expect the results to be very similar to those of the MSSM.

Similarly to the gluino-LSP scan, the upper limit on the LSP mass is set by requiring $m_{\tilde{t}_1} < m_{\tilde{t}_2}$. For the $\lambda_L$ case, mixing in the neutralino sector leads to a maximum achievable value of $m_{\tilde{\chi}^0_1}$ which lies somewhat below $m_{\tilde{t}_2}$. In the $\lambda_S$ scenario, realistic parameter points are not possible with $m_{\tilde{t}_1} \lesssim 400$ GeV if $m_{\tilde{t}_2} \approx m_{\tilde{b}_1}$ is small. The reason is that for small $m_{\tilde{b}_1}$, the radiative corrections to the Higgs boson mass are not large enough to correctly reproduce the LHC measurement (see e.g. Eq. (18)).

In all plots we again show, for each individual considered data point, the most sensitive analysis that has been used to calculate the confidence level of that particular point. However, we do not show the numerous individual signal regions (as we did for Fig. 2) to keep the amount of different values to a reasonable level.

We again observe that the choice of analysis responsible for the limit setting is strongly correlated with the branching ratio of the lightest stop and from Fig. 16 we expect four main regions of interest. These are respectively, direct decays of the stop into the LSP and an $a)$ on- or $b)$ off-shell top, $c)$ intermediate decays via charginos or $d)$ via neutral higgsinos. The thresholds for these regions often coincide with similar threshold for sbottom decays, as can be seen in Fig. 18. As an example the $b_1 \to t_{\chi^0_1}$ and the $t_1 \to t_{\chi^0_{2/3}}$ lie very close in the $\lambda_L$ scenario.

Using the branching ratio information, we can closely follow the explanations from the last section to understand the limits in Figs. 8-11. For stops lighter than the given chargino, only direct decays $t_1 \to t^{(*)}_{\chi^0_1}$ are kinematically allowed. $t\chi$-like analyses are therefore the most sensitive and lead to similar limits for $\lambda_L$ and $\lambda_S$, with the former being slightly stronger than the latter due to the lighter sbottoms in this model. In $\lambda_S$, a strip for $m_{\tilde{t}_1} - m_{\tilde{\chi}^0_1} < m_{\tilde{t}_1}$ cannot be excluded as the final state with an off-shell top is not observed by $t\chi$-like analyses and hard to distinguish from the SM background. Within $\lambda_L$, this region can still be explored since it is possible that the spectrum also contains a light $b_1$. This can be excluded via $\tilde{b} \to t_{\chi^0_1}$ specific selections in atlas.1404.2500 (see e.g. Figs. 10, 11).

For kinematically allowed chargino decays, a transition from ‘$t\chi$’ into ‘$b\chi$’ signal regions can be observed for increasing $m_{\tilde{t}_1}$, that is for larger stop-chargino splitting. As in the previous setup, $\lambda_L$ profits from the Higgsino fraction of the LSP and the generally lighter 3rd generation squarks. The highest sensitivities are reached via ‘$t\chi$’ final states in atlas.conf.2013.024. The highest sensitivity to the LSP mass can be reached when these final states set the limit, which can reach up to $(m_{\tilde{\chi}^0_{2/3}} \approx 325$ GeV. In $\lambda_S$, $b\chi$ signal regions dominate the limit earlier, which require lighter neutralinos to observe the intermediate chargino decay step. The experimental reach to the LSP mass is therefore smaller in these scenarios and of order 250 GeV.

As we further increase the stop masses, a maximum value of $m_{\tilde{t}_1}$ is reached. This stop sensitivity limit seems to depend on the chosen chargino mass and the considered $\lambda$ scenario and is rather independent of the LSP mass as long as it is light enough, that is for $m_{\tilde{\chi}^0_1} \lesssim 150$ GeV.

To better understand the parameter dependence, we chose to show results in the $m_{\tilde{t}_1}$-$m_{\tilde{\chi}^0_{2/3}}$-plane for a fixed, light LSP mass of 100 GeV in Fig. 12. We show the previously discussed thresholds for $t_1 \to b_{\chi^\pm}$ and $\tilde{t}_1 \to t_{\chi^0_{2/3}}$ and it can be seen that they can have an important impact on the sensitivity of the experimental analyses to the stop mass. Within $\lambda_L$, the upper limit on $m_{\tilde{t}_1}$ is almost constant at $\approx 700$ GeV for charginos above the $t_{\chi^0_{2/3}}$ threshold. This corresponds to similar limits from simplified $\tilde{t} \to t\tilde{L}$ topologies as in [40, 41]. The limit gets slightly weaker if the chargino threshold is passed, dropping by at most 50 GeV as soon as $b\chi$ signal regions dominate the limit. In Figs. 17, 19 and 21 we show the branching ratio distributions in the same plane and the same LSP mass as the results in Fig. 12. One observes that the mass values in our spectrum are such that the above behaviour coincides with the threshold for $b \to t_{\chi^0_{2/3}}$, which also explains why the $b\chi$-like signal regions become important within this region of parameter space.

As long as the higgsinos do not appear in the squark decay chains, $\lambda_S$ returns similar limits as the $\lambda_L$ scenario, for the same reasons discussed in the previous section. However, within this model one observes a sizeable weakening of the limits as soon as the intermediate chargino and NLSP higgsino decays open kinematically. Interestingly, the latter has a particularly negative impact on the result, as the experimental analyses are only weakly sensitive to parameter regions in $\lambda_S$ where $t_1 \to t_{\chi^0_{2/3}}$ is kinematically allowed. As discussed in Sec. 1C it is this decay chain which yields NMSSM-specific features in the final state topology: the decay of the higgsino NLSPs into the singlino LSP should create a sizeable excess of $h/H/A_1 \to bb$ final states. It seems, however, that none of the many distinct final states within the numerous analyses that CheckMATE contains is sufficiently sensitive to this topology. Thus, the existing $b\chi$-like limits are weakened due to reduced branching ratios after passing the NLSP higgsino threshold.

We therefore conclude that not only can many limits on natural NMSSM scenarios be derived from very similar topologies in natural MSSM studies, but we also find that regions of parameter space which produce NMSSM-exclusive final state features are not sufficiently covered by existing studies. Therefore, only weak limits on the NMSSM can be set within this region of parameter space which suffer under branching-ratio penalties.

V. CONCLUSION

In this study we explore the natural NMSSM to determine how the additional singlino can effect the LHC searches compared to the more studied MSSM case. To
Most Sensitive Analysis + $r_{\text{obs}}$ limit

FIG. 12. Observed 95% C.L. exclusion limit and most sensitive analysis per point for $m_{\tilde{t}^0} = 100$ GeV. Diagonal dashed (dashed-dotted) lines show shows the threshold for $t_1 \rightarrow b\tilde{\chi}^\pm_1$ ($t_1 \rightarrow b\tilde{\chi}^0_{2/3}$). Left: $\lambda_L$. Right: $\lambda_S$.

We find that, when constructing a realistic phenomenological model, the NMSSM-specific decay chains via intermediate heavy neutralinos often create an MSSM-like topology, $q_3 \rightarrow q_3\tilde{\chi}^0_1$ which can be preceded by $\tilde{g} \rightarrow q_3q_3$ if the gluino is light. If the branching ratio to these decay chains are large, the limits very closely follow those of the NMSSM-coupling $\lambda$. If it is large, all neutralinos have a sizeable higgsino fraction and direct decays into the lightest neutralino are significant. However, in case of small $\lambda$, the coupling of the squarks to the LSP is made small since it has a large singlino content. Therefore decays via intermediate charged and neutral higgsinos are preferred if kinematically allowed which lengthens the decay chains seen. In addition, since different decay modes may be competing with similar branching ratios, ‘asymmetric’ decay chains can often occur.

These longer decay chains can lead to weaker LHC bounds for two particular reasons. First of all, the ATLAS searches have more focussed on the MSSM specific signatures and consequently not been designed with these final states in mind. Secondly, the longer decay chains lead to a higher final state particle multiplicity but with each individual particle carrying smaller $p_T$. In addition the same effect reduces the final state $E_T^{\text{miss}}$ as observed in other studies with more complicated decay topologies e.g. [87]. On the other hand, additional final states, namely jets and leptons, can improve the sensitivity even though the invisible transverse momentum is reduced. Therefore an important conclusion of this study is that it is not obvious if the efficiency is smaller or larger in a particular NMSSM scenario simply by looking at the spectrum and decays. Instead it is crucial to test the model against a large number of searches covering various final state topologies.

Within this study we do test a large variety of different analyses but still only use one signal region to define the overall limit. In the models with extended and asymmetric decay chains (where we observe a weakening of the LHC limit), we expect the signal to populate a more varied number of signal regions than if the model predicted a single dominating decay chain. Therefore it may be expected that a combination of the sensitivities across all analyses can significantly enhance the limits but this is beyond the scope of this study.

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Appendix A: Mass Distributions

FIG. 13. Mass of the heavier stop and the lighter sbottom (which is very degenerate with the heavier sbottom) for a decoupled gluino and $m_{\tilde{\chi}^\pm} = 500$ GeV. Left: $\lambda_L$. Right: $\lambda_S$
Appendix B: Cross Section Distributions

FIG. 14. Total production cross section for gluinos, using $m_{\tilde{t}} = 500$ GeV, $m_{\tilde{\chi}_{\pm}^1} = 400$ GeV. Left: $\lambda_L$. Right: $\lambda_S$.

FIG. 15. Total production cross section for the third generation squarks for a decoupled gluino and $m_{\tilde{t}} = 500$ GeV. Left: $\lambda_L$. Right: $\lambda_S$.
FIG. 16. Most significant branching ratios of the lightest stop into the the singlino LSP, the higgsino NLSPs and the chargino for a decoupled gluino and $m_{\tilde{g}} = 500$ GeV. Left: $\lambda_L$. Right: $\lambda_S$. 
FIG. 17. Most significant branching ratios of the lightest stop into the singlino LSP, the higgsino NLSPs and the chargino for a decoupled gluino and $m_{\tilde{\chi}_1^0} = 100$ GeV. Left: $\lambda_L$. Right: $\lambda_S$. 
FIG. 18. Most significant branching ratios of the (mostly degenerate) sbottoms into the lightest stop, the higgsino NLSPs and the chargino for a decoupled gluino and $m_{\tilde{\chi}_1^\pm} = 500$ GeV. Left: $\lambda_L$. Right: $\lambda_S$. 

Appendix D: $b_{1/2}$ Branching Ratio Distributions
FIG. 19. Most significant branching ratios of the (mostly degenerate) sbottoms into the the lightest stop, higgsino NLSPs and the chargino for a decoupled gluino and $m_{\tilde{\chi}_1^0} = 100$ GeV. Left: $\lambda_L$. Right: $\lambda_S$. 

Appendix E: $\tilde{t}_2$ Branching Ratio Distributions

FIG. 20. Most significant branching ratios of the heavier stop into the the singlino LSP, the higgsino NLSPs and the chargino for a decoupled gluino and $m_{\tilde{g}} = 500$ GeV. Left: $\lambda_L$. Right: $\lambda_S$
FIG. 21. Most significant branching ratios of the heavier stop into the the singlino LSP, the higgsino NLSPs and the chargino for a decoupled gluino and $m_{\tilde{\chi}^0_1} = 100$ GeV. Left: $\lambda_L$. Right: $\lambda_S$. 
FIG. 22. Most significant branching ratios of the higgsino-like neutralinos for a decoupled gluino and $m_{\tilde{t}_1} = 500$ GeV. Left: $\lambda_L$. Right: $\lambda_S$. 

Appendix F: $\tilde{\chi}^0_{2/3}$ Branching Ratio Distributions
FIG. 23. Most significant branching ratios of the higgsino-like neutralinos for a decoupled gluino and $m_{\tilde{\chi}_1^0} = 100$ GeV. Left: $\lambda_L$. Right: $\lambda_S$. 
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