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Higgs bosons in the NMSSM with exact and slightly broken PQ-symmetry

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Abstract

We explore the Higgs sector of the NMSSM in the limit when the Peccei–Quinn symmetry is exact or only slightly broken. In this case the Higgs spectrum has a hierarchical structure which is caused by the stability of the physical vacuum. We find a strong correlation between the parameters of the NMSSM if $\kappa = 0$ or $\kappa \lesssim \lambda^2$. It allows one to distinguish the NMSSM with exact or softly broken PQ-symmetry from the MSSM even when extra scalar and pseudoscalar Higgs states escape direct detection.

1. Introduction

Nowadays the simplest supersymmetric (SUSY) extension of the Standard Model (SM), the so–called Minimal Supersymmetric Standard Model (MSSM), is possibly the best motivated model beyond the SM. Indeed the quadratic divergences, that lead to the hierarchy problem in the SM, are naturally canceled in supersymmetric theories. By making supersymmetry local (supergarvity) a partial unification of gauge interactions with gravity can be achieved. The remarkable coincidence of gauge coupling constants at the high energy scale $M_X \sim 10^{16}$GeV obtained in the framework of the MSSM allows one to embed the simplest SUSY extension of the SM into Grand Unified and superstring theories.

The stabilization of the mass hierarchy in the MSSM does not provide any explanation for the origin of the electroweak scale, and therefore a minimal SUSY model should not know about the electroweak scale before symmetry breaking. However, the MSSM superpotential contains one bilinear term $\mu(\hat{H}_1\epsilon\hat{H}_2)$ which is present before supersymmetry is broken; from dimensional considerations one would naturally expect the parameter μ to be either zero or the Planck scale, but in order to provide the correct pattern of electroweak symmetry breaking, μ is required to be of the order of the electroweak scale. Thus minimal SUSY has its own "hierarchy" problem, known as the μ-problem.

The most elegant solution of the μ–problem naturally appears in the framework of superstring–inspired E_6 models where the bilinear terms are forbidden by gauge symmetry. In general these models contain a few pairs of the Higgs doublets and a few singlet fields S_i. Assuming that only one pair of Higgs doublets and one singlet survive to low energies the superpotential of the Higgs sector takes the form $\lambda\hat{S}(\hat{H}_1\epsilon\hat{H}_2)$. The considered model includes only one additional singlet field and almost the same number of parameters as the MSSM. For this reason it can be regarded as the simplest extension of the MSSM. As a result of spontaneous symmetry breakdown at the electroweak scale the superfield $\hat{S}$ gets a non-zero vacuum expectation value ($\langle S\rangle \equiv s/\sqrt{2}$) and an effective μ-term ($\mu = \lambda s/\sqrt{2}$) of the required size is automatically generated.

The model discussed above possesses a $SU(2)\times[U(1)]^2$ global symmetry. Being broken by the vacuum an extended global symmetry leads to the appearance of a massless CP-odd spinless particle in the Higgs boson spectrum which is a Peccei-Quinn (PQ) axion. The usual way to avoid this axion is to introduce a term cubic in the new singlet superfield $\hat{S}$ in the superpotential that explicitly breaks the additional $U(1)$ global symmetry. The superpotential of the Higgs sector of the obtained model, which is the so–called Next–to–Minimal Supersymmetric Standard Model (NMSSM), is given by [1]

$$W_H = \lambda\hat{S}(\hat{H}_1\epsilon\hat{H}_2) + \frac{1}{3}\kappa\hat{S}^3\,. \tag{1}$$

In this paper we study the Higgs sector of the NMSSM. In Section 2 we discuss the MSSM limit of the NMSSM. In Section 3 we investigate the spectrum and couplings of the Higgs bosons in the exact PQ–symmetry limit of the NMSSM where $\kappa = 0$. The scenario with a slightly broken PQ–symmetry, where κ is small, is considered in Section 4. The results are summarized in Section 5.

2. The MSSM limit of the NMSSM

The Higgs sector of the NMSSM includes two Higgs doublets $H_{1,2}$ and one singlet field S. The potential energy of the Higgs field interaction can be written as a sum

$$
\begin{aligned}
V &= V_F + V_D + V_{soft} + \Delta V\,,\\
V_F &= \lambda^2|S|^2(|H_1|^2+|H_2|^2) + \lambda^2|(H_1\epsilon H_2)|^2 + \lambda\kappa\left[S^{*2}(H_1\epsilon H_2)+h.c.\right] + \kappa^2|S|^4\,,\\
V_D &= \frac{g^2}{8}\left(H_1^+\sigma_a H_1 + H_2^+\sigma_a H_2\right)^2 + \frac{g'^2}{8}\left(|H_1|^2-|H_2|^2\right)^2\,,\\
V_{soft} &= m_1^2|H_1|^2 + m_2^2|H_2|^2 + m_S^2|S|^2 + \left[\lambda A_\lambda S(H_1\epsilon H_2) + \frac{\kappa}{3}A_\kappa S^3 + h.c.\right]\,,
\end{aligned}
\tag{2}
$$

where $H_1^T = (H_1^0,\, H_1^-)$, $H_2^T = (H_2^+,\, H_2^0)$ and $(H_1\epsilon H_2) = H_2^+H_1^- - H_2^0H_1^0$. At the tree level the Higgs potential (2) is described by the sum of the first three terms. V_F and V_D are the F and D terms. Their structure is fixed by the NMSSM superpotential (1) and the electroweak gauge interactions in the common manner. The last term in Eq.(2), ΔV, corresponds to the contribution of loop corrections.

The parameters g and g' are $SU(2)$ and $U(1)$ gauge couplings respectively ($g_1 = \sqrt{5/3}g'$), which are known precisely. The couplings g, g', λ and κ do not violate supersymmetry. The soft supersymmetry breaking terms are collected in V_{soft}. At the tree level the set of soft SUSY breaking parameters involves soft masses m_1^2, m_2^2, m_S^2 and trilinear couplings A_κ, A_λ. The inclusion of loop corrections draw into the analysis many other soft SUSY breaking terms that define the masses of different superparticles. All parameters listed above are here assumed to be real. In general, complex values of λ, κ and soft couplings induce CP–violation, and we here restrict our consideration to the part of the NMSSM parameter space where CP is conserved.

At the physical minimum of the potential (2) the neutral components of the Higgs doublets H_1 and H_2 develop vacuum expectation values v_1 and v_2 breaking the electroweak symmetry down to $U(1)$. Instead of v_1 and v_2 it is more convenient to use $\tan\beta = \dfrac{v_2}{v_1}$ and $v = \sqrt{v_1^2+v_2^2}$, where v=246 GeV in the physical vacuum.

To start with let us specify the transition from the NMSSM to the minimal SUSY model which has been studied thoroughly. Because the strength of the interaction of the extra singlet fields with other bosons and fermions is determined by the size of the coupling λ in the superpotential (1) the MSSM sum rules for the Higgs masses and couplings are reproduced when λ tends to be zero. It implies that the vacuum expectation value of the singlet field should grow with decreasing λ as M_Z/λ to ensure the correct breakdown of the electroweak symmetry. The increasing of s can be achieved either by decreasing κ or by raising m_S^2 and A_κ. Since there is no natural reason why m_S^2 and A_κ should be very large while all other soft SUSY breaking terms are left in the TeV range, the values of λ and κ are obliged to go to zero simultaneously so that their ratio remains unchanged.

Since, in the MSSM limit of the NMSSM, mixing between singlet states and neutral components of the Higgs doublets vanish, the structures of the Higgs mass matrices are simplified. This allows one to obtain simple approximate solutions for the Higgs masses. The NMSSM Higgs sector involves two charged states with masses

$$
m_{H^\pm}^2 \approx m_A^2 + M_W^2\,, \tag{3}
$$

where $M_W = \dfrac{g}{2}v$ is the mass of the charged W–boson while

$$
m_A^2 = \frac{4\mu^2}{\sin^2 2\beta}\left(x - \frac{\kappa}{2\lambda}\sin 2\beta\right)\,, \qquad x = \frac{1}{2\mu}\left(A_\lambda + 2\frac{\kappa}{\lambda}\mu\right)\sin 2\beta\,. \tag{4}
$$

Also there are five neutral fields in the Higgs spectrum. If CP is conserved then two of them are CP–odd with masses

$$m_{A_1}^2 \approx -3\frac{\kappa}{\lambda}A_\kappa\mu\,, \qquad\qquad m_{A_2}^2 \approx m_A^2\,, \tag{5}$$

whereas three others are CP–even with masses

$$\begin{aligned} m_{h_1}^2 &\approx 4\frac{\kappa^2}{\lambda^2}\mu^2 + \frac{\kappa}{\lambda}A_\kappa\mu + \frac{\lambda^2 v^2}{2}x\sin^2 2\beta - \frac{2\lambda^2 v^2\mu^2(1-x)^2}{M_Z^2\cos^2 2\beta}\,, \\ m_{h_2,h_3}^2 &\approx \frac{1}{2}\left[m_A^2 + M_Z^2 \mp \sqrt{(m_A^2+M_Z^2)^2 - 4m_A^2 M_Z^2\cos^2 2\beta}\right]\,, \end{aligned} \tag{6}$$

where $M_Z = \dfrac{\bar{g}}{2}v$ is a Z–boson mass and $\bar{g} = \sqrt{g^2+g'^2}$. In Eqs.(3)–(6) we ignore the contribution of loop corrections to the Higgs masses. The terms of the order of $O(\lambda^2 v^2)$ and $O(\lambda\kappa v^2)$ are also omitted. The only exception is for the mass of the CP–even singlet field h_1 where we retain the two last terms proportional to $\lambda^2 v^2$, since they become significant when κ becomes very small compared to λ.

For appreciable values of κ/λ the Higgs spectrum presented above depends on four variables: m_A, $\tan\beta$, $\dfrac{\kappa}{\lambda}\mu$ and A_κ. The masses of MSSM–like Higgs bosons ($m_{H^\pm}, m_{A_2}, m_{h_2}$ and m_{h_3}) are defined by m_A and $\tan\beta$. As in the minimal SUSY model they grow if m_A is increased. At large values of m_A ($m_A^2 >> M_Z^2$) the masses of the charged, one CP-odd and one CP-even states are almost degenerate, while the SM-like Higgs boson mass attains its theoretical upper bound $M_Z|\cos 2\beta|$. Loop corrections from the top quark and its superpartners raise this upper limit up to $130-135\,\mathrm{GeV}$. The experimental constraints on the SUSY parameters obtained in the MSSM remain valid in the the NMSSM with small κ and λ. For example, non-observation of any neutral Higgs particle at the LEPII restricts $\tan\beta$ and m_A from below.

The combination of the NMSSM parameters $\dfrac{\kappa}{\lambda}\mu$ set the mass scale of singlet fields (m_{h_1} and m_{A_1}). Decreasing κ reduces their masses so that for $\dfrac{\kappa}{\lambda} \ll 1$ they can be the lightest particles in the Higgs boson spectrum. The parameter A_κ occurs in the masses of extra scalar and pseudoscalar with opposite sign, and is therefore responsible for their splitting. Too large a value of $|A_\kappa|$ pulls the mass-squared of either singlet scalar or singlet pseudoscalar below zero destabilizing the vacuum. An even stronger lower constraint on A_κ is found from the requirement that the vacuum be the global minimum. Together these requirements constrain A_κ and consequently the ratio $(m_{A_1}^2/m_{h_1}^2)$ from below and above

$$-3\left(\frac{\kappa}{\lambda}\mu\right)^2 \le A_\kappa\cdot\left(\frac{\kappa}{\lambda}\mu\right) \le 0\,, \qquad\qquad 0 \le \frac{m_{A_1}^2}{m_{h_1}^2} \le 9\,. \tag{7}$$

The main features of the NMSSM Higgs spectrum discussed above are retained when the couplings λ and κ increase. In this case Eqs.(3)–(6) provide some insight into the mass hierarchy of the NMSSM Higgs sector and qualitatively describe the dependence of the Higgs masses with respect to the variations of m_A, κ/λ, A_κ, μ and $\tan\beta$.

3. NMSSM with $\kappa = 0$

The analysis of the MSSM limit of the NMSSM reveals one of the main impediments in the study of the NMSSM Higgs sector — the large number of independent parameters. Indeed even in the limit κ, $\lambda \to 0$, when the number of variables parameterizing the spectrum at the tree level reduces drastically, the masses of the extra Higgs states take arbitrary values. Therefore it seems very attractive to take a step back to the simplest extension of the MSSM when $\kappa = 0$. Since the self interaction of the singlet fields no longer appears in the superpotential nor in the Higgs effective potential, there are only 4 parameters defining the masses and couplings of the Higgs bosons at the tree-level: $\lambda, \mu, \tan\beta$ and m_A (or x). For $\tan\beta \le 2.5$ and small values of λ ($\lesssim 0.1$) the predominant part of the NMSSM parameter space is excluded by unsuccessful Higgs searches; although the lightest Higgs boson may partially decouple and not be seen, the second lightest scalar would be SM-like and visible. Furthermore, non-observation of charginos at

LEPII restricts the effective μ-term from below: $|\mu| \geq 90 - 100\,\mathrm{GeV}$. Combining these limits one gets a useful lower bound on m_A at the tree level:

$$m_A^2 \gtrsim 9M_Z^2 x\,. \tag{8}$$

When $\lambda \to 0$ the Higgs boson masses are closely approximated by Eq.(3) and Eqs.(5)–(6) where κ and A_κ must be taken to be zero. In the considered limit the mass of the lightest CP–odd Higgs boson, which is predominantly a singlet pseudoscalar, vanishes. This is a manifestation of the enlarged $SU(2)\times[U(1)]^2$ global symmetry of the Lagrangian; the extra $U(1)$ (Peccei–Quinn) symmetry is spontaneously broken giving rise to a massless Goldstone boson (axion) [2]. The Peccei–Quinn (PQ) symmetry and its axion allows one to avoid the strong CP problem, eliminating the θ–term in QCD [3].

The lightest CP–even Higgs boson is also mainly a singlet field. As evident from Eq.(6), at large values of $\tan\beta$ or μ the mass–squared of the lightest Higgs scalar becomes negative if the auxiliary variable x differs too much from unity. Therefore vacuum stability localizes the auxiliary variable x to a rather narrow range

$$1 - \frac{M_Z|\cos 2\beta|}{m_A^0} < x < 1 + \frac{M_Z|\cos 2\beta|}{m_A^0}\,, \tag{9}$$

where $m_A^0 = 2\mu/\sin 2\beta$. For example, at $\mu = 100\,\mathrm{GeV}$ and $\tan\beta = 3$ the squared mass of the lightest Higgs scalar remains positive only if x lies between 0.78 to 1.22. From the definition of m_A, Eq.(4), we see that the tight bounds on the auxiliary variable x constrain the m_A to the vicinity of $\mu\tan\beta$, which is much larger than the Z-boson mass. As a result, a mass splitting occurs, where the heaviest CP-odd, CP-even Higgs and charged Higgs bosons have a mass rather close to $\mu\tan\beta$, while the SM–like Higgs h_2 has a mass of the order of M_Z.

Increasing the value of λ increases the lightest Higgs scalar mass and mixings between the MSSM–like Higgs bosons and singlet states. As before the masses of the heaviest states are almost degenerate and close to $m_A \approx \mu\tan\beta$. At the tree level the masses of the lightest and second lightest Higgs scalars vary within the limits [4]:

$$\begin{array}{rcccl} 0 & \leq & m_{h_1}^2 & \leq & \dfrac{\lambda^2}{2}v^2 x\sin^2 2\beta\,, \\ M_Z^2\cos^2 2\beta + \dfrac{\lambda^2}{2}v^2\sin^2 2\beta & \leq & m_{h_2}^2 & \leq & M_Z^2\cos^2 2\beta + \dfrac{\lambda^2}{2}v^2(1+x)\sin^2 2\beta\,. \end{array} \tag{10}$$

In Eq.(10) the value of λ must be smaller than about 0.7 to prevent the appearance of Landau pole up below the GUT scale M_X. Although the masses of the lightest Higgs scalars rise with growing λ the mass hierarchy in the Higgs spectrum is preserved, i.e. $m_{h_1}, m_{h_2} \ll m_A$. The couplings of the lightest CP–even Higgs states to the Z pair (g_{ZZi}) and to the axion and Z–boson (g_{ZA_1i}) obey the sum rules [4]

$$R_{ZZ1}^2 + R_{ZZ2}^2 \approx 1\,, \tag{11}$$

$$R_{ZA_11}^2 + R_{ZA_12}^2 \approx \frac{1}{4}\left(\frac{\lambda v}{m_A^0}\right)^4 \cos^2 2\beta\,. \tag{12}$$

where R_{ZZi} and R_{ZA_1i} are normalized couplings defined by $g_{ZZi} = \dfrac{\bar{g}}{2}M_Z \times R_{ZZi}$ and $g_{ZA_ji} = \dfrac{\bar{g}}{2}\times R_{ZA_ji}$.

Searches for massless pseudoscalars and light scalar particles at accelerators as well as the study of their manifestations in astrophysics and cosmology rule out almost the entire parameter space of the NMSSM with $\kappa = 0$. A particularly stringent constraint emerges from stellar-evolution [5]–[6]. Since axions interact with electrons, nucleons and photons with strength inversely proportional to the axion decay coupling $f_a \sim s$, they are produced during the process of star cooling. To evade the modification of stellar-evolution beyond observational limits one must impose a lower limit on f_a and $s > 10^9\,\mathrm{GeV}$ [7].

For large values of $\tan\beta$ the restrictions on s are even stronger. The axion is accompanied by the lightest scalar Higgs boson (saxion) which has a mass less than $10\,\mathrm{KeV}$ for $\tan\beta > 10$ and $\mu < 200\,\mathrm{GeV}$, see Eq.(10). This light scalar can be also produced during the cooling of globular–cluster stars significantly affecting their evolution if the scalar–electron coupling g_{Xe} is above $1.3\cdot 10^{-14}$ [6]. Since $g_{h_1e} \sim m_e/s$ this translates into a lower bound on the scale of PQ–symmetry breaking $f_a \sim s > 10^{11}\,\mathrm{GeV}$. Cold dark matter is composed of both axions and saxions while the supersymmetric partner to the axion, the axino (the lightest neutralino), is so light that it does not contribute [8].

The constraints on the vacuum expectation value of the singlet field restrict λ to be less then $10^{-6}\,(10^{-9})$ for $s > 10^9\,\mathrm{GeV}(10^{11}\,\mathrm{GeV})$. The smallness of λ could be caused by certain discrete and gauge symmetries which forbid the operator $\lambda S(H_1H_2)$ at the renormalizable level but which permit a similar non–renormalizable operator involving additional singlets, resulting in an effective λ proportional to the ratio of the vacuum expectation values of the new singlet fields and the Planck scale [9]. It has also been shown that the interactions of S with other extra singlet fields result in the stabilization of the Higgs scalar potential which otherwise has a direction unbounded from below when $\kappa = 0$ and $m_S^2 < 0$. Moreover, this axion could play the role of an inflaton field [10].

For tiny values of λ, the decoupling of new singlet states prevents their observation at future colliders. Thus if the NMSSM with unbroken global PQ–symmetry is realized in nature, only MSSM–like Higgs bosons will be discovered in the near future. However, the strong correlation between the masses of the heaviest Higgs bosons and $\mu\tan\beta$ revealed in the tree level analysis (see Eq.(9)) does not take place in the MSSM but must be observed here(see also [11]). The inclusion of loop corrections does not change the qualitative pattern of the Higgs spectrum and does not enlarge the allowed range of x (or m_A). Loop corrections only slightly shift the admissible range of the variable x which shrinks with increasing μ and $\tan\beta$ in compliance with Eq.(9). In order to demonstrate the correlation between m_A and $\mu\tan\beta$ we examined 10^6 different scenarios, with m_A and $\tan\beta$ chosen randomly between 0 to 6 TeV and 3 to 30 respectively. We calculated the one-loop mass spectrum and, for every scenario with a stable vacuum, plotted a single point on the m_A–$\tan\beta$ plane of Fig.1. We discarded scenarios with unstable vacua. It is immediately evident that the physically acceptable scenarios all lie within a small area around $m_A \approx \mu\tan\beta$ [7]. Since the positivity or negativity of $m_{h_1}^2$ is independent of s (or λ), the NMSSM with $\kappa = 0$ is ruled out for all large values of the singlet expectation value if after measuring μ and $\tan\beta$ at future accelerators, the heavy pseudoscalar mass is not found to lie close to $\mu\tan\beta$. Alternatively, if the mass prediction were found to hold, it would provide an indirect evidence for the PQ–symmetric NMSSM as a solution to the strong CP problem and for the axion and saxion as a source of dark matter.

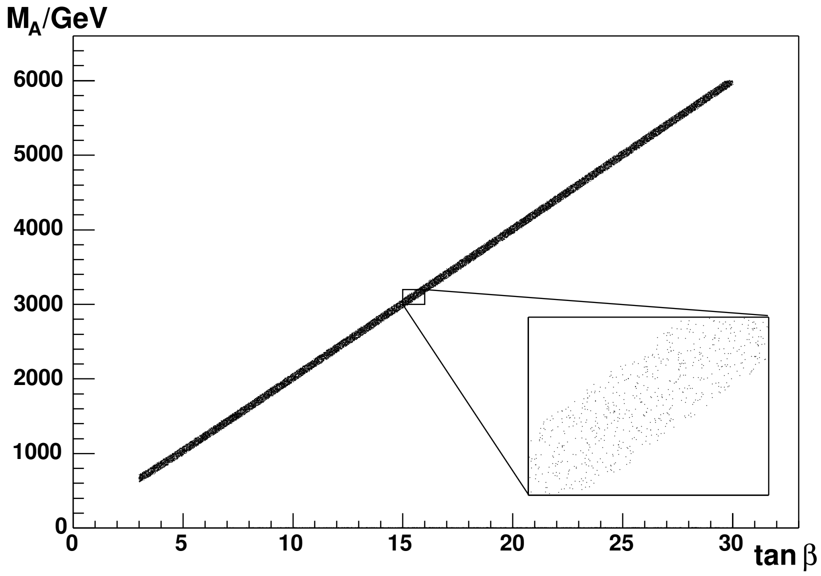

Figure 1: *The distribution of scenarios with physically acceptable vacua, with M_A chosen randomly between* 0 *and* 6 *TeV,* $\tan\beta$ *chosen randomly between* 3 *and* 30 *and* $\mu = 200\,GeV$. *The blow-up allows individual scenario points to be seen.*

4. Slight breaking of the PQ–symmetry

If one wants to avoid the introduction of a new intermediate scale that arises in the NMSSM with $\kappa = 0$ when the astrophysical limits on the couplings of the lightest scalar and pseudoscalar particles are applied, one must break the Peccei–Quinn symmetry. For a discussion of the possible origins of this symmetry breaking in the NMSSM see Refs.[11]-[12]. Here we assume that the violation of the Peccei–Quinn symmetry is caused by a non–zero value of κ. Moreover we restrict our consideration to small values of κ when the PQ–symmetry is only slightly broken. To be precise we consider values of κ that do not greatly change the vacuum energy density:

$$\kappa \lesssim \lambda^2 . \tag{13}$$

If $\kappa \gg \lambda^2$ then the terms $\kappa^2|S|^4$ and $\dfrac{\kappa}{3}A_\kappa S^3$ in the Higgs effective potential (2) become much larger than $|\mu|^4 \sim M_Z^4$ increasing the absolute value of the vacuum energy density significantly.

For small values of λ the approximate formulae Eqs.(5)–(6) obtained in section 2 remain valid. However, breaking the PQ–symmetry gives the lightest CP–odd Higgs an appreciable mass that allows it to escape the strong astrophysical constraints previously outlined. One must ensure that the value of κ is large enough for the lightest scalar and pseudoscalar to escape the exclusion limits of LEP, but it is still physically reasonable to only slightly break the PQ–symmetry, as defined by Eq.(15).

Indeed, for the appreciable values of κ and λ this slight breaking of the Peccei–Quinn symmetry may arise naturally from their renormalization group (RG) flow from M_X to M_Z [14]. While the values of the Yukawa couplings at the Grand Unification scale grow, the region where the solutions of the RG equations are concentrated at the electroweak scale shrinks and they are focused near the quasi fixed point [13]:

$$h_t(M_t) \approx 1.103, \qquad \lambda(M_t) \approx 0.514, \qquad \kappa(M_t) \approx 0.359 . \tag{14}$$

This point appears as a result of intersection of the Hill-type effective surface with the invariant line that connects the stable fixed point in the strong Yukawa coupling limit with the infrared fixed point of the NMSSM renormalization group equations. The requirement of perturbativity up to the Grand Unification scale provides stringent restrictions on the values of $\lambda(M_t)$ and $\kappa(M_t)$

$$\lambda^2(M_t) + \kappa^2(M_t) < 0.5 . \tag{15}$$

In order to obtain a realistic spectrum, one must of course include the leading one–loop corrections from the top and stop loops. We performed this exercise numerically [14] and present in Fig.2 the mass spectrum as a function of m_A for the parameters $\lambda = 0.3$, $\kappa = 0.1$, $\mu = 157\,\mathrm{GeV}$, $\tan\beta = 3$ and $A_\kappa = -60\,\mathrm{GeV}$. One sees that most of the structure outlined above is retained. The heaviest scalar and pseudoscalar are approximately degenerate with the charged Higgs boson and track m_A. The second lightest scalar is of the order of the Z-boson mass plus radiative corrections, mimicking the lightest scalar of the MSSM. The breaking of the PQ–symmetry has raised the masses of the lightest scalar and pseudoscalar to values which agree very well with the approximate expression of Eqs.(5)-(6)[1]. Also notice that the vacuum stability prevents having very high or very low values of m_A (or x) but now the allowed range has increased significantly, permitting m_A to substantially deviate from $\mu\tan\beta$.

One might expect that such a light Higgs scalar should already be ruled out by LEP, but this is not the case [15]. The reduced coupling to the Z-boson allows for a singlet like scalar substantially below the current SM LEP bounds. Indeed, LEP limits have been included in Fig.2 as a shaded area: for this parameter choice, values of m_A in the shaded region either provide a scalar Higgs boson which would have been seen at LEP or have an unstable vacuum. There is a substantial range in m_A, once more around the value $m_A \approx \mu\tan\beta$, where the Higgs scalar remains undetected. In this way, the mass hierarchy between the lighter Higgs bosons, around the electroweak scale, and the heavier Higgs bosons, at around $\mu\tan\beta$, is maintained. Since the coupling of the lightest scalar to the Z-boson must necessarily be suppressed in this region to avoid detection at LEP, the sum rule of Eq.(11) tells us that the couplings of the second lightest scalar will be similar to those of the lightest scalar in the MSSM. It is interesting to note that

[1] The agreement with tree-level expressions is good because the singlet nature of the new fields suppresses loops corrections.

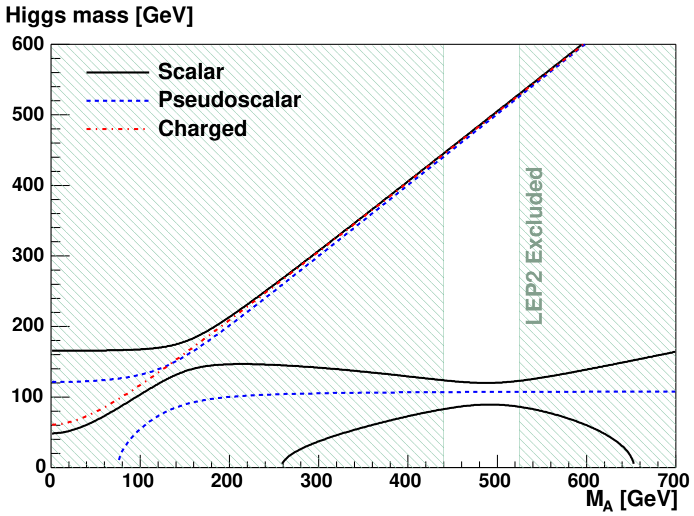

Figure 2: *The dependence of the Higgs boson masses on m_A for $\lambda = 0.3$, $\kappa = 0.1$, $\mu = 157\,GeV$, $\tan\beta = 3$ and $A_\kappa = -60\,GeV$. Solid, dashed and dashed–dotted curves correspond to the one–loop masses of the CP-even, CP-odd and charged Higgs bosons respectively. All masses are given in GeV.*

this light scalar would be very difficult to see at the LHC since it will principally decay hadronically, presenting a signal which is swamped by huge QCD backgrounds. According to Eq.(12), the coupling of the pseudoscalar to the Z-boson is always rather small. Nevertheless, if these light Higgs bosons could be seen, one would have a definitive signature of the NMSSM, even without observing the heavier states.

5. Conclusions

We have studied the Higgs sector of the NMSSM with exact and slightly broken Peccei–Quinn symmetry. In the PQ–symmetric NMSSM astrophysical observations exclude any choice of the parameters unless one allows s to be enormously large ($> 10^9 - 10^{11}$ GeV). These huge vacuum expectation values of the singlet field can be consistent with the electroweak symmetry breaking only if the coupling λ is extremely small $10^{-6} - 10^{-9}$. Such tiny values of λ may arise from non–renormalizable operators. In this limit the main contribution to the cold dark matter density comes from axion and saxion contributions while that of the lightest supersymmetric particle, the axino, is negligible.

If the PQ–symmetry is exact or only slightly broken, vacuum stability and LEP exclusion require parameters which cause a splitting in the NMSSM Higgs spectrum. The heaviest scalar, heaviest pseudoscalar and charged Higgs bosons are approximately degenerate with masses around $m_A \approx \mu \tan\beta$. The other three neutral states are considerably lighter. The masses of the lightest scalar and pseudoscalar, which are predominantly singlet fields, are governed by the combination of parameters $\dfrac{\kappa}{\lambda}\,\mu$. The SM–like Higgs boson mass remains at the electroweak scale.

In the limit of vanishing λ and κ the extra CP–even and CP–odd singlet states decouple from the rest of the spectrum and become invisible. However in the case of exact PQ–symmetry with $\kappa = 0$ (or very slightly broken PQ–symmetry with $\kappa \ll \lambda^2$) the stability of the physical vacuum constrains the masses of the heavy Higgs bosons in the vicinity of $m_A \approx \mu\,\tan\beta$. The strong correlation between m_A and $\mu \tan\beta$ coming from the dark sector of the NMSSM gives a unique "smoking gun" for distinguishing this model from the MSSM even if no extra Higgs states are discovered.

For appreciable values of λ and κ the slight breaking of the PQ–symmetry can be caused by the NMSSM renormalization group flow. Increasing λ increases the mixing between the light CP–even Higgs bosons, while increasing κ increases the masses of the predominantly singlet states. For small values of κ, one can have a light scalar Higgs boson which would not have been seen at LEP. Although the range of m_A allowed by vacuum stability increases significantly, one is still required to have $m_A \approx \mu \tan\beta$ in order to avoid the LEP constraints, leading to a mass splitting between the light and heavy Higgs bosons. Observing two light scalars and one pseudoscalar Higgs but no charged Higgs boson at future colliders would yield another opportunity to differentiate the NMSSM with a slightly broken PQ–symmetry from the MSSM even if the heavy Higgs states are inaccessible.

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