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A supersymmetric Stueckelberg $\mathrm{U}(1)$ extension of the MSSM

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Abstract: A Stueckelberg extension of the MSSM with only one abelian vector and one chiral superfield as an alternative to an abelian extension with Higgs scalars is presented. The bosonic sector contains a new gauge boson Z' which is a sharp resonance, and a new CP-even scalar, which combines with the MSSM Higgs bosons to produce three neutral CP-even massive states. The neutral fermionic sector has two additional fermions which mix with the four MSSM neutralinos to produce an extended 6×6 neutralino mass matrix. For the case when the LSP is composed mostly of the Stueckelberg fermions, the LSP of the MSSM will be unstable, which leads to exotic decays of sparticles with many leptons in final states. Prospects for supersymmetry searches and for dark matter are discussed.

Keywords: Gauge Symmetry, Beyond Standard Model, Supersymmetric Standard Model.

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The Stueckelberg mechanism [1] generates a mass for an abelian gauge boson in a gauge invariant, renormalizable way, but without utilizing the Higgs mechanism [2]. Recently, an extension of the Standard Model [3] was proposed [4], where the electroweak gauge group $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ was enlarged by an extra $\mathrm{U}(1)_X$ gauge sector, whose gauge field couples to an axionic scalar field in the way of a Stueckelberg coupling, giving rise to a massive neutral gauge boson Z'. This is the simplest realization of the Stueckelberg mechanism in a minimally extended Standard Model. The model predicts the presence of a sharp resonance in e^+e^- collision, which is distinctly different from other Z' extensions [5] that appear for instance in grand unified theories or string and brane models (but see also [6]). In this Letter we further extend this technique and obtain a Stueckelberg extension of the minimal supersymmetric Standard Model (StMSSM). Since the Stueckelberg lagrangian can only accommodate abelian gauge invariance it is mostly sufficient to concentrate on the abelian subsector consisting of the hypercharge $\mathrm{U}(1)_Y$ and the new $\mathrm{U}(1)_X$. The minimal set of degrees of freedom that has to be added to the MSSM consists of the abelian vector multiplet C with components (C_μ, λ_C, D_C) and the chiral multiplet S with components $(\chi, \rho + ia, F)$, which we call the Stueckelberg multiplet. Before supersymmetry breaking, the two combine into a single massive spin one multiplet, and mix with hypercharge and the 3-component of iso-spin, as we shall see. Later we will also include a hidden sector whose matter fields may carry charge under the $\mathrm{U}(1)_X$. For the Stueckelberg lagrangian we choose

$$\mathcal{L}_{\rm St} = \int d\theta^2 d\bar{\theta}^2 \left[M_1 C + M_2 B + S + \bar{S} \right]^2 , \tag{1}$$

where B is the $\mathrm{U}(1)_Y$ vector multiplet with components (B_μ, λ_B, D_B). The gauge transformations under $\mathrm{U}(1)_Y$ and $\mathrm{U}(1)_X$ are

$$\begin{aligned} \delta_Y B &= \Lambda_Y + \bar{\Lambda}_Y , \qquad & \delta_Y S &= -M_2 \Lambda_Y , \\ \delta_X C &= \Lambda_X + \bar{\Lambda}_X , \qquad & \delta_X S &= -M_1 \Lambda_X . \end{aligned} \tag{2}$$

The quantities M_1, M_2 are "topological" [7] input parameters of the model. We define C in Wess-Zumino gauge as

$$C = -\theta\sigma^\mu\bar{\theta} C_\mu + i\theta\theta\bar{\theta}\bar{\lambda}_C - i\bar{\theta}\bar{\theta}\theta\lambda_C + \frac{1}{2}\theta\theta\bar{\theta}\bar{\theta} D_C , \tag{3}$$

and similarly B, while S is

$$S = \frac{1}{2}(\rho + ia) + \theta\chi + i\theta\sigma^\mu\bar{\theta}\frac{1}{2}(\partial_\mu\rho + i\partial_\mu a) + \theta\theta F + \frac{i}{2}\theta\theta\bar{\theta}\bar{\sigma}^\mu\partial_\mu\chi + \frac{1}{8}\theta\theta\bar{\theta}\bar{\theta}(\Box\rho + i\Box a) . \tag{4}$$

Its complex scalar component contains the axionic pseudo-scalar a, which is the analogue of the real pseudo-scalar that appears in the non-supersymmetric version in [4]. We write $\mathcal{L}_{\rm St}$ in component notation (e.g. [8])

$$\begin{aligned} \mathcal{L}_{\rm St} = & -\frac{1}{2}(M_1 C_\mu + M_2 B_\mu + \partial_\mu a)^2 - \frac{1}{2}(\partial_\mu\rho)^2 - i\chi\sigma^\mu\partial_\mu\bar{\chi} + 2|F|^2 + \\ & + \rho(M_1 D_C + M_2 D_B) + \bar{\chi}(M_1\bar{\lambda}_C + M_2\bar{\lambda}_B) + \chi(M_1\lambda_C + M_2\lambda_B) . \end{aligned} \tag{5}$$

For the gauge fields we add the kinetic terms

$$\mathcal{L}_{\rm gkin} = -\frac{1}{4}(B_{\mu\nu}B^{\mu\nu} + C_{\mu\nu}C^{\mu\nu}) - i\lambda_B\sigma^\mu\partial_\mu\bar{\lambda}_B - i\lambda_C\sigma^\mu\partial_\mu\bar{\lambda}_C + \frac{1}{2}(D_B^2 + D_C^2) \tag{6}$$

with $C_{\mu\nu} = \partial_\mu C_\nu - \partial_\nu C_\mu$, and similarly for B. For the matter fields (quarks, leptons, Higgs scalars, plus hidden sector matter) chiral superfields with components (f_i, z_i, F_i) are introduced, defined similar to S but f_i with extra factor $\sqrt{2}$ and z_i without the extra $\frac{1}{2}$ of eq. (4). Their lagrangian is standard,

$$\begin{aligned}\mathcal{L}_{\rm matt} \;=\; & -|D_\mu z_i|^2 - if_i\sigma^\mu D_\mu\bar{f}_i - \sqrt{2}\left(ig_Y Y z_i\bar{f}_i\bar{\lambda}_B + ig_X Q_X z_i\bar{f}_i\bar{\lambda}_C + \text{h.c.}\right) + \\ & + g_Y D_B(\bar{z}_i Y z_i) + g_X D_C(\bar{z}_i Q_X z_i) + |F_i|^2\,,\end{aligned} \tag{7}$$

where (Y, g_Y) and (Q_X, g_X) are the charge operators and coupling constants of hypercharge and $\mathrm{U}(1)_X$, and $D_\mu = \partial_\mu + ig_Y Y B_\mu + ig_X Q_X C_\mu$ the gauge covariant derivative. One further has the freedom to add Fayet-Iliopoulos terms

$$\xi_B D_B + \xi_C D_C\,, \tag{8}$$

which leads to

$$-D_B = \xi_B + M_2\rho + g_Y\sum_i \bar{z}_i Y z_i\,, \qquad -D_C = \xi_C + M_1\rho + g_X\sum_i \bar{z}_i Q_X z_i\,. \tag{9}$$

We find it convenient to replace the two-component Weyl-spinors by real four-component Majorana spinors, which we label as $\psi_S = (\chi_\alpha, \bar{\chi}^{\dot{\alpha}})^T$, and $\lambda_X = (\lambda_{C\alpha}, \bar{\lambda}_C^{\dot{\alpha}})$ and $\lambda_Y = (\lambda_{B\alpha}, \bar{\lambda}_B^{\dot{\alpha}})$ for the two gauginos, and similarly for the matter fields as well. Using identities

$$\begin{aligned}\chi\lambda_C + \bar{\chi}\bar{\lambda}_C &= \bar{\psi}_S\lambda_X\,, \\ \chi\lambda_C - \bar{\chi}\bar{\lambda}_C &= \bar{\psi}_S\gamma_5\lambda_X\,, \\ \chi\sigma^\mu\partial_\mu\bar{\chi} - (\partial_\mu\chi)\sigma^\mu\bar{\chi} &= \bar{\psi}_S\gamma^\mu\partial_\mu\psi_S\,,\end{aligned} \tag{10}$$

the total lagrangian of our extension of the MSSM takes the form

$$\mathcal{L}_{\rm StMSSM} = \mathcal{L}_{\rm MSSM} + \Delta\mathcal{L}_{\rm St} + \Delta\mathcal{L}_{\rm hidden}\,, \tag{11}$$

with

$$\begin{aligned}\Delta\mathcal{L}_{\rm St} \;=\; & -\frac{1}{2}(M_1 C_\mu + M_2 B_\mu + \partial_\mu a)^2 - \frac{1}{2}(\partial_\mu\rho)^2 - \frac{1}{2}(M_1^2 + M_2^2)\rho^2 - \\ & -\frac{i}{2}\bar{\psi}_S\gamma^\mu\partial_\mu\psi_S - \frac{1}{4}C_{\mu\nu}C^{\mu\nu} - \frac{i}{2}\bar{\lambda}_X\gamma^\mu\partial_\mu\lambda_X + M_1\bar{\psi}_S\lambda_X + M_2\bar{\psi}_S\lambda_Y - \\ & -\sum_i\Big[\,|D_\mu z_i|^2 - |D_\mu z_i|^2_{C_\mu=0} + \rho\Big(g_Y M_2(\bar{z}_i Y z_i) + g_X M_1(\bar{z}_i Q_X z_i)\Big) + \\ & \qquad + \frac{1}{2}g_X C_\mu\bar{f}_i\gamma^\mu Q_X f_i + \sqrt{2}g_X\left(iz_i Q_X\bar{f}_i\lambda_X + \text{h.c.}\right)\Big] \\ & -\rho(M_2\xi_B + M_1\xi_C) - \frac{1}{2}\Big[\xi_C + g_X\sum_i \bar{z}_i Q_X z_i\Big]^2\,.\end{aligned} \tag{12}$$

Here and in the following we assume, for simplicity, that all hidden sector fields are neutral under the MSSM gauge group, and vice versa, that all fields of the MSSM are neutral under the new $U(1)_X$.[1]

To clarify the particular properties of the StMSSM let us briefly compare to other extensions of the MSSM with an extra $U(1)_X$ gauge field, but which use Higgs fields to generate its mass. Most of these are immediately distinct, since they involve direct couplings between the new gauge field and the MSSM, i.e. the new gauge boson is not neutral under hypercharge and iso-spin. This imposes much stronger bounds on its mass, usually in the range of 1 TeV or larger. Such couplings arise for instance in left-right symmetric unified models after breaking to the electro-weak gauge group, e.g. in [9, 10]. These models usually involve many more degrees of freedom than the minimal Stueckelberg $U(1)_X$ extension we consider here. The Higgs model that would actually come closest to producing the identical effect as the StMSSM, a mass for a single extra abelian U(1) gauge boson, would consist of adding the $U(1)_X$ to the MSSM plus a Higgs chiral multiplet with charges under hypercharge and $U(1)_X$, say both charges +1, but otherwise neutral. Its action would be given by a single copy of (7), from which one can read that a vacuum expectation value v_H for the scalar component of the Higgs multiplet produces all the terms in (5) that involve M_1 or M_2, with the replacement $M_1 \to g_X v_H$ and $M_2 \to g_Y v_H$, though the total lagrangians do not match completely. This similarity is, however, deceiving, since a single charged chiral multiplet with a lagrangian of the standard form (7) and its implied gauge invariance, would contribute to the ABJ gauge anomaly via the usual triangle diagram and spoil anomaly cancellation for the hypercharge. Therefore, one is forced to add at least one more Higgs multiplet of opposite charge assignments to cancel the anomaly. For the Stueckelberg multiplet with its unusual gauge transformation (2) this problem does not arise, since there are no trilinear couplings of the form $g_Y B_\mu \chi \sigma^\mu \bar{\chi}$ in (5), and the Stueckelberg fermion χ has zero charge and does not contribute in a triangle diagram. As a conclusion, the minimal anomaly-free supersymmetric abelian Higgs model, which would be closest to the StMSSM, differs from the latter already at the level of the number of degrees of freedom. We also note in passing that because of eq. (2) the chiral superfield S cannot appear in the superpotential unlike the usual abelian extensions with Higgs scalars (e.g. a term Sh_1h_2 is not allowed here; compare to [11]). Thus the StMSSM appears really unique, not only in its theoretical foundation, but also in its predictions.

In addition to the soft supersymmetry breaking termns of the MSSM we also add a soft mass $\tilde{m}_X$ for the new neutral gaugino λ_X. In principle, one could also allow soft mass terms for ρ and ψ_S, but we leave them out as they are not crucial to our discussion. Finally, one has to add gauge fixing terms similar to the R_ξ gauge, which remove the cross-terms $M_2 B_\mu \partial^\mu a$ etc. together with the usual ones involving the Higgs doublets, see [8, 12]. This completes the lagrangian of our model. As is typically done for the MSSM, we further assume that the FI parameters ξ_B and ξ_C give subdominant contributions relative to other sources of supersymmetry breaking. In fact, for the purpose of the present analysis, we

[1]Therefore no coupling $g_Y B_\mu \bar{f}_i \gamma^\mu Y f_i$ appears in (11), and further the interaction $g_X C_\mu \bar{f}_i \gamma^\mu Q_X f_i$ vanishes for the fermions of the MSSM, and only involves hidden sector fermions. Under these assumptions $\Delta\mathcal{L}_{\text{hidden}}$ will not be relevant for our discussion.

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let ξ_B and ξ_C vanish. We will discuss the modifications due to non-zero FI parameters elsewhere. Note that if $\xi_B = \xi_C = 0$ at the tree-level, there is no contribution to these terms from loop diagrams, since $\mathrm{U}(1)_Y$ and $\mathrm{U}(1)_X$ are both anomaly-free [13].

We first concentrate on the neutral vector bosons, ordered (C_μ, B_μ, A^3_μ), where A^3_μ is the 3-component of the iso-spin. By giving two out of the three vector bosons masses, the Stueckelberg axion a plus one CP-odd component of the Higgs scalars decouple after gauge fixing, and disappear from the physical spectrum. In order to avoid a mass term for the photon it is required that the expectation values for all the scalars charged under $\mathrm{U}(1)_X$ vanish. Thus, we demand all $\langle \bar{z}_i Q_X z_i \rangle = 0$. This part of the supersymmetric model is then identical to the non-supersymmetric version, that was the subject of [4]. After spontaneous electro-weak symmetry breaking the 3×3 neutral vector boson mass matrix takes the form

$$\left[\begin{array}{c|cc} M_1^2 & M_1 M_2 & 0 \\ \hline M_1 M_2 & M_2^2 + \frac{1}{4}g_Y^2 v^2 & -\frac{1}{4}g_Y g_2 v^2 \\ 0 & -\frac{1}{4}g_Y g_2 v^2 & \frac{1}{4}g_2^2 v^2 \end{array}\right] . \tag{13}$$

Here, $v = 2M_{\mathrm{W}}/g_2 = (\sqrt{2}G_F)^{-\frac{1}{2}}$, M_{W} being the mass of the W-boson, G_F the Fermi constant. The matrix has a zero eigen value which corresponds to the photon, and two massive eigen states, the Z and Z$'$ bosons. As was pointed out in [4], it is most convenient to use the two quantities $M = (M_1^2 + M_2^2)^{1/2}$ and $\delta = M_2/M_1$ to parametrize the extension of the Standard Model. Experimental bounds then only impose $\delta < 10^{-2}$, and $M > 150\,\mathrm{GeV}$, which makes the Z$'$ rather light, and a very sharp resonance in e^+e^- annihilation.[2] For further details on the gauge boson sector, see [4].

We next turn to the scalars of the StMSSM. The total scalar potential involves the two Higgs doublets h_i and ρ, and is given by

$$\mathcal{V} = \frac{1}{2}(m_1^2 - \rho g_Y M_2)|h_1|^2 + \frac{1}{2}(m_2^2 + \rho g_Y M_2)|h_2|^2 + (m_3^2 \epsilon_{ij} h_i h_j + \mathrm{h.c.}) + \frac{1}{2}(M_1^2 + M_2^2)\rho^2 + \mathcal{V}_D \,, \tag{14}$$

where $\mathcal{V}_D$ is the standard MSSM D-term potential that follows immediately from (7). Further, for $i = 1, 2$ we defined $m_i^2 = m_{h_i}^2 + |\mu|^2$, and $m_3^2 = |\mu B|$, μ being the Higgs mixing parameter and B the soft bilinear coupling. To introduce (real) expectation values for the neutral components of the Higgs fields and ρ we replace $h_1^0 \to (v_1 + h_1^0)/\sqrt{2}$, $h_2^0 \to (v_2 + h_2^0)/\sqrt{2}$, with $\tan(\beta) = v_2/v_1$, and $\rho \to v_\rho + \rho$, with

$$g_Y M_2 v_\rho = \frac{M_{\mathrm{W}}^2 M_2^2}{m_\rho^2} \tan^2(\theta_W) \cos(2\beta) \,, \tag{15}$$

where $m_\rho^2 = M_1^2 + M_2^2 = M^2$, and θ_W is the unmodified weak mixing angle with $\tan(\theta_W) = g_Y/g_2$. Due to $|g_Y M_2 v_\rho| < 10^{-4} M_Z^2$, the modification

$$\frac{(g_2^2 + g_Y^2)(v_1^2 + v_2^2)}{8} = \frac{m_1^2 - m_2^2 \tan^2(\beta)}{\tan^2(\beta) - 1} + \frac{g_Y M_2 v_\rho}{\cos(2\beta)} \tag{16}$$

[2] In models with Stueckelberg couplings that derive from string theory, the mass scale M is naturally tied to the string scale, and thus M could lie in the electro-weak range for a string scale in the TeV region. Furthermore, as stressed in [14], the Stueckelberg mass term in D-brane models depends on moduli and can in principle be of any size. Of course, even when M becomes really large there would still be effects of the extra Stueckelberg $\mathrm{U}(1)_X$ as long as M_2/M_1 remains finite and fixed.

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of the electro-weak symmetry breaking constraints is unimportant. The StMSSM does not at all affect the mass of the CP-odd neutral scalar in the MSSM, which is

$$m_A^2 = -\frac{m_3^2}{\sin(\beta)\cos(\beta)} \,. \tag{17}$$

The three CP-even neutral scalars (h_1^0, h_2^0, ρ) mix with mass^2 matrix (s_β, $c_\beta = \sin(\beta)$, $\cos(\beta)$, $t_\theta = \tan(\theta_W)$)

$$\left[\begin{array}{cc|c} M_0^2 c_\beta^2 + m_A^2 s_\beta^2 & -(M_0^2 + m_A^2) s_\beta c_\beta & -t_\theta c_\beta M_\mathrm{W} M_2 \\ -(M_0^2 + m_A^2) s_\beta c_\beta & M_0^2 s_\beta^2 + m_A^2 c_\beta^2 & t_\theta s_\beta M_\mathrm{W} M_2 \\ \hline -t_\theta c_\beta M_\mathrm{W} M_2 & t_\theta s_\beta M_\mathrm{W} M_2 & m_\rho^2 \end{array}\right]$$

where $M_0^2 = (g_2^2 + g_Y^2)(v_1^2 + v_2^2)/4 \simeq M_\mathrm{Z}^2$. One can organize the three eigen states as (H_1^0, H_2^0, ρ_S) such that in the limit $M_2/M_1 \to 0$, $(H_1^0, H_2^0, \rho_S) \to (H^0, h^0, \rho)$, where H^0 and h^0 are the heavy and the light CP-even neutral Higgs of the MSSM. The new real scalar ρ_S is dominantly ρ, but it also carries small components of H^0 and h^0. Although the off-diagonal terms proportional to $M_\mathrm{W} M_2$ can be larger than $(100\,\mathrm{GeV})^2$, the corrections to the mass eigen states through mixing are still under control, since the ratio $M_\mathrm{W} M_2/m_\rho^2$ remains small. For a very low mass scale $M_1 \sim 10^2\,\mathrm{GeV}$, ρ_S can be directly produced in the $J^{\mathrm{CP}} = 0^+$ channel. Thus there should be three resonances in the $J^{\mathrm{CP}} = 0^+$ channel in e^+e^- collisions in contrast to just two for the MSSM case. The decay of ρ_S into visible MSSM fields will be dominantly into $t\bar{t}$ (or $b\bar{b}$ if $m_{\rho_S} < 2m_t$) through the admixture of H^0 and h^0. The partial decay width can be estimated

$$\begin{aligned} \Gamma(\rho_S \to t\bar{t}\,) &= \frac{3m_{\rho_S}}{8\pi} \left[\frac{m_t S_{32}}{\sqrt{2} M_\mathrm{W} \sin(\beta)}\right]^2 \sqrt{1 - \frac{4m_t^2}{m_{\rho_S}^2}}\,, \\ \Gamma(\rho_S \to b\bar{b}\,) &= \frac{3m_{\rho_S}}{8\pi} \left[\frac{m_b S_{31}}{\sqrt{2} M_\mathrm{W} \sin(\beta)}\right]^2 \sqrt{1 - \frac{4m_b^2}{m_{\rho_S}^2}}\,. \end{aligned} \tag{18}$$

Here S_{ij} are the elements of the rotation matrix that diagonalizes the Higgs mass^2 matrix. One estimates that S_{32} and S_{31} are $\mathcal{O}(M_2/M_1) \sim 0.01$, and thus the ρ_S decay width will be in the range of MeV or less, similar to that of Z′ [4]. We estimate the production cross section for this new state to be not more than a few percent at best, of what one has for the other two CP even Higgs bosons of MSSM. Such a sharp resonance can escape detection unless a careful search is carried out [15]. The total decay width of ρ_S will be broadened, if it can decay into hidden sector matter through the much larger coupling $g_X M_1 \rho(\bar{z}_i Q_X z_i)$ in eq. (11). The StMSSM also modifies the D-term contribution to squark and slepton masses through the term $g_Y M_2 v_\rho(\bar{z}_i Y z_i)$ in (11), which is typically negligible due to eq. (15).

Finally, we come to the neutral fermions of the StMSSM. Instead of four neutralinos in the MSSM we now have six, consisting of the three gauginos, the two Higgsinos $\tilde{h}_i$, and the extra Stueckelberg fermion ψ_S, which we order as $(\psi_S, \lambda_X, \lambda_Y, \lambda_3, \tilde{h}_1, \tilde{h}_2)$. After spontaneous electro-weak symmetry breaking the 6×6 neutralino mass matrix in the above basis is given by

$$\left[\begin{array}{cc|cccc} 0 & M_1 & M_2 & 0 & 0 & 0 \\ M_1 & \tilde{m}_X & 0 & 0 & 0 & 0 \\ \hline M_2 & 0 & \tilde{m}_1 & 0 & -c_1M_0 & c_2M_0 \\ 0 & 0 & 0 & \tilde{m}_2 & c_3M_0 & -c_4M_0 \\ 0 & 0 & -c_1M_0 & c_3M_0 & 0 & -\mu \\ 0 & 0 & c_2M_0 & -c_4M_0 & -\mu & 0 \end{array}\right], \tag{19}$$

where $c_1 = c_\beta s_\theta$, $c_2 = s_\beta s_\theta$, $c_3 = c_\beta c_\theta$, $c_4 = s_\beta s_\theta$. We label the mass eigen states as $(\tilde{\chi}_a^0,\, \tilde{\chi}_5^0,\, \tilde{\chi}_6^0)$, $a = 1,2,3,4$, such that in the limit $\delta = M_2/M_1 \to 0$, $\tilde{\chi}_a^0$ are the four eigen states of the MSSM mass matrix in the lower right-hand corner, with $m_{\tilde{\chi}_1^0} < m_{\tilde{\chi}_2^0} < m_{\tilde{\chi}_3^0} < m_{\tilde{\chi}_4^0}$, plus the eigen states $\tilde{\chi}_5^0, \tilde{\chi}_6^0$ such that

$$m_{\tilde{\chi}_5^0},\, m_{\tilde{\chi}_6^0} = \sqrt{M_1^2 + \frac{1}{4}\tilde{m}_X^2} \pm \frac{1}{2}\tilde{m}_X + \mathcal{O}(\delta)\,, \qquad m_{\tilde{\chi}_5^0} \geq m_{\tilde{\chi}_6^0}\,.$$

As long as $m_{\tilde{\chi}_6^0} > m_{\tilde{\chi}_1^0}$, the lightest neutralino of the MSSM, $\tilde{\chi}_1^0$, will still function as the LSP of StMSSM. However, when $m_{\tilde{\chi}_6^0} < m_{\tilde{\chi}_1^0}$, $\tilde{\chi}_{\rm St}^0 = \tilde{\chi}_6^0$ becomes the LSP and (with R-parity conservation) a dark matter candidate. A numerical analysis shows that this can easily be the case for a wide range of parameters. This modifies completely the analysis of the decay channels for supersymmetric particles into the LSP, since the couplings of $\tilde{\chi}_{\rm St}^0$ to the visible (and to the hidden) matter are quite different than those of $\tilde{\chi}_1^0$. Aside from the issue of dark matter, the supersymmetric signals at particle colliders would be drastically modified, and the usual missing energy signals no longer apply. Indeed $\tilde{\chi}_1^0$ would itself be unstable to decay into $\tilde{\chi}_{\rm St}^0$ by a variety of channels, such as

$$\tilde{\chi}_1^0 \;\to\; l_i \bar{l}_i \tilde{\chi}_{\rm St}^0\,, \qquad q_j \bar{q}_j \tilde{\chi}_{\rm St}^0\,, \qquad Z\tilde{\chi}_{\rm St}^0\,, \tag{20}$$

where $i(j)$ are lepton (quark) flavors. The decay lifetime of $\tilde{\chi}_1^0$ is highly model dependent, involving the parameters of the Stueckelberg sector, i.e., M_2/M_1, as well as of the MSSM. An estimate using bounds on M_2/M_1 from [4] gives $\tau_{\tilde{\chi}_1^0} \sim 10^{-(19\pm3)}$ s, which implies that $\tilde{\chi}_1^0$ will decay in the detection chamber if $m_{\tilde{\chi}_6^0} < m_{\tilde{\chi}_1^0}$. In this circumstance the signatures for the detection of supersymmetry change drastically as discussed below.

Because the direct coupling between $\tilde{\chi}_{\rm St}^0$ and visible matter is weaker by M_2/M_1 than of the MSSM neutralinos $\tilde{\chi}_a^0$, sfermions $\tilde{f}_j$ will first decay dominantly into the MSSM neutralinos, i.e., $\tilde{f}_j \to f_j + \tilde{\chi}_a^0$. This is then followed by the decay of the $\tilde{\chi}_a^0$ with the chain ending with $\tilde{\chi}_{\rm St}^0$ in the end product. Typically this will lead to multi particle final states often containing many leptons. For example, the lightest slepton decay can result in a trileptonic final state

$$\tilde{l}^- \to l^- + \tilde{\chi}_1^0 \to l^- + \begin{cases} l_i^- l_i^+ + \{\tilde{\chi}_{\rm St}^0\} \\ q_j \bar{q}_j + \{\tilde{\chi}_{\rm St}^0\} \end{cases},$$

where $\{\tilde{\chi}_{\rm St}^0\}$ is the missing energy. A similar situation arises in the decay of the light charginos $\tilde{\chi}_1^\pm$ into charged leptons

$$\tilde{\chi}_1^- \to l^- + \tilde{\chi}_1^0 + \bar{\nu} \to l^- + \begin{cases} l_i^- l_i^+ + \{\tilde{\chi}_{\rm St}^0 + \bar{\nu}\} \\ q_j \bar{q}_j + \{\tilde{\chi}_{\rm St}^0 + \bar{\nu}\} \end{cases}.$$

As another example, the decay of a squark or a gluino will necessarily allow charged leptons in the final states,

$$\begin{aligned} \tilde{q} &\to q l \bar{l} + \{\tilde{\chi}^0_{\rm St}\}\,, \\ \tilde{g} &\to q\bar{q} l \bar{l} + \{\tilde{\chi}^0_{\rm St}\}\,. \end{aligned} \tag{21}$$

Using the above chain of decays, the processes $p + \bar{p} \to \tilde{\chi}^0_1 + \tilde{\chi}^0_1 + X, \tilde{\chi}^{\pm}_1 + \tilde{\chi}^0_1 + X, \tilde{\chi}^{\pm}_1 + \tilde{\chi}^0_2 + X, \tilde{\chi}^{\pm}_1 + \tilde{\chi}^{\mp}_1 + X$ at the Tevatron collider would lead to multi particle final states often with many leptons, and similar phenomena will occur at the LHC. Thus the conventional signal for supersymmetry in supergravity unified models where the decay of an off-shell W-boson leads to a trileptonic signal [16], $\mathrm{W}^{*-} \to \tilde{\chi}^-_1 \tilde{\chi}^0_2 \to l^- l_i \bar{l}_i +$missing energy, is replaced by a purely leptonic final state, which has seven leptons and missing energy. The decay branching ratios of these are model dependent and we leave a more complete investigation to a future work [17].

However, the preceding analysis is already sufficient to demonstrate, that in the specific scenario considered, where $m_{\tilde{\chi}^0_6} < m_{\tilde{\chi}^0_1}$, the supersymmetric signals at the Tevatron and at the LHC are drastically altered. A similar situation will hold for the supersymmetric signals at a linear collider. Here, the process $e^+e^- \to \tilde{\chi}^{\pm}_1 + \tilde{\chi}^{\mp}_1$ would lead to a purely leptonic final state with six leptons and missing energy which would provide signatures for this kind of a StMSSM scenario. Since the nature of physics beyond the standard model is largely unknown, it is imperative that one considers all viable scenarios, including the one discussed here, in the exploration of new physics.

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