The motivations for the NMSSM are reviewed, and possible unconventional signals for Higgs and sparticle production at the LHC are discussed. In the presence of a light pseudoscalar, the SM-like Higgs scalar can decay dominantly into a 4-tau final state. In the fully constrained NMSSM with mSUGRA-like soft SUSY breaking terms, the correct dark matter relic density is obtained for a singlino-like LSP which modifies considerably all sparticle decay chains.

1. The NMSSM

The Next-to-Minimal Supersymmetric Standard Model (NMSSM) addresses the so-called \( \mu \)-problem of the MSSM [1], whose origin we describe below:

Any supersymmetric extension of the Standard Model (SM) generalizes in a unique way – as dictated by supersymmetry – the interactions involving dimensionless gauge- and Yukawa couplings, whereas the electroweak scale originates from the softly supersymmetry breaking mass terms and trilinear interactions (of the order \( M_{\text{SUSY}} \)). The MSSM, however, requires the introduction of a supersymmetric (SUSY) mass term for the Higgs multiplets, the so-called \( \mu \)-term: both complex Higgs scalars \( H_u \) and \( H_d \) of the MSSM have to be components of chiral superfields which contain, in addition, fermionic \( SU(2) \)-doublets \( \psi_u \) and \( \psi_d \). Some of the \( SU(2) \)-components of \( \psi_u \) and \( \psi_d \) are electrically charged. Together with with the fermionic superpartners of the \( W^\pm \) bosons, they constitute the so-called chargino sector (two charged Dirac fermions) of the SUSY extension of the SM. Due to the fruitless searches for a chargino at LEP, the lighter chargino has to have a mass above \( \sim 100 \) GeV [2].

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Analysing the chargino mass matrix, this lower limit implies that a Dirac mass \( \mu \) for \( \psi_u \) and \( \psi_d \) – for arbitrary other parameters – has to satisfy the constraint \( |\mu| \gtrsim 100 \) GeV. A Dirac mass term is not among the soft supersymmetry breaking mass terms, hence \( \mu \) has to be a supersymmetric mass term for the Higgs multiplets.

In addition, an analysis of the Higgs potential shows that a non-vanishing soft SUSY breaking term \( B\mu H_u H_d \) is a necessary condition so that both neutral components of \( H_u \) and \( H_d \) are non-vanishing at the minimum. This, in turn, is required in order to generate masses both for up-type quarks and down-type quarks by the Higgs mechanism. The numerical value of \( B\mu \) should be roughly of the order of the electroweak scale (\( M_Z^2 \)).

However, \( |\mu| \) must not be too large: the Higgs potential must be unstable at its origin \( H_u = H_d = 0 \) in order to generate the electroweak symmetry breaking. Whereas negative soft SUSY breaking mass terms for \( H_u \) and \( H_d \) of the order of the SUSY breaking scale \( M_{\text{SUSY}} \) can generate such a desired instability, the \( \mu \)-induced masses squared for \( H_u \) and \( H_d \) are always positive, and must not dominate the negative soft SUSY breaking mass terms. Consequently the \( \mu \) parameter must obey \( |\mu| \lesssim M_{\text{SUSY}} \). Hence, both “natural” values \( \mu = 0 \) and \( \mu \) very large (\( \sim M_{\text{GUT}} \) or \( \sim M_{\text{Planck}} \)) are ruled out, and the need for an
explanation of $\mu \approx M_{\text{SUSY}}$ is the $\mu$-problem.

Within the NMSSM, $\mu$ is generated in a way similar to the quark- and lepton masses with the help of a vacuum expectation value (vev) of a scalar field: one introduces a Yukawa coupling $\lambda$ of the higgsinos $\psi_u$ and $\psi_d$ to a scalar field $S$ (a gauge singlet, since the $\mu$-parameter carries no gauge quantum numbers), and arranges that the vev $\langle S \rangle$ is of the order of $M_{\text{SUSY}}$. This is easy to obtain with the help of soft SUSY breaking negative masses squared (or trilinear couplings) of the order of $M_{\text{SUSY}}$ for $S$; then, $M_{\text{SUSY}}$ is the only scale in the theory. In this sense, the NMSSM is the simplest supersymmetric extension of the SM in which the weak scale is generated by the supersymmetry breaking scale $M_{\text{SUSY}}$ only.

It should be mentioned that additional attractive features of the MSSM — the unification of the running gauge couplings at $M_{\text{GUT}} \sim 10^{16}$ GeV and the possibility to explain the dark matter relic density — remain unchanged.

### 1.1. The NMSSM Superpotential and Particle Content

The superpotential $W_{\text{NMSSM}}$ of the NMSSM depends on the additional gauge singlet superfield $S$, and is obtained from the superpotential $W_{\text{MSSM}}$ of the MSSM by the following substitutions:

$$W_{\text{MSSM}} = \mu H_u H_d + \ldots$$

$$\rightarrow W_{\text{NMSSM}} = \lambda S H_u H_d + \frac{1}{3} \mu S^3 + \ldots$$

and similarly for the soft SUSY breaking terms:

$$B_{\mu} H_u H_d + \ldots \rightarrow \lambda A_\lambda S H_u H_d + \frac{1}{3} \kappa A_\kappa S^3 + \ldots$$

The term $\sim \kappa$ in $W_{\text{NMSSM}}$ serves to stabilize the potential for the scalar singlet field $S$; a possibly negative SUSY breaking mass term $m_3^2 |S|^2$ (together with $\frac{1}{3} \kappa A_\kappa S^3 + \text{h.c.}$) can easily generate a vev $\langle S \rangle \sim M_{\text{SUSY}}$. Comparing $W_{\text{NMSSM}}$ and $W_{\text{MSSM}}$ one finds an effective $\mu$-term with $\mu_{\text{eff}} = \lambda \langle S \rangle$.

Discarding the Goldstone bosons, the particle content of the Higgs sector of the NMSSM consists in 3 CP-even neutral scalars $H_i$, 2 CP-odd neutral scalars $A_i$, 1 charged Higgs scalar $H^\pm$, and — together with the neutral gauginos — 5 neutralinos $\chi_i^0$. (The singlet-like states $H_S$, $A_S$ and $\chi_3^0$ mix with the $SU(2)$-doublets and the neutral gauginos; the decomposition of the eigenstates depend on the parameters $\lambda$, $\kappa$, $A_\lambda$, and $A_\kappa$.)

It is important to note that the lightest CP-odd scalar $A_1$ can be quite light, notably in the case of an approximate $R$-symmetry in the Higgs sector ($A_\lambda, A_\kappa \rightarrow 0$) or Peccei-Quinn-symmetry ($\kappa \rightarrow 0$), in which case it plays the role of a (pseudo-) Goldstone boson.

### 2. Phenomenological Aspects of the NMSSM

As a result of the new states and mixings in the Higgs and neutralino sectors, the corresponding phenomenology of the NMSSM can differ considerably from the MSSM. Already the LEP constraints from Higgs boson searches have to be interpreted anew.

#### 2.1. Lessons/Hints from LEP

The results of the four LEP experiments searching for a Higgs scalar decaying into $H \rightarrow b\bar{b}$, $\tau^+\tau^-$ (assuming SM branching fractions) have been combined by the LEP-Higgs Working Group [4] and are shown in Fig. 1. There, $\xi$ denotes the reduced coupling of a Higgs scalar to the $Z$ boson (compared to the coupling of the SM Higgs scalar), $\xi \equiv g_{HZZ}/g_{HZZ}^{SM}$. Shown are upper bounds on $\xi^2$ as function of a scalar Higgs mass $m_H$.

One can note a light excess of events for $m_H \sim 95 - 100$ GeV (of $\sim 2.3 \sigma$ statistical significance), which is difficult to explain in the SM. The NMSSM offers two possible explanations for this excess of events: i) a Higgs scalar with a mass of $\sim 95 - 100$ GeV can have a reduced coupling to the $Z$ boson ($\xi \lesssim 0.4 - 0.5$) due to its large singlet component; or ii) a Higgs scalar with a mass of $\sim 95 - 100$ GeV can have a reduced branching ratio into $b\bar{b}$, $\tau^+\tau^-$, since it decays dominantly into a pair of light CP-odd scalars with a $\text{BR}(H \rightarrow A_1 A_1) \sim 80 - 90 \%$. In the latter case, the coupling of $H$ to $Z$ bosons can be SM-like. In [4], it has been argued that this scenario allows to alleviate the “little finetuning problem”
of supersymmetric extensions of the SM (since $m_H > 114$ GeV is not required).

However, the LEP experiments have also searched for $H \to A_1 A_1 \to b\bar{b}$ [3], and the constraints are very strong for $m_H \sim 95 - 100$ GeV. On the other hand, if $M_{A_1}$ is below the $b\bar{b}$ threshold of $\sim 10.5$ GeV, $A_1$ would decay dominantly into $\tau^+ \tau^-$. LEP constraints on $H \to A_1 A_1 \to 4\tau$ impose no bounds for $m_H \sim 95 - 100$ GeV [3]; hence $M_{A_1} \lesssim 10.5$ GeV is an attractive scenario.

2.2. Lessons/Hints from radiative $\Upsilon$ decays

For $M_{A_1} \lesssim 10.5$ GeV, decays $\Upsilon(nS) \to A_1 + \gamma$ are kinematically possible. The corresponding branching ratios depend on $M_{A_1}$, and the coupling $g_{A_1 b\bar{b}}$ of $A_1$ to $b$-quarks. It is useful to define a reduced coupling $X_d \equiv g_{A_1 b\bar{b}}/g_{H b\bar{b}}^{SM}$, and for large $\tan \beta$ one can obtain $X_d > 1$ in the NMSSM.

In the range $M_{A_1} \lesssim 9$ GeV where $A_1 \to \tau\tau$ (or $A_1 \to \mu\mu$) would be dominant, the CLEO collaboration has searched for $\Upsilon(1S) \to A_1 + \gamma \to \tau\tau + \gamma$ decays [8]. No signal has been observed, which implies upper bounds on $X_d$ dependent on $M_{A_1}$.

The BABAR collaboration has recently observed an $\eta_b$-like state with a mass of $\sim 9.39$ GeV in $\Upsilon(2S)$ and $\Upsilon(3S)$ radiative decays [7,8]. For large $g_{A_1 b\bar{b}}$, the $\eta_b(nS) b\bar{b}$ bound states ($n = 1, 2, 3$) can mix with $A_1$, since these states have the same quantum numbers $[9,10,11,12]$, and the mixing elements/eigenvalues of the $\eta_b(nS) - A_1$ mass matrix can be determined in terms of $X_d$ and $M_{A_1}$ [10,11,12].

The measured mass by BABAR of the $\eta_b$-like state must correspond to one of these eigenvalues. On the one hand, this implies that $M_{A_1}$ cannot be equal to $9.39$ GeV, if the mixing (which is proportional to $X_d$) is large: for a large mixing, an eigenvalue of a $2 \times 2$ matrix (of the $\eta_b(1S) - A_1$ system) cannot coincide with one of its diagonal elements. This reasoning implies an upper bound on $X_d$ for $M_{A_1}$ near $9.39$ GeV. In Fig. 2 we show upper bounds on $X_d$ in the NMSSM from CLEO (black), the muon anomalous magnetic moment $a_\mu$ (blue), $B$-physics (green) and constraints due to the measured $\eta_b(1S)$ mass by Babar as a red line (from [11]).

On the other hand, the average value for the hyperfine splitting $E_{hfs}(1S) = m_\Upsilon(1S) - m_{\eta_b(1S)}$ found by BABAR [8] is somewhat large:

$$E_{hfs}(1S) = 69.9 \pm 3.1 \text{ MeV} \quad (3)$$

This result can be compared to recent predictions from perturbative QCD: $E_{hfs}(1S) = 44\pm11$ MeV [13] and $E_{hfs}(1S) = 39\pm14$ MeV [14].

Whereas an explanation for the discrepancy between [3] and perturbative QCD is not excluded at present, the difference might be ascribed to the mixing of the $\eta_b(1S)$ state with a light pseudoscalar Higgs $A_1$, if $M_{A_1}$ is somewhat above $9.4$ GeV. (For $M_{A_1} > 9$ GeV, $X_d$ is bounded from above just by the red line in Fig. 2.) The assumption that the $\eta_b(1S)$-$A_1$ mixing explains the discrepancy between [3] and perturbative QCD, allows to determine $X_d$ and hence all eigenvalues and mixing angles in the $\eta_b(nS)-A_1$ system as function of $M_{A_1}$ [12]. Then, the masses of the states interpreted as $\eta_b(2S)$ and $\eta_b(3S)$ can also be modified. Furthermore, all $\eta_b(nS)$ states can acquire non-negligible branching ratios into

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig1.pdf}
\caption{Upper bound on $\xi^2$ as function of a scalar Higgs mass $m_H$, where $\xi$ denotes the coupling of the Higgs scalar to $Z$ bosons (normalized w.r.t. the SM Higgs boson); from [3].}
\end{figure}
Figure 2. Upper bounds on $X_d$ versus $M_{A_1}$ in the NMSSM. Indicated are constraints from $B$-physics as a green dashed line, constraints from $a_\mu$ as a blue dashed line, the bounds from CLEO on $BR (\Upsilon \to \gamma \tau \tau)$ as a black line and constraints due to the measured $\eta_b(1S)$ mass by Babar as a red line; from [11].

$\tau^+ \tau^-$ due to their mixing with $A_1$.

In Fig. 3 (from [12]), the masses of all 4 physical states (denoted by $\eta_i$, $i = 1 \ldots 4$) as functions of $M_{A_1}$ are shown together with the error bands; by construction, $m_{\eta_1} \equiv 9.39$ GeV and for clarity the assumed values for $m_{\eta_i(0)}$ (before mixing) are indicated as horizontal dashed lines. For $M_{A_1}$ not far above 9.4 GeV the effects of the mixing on the states $\eta_0^0(2S)$ and $\eta_0^0(3S)$ are negligible, but for larger $M_{A_1}$ the spectrum can differ considerably from the expectations from QCD-based quark models.

Future experiments at $B$ factories could verify this scenario in different ways: through the spectrum of the $\eta_b(nS)$-$A_1$-system beyond the $\eta_b(1S)$-state seen by BABAR, and/or through a violation of lepton universality in inclusive $\Upsilon$ decays $\Upsilon \to l^+ l^- + X$: from $\Upsilon \to \gamma + A_1 \to \gamma + \tau^+ \tau^-$ one would expect an excess in the $\tau^+ \tau^-$ final state [15].

2.3. Higgs Searches at the LHC with dominant $H \to A_1 A_1$ decays

If a SM-like Higgs boson $H$ decays dominantly into a pair of light pseudoscalars with $M_{A_1} < 10.5$ GeV (such that $A_1$ decays dominantly into $\tau^+ \tau^-$), the detection of $H$ will be quite difficult at the LHC: the final state from $H \to A_1 A_1 \to 4\tau$ will contain four $\tau$ leptons, which escape undetected and make it difficult to observe peaks in invariant masses corresponding to $A_1$ or $H$; two $\tau$ leptons at a time (originating from one $A_1$ boson) will be nearly collinear, and the average $p_T$ of the particles in the final state is quite low. Important backgrounds originate from $\Psi$ production and heavy flavour jets.

Up to now, the following proposals for $H$ searches at the LHC in this scenario have been made:

(i) In [16] it has been proposed to consider diffractive Higgs production ($pp \to pp + H$) in order to be sensitive to $H \to 4\tau$, which requires to install additional forward detectors. Using a track-based analysis in which all events with more than 6 tracks in the central region are discarded,
a viable signal seems possible after accumulating 300 fb$^{-1}$ of integrated luminosity.

(ii) Proposals for signals and cuts appropriate for the $A_1 A_1 \to 4\tau \to 2\mu + 2$ jets final state have been made in [17]; with 100 fb$^{-1}$ of integrated luminosity, the expected rates after cuts are $\sim 8 \times 10^3$ from $H$ production via vector boson fusion, and $\sim 10^3$ from $H$ production via Higgs Strahlung ($W^\pm \to H + W^\pm$) where one can trigger on a lepton from $W^\pm$ decays.

(iii) In [18], the subdominant $H \to A_1 A_1 \to 2\tau 2\mu$ final state (with $2\mu$ from direct $A_1$ decays) was discussed: in spite of the small branching fraction it was argued that, for $M_H \sim 102$ GeV and with $H$ being produced via gluon-gluon fusion, the Tevatron can see a signal over the background for an integrated luminosity $L\sim 10$ fb$^{-1}$, and the LHC already for $L\sim 1$ fb$^{-1}$.

Further details of current ATLAS and CMS studies of benchmark scenarios including the $H \to AA \to 4\tau$ final state can be found in [19].

3. The constrained NMSSM

As in the MSSM, one can assume universal soft SUSY breaking masses and trilinear couplings at the GUT or Planck scale, which is motivated by a gravitational (flavour blind) origin for these terms in minimal supergravity. Then, all soft SUSY breaking terms (including those involving the singlet in the NMSSM) are specified by universal gaugino masses $M_{1/2}$, universal scalar masses $m_0$ and universal trilinear couplings $A_0$ at a large scale. The constrained NMSSM (cNMSSM) has the same number of free parameters as the constrained MSSM (cMSSM): the parameters $\mu$, $B$ of the cMSSM are replaced by the Yukawa couplings $\lambda$, $\kappa$ of the cNMSSM.

The allowed ranges of the parameters $M_{1/2}$, $m_0$ and $A_0$ in the cMSSM have been widely discussed in the literature; it turns out that these ranges are very different in the cNMSSM [20-21]: in order to obtain a non-vanishing vev $\langle S \rangle \neq 0$ in the NMSSM, the soft SUSY breaking mass $m_3^2$ must not be large and positive. Since $m_3^2$ is hardly renormalized between the GUT and the weak scale, the same condition applies to $m_0^2$ in the cNMSSM. In the cMSSM, small values of $m_0$ (compared to $M_{1/2}$) lead to a charged stau ($\tilde{\tau}$) LSP, which is ruled out. In the cNMSSM, an additional singlino-like neutralino $\chi^0_1$ can have a mass below the $\tau$ mass and be the true LSP. In order to give the correct (not too large) dark matter relic density, $M_{\chi^0_1}$ must be just a few GeV below $M_\tau$, such that the $\chi^0_1$ relic density can be reduced via co-annihilation with the $\tilde{\tau}$ NLSP. The latter condition fixes $A_0$ in terms of $M_{1/2}$. Finally LEP constraints on the Higgs sector lead to an upper bound on $\lambda$ of $\lambda \lesssim 0.02$ [20-21].

Hence, the full Higgs and sparticle spectrum depends essentially on $M_{1/2}$ only. A preferred range for $M_{1/2}$ can be obtained by computing the muon anomalous magnetic moment $a_\mu$ as a function of $M_{1/2}$, and requiring that the supersymmetric contributions explain the deviation $\delta a_\mu \sim 3 \times 10^{-9}$ between the result of the E821 experiment at BNL [22] and the SM [23]. The result for $\delta a_\mu^{\text{SUSY}}$ is shown in Fig. 3 according to which values for $M_{1/2} \lesssim 1$ TeV are preferred, with $M_{1/2} \approx 500$ GeV within 1 $\sigma$.

Choosing $m_0 = 0$ for simplicity, the stau and neutralino spectrum as a function of $M_{1/2}$ in the cNMSSM is shown in Fig. 4. One finds

![Figure 4](image-url)
that, for $M_{1/2} \lesssim 400$ GeV, the lightest stau mass would fall below $\sim 100$ GeV and violate LEP constraints; hence one must require $M_{1/2} \gtrsim 400$ GeV. (As discussed above, the $\chi^0_1$ mass is just below the lightest stau mass in order to give the correct dark matter relic density.)

The Higgs spectrum as a function of $M_{1/2}$ in the cNMSSM is shown in Fig. 6. The heavy states $H_3$, $A_2$ and $H^\pm$ form a practically degenerate $SU(2)$ doublet, and the lighter CP-odd scalar $A_1$ is too heavy to be produced in $H_1$ decays in the cNMSSM. The nature of the lightest CP-even Higgs scalar depends on $M_{1/2}$: for $M_{1/2} \lesssim 640$ GeV, $H_1$ is dominantly singlet-like which allows to satisfy LEP constraints even for a mass well below 114 GeV. The next-to-lightest state $H_2$, for $M_{1/2} \lesssim 640$ GeV, is SM-like with a mass just above 114 GeV. For $M_{1/2} \gtrsim 640$ GeV, the situation is reversed: here $H_2$ is dominantly singlet-like, and the lightest CP-even Higgs scalar $H_1$ is SM-like with a mass up to 120 GeV for $M_{1/2} \to 1.5$ TeV.

An interesting scenario is possible for $M_{1/2} \sim 570$ GeV, where $M_{H_1} \sim 100$ GeV and the reduced coupling $\xi_{H_1 ZZ}$ of the singlet-like $H_1$ (for not too small $\lambda \sim 0.005$) is $\xi_{H_1 ZZ} \sim 0.3$ [21]: in this case a production of $H_1$ at LEP could explain the excess of events at LEP as discussed before. Simultaneously, $H_2$ with its mass just above 114 GeV could generate an even slighter excess of events in this mass range, which has equally been observed [3].

The consequences of the nearly pure singlino-like $\chi^0_1$ LSP in the cNMSSM on sparticle decay chains are important: first, each sparticle decay cascade will end up in the next-to-lightest $R$-odd particle, the charged $\tilde{\tau}$ in the present case. Only subsequently the $\tilde{\tau}$ will decay as $\tilde{\tau} \to \tau + \chi^0_1$, leading to at least one $\tau$ lepton per sparticle. Notably, if the $\tilde{\tau} - \chi^0_1$ mass difference is small and/or $\lambda$ is very small (in which case the couplings to $\chi^0_1$ become tiny), the $\tilde{\tau}$ width can become so small that the $\tilde{\tau}$ decay length becomes $\mathcal{O}(\text{mm-cm})$ [20,21]. Such displaced vertices can also appear in scenarios with gauge mediated supersymmetry breaking where, however, the decay lengths are rather of $\mathcal{O}(\text{m})$. Hence, dedicated simulations of the sparticle signatures at colliders in the cNMSSM will be required.

Figure 5. The stau and neutralino spectrum as a function of $M_{1/2}$ in the cNMSSM; from [20].

Figure 6. The Higgs spectrum as a function of $M_{1/2}$ in the cNMSSM; from [20].
4. Summary and Conclusions

Assuming that the SUSY breaking scale $M_{\text{SUSY}}$ generates the weak scale $\sim M_Z$ (i.e. that no dimensionful terms as $\mu$ are present in the superpotential), the NMSSM is the most natural supersymmetric extension of the Standard Model.

Its simplest version, the cNMSSM (with soft SUSY breaking terms from mSUGRA), could explain the $2.3 \sigma$ excess of events in Higgs searches at LEP, in contrast to the cMSSM. Due to $m_0 \ll M_1/2$, its sparticle spectrum would be very different from the cMSSM: the LSP $\chi_1^0$ is singlino-like, but every sparticle decay chain contains a $\tilde{\tau}$ decaying, in turn, into $\tau + \chi_1^0$. This last decay can lead to macroscopically displaced vertices.

In the general NMSSM, a light CP-odd scalar $A_1$ with $M_{A_1} < 10.5$ GeV could (i) explain the $2.3 \sigma$ excess of events at LEP; (ii) alleviate the “little finetuning problem” of supersymmetric extensions of the SM, and (iii) explain a low $\eta_b$ mass as measured by BABAR. But: this scenario would constitute a real challenge for Higgs searches both at the Tevatron and at the LHC!

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