

Higgs Inflation at the LC

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Abstract

Most cosmological models of inflation are far away from providing a smoking gun at low energies. A model of Higgs inflation in the Next-to-Minimal Supersymmetric Standard Model, however, changes the NMSSM phenomenology drastically and may be well distinguished from the pure NMSSM or MSSM at a future Linear Collider. We point out certain differences of the inflationary model to the ordinary NMSSM and discuss the Higgs and neutralino/chargino sector in particular to identify the smoking gun of inflation at electroweak energies.

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1 Introduction

Cosmological models that can be tested in the laboratory are typically very rare. A smoking gun at low energies of a model acting at a very large scale like the Planck scale requires a precision machine like a Linear Collider. Precise electroweak observations may then distinguish between an ordinary extension of the Standard Model (SM) or an extension which simultaneously grasps cosmological problems. Early universe inflation is to be seen as a cosmological fact which has to be addressed. If it is addressed in a way that interrelates the Planck-scale physics with Fermi-scale physics, such a model will most probably modify the Higgs sector of the SM. On an economic basis, employing the SM Higgs field as the inflaton field of cosmology, such a model can be called minimal. While Higgs inflation in the SM tends to become “unnatural” towards the high scales, see Ref. [1], a more viable candidate is the scale-free Supersymmetric Standard Model. The scale-free model requires compared to the Minimal Supersymmetric Standard Model (MSSM) an additional singlet superfield and thus is known as the Next-to-Minimal Supersymmetric Standard Model (NMSSM).

The Higgs sector of the NMSSM is characterised by the \mathbb{Z}_3 -invariant superpotential

$$\mathcal{W}_{\text{Higgs}} = \lambda \hat{S} \hat{H}_u \cdot \hat{H}_d + \frac{\kappa}{3} \hat{S}^3, \quad (1)$$

where \hat{H}_u and \hat{H}_d are the $SU(2)_L$ doublet Higgs superfields and \hat{S} the singlet superfield. The scalar components of the doublet fields decompose as

$$H_u = \begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix}, \quad H_d = \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix}, \quad (2)$$

such that $H_u \cdot H_d = H_u^+ H_d^- - H_u^0 H_d^0$. This superpotential has an accidental \mathbb{Z}_3 -invariance under the transformation

$$\hat{\Phi} \rightarrow e^{i\frac{2\pi}{3}k} \hat{\Phi}, \quad k \in \mathbb{Z}, \quad \text{for } \hat{\Phi} = \hat{H}_u, \hat{H}_d, \hat{S},$$

such that only trilinear terms are allowed. Especially the μ -term of the MSSM, $\mu H_u \cdot H_d$, is forbidden if the \mathbb{Z}_3 symmetry is imposed. After electroweak symmetry breaking, however, the singlet scalar acquires a vacuum expectation value (v_{eff}) and dynamically induces a μ -term via the λ coupling, which plays the role of an effective higgsino mass term:

$$\mu_{\text{eff}} = \lambda \langle S \rangle. \quad (3)$$

Non-minimal coupling in Canonical Superconformal Supergravity The implementation of Higgs inflation in superconformal theories follows the non-minimal coupling of the Higgs field content to supergravity, as suggested by Ref. [1], and comes with a single dimensionless and holomorphic coupling $X(\hat{\Phi})$:

$$\mathcal{L} = -6 \int d^2 \theta \mathcal{E} \left[R + X(\hat{\Phi})R - \frac{1}{4} (\bar{\mathcal{D}}^2 - 8R) \hat{\Phi}^\dagger \hat{\Phi} + \mathcal{W}(\hat{\Phi}) \right] + \text{h. c.} + \dots, \quad (4)$$

where \mathcal{E} is the vierbein multiplet, R the curvature multiplet and $\bar{\mathcal{D}}$ a covariant derivative. The chiral superfields $\hat{\Phi}$ shall be any of the fields H_u, H_d or S . A realization of such a non-minimal coupling involving the doublet Higgs fields only can be found to be

$$X = \chi \hat{H}_u \cdot \hat{H}_d, \quad (5)$$

with a numerical factor χ . Note that this term breaks the \mathbb{Z}_3 symmetry of the NMSSM and the superconformal symmetry.

The addition of the superconformal symmetry breaking term changes the frame function in Jordan frame supergravity and affects the Kähler potential in such a way that the NMSSM superpotential gets modified [2,3]. In Planck units ($M_p = 1$), the frame function $\Omega = \hat{\Phi}_i^* \hat{\Phi}_i - 3$ gets extended by the χ -term to

$$\Omega_\chi = \Omega - \frac{3}{2} (X(\hat{\Phi}) + \text{h. c.}), \quad (6)$$

and similarly the Kähler potential

$$\mathcal{K} = -3 \log(-\Omega/3) \rightarrow \mathcal{K}_\chi = -3 \log(-\Omega_\chi/3). \quad (7)$$

In the canonical superconformal supergravity (CSS) model, the frame function is explicitly given by [3]

$$\Omega_{\text{CSS}} = -3 + |\hat{S}|^2 + |\hat{H}_u|^2 + |\hat{H}_d|^2 + \frac{3}{2} \chi (\hat{H}_u \cdot \hat{H}_d + \text{h. c.}). \quad (8)$$

In order to have successful inflation in the NMSSM, however, a stabilisator term $\zeta(\hat{S}\hat{S})^2$ has to be added [3], which disappears from the low-energy phenomenology (Planck-suppressed).

The χ -term breaks a continuous \mathcal{R} symmetry and its discrete \mathbb{Z}_3 subgroup at dimension six $\sim \chi \frac{\lambda^2 h^6}{M_p^2}$. Much below the Planck scale, the additional term induces a correction in the

superpotential

$$\mathcal{W}_{\text{eff}} \rightarrow \mathcal{W} e^{\frac{3}{2}\chi H_u \cdot H_d / M_P^2} \approx \mathcal{W} + \frac{\langle \mathcal{W}_{\text{hid}} \rangle}{M_P^2} \frac{3}{2} \chi H_u \cdot H_d \equiv \mathcal{W} + \frac{3}{2} \chi m_{3/2} H_u \cdot H_d, \quad (9)$$

where the *vev* of the hidden sector superpotential can be related to the gravitino mass scale

$$m_{3/2} \approx \frac{\langle \mathcal{W}_{\text{hid}} \rangle}{M_P^2}. \quad (10)$$

Effectively, the superpotential of the NMSSM gets modified by an additional μ -like term,

$$\mathcal{W}_{\text{iNMSSM}} = \lambda S H_u \cdot H_d + \frac{\kappa}{3} S^3 + \mu H_u \cdot H_d, \quad (11)$$

with $\mu = \frac{3}{2} \chi m_{3/2}$. Thus, the effective higgsino mass term of the NMSSM Eq. (3) gets shifted by the contribution from the non-minimal coupling to supergravity leading to inflation as

$$\mu'_{\text{eff}} = \lambda \langle S \rangle + \frac{3}{2} \chi m_{3/2} = \mu_{\text{eff}} + \mu. \quad (12)$$

The low-energy smoking gun of Higgs inflation in the superconformal setup appears to be the NMSSM extended with an MSSM-like μ -term and can be quite well distinguished from either the pure MSSM or NMSSM as will be discussed in the following. We refer to this model setup as the inflationary NMSSM, or short iNMSSM.

2 A short introduction to the iNMSSM

We consider the NMSSM extended with the additional μ -term as described above only. Its presence can be motivated from a non-minimal coupling to supergravity and a proceeding transformation in the Kähler potential in such a way that only the term $\mu H_u \cdot H_d$ is present in the superpotential with $\mu = \frac{3}{2} \chi m_{3/2}$. Cosmological observations require the size of this non-minimal coupling to be

$$\chi \simeq 10^5 \lambda. \quad (13)$$

The size of the μ -term is then mainly given by the gravitino mass $m_{3/2}$ and the λ coupling, which we will assume to be $\mathcal{O}(0.1)$ in order to have sizeable NMSSM effects. Generically, we also assume $\mu \sim \mathcal{O}(1 \text{ TeV})$, which in combination requires rather light gravitinos of $m_{3/2} \sim 10 \text{ MeV}$. This might cause the cosmological gravitino problem, see Ref. [4]. Over-

abundance of gravitino dark matter, however, can be constrained by constraining the reheating temperature after inflation [5]. We assume that the details of the inflationary model can accommodate this problem as outlined in [3].

Soft SUSY breaking in the iNMSSM and the Higgs potential The additional \mathbb{Z}_3 -breaking μ -term may generate an additional soft Supersymmetry (SUSY) breaking bilinear parameter in a similar manner as the Higgs B_μ -term exists in the MSSM. All further \mathbb{Z}_3 -breaking soft SUSY breaking terms that are in general allowed, see [6] and [7], are assumed to be suppressed [8] and cannot be generated to sizeable amount by radiative corrections. Therefore, the soft SUSY breaking potential can be summarised as the usual trilinear terms of the NMSSM $\sim A_\lambda, A_\kappa$ plus the bilinear term from the non-minimal coupling to supergravity:

$$V_{\text{soft}} = \lambda A_\lambda S H_u \cdot H_d + \frac{1}{3} \kappa A_\kappa S^3 + \frac{3}{2} B_\mu \chi m_{3/2} (H_u \cdot H_d + \text{h. c.}). \quad (14)$$

The scalar potential for the two doublet and one singlet Higgs fields is defined according to the rules of SUSY and consists of the F - and D -terms as well as the contribution from SUSY breaking. In comparison to the \mathbb{Z}_3 -symmetric NMSSM, we have the additional μ -term appearing as mass term for the doublet fields and the bilinear soft breaking term $\sim H_u \cdot H_d$. The full Higgs potential is thus given by

$$\begin{aligned} V_{\text{Higgs}} = & \left[m_{H_d}^2 + (\mu + \lambda S)^2 \right] |H_d|^2 + \left[m_{H_u}^2 + (\mu + \lambda S)^2 \right] |H_u|^2 + m_S^2 |S|^2 \\ & + \frac{2}{3} \kappa A_\kappa S^3 + \left[\kappa S^2 + \lambda H_u \cdot H_d \right]^2 + 2 (B_\mu \mu + \lambda A_\lambda S) H_u \cdot H_d \\ & + \frac{g_1^2 + g_2^2}{8} (|H_d|^2 - |H_u|^2)^2 + \frac{g_2^2}{2} |H_d^\dagger H_u|^2, \end{aligned} \quad (15)$$

where we assume all parameters to be real.

The potential of Eq. (15) can easily provide the observed phenomenology of electroweak symmetry breaking meaning Higgs $vevs$ of the neutral doublet Higgs field components, $\langle H_u^0 \rangle = v_u$ and $\langle H_d^0 \rangle = v_d$, with $v_u^2 + v_d^2 = v^2 = (174 \text{ GeV})^2$ and $\tan \beta = v_u/v_d$ a free parameter; additionally, the singlet vev generates the effective μ -term of the NMSSM, $\mu_{\text{eff}} = \lambda \langle S \rangle = \lambda v_s$. In order to do so, the soft SUSY breaking masses $m_{H_d}^2$, $m_{H_u}^2$ and m_S^2 are adjusted in such a way, that the minimisation conditions

$$\left. \frac{\partial V_{\text{Higgs}}}{\partial H_d^0} \right|_{\text{vev}} = 2m_{H_d}^2 v_d + \dots, \quad \left. \frac{\partial V_{\text{Higgs}}}{\partial H_u^0} \right|_{\text{vev}} = 2m_{H_u}^2 v_u + \dots, \quad \left. \frac{\partial V_{\text{Higgs}}}{\partial S} \right|_{\text{vev}} = 2m_{H_d}^2 v_s + \dots \quad (16)$$

are fulfilled. Solving for the mass parameters is trivial and we obtain

$$m_{H_d}^2 = -(\mu + \mu_{\text{eff}})^2 - v^2 \lambda^2 \sin^2 \beta - \frac{1}{2} M_Z^2 \cos(2\beta) + a_1 \tan \beta, \quad (17a)$$

$$m_{H_u}^2 = -(\mu + \mu_{\text{eff}})^2 - v^2 \lambda^2 \cos^2 \beta + \frac{1}{2} M_Z^2 \cos(2\beta) + a_1 \cot \beta, \quad (17b)$$

$$m_S^2 = a_4 - a_5 - a_7 - v^2 \lambda^2 - 2\mu_{\text{eff}}^2 \left(\frac{\kappa}{\lambda} \right)^2, \quad (17c)$$

where $M_Z^2 = \frac{g_1^2 + g_2^2}{2} v^2$ and the abbreviations a_i are defined in Appendix A.

The electroweak breaking conditions have to be taken with great care, since the potential possesses multiple minima and even if *one particular* minimum is selected to be the electroweak (*desired*) vacuum by the above definition, there might be other minima deeper than the desired vacuum and thus the *true* vacuum, i. e. the global minimum, is not the desired one anymore. At tree-level, the minimisation can only be done numerically; at the loop-level, the situation even gets worse and one has to guess suitable starting values for the numerical routines, which have the potential to miss several of the minima. We take those constraints at the tree-level seriously and therefore exclude parameter points leading to a non-standard true vacuum of the theory. Typically, this global minimum appears at larger *vevs* for the fields and thus gets more easily selected by the cosmological history of the universe after inflation [9]. Since we start with *vevs* shortly below the Planck scale after inflation ends, the universe while cooling down may get stuck in the higher scale vacuum. If it is a local minimum, one should consider the tunneling to the desired one. Typically, however, the larger *vev* vacuum appears to be deeper than the desired vacuum.

Besides the fact that there are multiple vacua implying alternative *vevs* in the Higgs potential, the Higgs mass matrices (defined in Appendix A) show tachyonic states at the tree-level depending on the input parameters. Tachyonic states have negative masses squared and simply invalidate the electroweak expansion point because the potential at that point appears to be a local maximum (the “minimisation” conditions are rather conditions for stationary points and may also result in maxima or saddle points) and thus pointing towards the deeper minimum in the tachyonic direction. Actually, radiative corrections may lift up the potential in some cases leading to rather light instead of tachyonic states. We take these constraints nevertheless seriously and exclude tachyonic parameter configurations irrespective whether radiative corrections lift the masses up or not.

The tachyonic constraints already confine clear portions of parameter space that remain valid. In addition, vacuum stability considerations exclude additional parts at the borderline.

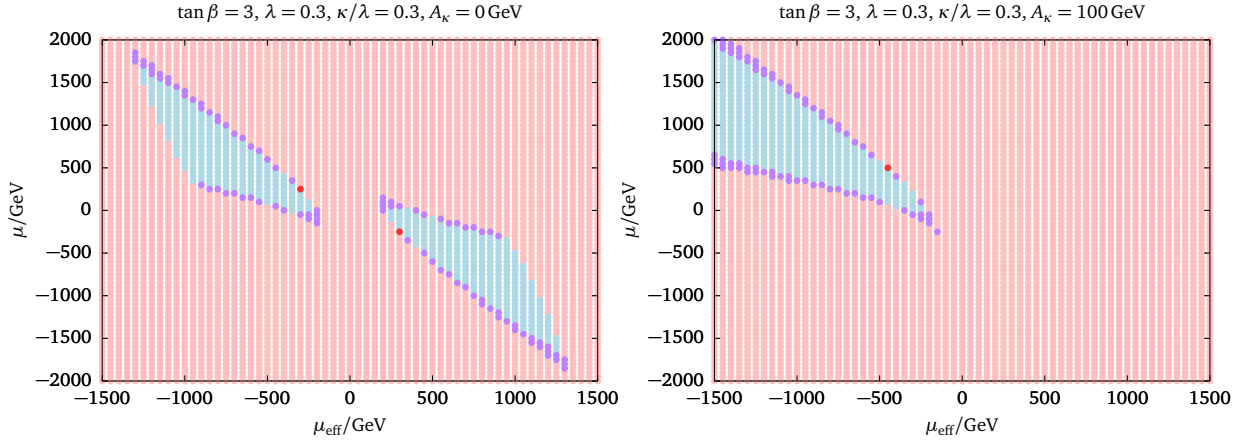


Figure 1: Mainly tachyonic constraints (pink points) confine the allowed parameter space (light blue). In addition, there exist other global minima than the electroweak v_{evs} with a mostly long-lived desired vacuum (purple) and rarely short-lived configurations (red). On the left panel, A_κ is taken to be 0 GeV, where on the right panel $A_\kappa = 100$ GeV. The positive value of A_κ disallows the right wing for positive μ_{eff} that was allowed for vanishing A_κ (in this case, there is a reflection symmetry). Moreover, there is a clear correlation between the allowed signs of μ and μ_{eff} , which in most cases have to differ unless μ appears to be small. The other sign of A_κ reversed the situation.

3 Electroweak phenomenology of the iNMSSM

The phenomenology of the iNMSSM at the electroweak scale deviates significantly from the usual NMSSM. On the one hand, the number of states remain the same which may look like the same phenomenology. On the other hand, the dependence on certain parameters appears to be very different and the additional μ -term changes the interpretation of the higgsino mass parameter as well as the functional dependence of the Higgs masses on it.

First of all, the tachyonic selection rule excludes large as well as very small ($\lesssim \sqrt{2}\nu$) values of μ_{eff} . Moreover, both values μ and μ_{eff} appear to be correlated. This can be seen from Figure 1. The tachyonic boundaries can be easily understood from a look at the mass matrices, see Appendix A, where the small μ_{eff} value sets A_λ to be large, which sits on the off-diagonal elements and thus is responsible for a large mixing which potentially drives one state negative. Similarly, if the combination $\mu + \mu_{\text{eff}}$ appears to be large; therefore same signs of μ and μ_{eff} are excluded in most cases. The trilinear soft SUSY breaking parameter A_κ mainly influences the pseudoscalar singlet-like state. If this one appears to be tachyonic for small A_κ , larger values of this parameter have the ability to lift this mass up and open up parameter space that is excluded with a vanishing A_κ . This can be clearly seen in the comparison of the allowed and excluded parameter space in the μ - μ_{eff} -plane shown in Figure 1.

The effect of both μ and μ_{eff} on the tachyonicity of states can be seen from the dependence

of the Higgs spectra on these parameters. In Figure 2, we show the functional dependence of the two lightest scalar and the lightest pseudoscalar masses on μ_{eff} for several values of μ . The heavy states are mainly dominated by the input $m_{H^\pm} = 800$ GeV.

A precise knowledge of the Higgs sector in the NMSSM hence allows to distinguish between the pure \mathbb{Z}_3 -symmetric NMSSM and the inflation-inspired iNMSSM with the additional \mathbb{Z}_3 -breaking μ -term. So far, we have not considered the additional soft SUSY breaking bilinear and kept it zero. In combination with a measurement of the neutralino sector, which in contrast rather mimics the NMSSM, there is a clear smoking gun of inflation that can be detected at an electroweak precision machine like a future Linear Collider. The Higgs spectrum varies severely with varying μ as we show in Figure 3. Here, we compare the light pseudoscalar case with vanishing A_κ with the heavier scenario where $A_\kappa = 100$ GeV for illustrative reasons. While the light pseudoscalar mass is lifted up mainly by the amount of A_κ , the tachyonic state for $\mu = 1000$ GeV gets non-tachyonic and the scalar spectrum only changes marginally. The heavy states are merely fixed by the input value of the charged Higgs mass $m_{H^\pm} = 800$ GeV.

The electroweakino sector is defined and briefly described in Appendix A, where it can be seen from the neutralino mass matrix that the singlino mass is governed by $\frac{\kappa}{\lambda}\mu_{\text{eff}}$, where the higgsino mass is determined by $\mu + \mu_{\text{eff}}$. Thus, a small higgsino mass, and therefore especially also a small charged higgsino mass, which is preferably detectable at a Linear Collider, is somewhat naturally selected in the iNMSSM where μ and μ_{eff} have to have opposite signs and rather the same magnitude. Such a cancellation, however, if μ is significantly large, tends to produce a heavy singlino in the iNMSSM in contrast to the NMSSM. This effect can be removed by adjusting the ratio κ/λ in such a way that both singlino and higgsino masses scale the same with μ . By this redefinition, however, if λ is kept fixed, the value of κ changes dramatically. While the electroweakino sector may look the same as in the NMSSM even in the presence of a large μ -term, the (pseudo)scalar sector still has a strong dependence on the additional μ -term which is shown in Figure 4.

All the considerations and predictions presented above still depend on other, previously suppressed parameters. For the illustrative purpose, these additional parameters have been fixed to some values, variation of them also changes the structure of the plots shown in this talk. Especially the top/stop sector enters the determination of the SM-like Higgs mass, where we show mass contours in an interesting slice of parameter space in the following. The Higgs mass predictions contain the full iNMSSM one-loop and leading two-loop contributions, where the stop contribution in all cases was fixed to be sizeable and beyond the

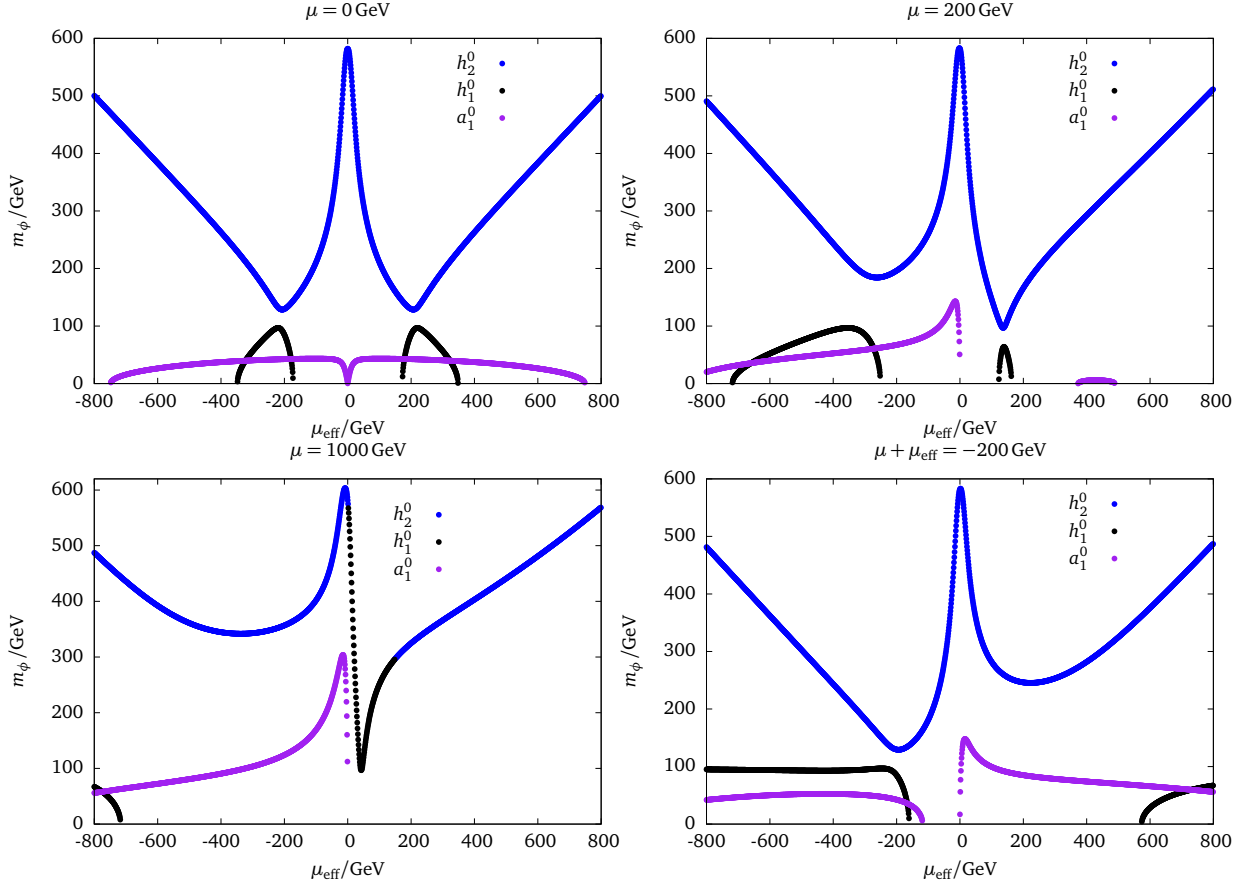


Figure 2: The spectra of the lightest states and how they vary with varying μ_{eff} , exemplarily for some choice of parameters. In each plot, the μ value is fixed to a given value, where the lower right plot has a fixed sum $\mu + \mu_{\text{eff}} = -200$ GeV. It can be clearly seen which intervals are allowed (those with all three states appearing in the plot; where one or more are missing, these are tachyonic). In the case with $\mu = 1000$ GeV h_1^0 and h_2^0 apparently change their role which is due to the fact that the absolute value of the tachyonic state grows above the corresponding value of h_2^0 . The scenario with $\mu = 200$ GeV shows the feature that the tachyonic exclusions are exclusive in the sense that one tachyonic state (scalar or pseudoscalar) is enough to exclude the spectrum. Here, both the lightest scalar and pseudoscalar have some small interval for positive μ_{eff} where they are non-tachyonic but the respective other one is and thus all the range for positive μ_{eff} is excluded (where a_1^0 turns tachyonic the first time for growing μ_{eff}) and additionally already the light scalar mass gets tachyonic at larger negative values of μ_{eff} . This artefact can be also seen in the region plot of Figure 1.

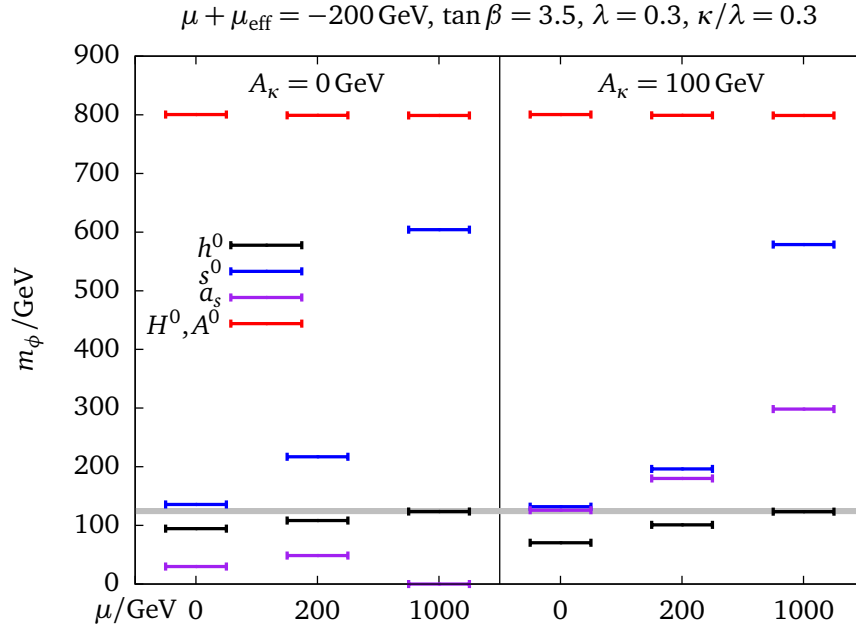


Figure 3: The Higgs spectra change with different choice of μ from 0, 200 to 1000 GeV, where the black line corresponds to the SM-like state. The grey band around 125 GeV shows the experimentally favoured region with an error of ± 3 GeV. The effective higgsino parameter was fixed to a value of $\mu + \mu_{\text{eff}} = -200$ GeV. On the right side, with respect to the left side, the A_κ contribution is risen from 0 to 100 GeV which lifts the pseudoscalar singlet mass of a_s up and turns the tachyonic point at $A_\kappa = 0$ GeV and $\mu = 1000$ GeV non-tachyonic with a rather large a_s mass.

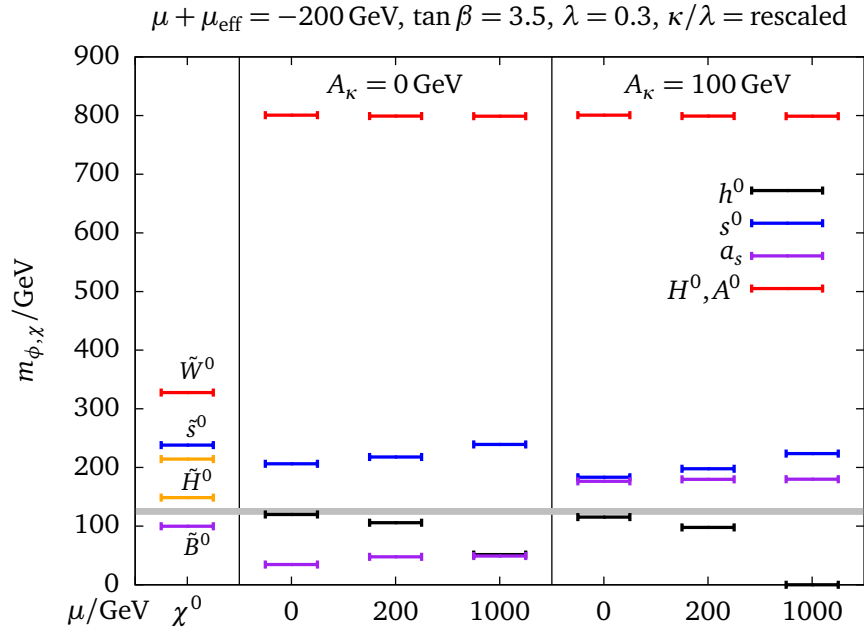


Figure 4: While the neutralino spectrum stays invariant under change of μ , where the sum $\mu + \mu_{\text{eff}} = -200 \text{ GeV}$ is fixed and the ratio κ/λ rescaled in such a way to keep the singlino mass fixed as well, the scalar and pseudoscalar Higgs spectrum still change with varying μ .

direct limits on stop searches in such a way that $m_{\tilde{t}} = 2 \text{ TeV}$ and the mixing was chosen to be $A_t = 2m_{\tilde{t}}$. This particular choice, of course, can and has to be adjusted in a precision analysis. Moreover, the influence of the other input parameters as $\tan \beta$, λ , κ , and to some extent A_{κ} , still has to be tackled down in order to clearly determine the precision needed to distinguish two different scenarios of the NMSSM and the iNMSSM generating similar spectra. In addition, the electroweak phenomenology also involves production and decay rates of the Higgs states and thus one has an additional handle to distinguish the two models. In any case, a precise measurement of the electroweak sector at a future collider will give clear insights whether there is a smoking gun of inflation at the Linear Collider or not. This will be discussed in a forthcoming publication [10].

We have discussed above that an interesting slice of parameter space is defined by the sum of the two μ -terms, $\mu + \mu_{\text{eff}}$, and the ratio κ/λ . The couplings λ and κ are known to run into a Landau pole below the GUT scale in the NMSSM, and the same is true for the iNMSSM since the additional μ -term does not change the running. This non-perturbativity can be avoided, if λ and κ are taken to be constrained by $\lambda^2 + \kappa^2 \lesssim 0.5$, which will be always the case in the region plots shown in the following. The phenomenologically interesting regions are those where the SM-like Higgs state can accommodate for the observed 125 GeV,

the branching ratios are in those cases SM-like as well. There is an experimental exclusion from direct higgsino searches which constrains the chargino mass to be $\lesssim 94$ GeV; this bound approximately transfers to $\mu + \mu_{\text{eff}}$.

Parameter scans and results The parameter space of the iNMSSM gets enlarged by two dimensions (μ and B_μ) with respect to the NMSSM. Additionally, the new parameters may invalidate certain allowed regions of the NMSSM that become e. g. tachyonic once a sufficiently large μ parameter is turned on. On the other hand, as we have seen, there might be cancellations between μ and μ_{eff} . The surviving parameter space appears to be rather constrained, which allows to have some clear predictions, especially on the hierarchy of Higgs boson masses. Unfortunately, as there are many parameters available, modification of one of these where the others remain fixed, relax the constraints and thus diminish the predictability. This, however, comes along with a different phenomenology and thus clearly distinguish different scenarios.

One crucial parameter, as already discussed above, is given by A_κ which controls the mass of the singlet pseudoscalar state. Low values of A_κ produce a rather light state, heavier masses can be generated by lifting A_κ up and simultaneously removing tachyons from the spectrum. There appears to be a larger fraction in the parameter space allowed, if one takes a look at the κ/λ vs. $\mu + \mu_{\text{eff}}$ slice. This is shown in the samples of Figure 5. Comparing the cases with $A_\kappa = 0$ and 100 GeV, the effect of opening up excluded tachyonic parameter space can be clearly seen for the cost of a heavier singlet pseudoscalar (green contours). The other parameters only have a minor effect, so $\tan\beta$ is enhanced from 2.5 to 3.5 from the second to the third row of Figure 5 and simultaneously λ reduced from 0.6 to 0.3, where $\mu = 1$ TeV in all cases. The single plot on top of Figure 5 illustrates the “NMSSM-limit” with vanishing μ . Here, apparently only a very constrained region is allowed (note that $A_\kappa = 0$ GeV) and the singlet-like pseudoscalar can be rather light. The grey bands show a rough experimental exclusion on the higgsino mass given by the LEP-limit on the chargino mass $m_{\chi_1^\pm} > 94$ GeV [11]. The chargino mass in the iNMSSM is mainly given by $\mu + \mu_{\text{eff}}$, see Appendix A, up to small modifications from the mixing.

Figure 5 also reveals information about the vacuum structure of the scanned points: light blue points denote an absolutely stable electroweak vacuum, where tachyonic states (at the tree-level!) appear in the pink points. Interestingly, the allowed regions can also easily accommodate for a 125 GeV SM-like Higgs state, where we added a uniform stop contribution as discussed above. Unstable or metastable desired vacua are coded in purple (long-lived)

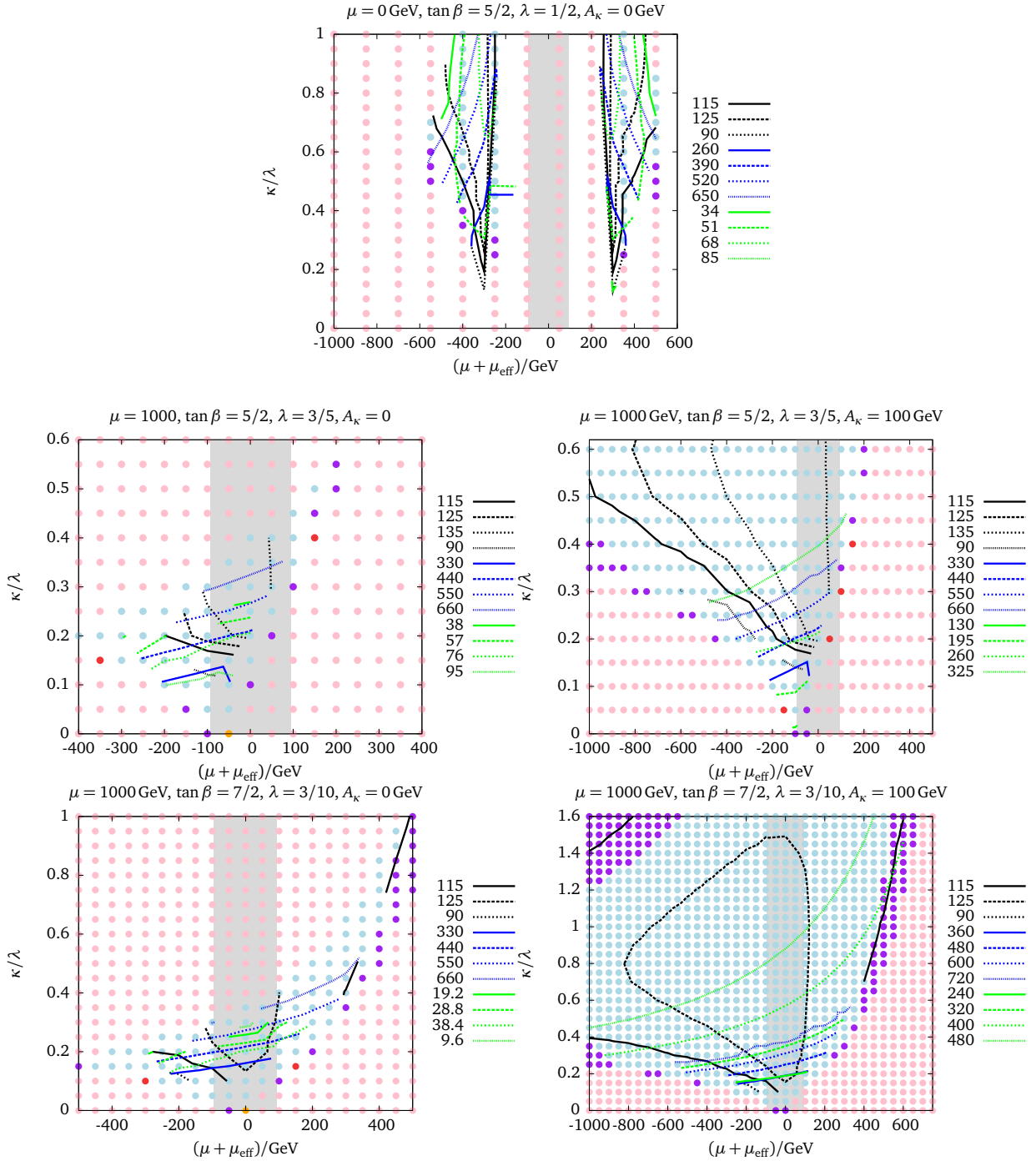


Figure 5: This selection shows some typical features of the iNMSSM and NMSSM. The upper plot has the situation with a vanishing μ and A_κ which is symmetric around $\mu_{\text{eff}} = 0 \text{ GeV}$. The lower plots have a larger $\mu = 1 \text{ TeV}$ and show the effect of A_κ on the allowed space and the Higgs masses. The contour lines give lines of equal masses for the SM-like Higgs (black), the singlet-like scalar (blue) and singlet-like pseudoscalar (green) which can be light for small A_κ . The colour code for the points describes the stability of the electroweak ground state: light blue points have a global electroweak vacuum, purple ones a long-lived and red points a short-lived desired vacuum. Tachyonic points are shown in pink, where the orange points do not fulfil the NMSSM constraint $A_\kappa^2 > 9m_S^2$ for a non-vanishing singlet vev .

and red (short-lived). We briefly describe in Appendix B how we estimate the life-time. In the NMSSM there exists a bound on A_κ ,

$$A_\kappa^2 > 9m_S^2, \quad (18)$$

relating the trilinear soft SUSY breaking singlet coupling with the soft SUSY breaking singlet mass. This constraint is needed to generate a sufficiently large singlet $\nu\nu$ and therefore higgsino mass. In the presence of the \mathbb{Z}_3 -breaking μ -term, this unequation does not have to be necessarily fulfilled. Eq. (18) can be easily derived from the singlet-only potential with the requirement that the minimum $\langle S \rangle \neq 0$ is the true vacuum and thus a non-vanishing singlet $\nu\nu$ is generated. This is needed in the \mathbb{Z}_3 -invariant NMSSM to produce the correct electroweak phenomenology. In the iNMSSM, however, the \mathbb{Z}_3 -breaking MSSM-like μ -term is generated by the non-minimal coupling to supergravity and related to the scale of SUSY breaking and the gravitino mass. If both μ and μ_{eff} are present, there can be cancellations since they have to have different signs and hence a small higgsino mass still can be valid even if both μ parameters are in the TeV range.

Constraints on B_μ The iNMSSM has in addition to the superpotential parameter one more soft SUSY breaking term, the bilinear B_μ -term, which has been ignored to far in the discussion above. It turns out that it cannot be arbitrarily large anyway and thus there are good reasons to keep it small. If it is non-zero, the effect is merely under control as the contribution from B_μ grows linearly with μ (note that it appears as μB_μ in the soft breaking potential). Together with A_λ it influences the charged Higgs mass and therefore, in our approach where we treat A_λ for m_{H^\pm} as input, it enters the determination of A_λ , see Eq. (24) and Appendix A.

Unfortunately, the role of B_μ is less clear than compared to the MSSM where it can easily be replaced by the pseudoscalar mass m_A . However, its impact on the Higgs boson masses is very well-defined as it always enters in sum with $\mu_{\text{eff}} \frac{\kappa}{\lambda}$ and $\mu_{\text{eff}} A_\lambda$. Thus, it might be absorbed in A_λ , which nevertheless appears also at different places. Treating the charged Higgs mass m_{H^\pm} as input and solving for A_λ , μB_μ enters the determination of A_λ . For too large values of B_μ , tachyonic states are generated again. There is, however, a valley that allows for non-tachyonic states even for large but negative B_μ values as can be seen from Figure 6. It has nevertheless the power to destabilize the desired electroweak ground state of the theory as such large values of B_μ induce a global minimum different from the standard vacuum. The desired but local vacuum now appears to be rather short-lived with respect to the life-time of the universe, which is depicted by the red points of Figure 6. In the boundary region, the

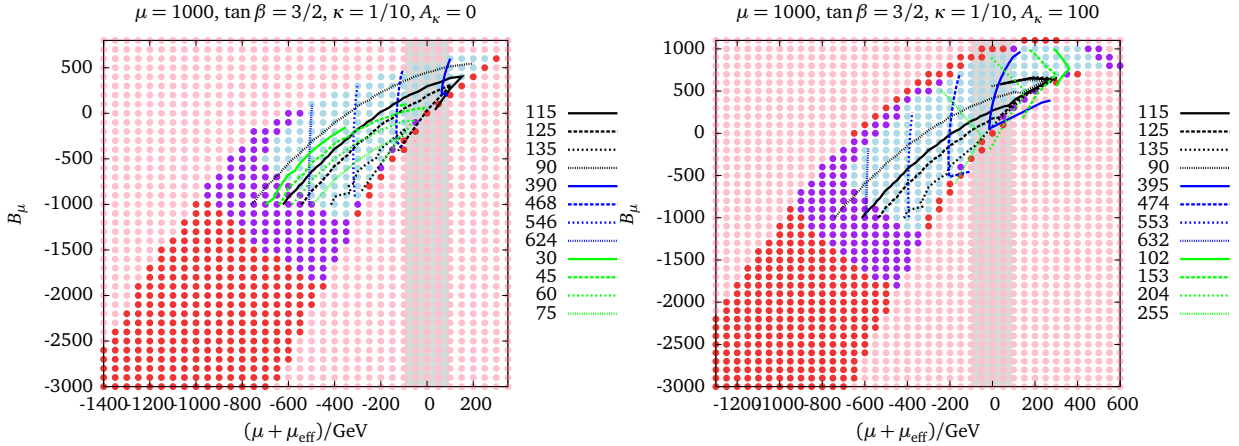


Figure 6: The color code of the points and lines is as in Figure 5. Now we show for fixed values of $\kappa = 0.1$ and $\lambda = 0.5$ a different slice in parameter space for non-vanishing B_μ . In this region, the influence from A_κ is apparently very mild (left $A_\kappa = 0$ GeV, right $A_\kappa = 100$ GeV). For too large negative values of B_μ , the electroweak vacuum gets severely short-lived (red points) with a broad band of long-lived desired vacua in between (purple).

electroweak vacuum is sufficiently long-lived (purple points). The scalar singlet mass (blue lines) appears to be rather independent of B_μ , where the pseudoscalar singlet (green) shows a striking behaviour depending on A_κ . The mass of the SM-like Higgs boson is very much aligned with the allowed valley and explains very well the tachyonic boundary (together with the pseudoscalar singlet as can be already seen from Figure 2, where $B_\mu = 0$ GeV but the two light states running tachyonic at different places).

4 Conclusions

We have presented the electroweak phenomenology of an inflation-inspired NMSSM as first discussed in Refs. [1–3]. We briefly summarized the idea of Higgs inflation in the superconformal sector and showed how the non-minimal coupling of the Higgs sector to supergravity shows up in the effective low-energy superpotential. The remaining model can be described by the NMSSM augmented with an MSSM-like μ -term $\mu H_u \cdot H_d$, which breaks the accidental \mathbb{Z}_3 invariance of the NMSSM. Additionally, a soft SUSY breaking term μB_μ is generated and has to be taken into account. The rules of supergravity dictate μ and together with the effective μ -term of the NMSSM arising from the singlet v_{ev} , both sum up to an effective higgsino mass $\mu + \mu_{\text{eff}}$. This combination plays an important role for the phenomenology, especially since the signs of both contributions appear to be anticorrelated and thus a natural cancella-

tion among those fundamentally different contributions to the higgsino mass appears. Thus, scenarios with light higgsinos but heavy singlinos exist and a precise knowledge of the Higgs and electroweakino sector as it might be achieved at a future Linear Collider helps to clearly distinguish this model from the ordinary NMSSM. A smoking gun of the inflationary remnant exists as a footprint in the electroweak spectrum.

We have extensively discussed the influence of the model parameters on the Higgs masses and how tachyonic states are generated at the tree-level. Tachyonic masses invalidate the expansion point in such a way that the electroweak point appears to be a local maximum instead of a minimum and thus the tachyonic direction points towards the global minimum. In addition, the iNMSSM as well as the NMSSM may reveal several vacua out of which the desired vacuum appears to be a false vacuum. A numerical analysis minimising the scalar potential finds the global minimum of the theory, which in some cases is the electroweak vacuum in others not. If the desired vacuum is a local minimum, vacuum decay rates have been estimated to compare the life-time of the false vacuum with the life-time of the universe. Only in the case of large and negative B_μ values, reasonable amounts of short-lived vacua have been found.

Higgs inflation embedded into a superconformal framework appears to be distinguishable at low energies from the common SUSY models beyond the SM. The iNMSSM needs an additional singlet as the NMSSM; the spectrum, however, appears to be different and cannot be matched to the parameters of the NMSSM. The model is also different from the MSSM, in which Higgs inflation cannot be accommodated.

The results presented in this talk are going to be discussed in more detail in a forthcoming publication [10].

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A Higgs boson and neutralino/chargino mass matrices

We define the Higgs mass matrices via the second derivatives of the potential, where we distinguish between scalar and pseudoscalar neutral states by the decomposition

$$H_u = \begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix} = \begin{pmatrix} \phi_u^+ \\ v_u + \frac{1}{\sqrt{2}}(\sigma_u + i\phi_d) \end{pmatrix}, \quad H_u = \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix} = \begin{pmatrix} v_d + \frac{1}{\sqrt{2}}(\sigma_d + i\phi_d) \\ \phi_d^- \end{pmatrix}, \quad (19)$$

$$S = v_s + \frac{1}{\sqrt{2}}(\sigma_s + i\phi_s).$$

The mass matrices for the scalar and pseudoscalar states \mathcal{M}_S^2 and \mathcal{M}_P^2 , respectively, are then given by the expressions

$$\mathcal{M}_S^2 = \begin{pmatrix} M_Z^2 c_\beta^2 + a_1 t_\beta & (2v^2 \lambda^2 - M_Z^2) c_\beta s_\beta - a_1 & a_2 c_\beta - a_3 s_\beta \\ * & M_Z^2 s_\beta^2 + a_1/t_\beta & a_2 s_\beta - a_3 c_\beta \\ * & * & a_4 + a_5 \end{pmatrix}, \quad (20a)$$

$$\mathcal{M}_P^2 = \begin{pmatrix} a_1 t_\beta & a_1 & -a_6 s_\beta \\ * & a_1/t_\beta & -a_6 c_\beta \\ * & * & a_4 - 3a_5 - 2a_7 \end{pmatrix}, \quad (20b)$$

with $s_\beta = \sin \beta$, $c_\beta = \cos \beta$, $t_\beta = \tan \beta$ and where the abbreviations a_i are

$$a_1 = B_\mu \mu + \mu_{\text{eff}} \left(\frac{\kappa}{\lambda} \mu_{\text{eff}} + A_\lambda \right), \quad (21a)$$

$$a_2 = 2v\lambda(\mu + \mu_{\text{eff}}), \quad (21b)$$

$$a_3 = v\lambda \left(2\frac{\kappa}{\lambda} \mu_{\text{eff}} + A_\lambda \right), \quad (21c)$$

$$a_4 = \frac{1}{\mu_{\text{eff}}} \left[v^2 \lambda^2 c_\beta s_\beta \left(\frac{\kappa}{\lambda} \mu_{\text{eff}} + A_\lambda \right) - v^2 \lambda^2 \mu \right], \quad (21d)$$

$$a_5 = 4 \left(\frac{\kappa}{\lambda} \right)^2 \mu_{\text{eff}}^2 + \frac{\kappa}{\lambda} \left[\mu_{\text{eff}} A_\kappa - v^2 \lambda^2 c_\beta s_\beta \right], \quad (21e)$$

$$a_6 = v\lambda \left(2\frac{\kappa}{\lambda} \mu_{\text{eff}} - A_\lambda \right), \quad (21f)$$

$$a_7 = -6 \left(\frac{\kappa}{\lambda} \right)^2 \mu_{\text{eff}}^2. \quad (21g)$$

Note, that the pseudoscalar mass matrix comprises one massless state, the Goldstone mode. The charged Higgs mass matrix is given by

$$\mathcal{M}_C^2 = [(M_W^2 - v^2\lambda^2)c_\beta s_\beta + a_1] \begin{pmatrix} t_\beta & 1 \\ 1 & 1/t_\beta \end{pmatrix}, \quad (22)$$

with the W boson mass $M_W^2 = \frac{1}{2}g_2^2 v^2$ and the eigenvalue given by

$$m_{H^\pm} = M_W^2 - v^2\lambda^2 + \frac{a_1}{c_\beta s_\beta}, \quad (23)$$

which can be used to eliminate A_λ as a free parameter for the sake of the charged Higgs boson mass m_{H^\pm} as input value (for the numerical analyses presented in this talk, we used all over the scenarios $m_{H^\pm} = 800$ GeV), such that

$$A_\lambda = \frac{c_\beta s_\beta}{\mu_{\text{eff}}} (m_{H^\pm}^2 - M_W^2 + v^2\lambda^2) - \frac{B_\mu \mu}{\mu_{\text{eff}}} - \mu_{\text{eff}} \frac{\kappa}{\lambda}. \quad (24)$$

The mass matrices of charginos and neutralinos resemble very much the ordinary NMSSM, where the effective higgsino parameter is replaced by $\mu + \mu_{\text{eff}}$. However, the singlino mass is only governed by μ_{eff} , since the additional μ term couples the doublet superfields. Therefore, the neutralino mass matrix is given by

$$\mathcal{M}_{\chi^0} = \begin{pmatrix} M_1 & 0 & -M_Z s_w c_\beta & M_Z s_w s_\beta & 0 \\ * & M_2 & M_Z c_w c_\beta & -M_Z c_w s_\beta & 0 \\ * & * & 0 & -(\mu + \mu_{\text{eff}}) & -\lambda v s_\beta \\ * & * & * & 0 & -\lambda v c_\beta \\ * & * & * & * & 2\frac{\kappa}{\lambda} \mu_{\text{eff}} \end{pmatrix}, \quad (25)$$

where M_1 and M_2 are the gaugino masses for the $U(1)_Y$ and $SU(2)_L$ gauginos, respectively, and the weak mixing angle θ_w enters via $\tan \theta_w = s_w/c_w = g_1/g_2$. Apparently, the neutralino spectrum in the iNMSSM can be rescaled via the ratio $\frac{\kappa}{\lambda}$ in such a way to match the NMSSM neutralino spectrum for a given higgsino mass $\mu + \mu_{\text{eff}}$.

The chargino mass matrix is given by

$$\mathcal{M}_{\chi^\pm} = \begin{pmatrix} M_2 & \sqrt{2}M_W s_\beta \\ \sqrt{2}M_W c_\beta & \mu + \mu_{\text{eff}} \end{pmatrix}. \quad (26)$$

B Vacuum tunneling

We briefly describe our estimate on the tunneling rates in the case where the desired vacuum appears to be a false vacuum. The electroweak input parameters determine the position and depth of the local minimum, where the global minimum and true vacuum is found by numerical minimisation of the tree-level potential. In general, the true vacuum has *vevs* $\langle H_u^0 \rangle \neq v_u$, $\langle H_d^0 \rangle \neq v_d$ and $\langle S \rangle \neq v_s$. We approximate the potential barrier between the two minima by a one-dimensional quartic potential

$$V(\phi) = g \phi^4 - a \phi^3 + b \phi^2 + c \phi + d, \quad (27)$$

which allows for an exact solution of the bounce action [12]. This is given by

$$B = \frac{\pi^4}{3g} (2 - \delta)^{-3} [\alpha_1 \delta + \alpha_2 \delta^2 + \alpha_3 \delta^3], \quad (28)$$

with $\delta = 8g^2 b/a^2$ and $\alpha_{1,2,3}$ numerical coefficients. By comparison of the decay rate of the false vacuum per unit volume [13]

$$\Gamma/V = Ae^{-B/\hbar} [1 + \mathcal{O}(\hbar)], \quad (29)$$

one estimates bounce actions $B \gtrsim 400$ to be sufficiently long-lived. The prefactor A is difficult to calculate and usually approximated by the height of the potential or the electroweak scale, $A \sim (100 \text{ GeV})^4$, where the error enters only logarithmically the decay time.

The interpolation between the two minima is done by a straight line, where the electroweak point is shifted to the origin. Therefore, in the expression of the neutral Higgs potential, we have the replacement

$$V(\phi) = V \left(H_u^0 = v_u + (V_u - v_u) \frac{\phi}{\sqrt{2}}, H_d^0 = v_d + (V_d - v_d) \frac{\phi}{\sqrt{2}}, S = v_s + (V_s - v_s) \frac{\phi}{\sqrt{2}} \right), \quad (30)$$

with the “true” *vevs* V_u , V_d and V_s . The factor $1/\sqrt{2}$ is employed to keep the ϕ -field canonically normalised. This way, the field ϕ interpolates between the desired, false vacuum ($\phi = 0$) and the true vacuum ($\phi = 1$). It is, however, more convenient to keep ϕ dimensionful and thus the coefficients of the potential in Eq. (27) of the same order of magnitude as the original

coefficients. Therefore, we use a normalised field in the one-field potential, $V(\bar{\phi})$, with

$$\bar{\phi} = \frac{\phi}{\sqrt{(V_u - v_u)^2/2 + (V_d - v_d)^2/2 + (V_s - v_s)^2/2}}. \quad (31)$$

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