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# LHC signatures of the constrained exceptional supersymmetric standard model

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We discuss two striking Large Hadron Collider (LHC) signatures of the constrained version of the exceptional supersymmetric standard model, based on a universal high-energy soft scalar mass $m_0$, soft trilinear coupling $A_0$ and soft gaugino mass $M_{1/2}$. The first signature we discuss is that of light exotic color triplet charge $1/3$ fermions, which we refer to as $D$-fermions. We calculate the LHC production cross section of $D$-fermions, and discuss their decay patterns. Secondly we discuss the $E_6$ type $U(1)_N$ spin-1 $Z'$ gauge boson and show how it may decay into exotic states, increasing its width and modifying the line shape of the dilepton final state. We illustrate these features using two representative exceptional supersymmetric standard model benchmark points, including an ''early LHC discovery'' point, giving the Feynman rules and numerical values for the relevant couplings in order to facilitate further studies.

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## I. INTRODUCTION

Last year the LHC experiments started to collect data. We expect that the LHC will shed light on the physics beyond the standard model (SM), the origin of dark matter and the mechanism of electroweak symmetry breaking in the near future. However it may take some time for the LHC experiments to discover the Higgs boson if it is light. On the other hand the LHC can relatively quickly discover new colored particles and a $Z'$ if these states are kinematically accessible. In this article we study the production and decay signatures of the $Z'$ and exotic color triplet charge $1/3$ fermions, which we refer to as $D$-fermions, that naturally appear within well a motivated supersymmetric extension of the SM known as the exceptional supersymmetric standard model ($E_6$SSM) [1].

Softly broken supersymmetry (SUSY) provides a very attractive framework for physics beyond the standard model, in which the hierarchy problem is solved and the unification of gauge couplings can be realized [2]. Despite these attractive features, the minimal supersymmetric standard model (MSSM) suffers from the $\mu$ problem. The superpotential of the MSSM contains the bilinear term $\mu H_d H_u$, where $H_{d,u}$ are the two Higgs doublet superfields[1] whose scalar components develop vacuum expectation values (VEVs) at the weak scale and $\mu$ is the supersymmetric Higgs mass parameter which can be present before SUSY is broken. One naturally expects $\mu$ to be the order of the Planck mass or to be zero, having been forbidden by some symmetry, whereas phenomenologically, to achieve correct electroweak symmetry breaking (EWSB), $\mu$ is required to be in the TeV region.

A very elegant solution to this problem is to generate an effective $\mu$-term from an interaction, $\lambda S H_d H_u$, between the usual Higgs doublets and a new Higgs singlet superfield $S$, whose scalar component develops a low-energy VEV. However, although an extra singlet superfield $S$ seems like a minor modification to the MSSM, which does no harm to gauge coupling unification, its introduction leads to an additional accidental global $U(1)$ (Peccei-Quinn [3]) symmetry which will result in a weak-scale massless axion when it is spontaneously broken by the VEV $\langle S \rangle$ [4].

To avoid this one can promote the Peccei-Quinn symmetry to an Abelian $U(1)'$ gauge symmetry [5]. The troublesome would-be axion is then eaten by the new $U(1)'$ gauge boson to give a massive $Z'$ at the TeV scale. An extra $U(1)'$ gauge group can also be motivated within the framework of grand unified theories (GUTs), arising as the relic of the breakdown of the unified gauge group. For example, an $E_6$ GUT symmetry can be broken to the rank-5 subgroup $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)'$ where in general $U(1)' = U(1)_\chi \cos\theta + U(1)_\psi \sin\theta$ [6], and the two anomaly-free $U(1)_\psi$ and $U(1)_\chi$ symmetries originate from the breakings $E_6 \to SO(10) \times U(1)_\psi$, $SO(10) \to SU(5) \times U(1)_\chi$. For a review see [7]and for a discussion of the latest Tevatron and early LHC $Z'$ mass limits see [8].

With additional Abelian gauge symmetries it is also important to ensure the cancellation of anomalies. This fits very nicely into the framework of an $E_6$ GUT since, for any $U(1)'$ that is a subgroup of $E_6$, anomalies are cancelled automatically if the low-energy spectrum constitutes a complete 27-plet.

Furthermore, within the class of $E_6$ models, there is a unique choice of Abelian gauge group that allows zero charges for right-handed neutrinos and thus large Majorana masses and a high-scale seesaw mechanism. This is the $U(1)_N$ gauge symmetry given by $\theta = \arctan\sqrt{15}$ which is

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[1]Note that we will not put hats on the superfields.

 

naturally achieved by GUT-scale Higgses which develop VEVs in the "right-handed neutrino" component. The choice of $U(1)_N$ gauge group coupled with complete 27-plets of matter at low energy defines the $\mathrm{E}_6$SSM [1].

The right-handed neutrinos acquire heavy Majorana masses and may play a role in the early Universe by decaying unequally into final states with lepton number $L = \pm 1$, creating a cosmological lepton asymmetry. Because the Yukawa couplings of the new exotic particles of the model are not constrained by the neutrino oscillation data, substantial values of $CP$-violating lepton asymmetries can be induced even for a relatively small mass of the lightest right-handed neutrino ($M_1 \sim 10^6$ GeV) so that successful thermal leptogenesis may be achieved without encountering any gravitino problem [9].

The extra $U(1)_N$ gauge symmetry survives to low energies and forbids a bilinear term $\mu H_d H_u$ in the superpotential but allows the interaction $\lambda S H_d H_u$. At the electroweak (EW) scale, the scalar component of the SM-singlet superfield $S$ acquires a nonzero VEV, $\langle S \rangle = s/\sqrt{2}$, breaking $U(1)_N$ and yielding an effective $\mu = \lambda s/\sqrt{2}$ term. Thus the $\mu$ problem in the $\mathrm{E}_6$SSM is solved in a similar way to the next-to-minimal supersymmetric standard model [10], but without the accompanying problems of singlet tadpoles or domain walls.

Recently we discussed a constrained version of the $\mathrm{E}_6$SSM (c$\mathrm{E}_6$SSM), based on a universal high-energy soft scalar mass $m_0$, soft trilinear coupling $A_0$ and soft gaugino mass $M_{1/2}$ [11–13]. We proposed a number of benchmark points, and calculated the SUSY and exotic spectrum which we found to have the following characteristics:

(i) a spin-1 $Z'_N$ gauge boson of mass around 1–2 TeV;

(ii) light gauginos including a light gluino of mass $\sim M_3$ (typically 350–650 GeV), a light winolike neutralino and chargino pair of mass $\sim M_2$ (typically 100–200 GeV), and a light binolike neutralino of mass $\sim M_1$ (typically 60–120 GeV), where $M_i$ are the low-energy gaugino masses, which are typically driven small compared to the effective $\mu$ parameter (typically 700–1400 GeV) by renormalization group (RG) running;

(iii) heavier sfermions (typically 800–1600 GeV), except for the lightest stop which may be 500–800 GeV;

(iv) possibly light exotic color triplet charge $1/3$ $D$-fermions, with masses controlled by independent Yukawa couplings enabling them to be as light as the Tevatron limit of about 300 GeV.

In this paper, motivated by the light spectrum above, we consider it urgent and timely to discuss two of the most characteristic and striking LHC signatures of the c$\mathrm{E}_6$SSM in considerably more detail than was done in [11,12]. First, we discuss the $U(1)_N$ spin-1 $Z'$ gauge boson (referred to as $Z'_N$) and show how it may decay into exotic states, including the exotic $D$-fermions and singlinos. This increases its width compared to that for SM fermion decays only, making its line shape more easily observed. Second, we calculate the LHC production cross section of exotic $D$-fermions and discuss their decay patterns. We illustrate these features by considering two of the benchmark points previously proposed in some detail. Crucially, we also give the numerical Feynman rules which will enable further studies (e.g., by experimentalists) to be performed.

We note that the phenomenology of $D$-fermions has also been discussed the general framework of $E_6$ models in [14], but not specifically for the c$\mathrm{E}_6$SSM which provides a more predictive framework via the use of benchmark points.

The layout of the rest of the paper is as follows. In Sec. II we review the c$\mathrm{E}_6$SSM. Sec. III discusses the LHC predictions of the c$\mathrm{E}_6$SSM illustrated through two benchmark points. Sec. IV concludes the paper. We then have one Appendix, where the numerical Feynman rules utilized in this work are given.

## II. THE CONSTRAINED $\mathrm{E}_6$SSM

The $\mathrm{E}_6$SSM is a supersymmetric model with three generations of complete 27 multiplets of matter and a low-energy gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_N$, where the $U(1)_N$ is specified by the charges given in Table I and the combination $U(1)_\chi \cos\theta + U(1)_\psi \sin\theta$, with $\theta = \arctan\sqrt{15}$.

The $27_i$ of $E_6$, each containing a quark and lepton family, decompose under the $SU(5) \times U(1)_N$ subgroup of $E_6$ as follows:

$$27_i \rightarrow (10, 1)_i + (5^*, 2)_i + (5^*, -3)_i + (5, -2)_i + (1, 5)_i + (1, 0)_i. \tag{1}$$

The first and second quantities in the brackets are the $SU(5)$ representation and extra $U(1)_N$ charge while $i$ is a family index that runs from 1 to 3. From Eq. (1) we see that, in order to cancel anomalies, the low-energy (TeV-scale)

TABLE I. The $U(1)_Y$ and $U(1)_N$ charges of matter fields in the $\mathrm{E}_6$SSM, where $Q_i^N$ and $Q_i^Y$ are here defined with the correct $E_6$ normalization factor required for the RG analysis.

| | $Q$ | $u^c$ | $d^c$ | $L$ | $e^c$ | $N^c$ | $S$ | $H_2$ | $H_1$ | $D$ | $\bar{D}$ | $H'$ | $\bar{H}'$ |
|---|---|---|---|---|---|---|---|---|---|---|---|---|---|
| $\sqrt{\frac{5}{3}}Q_i^Y$ | $\frac{1}{6}$ | $-\frac{2}{3}$ | $\frac{1}{3}$ | $-\frac{1}{2}$ | 1 | 0 | 0 | $\frac{1}{2}$ | $-\frac{1}{2}$ | $-\frac{1}{3}$ | $\frac{1}{3}$ | $-\frac{1}{2}$ | $\frac{1}{2}$ |
| $\sqrt{40}Q_i^N$ | 1 | 1 | 2 | 2 | 1 | 0 | 5 | $-2$ | $-3$ | $-2$ | $-3$ | 2 | $-2$ |

spectrum must contain three extra copies of $5^* + 5$ of $SU(5)$ in addition to the three quark and lepton families in $5^* + 10$. To be precise, the ordinary SM families which contain the doublets of left-handed quarks $Q_i$ and leptons $L_i$, right-handed up- and down-quarks ($u_i^c$ and $d_i^c$) as well as right-handed charged leptons, are assigned to $(10, 1)_i + (5^*, 2)_i$. Right-handed neutrinos $N_i^c$ should be associated with the last term in Eq. (1), $(1, 0)_i$. The next-to-last term in Eq. (1), $(1, 5)_i$, represents SM-singlet fields $S_i$ which carry nonzero $U(1)_N$ charges and therefore survive down to the EW scale. The three pairs of $SU(2)$-doublets ($H_i^d$ and $H_i^u$) that are contained in $(5^*, -3)_i$ and $(5, -2)_i$ have the quantum numbers of Higgs doublets, and we shall identify one of these pairs with the usual MSSM Higgs doublets, with the other two pairs being "inert" Higgs doublets which do not get VEVs. The other components of these $SU(5)$ multiplets form color triplets of exotic fermions $D_i$ and $\bar{D}_i$ with electric charges $-1/3$ and $+1/3$ respectively. The matter content and correctly normalized Abelian charge assignment are in Table I.

If there are only complete matter multiplets at low energy, the gauge couplings do not unify in a single step. Therefore one can either proceed with two-step unification, leading to unification at the string scale [15], or add incomplete multiplets.

In this paper we follow the latter path and require a further pair of superfields $H'$ and $\bar{H}'$ with a mass term $\mu' H' \bar{H}'$ from incomplete extra $27'$ and $\overline{27}'$ representations surviving to low energies. Anomaly cancellation is still guaranteed since $H'$ and $\bar{H}'$ originate from the $27'$ and $\overline{27}'$ supermultiplets. Previous analysis reveals that the unification of the gauge couplings in the $\mathrm{E}_6$SSM can be achieved for any phenomenologically acceptable value of $\alpha_3(M_Z)$, consistent with the measured low-energy central value, unlike in the MSSM which requires significantly higher values of $\alpha_3(M_Z)$, well above the central measured one [16].

The superpotential of the $\mathrm{E}_6$SSM contains many Yukawa couplings, including interactions between the SM singlets, $S_i$ to both the three generations of Higgs-like fields and the new exotic $D$-fermion fields, as well as interactions between the exotic $D$-fermions and inert Higgs fields with ordinary matter (the first two generations of the Higgs-like fields), which are new in comparison to the SM. Since some of these new interactions violate baryon number conservation and induce nondiagonal flavor transitions there should be some symmetry structure suppressing or forbidding the dangerous terms. A structure to do this can arise from a family symmetry at the GUT scale [17].

In the scenarios considered in this paper, following previous work [1,11,12], to suppress baryon number-violating and flavor-changing processes we postulate a $Z_2^H$ symmetry under which all superfields except one pair of $H_i^d$ and $H_i^u$ (say $H_d \equiv H_3^d$ and $H_u \equiv H_3^u$) and one SM-type singlet field ($S \equiv S_3$) are odd. The $Z_2^H$ symmetry reduces the number of the Yukawa interactions, and together with a further assumed hierarchical structure of the Yukawa interactions, we can simplify the form of the $\mathrm{E}_6$SSM superpotential substantially. Keeping only Yukawa interactions whose couplings are allowed to be of order unity leaves us with the following phenomenologically viable superpotential

$$W_{\mathrm{E}_6\mathrm{SSM}} \simeq \lambda S(H_d H_u) + \lambda_\alpha S(H_\alpha^d H_\alpha^u) + \kappa_i S(D_i \bar{D}_i) + h_t(H_u Q)t^c + h_b(H_d Q)b^c + h_\tau(H_d L)\tau^c + \mu'(H'\bar{H}'), \quad (2)$$

where $\alpha, \beta = 1, 2$ and $i = 1, 2, 3$, and where the superfields $L = L_3$, $Q = Q_3$, $t^c = u_3^c$, $b^c = d_3^c$ and $\tau^c = e_3^c$ belong to the third generation and $\lambda_i$, $\kappa_i$ are dimensionless Yukawa couplings with $\lambda \equiv \lambda_3$. Since the right-handed neutrino has no charge under the $U(1)_N$ gauge symmetry, nor under the SM gauge group, we assume that all right-handed neutrinos are relatively heavy so that they can be integrated out. The $SU(2)_L$ doublets $H_u$ and $H_d$, and singlet $S$, which are even under the $Z_2^H$ symmetry, now play the role of Higgs fields, generating the masses through EWSB, while the other generations of these Higgs-like fields remain inert. The $H_u$ and $H_d$ fields provide masses to the up-type and down-type quarks and leptons, respectively, just as in the MSSM, while $S$, which must acquire a large VEV to induce sufficiently large masses for the $Z_N'$ boson, also give masses to the exotic $D$-fermions and inert Higgs bosons from Yukawa interactions, $\lambda_\alpha S(H_\alpha^d H_\alpha^u)$ and $\kappa_i S(D_i \bar{D}_i)$. The couplings $\lambda_i$ and $\kappa_i$ should be large enough to ensure the exotic fermions are sufficiently heavy to avoid conflict with direct particle searches at present and past accelerators. One generation of the new Yukawa couplings (chosen to be the third generation) should also be large enough so that the evolution of the soft scalar mass $m_S^2$ of the singlet field $S$ results in negative values of $m_S^2$ at low energies, triggering the breakdown of the $U(1)_N$ symmetry.

However, the $Z_2^H$ can only be approximate since under an exact $Z_2^H$ decays of the exotic particles would be forbidden. Therefore, while Eq. (2) does not induce any proton decay, some suppressed couplings can, and so to prevent rapid proton decay in the $\mathrm{E}_6$SSM we should still introduce a discrete symmetry to play the role of $R$ parity in the MSSM. We give two examples of possible symmetries that can achieve that.

If $H_i^d$, $H_i^u$, $S_i$, $D_i$, $\bar{D}_i$ and the quark superfields ($Q_i, u_i^c, d_i^c$) are even under a discrete $Z_2^L$ symmetry while the lepton superfields ($L_i, e_i^c, N_i^c$) are odd (Model I) then the allowed superpotential is invariant with respect to a $U(1)_B$ global symmetry. The exotic $\bar{D}_i$ and $D_i$ are then identified as diquark and antidiquark, i.e., $B_D = -2/3$ and $B_{\bar{D}} = 2/3$. An alternative possibility is to assume that the exotic quarks $D_i$ and $\bar{D}_i$ as well as lepton superfields are all odd under $Z_2^B$ whereas the others remain even. In this case

(Model II) the $\bar{D}_i$ and $D_i$ are leptoquarks [1]. With both of these symmetries the MSSM particle content behaves like it does under $R$ parity, with the subset of particles present in the standard model and Higgs (and also inert Higgs) bosons being even under this generalized $R$ parity, while their supersymmetric partners are odd and therefore, as usual, must be pair-produced, and upon decaying will always give rise to a stable lightest supersymmetric particle (LSP). However the exotic $D$-fermions are odd and so must be pair-produced and will decay into an LSP, while their scalar superpartners are even and can be singly produced.

After $U(1)_N$ and EW symmetry breaking the Higgs fields, $H_u$, $H_d$ and $S$, give a physical Higgs spectrum of three $CP$-even, one $CP$-odd and two charged states. Two of the $CP$-even Higgs bosons tend to be rather heavy, with one mass being close to the $Z'$ boson mass $M_{Z'}$ and the other almost degenerate with the $CP$-odd Higgs boson and the charged Higgs states. The remaining $CP$-even Higgs boson is always light irrespective of the SUSY-breaking scale, and has an upper bound on its mass, as in the MSSM and next-to-minimal supersymmetric standard model, but in the $E_6$SSM it can be heavier than 110–120 GeV even at tree level. In the two-loop approximation the lightest Higgs boson mass does not exceed 150–155 GeV [1,18]. However for the benchmarks considered in the constrained model defined below [11,12] the lightest Higgs mass was in the range 115–121 GeV, and the points we selected for the study in this paper have light Higgs masses just above the LEP bound.

While the simplified superpotential of the $E_6$SSM in Eq. (2) only has six more couplings than the MSSM superpotential, the soft breakdown of SUSY gives rise to many new parameters. The number of fundamental parameters can be reduced drastically though within a constrained version of the model. Constrained SUSY models imply that all soft scalar masses are set to be equal to $m_0$ at some high-energy scale $M_X$, taken here to be equal to the GUT scale, all gaugino masses $M_i(M_X)$ are equal to $M_{1/2}$, and trilinear scalar couplings are such that $A_i(M_X) = A_0$. Thus the c$E_6$SSM is characterized by the following set of Yukawa couplings, which are allowed to be of the order of unity, and universal soft SUSY-breaking terms

$$\lambda_i(M_X),\ \kappa_i(M_X),\ h_t(M_X),\ h_b(M_X),\ h_\tau(M_X),\ m_0,\ M_{1/2},\ A_0, \tag{3}$$

where $h_t(M_X)$, $h_b(M_X)$ and $h_\tau(M_X)$ are the usual $t$-quark, $b$-quark and $\tau$-lepton Yukawa couplings, and $\lambda_i(M_X)$, $\kappa_i(M_X)$ are the extra Yukawa couplings defined in Eq. (2). The universal soft scalar and trilinear masses correspond to an assumed high-energy soft SUSY-breaking potential of the universal form

$$V_{\text{soft}} = m_0^2 27_i 27_i^* + A_0 Y_{ijk} 27_i 27_j 27_k + \text{H.c.}, \tag{4}$$

where $Y_{ijk}$ are generic Yukawa couplings from the trilinear terms in Eq. (2) and the $27_i$ represent generic fields from Eq. (1), and, in particular, those which appear in Eq. (2). In previous analyses we always set $m_0^2$ positive for correct EWSB and to simplify the analysis assume that all parameters in Eq. (3) are real and $M_{1/2}$ is positive. The set of c$E_6$SSM parameters in Eq. (3) should in principle be supplemented by $\mu'$ and the associated bilinear scalar coupling $B'$. However, since $\mu'$ is not constrained by the EWSB and the term $\mu' H' \bar{H}'$ in the superpotential is not suppressed by $E_6$, the parameter $\mu'$ was assumed to be $\sim$10 TeV so that $H'$ and $\bar{H}'$ decoupled from the rest of the particle spectrum. As a consequence the parameters $B'$ and $\mu'$ are irrelevant for the analysis [11,12].

In addition several of the parameters specified above are fixed by experimental measurements and the RG flow. This means that the particle spectrum and many phenomenological aspects of the model can be determined from only eight free parameters, which in previous analyses have been taken to be[2] $\{\lambda_i, \kappa_i, s, \tan\beta\}$, which can be compared to the cMSSM with $\{m_0, M_{1/2}, A, \tan\beta, \text{sign}(\mu)\}$, and could be reduced further by considering scenarios with some Yuakawa coupling universality or other well-motivated relations between the Yukawa couplings at the GUT scale.

To calculate the particle spectrum within the c$E_6$SSM a private spectrum generator has been written, based on some routines and the class structure of SOFTSUSY 2.0.5 [21]and employing two-loop RG equations (RGEs) for the gauge and Yukawa couplings together with two-loop RGEs for $M_a(Q)$ and $A_i(Q)$ as well as one-loop RGEs for $m_i^2(Q)$, where $Q$ is the renormalization scale. The details of the procedure we followed, including the RGEs for the $E_6$SSM and the experimental and theoretical constraints can be found in [11,12].

## III. LHC SIGNATURES OF THE c$E_6$SSM

### A. Benchmark spectra and couplings

In previous publications we presented a set of "early discovery" benchmark points which should be discovered using first LHC data and a set of slightly heavier ("late discovery") benchmarks to illustrate the wider range of possible c$E_6$SSM scenarios which could be discovered at the LHC. Here we select two of these points for a more detailed phenomenological study, focusing on the $Z_N'$ and the new exotic colored states. For this we have chosen the "early discovery" benchmark C (BMC) and a heavier, qualitatively different benchmark 4 (BM4). The mass spectra for these are given in Table II.

These spectra both exhibit the characteristic c$E_6$SSM signature of a heavy sfermion sector, with light gauginos. Previously we observed that in the c$E_6$SSM $m_0 \gtrsim M_{1/2}$ for all phenomenologically viable points [11,12]. Additionally we discovered that the low-energy gluino mass parameter

[2]Note that $m_0$, $M_{1/2}$, and $A_0$ have been replaced by $v$, $\tan\beta$ and $s$ through the EWSB conditions, in a similar manner to the way $|\mu|$ and $B$ are traded for $\tan\beta$ and $v$ in the MSSM.

TABLE II. Parameters for the "early discovery" benchmark point C (left) (from [11]) and "late discovery" benchmark point 4 (from [12]).

| | BMC | BM4 |
|---|---|---|
| $\tan\beta$ | 10 | 30 |
| $\lambda_3(M_X)$ | $-0.378$ | $-0.38$ |
| $\lambda_{1,2}(M_X)$ | 0.1 | 0.1 |
| $\kappa_3(M_X)$ | 0.42 | 0.16 |
| $\kappa_{1,2}(M_X)$ | 0.06 | 0.16 |
| $s$ [TeV] | 2.7 | 5.0 |
| $M_{1/2}$ [GeV] | 388 | 725 |
| $m_0$ [GeV] | 681 | 1074 |
| $A_0$ [GeV] | 645 | 1726 |
| $m_{\tilde{D}_1}(3)$ [GeV] | 1465 | 312 |
| $m_{\tilde{D}_2}(3)$ [GeV] | 2086 | 2623 |
| $\mu_D(3)$ [GeV] | 1747 | 1612 |
| $m_{\tilde{D}_1}(1,2)$ [GeV] | 520 | 312 |
| $m_{\tilde{D}_2}(1,2)$ [GeV] | 906 | 2623 |
| $\mu_D(1,2)$ [GeV] | 300 | 1612 |
| $\lvert m_{\chi_6^0}\rvert$ [GeV] | 1054 | 1950 |
| $m_{h_3} \simeq M_{Z'}$ [GeV] | 1021 | 1889 |
| $\lvert m_{\chi_5^0}\rvert$ [GeV] | 992 | 1832 |
| $m_S(1,2)$ [GeV] | 1001 | 1732 |
| $m_{H_2}(1,2)$ [GeV] | 627 | 1117 |
| $m_{H_1}(1,2)$ [GeV] | 459 | 220 |
| $\mu_{\tilde{H}}(1,2)$ [GeV] | 233 | 491 |
| $m_{\tilde{u}_1}(1,2)$ [GeV] | 911 | 1557 |
| $m_{\tilde{d}_1}(1,2)$ [GeV] | 929 | 1595 |
| $m_{\tilde{u}_2}(1,2)$ [GeV] | 929 | 1595 |
| $m_{\tilde{d}_2}(1,2)$ [GeV] | 964 | 1664 |
| $m_{\tilde{e}_2}(1,2,3)$ [GeV] | 849 | 1427 |
| $m_{\tilde{e}_1}(1,2,3)$ [GeV] | 765 | 1254 |
| $m_{\tilde{\tau}_2}$ [GeV] | 845 | 1363 |
| $m_{\tilde{\tau}_1}$ [GeV] | 757 | 1102 |
| $m_{\tilde{b}_2}$ [GeV] | 955 | 1491 |
| $m_{\tilde{b}_1}$ [GeV] | 777 | 1193 |
| $m_{\tilde{t}_2}$ [GeV] | 829 | 1248 |
| $m_{\tilde{t}_1}$ [GeV] | 546 | 837 |
| $\lvert m_{\chi_3^0}\rvert \simeq \lvert m_{\chi_4^0}\rvert \simeq \lvert m_{\chi_2^\pm}\rvert$ [GeV] | 674 | 1343 |
| $m_{h_2} \simeq m_A \simeq m_{H^\pm}$ [GeV] | 963 | 998 |
| $m_{h_1}$ [GeV] | 115 | 114 |
| $m_{\tilde{g}}$ [GeV] | 353 | 642 |
| $\lvert m_{\chi_1^\pm}\rvert \simeq \lvert m_{\chi_2^0}\rvert$ [GeV] | 109 | 206 |
| $\lvert m_{\chi_1^0}\rvert$ [GeV] | 61 | 116 |

$M_3$ is driven to be smaller than $M_{1/2}$ by RG running, due to the much larger (super)field content of the $E_6$SSM in comparison to the MSSM (three $27'$s instead of three $16'$s). This implies that the low-energy gaugino masses are all less than $M_{1/2}$ in the c$E_6$SSM, being given by roughly[3] $M_3 \sim 0.7M_{1/2}, M_2 \sim 0.25M_{1/2}, M_1 \sim 0.15M_{1/2}$. These two features imply that the sfermions of ordinary matter will always be heavier than the lightest gauginos, and the lightest SUSY states will include of a light gluino of mass $\sim M_3$, a light winolike neutralino and chargino pair of mass $\sim M_2$, and a light binolike neutralino of mass $\sim M_1$.

The heavier spectrum of BM4 is due to a significantly larger choice for the singlet vacuum expectation value, $s = \sqrt{2}\langle S\rangle = 5$ TeV, as opposed to $s = 2.7$ TeV in BMC. While substantial variation in the spectra can be produced by varying the new Yukawa couplings associated with exotic interactions, $\langle S\rangle$ is linked to the spectrum through the EWSB conditions and $U(1)_N$ $D$-terms, so choosing a particular value places restrictions on the masses and in general the larger $\langle S\rangle$ the heavier the spectrum.

The $U(1)_N$ gauge coupling, $g_1'$, is fixed by gauge coupling unification with the RG flow leading to $g_1'(M_Z) \approx g_1(M_Z)$. This means that $\langle S\rangle$ fixes the mass of the $Z_N'$, since $M_{Z'} \sim g_1'\langle S\rangle$ and this leads to $M_{Z'} = 1890$ GeV for BM4 and $M_{Z'} = 1021$ GeV for BMC, affecting the discovery potential at the LHC, as will be discussed later.

Another consequence of this is that the couplings to the $Z_N'$ are also highly constrained in this model since they are given by the gauge coupling and the $U(1)_N$ charges. Variation of these couplings between benchmark points comes only from mass mixing of gauge eigenstates, the scale dependence of $g_1'$, and two-loop running effects. This variation can be seen in the Appendix, where the $Z_N'$ Feynman rules are presented for our two benchmarks. However, despite this, there is still considerable room for different phenomenologies for a given $M_{Z'}$ (or equivalently $\langle S\rangle$), and this can also strongly impact on the Drell-Yan production cross section of the $Z_N'$.

For example, the exotic colored fermions can be light or heavy, since their masses are given by $\mu_{D_i} = \frac{1}{\sqrt{2}}\kappa_i s$, and if $\kappa$ universality is not assumed[4] it is possible to obtain two $\kappa_i(M_S)$ (where $M_S$ is the SUSY-breaking scale) small enough that the exotic fermions are just above their mass limit (300 GeV), as BMC illustrates. However the masses of the scalar partners to the exotic colored fermions also have soft mass contributions which tend to increase with $M_{Z'}$, and as a result only one of the two scalars can be light, and it is unlikely that both scalars will be available as $Z_N'$ decay modes. Nonetheless, even without small $\kappa_i$, it is still possible to have a light exotic sfermion due to large mixing, and this is demonstrated in BM4.

The inert Higgsino masses, $\mu_{H_\alpha} = \frac{1}{\sqrt{2}}\lambda_i s$, may also be light for a sufficiently small $\lambda_\alpha$ coupling and this is the case in BMC (and to a lesser extent BM4). However it is possible to also have all $\lambda_i$ large, giving Higgsinos of a TeV or above, and not available for the $Z_N'$ to decay into. The scalar inert Higgs masses can be very light depending on

[3]These should be compared to the corresponding low-energy values in the cMSSM, $M_3 \sim 2.7M_{1/2}$, $M_2 \sim 0.8M_{1/2}$, $M_1 \sim 0.4M_{1/2}$.

[4]At least one $\kappa_i$ coupling must be large to generate EWSB.

the particular parameters chosen. However, as with the exotic sfermions, due to the soft mass contribution, there is usually a hierachy between the inert Higgs bosons of a particular generation.

Both the inert and exotic colored states have large $U(1)_N$ charges which mean they can play an important role in $Z'_N$ phenomenology, as well as also producing interesting signatures from direct production.

All the sfermions of ordinary matter are rather heavy, with the sfermions in BM4 being substantially heavier than in BMC, arising from the influence of the larger $M_{Z'}$ in the EWSB conditions. The stops tend to be the lightest of the sfermions and due to large mixing the lightest stop is the only ordinary sfermion which can be really light with the possibility of being just above 400 GeV. In the two benchmarks here we have $m_{\tilde{t}_1} = 546$ GeV for BMC and $m_{\tilde{t}_1} = 837$ GeV for BM4, due as usual to the heavier $M_{Z'}$.

The light SUSY states that are always present in the spectrum include a light gluino $\tilde{g}$, two light neutralinos $\chi_1^0$, $\chi_2^0$, and a light chargino $\chi_1^\pm$. The lightest neutralino $\chi_1^0$ is essentially pure bino, while $\chi_2^0$ and $\chi_1^\pm$ are the degenerate components of the wino. Since these particles are composed primarily from states that do not couple to the $Z'_N$, they do not play a large role in the $Z'_N$ phenomenology. In addition there are other neutralinos $\chi_3^0$ and $\chi_4^0$ which are essentially pure Higgsino states and $\chi_5^0$ and $\chi_6^0$ associated with the third family singlino $\tilde{S}$ and the $\tilde{Z}'$ gaugino.

Nonetheless pair production of $\chi_2^0\chi_2^0$, $\chi_2^0\chi_1^\pm$, $\chi_1^\pm\chi_1^\mp$ and $\tilde{g}\,\tilde{g}$ should always be possible at the LHC irrespective of the $Z'$ mass.

The second lightest neutralino decaying through $\chi_2^0 \to \chi_1^0 + l\bar{l}$ would produce an excess in $pp \to l\bar{l}l\bar{l} + E_T^{\rm miss} + X$, where $X$ refers to any number of light quark/gluon jets, which could be observed at the LHC. While the focus of this paper is on exotics, so we do not present branching ratios for the MSSM-like states decaying, we note that since the second lightest and lightest neutralino are winolike and binolike states,, respectively, with squarks and sleptons and new exotic particles significantly heavier compared to the branching ratio $\mathrm{BR}(\chi_2^0 \to \chi_1^0 + l\bar{l})$, will be rather similar to MSSM scenarios with the same structure where $\mathrm{BR}(\chi_2^0 \to \chi_1^0 + l\bar{l})$ varies from 1.5% to 6% [22].

The rigid structure of the model also implies that gluinos can be relatively narrow states with width $\Gamma_{\tilde{g}} \propto M_{\tilde{g}}^5/m_{\tilde{q}}^4$, where the squarks are always significantly heavier that the gluino, leaving a width comparable to that of $W^\pm$ and $Z$ bosons. Because of the absence of lighter colored, r parity odd states, the gluinos can only decay in a three body decay involving a virtual squark, $\tilde{g} \to q\tilde{q}^* \to q\bar{q} + E_T^{\rm miss}$. Therefore gluino pair production will result in an appreciable enhancement of the cross section for $pp \to q\bar{q}q\bar{q} + E_T^{\rm miss} + X$.

Finally there are also two light inert singlinos not shown explicitly in Table II [the SUSY partners to the two families of inert singlet scalars $S(1, 2)$] whose masses are given by suppressed couplings that are assumed to be small enough so that they do not perturb the RG running of the other couplings. So, although these masses are not precisely fixed in previous analyses of the cE$_6$SSM spectrum, they are assumed to be very light. These particles then guarantee that there will be a substantial non-SM contribution to the $Z'_N$ width. However, if there is also other light exotic matter, then it can also make a significant contribution, as will be discussed later in the paper.

### B. LEP, Tevatron and LHC limits

The presence of light states (neutralinos, chargino and inert singlinos) in the E$_6$SSM particle spectrum raises serious concerns that they could have already been observed at the Tevatron and/or even earlier at LEP. For example, the light neutralino and chargino states could be produced at the Tevatron [23]. Recently, the CDF and D0 collaborations set a stringent lower bound on chargino masses using searches for SUSY with a trilepton final state [24]. These searches ruled out chargino masses below 164 GeV. However this lower bound on the chargino mass was obtained by assuming that the corresponding chargino and neutralino states decay predominantly into the LSP and a pair of leptons. In our case, the lightest neutralino and chargino states are expected to decay via virtual $Z$ and $W$ exchange, and then predominantly into the LSP and a pair of quarks. As a consequence the lower limit on the mass of charginos that is set by the Tevatron is not directly applicable to the benchmark scenarios that we consider here. Instead in our study we use the 95% confidence level lower limit on the chargino mass of about 100 GeV that was set by LEP II [25].

LEP experiments also set stringent constraints on the masses and couplings of neutral particles that interact with the $Z$ boson. Since inert singlinos have masses below $M_Z/2$, the $Z$ boson could decay into these states. However the couplings of these exotic states to the $Z$-boson are rather small due to their singlino nature [26]. Consequently their contribution to the $Z$ boson decay width and the corresponding branching ratios are negligible. Because of the small $Z$ couplings, the production of light inert singlinos at LEP was extremely suppressed, which allowed these states to escape detection at LEP.

Nevertheless the presence of light inert singlinos could lead to other phenomena which could be observed at LEP. In the case of BMC, $\chi_1^0\chi_1^0$ and $\chi_1^0\chi_2^0$ could be produced followed by their decay into inert singlino via virtual $Z$ exchange, resulting in $q\bar{q}q'\bar{q}'$ and missing energy in the final state. LEP has set limits on the cross section of $e^+e^- \to \chi_2^0\chi_1^0(\chi_1^+\chi_1^-)$ in the case where the subsequent decay is predominantly $\chi_2^0 \to q\bar{q}\chi_1^0$ $(\chi_1^\pm \to q\bar{q}'\chi_1^0)$ [27]. Unfortunately, these bounds are not directly applicable to our study, but they do demonstrate that it was difficult to observe

$$e^+e^- \to X + Y \to q\bar{q}q'\bar{q}' + \not{E}_T,$$

where $X$ and $Y$ are neutral particles, if the corresponding production cross section was 0.1–0.3 pb. In the case of the BMC the lightest and second lightest neutralinos have rather small couplings to the $Z$ boson. The corresponding relative couplings are of the order $(M_W/\mu)^2$. Since the selectron is also heavy in the considered scenario the production cross sections of $\chi_1^0\chi_1^0$ and $\chi_1^0\chi_2^0$ are suppressed by $O(\frac{1}{M^4})$ where $M \sim 700$–$800$ GeV. At LEP energies the cross sections of colorless particle production through $s$-channel $\gamma/Z$ exchange are typically a few picobarns, so the production cross sections of $\chi_1^0\chi_1^0$ and $\chi_1^0\chi_2^0$ in the case of BMC are expected to be of the order of $10^{-2}$ pb or even less. Thus BMC could not be ruled out by LEP experiments.

The Higgsino states are much heavier with the degenerate Higgsinos $\chi_{3,4}^0$ and $\chi_2^\pm$ having masses given by $\mu = \lambda s/\sqrt{2}$ in the range 675–830 GeV for both benchmark points considered. The remaining neutralinos are dominantly third-generation singlino and the gaugino partner of the $Z_N'$ with masses approximately given by $M_{Z'}$.

The Higgs spectrum for all the benchmark points contains a very light SM-like $CP$-even Higgs boson $h_1$ with a mass close to the LEP limit of 115 GeV, making it accessible to LHC or even Tevatron. The heavier $CP$-even Higgs $h_2$, the $CP$-odd Higgs $A_0$, and the charged Higgs $H^\pm$ are all closely degenerate with masses above 900 GeV making them difficult to discover. The remaining mainly singlet $CP$-even Higgs $h_3$ is closely degenerate with the $Z_N'$.

Tevatron, LEP and other experiments also set limits on the mass of the $Z_N'$ boson, $Z - Z_N'$ mixing and masses of exotic scalars ($\tilde{D}$). The direct searches at the Fermilab Tevatron ($p\bar{p} \to Z_N' \to l^+l^-$) exclude $Z_N'$ with mass below 892 GeV [8].[5] At the LHC, the $Z'$ boson that appears in the $E_6$-inspired models can be discovered if it has a mass below 4–4.5 TeV [29]. The determination of its couplings should be possible if $M_{Z'} \lesssim 2$–$2.5$ TeV [30]. The precision EW tests bound the $Z - Z'$ mixing angle to be around $(-1.5) - 0.7 \times 10^{-3}$ [31]. Recent results from Tevatron searches for dijet resonances [32]rule out scalar diquarks with mass less than 630 GeV. However, scalar leptoquarks may be as light as 300 GeV since at hadron colliders they are pair-produced through gluon fusion [33].

Recent SUSY searches at the LHC have substantially reduced the parameter space of the cMSSM. It is therefore worthwhile considering what we can infer from these searches regarding the parameter space of the $cE_6SSM$ and especially the benchmark points under consideration in this paper. To understand precisely how the model is constrained a detailed analysis, involving detector effects, and taking into account the signatures from the $cE_6SSM$ spectrum, which cannot be directly compared to the cMSSM, would be required.

While there are many differences between the two models, since the experimental searches are for squarks and gluinos we can try to map the $m_0$ and $M_{1/2}$ of the cMSSM to different values $m_0'$ and $M_{1/2}'$ which give similar squark and gluino masses in the $cE_6SSM$. The largest difference comes from the dramatically altered RG running of the gaugino sector. So as a rough approximation one can try to rescale $M_{1/2}$, using the RG coefficients given in the previous section. One observes that a particular value of $M_{1/2}$ in the $cE_6SSM$ gives the same gluino mass as a corresponding value of $M_{1/2}$ in the cMSSM (or mSUGRA) approximately 4 times smaller. In contrast the RG running does not dramatically alter the $m_0$ coefficients of squark masses, but due to other effects like large $U(1)_N$ $D$-terms we note that for points with the same $m_0$ in both models the squarks are significantly heavier, one should also rescale $m_0$ noting that a particular value of $m_0$ in the $cE_6SSM$ gives similar squark masses as a corresponding value of $m_0$ in the cMSSM (or mSUGRA) very roughly of order one and a half times larger. Thus as an extremely crude approximation $(m_0, M_{1/2})_{cE_6SSM} \to ((3/2)m_0, (1/4)M_{1/2})_{cMSSM}$ which underlines the $cE_6SSM$ prediction of relatively heavy squarks and relatively light gluinos. Note that this is the least sensitive region of the recent cMSSM analyses by CMS [19]and ATLAS [20] and suggests that these recent results do not restrict much of the previously phenomenologically viable parameter space.

For the benchmark points in question as a first rough estimate we can simply match the specific squark and gluino masses of the benchmark to those of the cMSSM and compare against the limits there. For BM4 we have $m_{\tilde{u}_2} = 1595$ GeV (the mass of the first- and second-generation left-handed squark) and $m_{\mathrm{gluino}} = 642$ GeV, which leaves the benchmark very far away from the experimental constraints presented in the cMSSM by ATLAS and CMS. However for BMC we have $m_{\tilde{u}_2} = 929$ GeV and $m_{\mathrm{gluino}} = 353$ GeV, which suggests the data used to constrain the cMSSM may have an impact on this point. In fact, ATLAS and CMS exclusion plots in imply that in the cMSSM such a light gluino was ruled out by LEP due to the lower bound on the mass of the lightest chargino. On the other hand the results of the calculations presented in Table II demonstrate that in the case of BMC all charged sparticles satisfy LEP constraints. This clearly indicates that the cMSSM and $cE_6SSM$ are extremely different models. For this reason the simplistic approach based on the matching of the squark and gluino masses does not

[5]Slightly weaker lower bound on the mass of the $Z_N'$ boson was obtained in [28]. Note that these bounds assume $Z_N'$ boson decays only into quarks and leptons. If the width increases by about a factor of 2 due to exotics and SUSY particles (as will be the case for the benchmarks studied in this paper) then this would reduce the branching ratio into charged leptons also by a factor of 2, which we estimate would reduce the mass limit quoted in from 892 GeV down to about 820 GeV.

really allow one to judge if a benchmark scenario in the $cE_6SSM$ is excluded when the squark and gluino masses of the corresponding cMSSM scenario are relatively close to the exclusion limits. Nevertheless such approach is not meaningless especially in the case of the ATLAS and CMS limits on the gluino mass. Indeed, when squarks are very heavy the gluino production rate is basically determined by the gluino couplings to gluon fields which are fixed by QCD in the leading approximation. Because the gluino production cross section decreases rapidly with increasing
gluino mass the lower limit on $m_{\rm gluino}$ in the $cE_6SSM$ cannot be much smaller than the corresponding limit in the cMSSM.

By means of direct comparison of BMC with CMS/ATLAS exclusion plots presented in one can establish that the MSSM benchmark point associated with BMC is either outside the exclusion region or very close to the ATLAS observed limit. However, also note that the experimental searches use specific choices of $\tan\beta$ and $A$ which are very far from the values for BMC and the $cE_6SSM$ mass splitting between squarks is very different compared to that of the cMSSM, due to altered RG running and large $U(1)_N$ $D$ terms. Thus without a thorough, detailed analysis, including the full $cE_6SSM$ spectrum and accounting for detector effects, it is not possible to determine whether borderline points like BMC are ruled out or not based on the 35 $\text{pb}^{-1}$ of data already used in these published analyses. It is therefore essential that such analysis is carried out soon on this question.

### C. Phenomenology

In this subsection, we focus on the phenomenology of the two benchmark points (BMC and BM4), in order to illustrate two of the most striking $cE_6SSM$ predictions: the exotic contributions to the heavy jet rate and the existence of a $Z'_N$ boson with an enhanced and resolvable width due to its additional decays into exotic states.

Table III presents the $cE_6SSM$ $Z'_N$ partial decay widths in all available channels. Apart from the leading SM decays into quarks ($q$) and leptons ($l$), one can notice, amongst the $cE_6SSM$ channels, the dominance of the decay into singlinos (collectively denoted by $\tilde{S}$), whose mass we have set at 10 and 30 GeV, for the two generations, respectively.[6] Next in line in order of importance are the exotic fermion (specifically, $D$-fermion, when open) and inert Higgsino ($\tilde{H}$) channels. The genuine SUSY contributions into gauginos ($\tilde{\chi}$) are never sizable while exotic scalars ($\tilde{D}$) and sfermions ($\tilde{f}$) count negligibly. At times (here for BM4), decays into inert Higgs ($H^{0/\pm}_{\alpha,i}$) states can also be tangible. Overall, non-SM contributions to the $cE_6SSM$ $Z'_N$ width are of order 100% for both benchmarks considered.

TABLE III. $Z'_N$ widths for the ''early discovery'' benchmark point C (left) (from [11]) and ''late discovery'' benchmark point 4 (from [12]). The index $i$ is summed over three families, the index $\alpha$ is summed over the two inert families of exotics while $j$ is summed over the light neutralino and chargino states. The leptonic branching ratio into $l^+l^-$ is given by $\text{Br}(l^+l^-) \approx 0.023$ for BMC and $\text{Br}(l^+l^-) \approx 0.028$ for BM4, as compared to the value calculated by ignoring the exotics and SUSY partners of $\text{Br}(l^+l^-) \approx 0.055$ in both cases. The Drell-Yan cross section may be defined in terms of two parameters $c_u$ and $c_d$ which are defined and discussed in [34]. In the limit where exotics and SUSY partners are ignored their values for this model are given by $c_u \approx 5.9 \times 10^{-4}$ and $c_d \approx 1.5 \times 10^{-3}$ [8]. Since $c_u$ and $c_d$ are both proportional to $\text{Br}(l^+l^-)$ they will therefore be reduced for both benchmarks due to the presence of exotics and SUSY partners by about a factor of 2 in each case. For BMC we find $c_u \approx 2.4 \times 10^{-4}$ and $c_d \approx 0.61 \times 10^{-3}$, while for BM4 we find $c_u \approx 3.0 \times 10^{-4}$ and $c_d \approx 0.75 \times 10^{-3}$.

| $Z'_N$ partial width [GeV] | BMC | BM4 |
|---|---|---|
| $\Gamma(Z'_N \to l^+l^-)$ ($l = e, \mu$ or $\tau$) | 0.41 | 0.77 |
| $\Sigma_l \Gamma(Z'_N \to \nu_l \bar{\nu}_l)$ (all neutrinos) | 0.87 | 1.64 |
| $\Sigma_l \Gamma(Z'_N \to l^+l^-, \nu_l \bar{\nu}_l)$ (all leptons) | 2.10 | 3.96 |
| $\Sigma_q \Gamma(Z'_N \to q\bar{q})$ (all quarks) | 5.31 | 10.08 |
| $\Sigma_i \Gamma(Z'_N \to D_i \bar{D}_i)$ (exotic fermions) | 3.49 | 0.00 |
| $\Sigma_\alpha \Gamma(Z'_N \to \tilde{H}_\alpha \tilde{H}_\alpha)$ (inert Higgsinos) | 3.09 | 5.19 |
| $\Sigma_\alpha \Gamma(Z'_N \to \tilde{S}_\alpha \tilde{S}_\alpha)$ (singlinos) | 4.05 | 7.63 |
| $\Sigma_i \Gamma(Z'_N \to \tilde{D}_i \tilde{D}_i)$ (exotic scalars) | 0.00 | 0.19 |
| $\Sigma_f \Gamma(Z'_N \to \tilde{f}\tilde{f})$ (sfermions) | 0.00 | 0.010 |
| $\Sigma_\alpha \Gamma(Z'_N \to H_\alpha H_\alpha)$ (inert Higgses) | 0.026 | 0.39 |
| $\Sigma_j \Gamma(Z'_N \to \tilde{\chi}_j \tilde{\chi}_j)$ (gauginos) | $6.50 \times 10^{-4}$ | $7.92 \times 10^{-5}$ |
| $\Gamma_{\rm tot}$ (all) | 18.07 | 27.45 |

The presence of light exotic particles and gauginos gives rise to nonstandard decays of the $Z'_N$ gauge boson. Indeed, exotic states, that originate from the $Z'_N$ decays, subsequently decay resulting in the four-fermion final states with and without missing energy. For example, the $Z'_N$ can decay into a pair of second lightest singlinos. Then second lightest singlino sequentially decays into the lightest one and a fermion-antifermion pair mainly via a virtual $Z$. Since lightest singlino is stable it leads to the missing energy in the final state. Because second lightest singlino tend to be relatively light it decays predominantly into light quarks and leptons. At the same time the decays of the $Z'_N$ into $D$-fermions (or inert Higgsinos) give rise to the final states that contain four third-generation fermions and missing energy as will be clarified later. Because $Z'_N$ is relatively heavy its decay products, which appear in the corresponding exotic final states, should have sufficiently high energies. Therefore some of them (in particular, charged leptons) might be observed at the LHC.

[6]Notice that their contribution to the total $Z'$ width is typically always about 30%, irrespectively of their actual mass, so long as the singlino masses remain within the boundaries established in [26], as space effects are minimal for the considered $Z'$ masses.

### *1. Benchmark C*

We now discuss the details of the "early discovery" BMC in Table II corresponding to a lighter spectrum first observable at the LHC with 7 TeV, then subsequently amenable to detailed study at 14 TeV.

*$Z'_N$ bosons.*—Figure 1 (top frame) shows the differential distribution in invariant mass of the lepton pair $l^+l^-$ (for one species of lepton $l = e$, $\mu$ or $\tau$) in Drell-Yan production at the LHC for $\sqrt{s} = 7$ TeV, assuming a sequential $Z'$ (that is, with the same mass as in the cE$_6$SSM but with SM-like couplings, i.e., no additional matter) as well as a cE$_6$SSM $Z'$ field with and without light exotic quarks and inert Higgsinos.[7]

This distribution is promptly measurable even at the lower-energy stage of the CERN collider with a high resolution and would enable one to not only confirm the existence of a $Z'$ state but also to establish the possible presence and nature of additional exotic matter, by simply fitting to the data the width of the $Z'$ resonance, its height at the resonance point, and its profile in the interference region with the SM channels ($\gamma$- and $Z$-mediated). In fact, for our choice of $\mu_{D_i}$, $\mu_{H_i}$ and $M_{Z'}$, the $Z'_N$ total width varies from $\approx 7$ GeV (in case of SM-only matter) to $\approx$ 18 GeV (in case of additional cE$_6$SSM matter). In particular, notice the different normalization around the $Z'$ resonance of the three curves in Fig. 1 (top frame).[8]

Another $Z'$ observable (alongside the cross section normalization and its line shape near and below the $Z'$ peak) which will be useful to access $Z'$ couplings is the forward-backward asymmetry (here denoted by AFB$_{l^+l^-}$). Figure 1 (middle frame) indeed shows a sizable difference in its shape (around the $Z'$ mass resonance, especially) between the cases of a sequential $Z'$ and a cE$_6$SSM $Z'_N$, albeit difficult to measure at the 7 TeV LHC (assuming 1 fb$^{-1}$ of total accumulated luminosity). Remarkably, the shape (and normalization) of AFB$_{l^+l^-}$ is essentially the same in the cE$_6$SSM irrespective of its particle content, so that the ability of accessing the $Z'$ couplings in such a model does not require a knowledge of its spectrum beforehand.

Figure 2 (top and middle frame) reinstates the above phenomelogical aspects at 14 TeV, with the added bonus of much larger event rates (by a factor of 6 or so around the $Z'$ peak) and luminosity (which could be up to 300 fb$^{-1}$ at the end of the collider lifetime).

[7]We have three generations of the exotic quarks but only two of inert Higgs. For convenience, in the legends of the plots we only refer to the former. Also note that we always include the other width contributions, according to Table III.

[8]Clearly, in order to perform such an exercise, the $Z'$ couplings to ordinary matter ought to have been previously established elsewhere, as a modification of the latter may well lead to effects similar to those induced by the additional matter present in our model. (Recall that in our model $Z'_N$ couplings to SM particles and exotic matter are simultaneously fixed.)

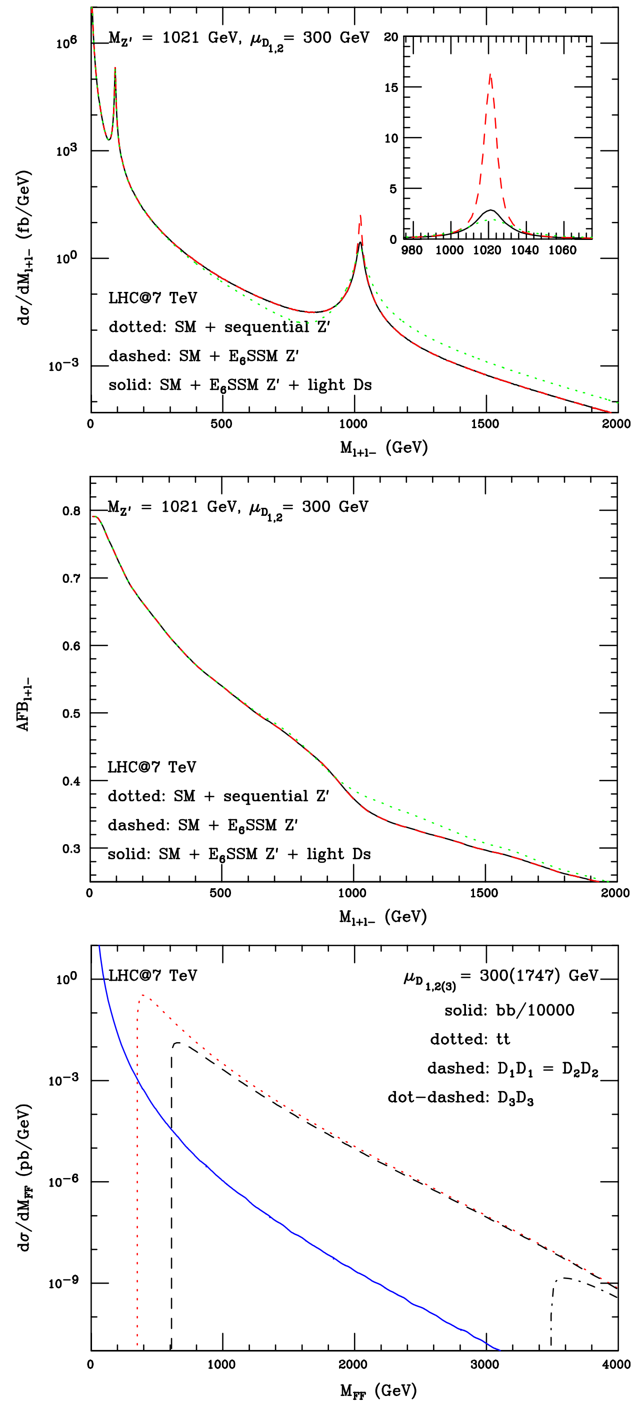


FIG. 1 (color online). Results for benchmark C at the 7 TeV LHC. *Top:* Differential cross sections for Drell-Yan production, with respect to the lepton pair invariant mass. *Middle:* Forward-backward asymmetries. *Bottom:* Production cross sections of exotic $D$-fermion pairs, in comparison to bottom- and top-quark pair production. The total production rates are $\sigma(D_1D_1) = \sigma(D_2D_2) = 3$ pb and $\sigma(D_3D_3) = 0.0005$ fb.

*Exotics.*—If exotic particles of the nature described here do exist at low scales, they could possibly be accessed through direct pair hadroproduction. However, as remarked in [1], the corresponding fully inclusive and differential cross

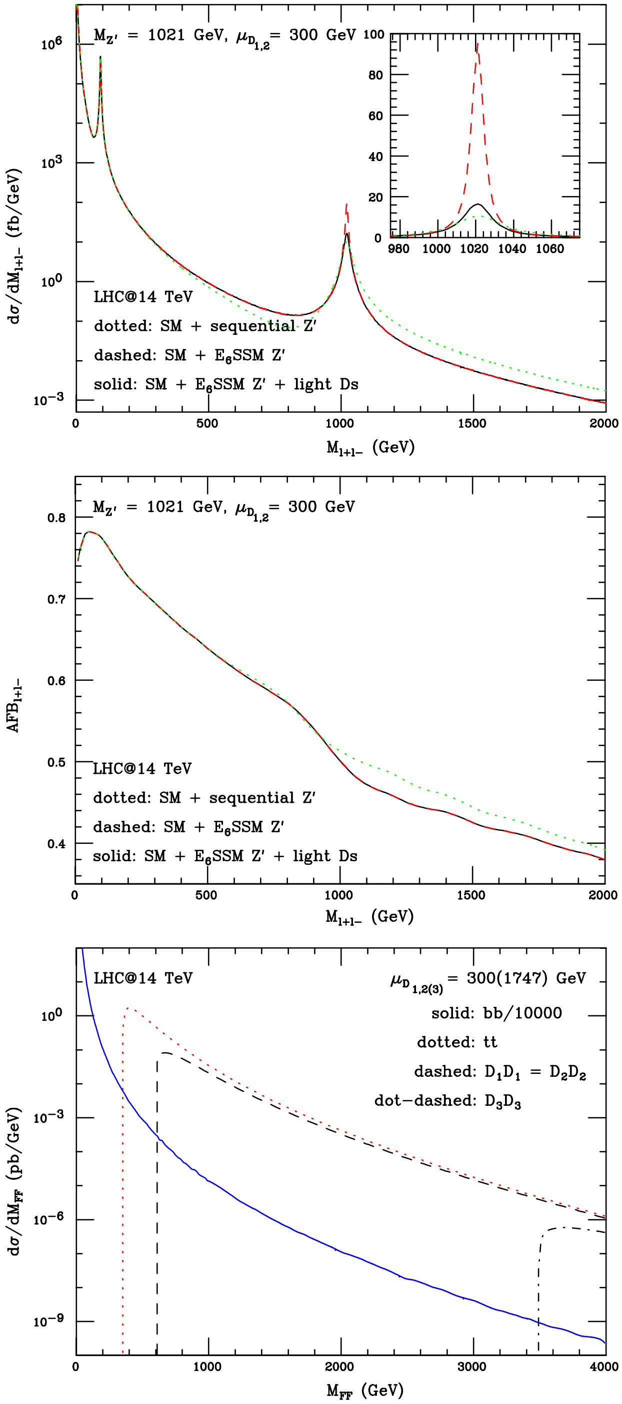


FIG. 2 (color online). Results for benchmark C at the 14 TeV LHC. *Top:* Differential cross sections for Drell-Yan production, with respect to the lepton pair invariant mass. *Middle:* Forward-backward asymmetries. *Bottom:* Production cross sections of exotic $D$-fermion pairs, in comparison to bottom- and top-quark pair production. The total production rates are $\sigma(D_1D_1) = \sigma(D_2D_2) = 25$ pb and $\sigma(D_3D_3) = 0.5$ fb.

sections are sufficient only in the case of exotic $D$-fermions (because they are pair-produced via QCD interactions) while inert Higgsinos most likely remain inaccessible (as their pair production is induced by EW interactions).

Therefore, we plot the production cross section of exotic $D$-fermion pairs, in comparison to those for bottom- and top-quark pair production, in the bottom frame of both Figs. 1 and 2, for an LHC with 7 and 14 TeV center-of-mass energy,, respectively, using CTEQ5L with $Q^2 = \hat{s}$. Although the detectable final states resulting from exotic $D$-fermion production do depend on the underlying nature of the exotic particles, we find that experimental signatures involve multijet states containing identifiable $b$-hadrons, whether produced via $t$-resonances or not, as we shall now discuss.

As outlined in [1], the lifetime and decay modes of the exotic $D$-fermions are determined by the operators that break the $Z_2^H$ symmetry. When $Z_2^H$ is broken significantly exotic fermions can produce a remarkable signature.[9] Since, according to our initial assumptions, the $Z_2^H$ symmetry is mostly broken by operators involving quarks and leptons of the third generation, the exotic $D$-fermions decay either via

$$\bar{D} \to t + \tilde{b}, \quad \bar{D} \to b + \tilde{t}, \quad D \to \bar{t} + \tilde{b}^*, \quad D \to \bar{b} + \tilde{t}^*,$$

if exotic $\bar{D}_i$ fermions are diquarks or via

$$D \to t + \tilde{\tau}, \quad D \to \tau^- + \tilde{t}, \quad \bar{D} \to \bar{t} + \tilde{\tau}^*, \quad \bar{D} \to \tau^+ + \tilde{t}^*,$$
$$D \to b + \tilde{\nu}_\tau, \quad D \to \nu_\tau + \tilde{b}, \quad \bar{D} \to \bar{b} + \tilde{\nu}_\tau^*, \quad \bar{D} \to \nu_\tau + \tilde{b}^*$$

if exotic $D$-fermions are leptoquarks. In general, sfermions decay into the corresponding fermion and a neutralino, so one expects that each diquark will decay into $t$- and $b$-quarks while a leptoquark will produce a $t$-quark and $\tau$-lepton in the final state with rather high probability. Thus the presence of light exotic $D$-fermions in the particle spectrum could result in an appreciable enhancement of the cross section of either $pp \to t\bar{t}b\bar{b} + X$ and $pp \to b\bar{b}b\bar{b} + X$ if exotic $D$-fermions are diquarks or $pp \to t\bar{t}\tau^+\tau^- + X$ and consequently $pp \to b\bar{b}\tau^+\tau^- + X$ if $D$-fermions are leptoquarks.[10]

Each $t$-quark decays into a $b$-quark while a $\tau$-lepton gives one charged lepton $l$ in the final state with a probability of 35%. Therefore both these scenarios would ultimately generate an excess in the $b$-quark production cross section. Thus the presence of exotic $D$-fermions alters the SM data samples involving $t\bar{t}$ production and decay as well as direct $b\bar{b}$ production.

[9] If $Z_2^H$ is only slightly broken, exotic quarks may live for a long time, and form compound states with ordinary quarks. This means that at future colliders it may be possible to study the spectroscopy of new composite scalar leptons or baryons. Also one can observe quasistable charged colorless fermions with zero lepton number.

[10] It is worth to remind the reader here that the production cross sections of $pp \to t\bar{t}b\bar{b} + X$ and $pp \to t\bar{t}\tau^+\tau^- + X$ in the SM are suppressed *at least* by a factor $(\frac{\alpha_s}{\pi})^2$ and $(\frac{\alpha_W}{\pi})^2$, respectively, as compared to the cross section of $t\bar{t}$ pair production (and, similarly, for $t$-quarks replaced by $b$-quarks).

Since the collider signatures associated with the $D$-fermions are so unique a detailed LHC analysis is required to establish the detectability of the corresponding processes at the LHC experiments. However, our results clearly show that, for the discussed parameter configuration, the position is favorable, as the product of production rates and branching ratios for these channels are typically larger than the expected four-body SM cross sections involving heavy quarks. For example, for BMC our estimations indicate that with 35 $\mathrm{pb}^{-1}$ of data the LHC experiments should have been produced about 200 pairs of $D$-fermions, giving serious reason to believe that BMC is on the edge of observability of the LHC experiments. Thus a detailed study of the detectability of the exotic $D$-fermions is an urgent necessity.

### *2. Benchmark point 4*

Having discussed in detail the phenomenology of the "early discovery" BMC, we now discuss BM4 in Table II, which represents the case of a heavier spectrum for the $cE_6SSM$, not necessarily discoverable at 7 TeV, hence dubbed "late discovery." BM4 is included in order to fairly show that the $cE_6SSM$ does not always lead to a light spectrum.

*$Z'_N$ bosons.*—For this parameter configuration, with a rather heavy $Z'_N$, cross sections are much smaller, beyond detectability at the 7 TeV LHC. We therefore only present results for the higher-energy stage of the CERN collider, in Fig. 3 (top and middle frame) for the $Z'$ line shape and forward-backward asymmetry. The pattern that emerges here is very much in line with that of the previous benchmark, albeit with reduced production rates overall. However, the $Z'_N$ should remain detectable at the 14 TeV LHC after full luminosity is collected (Also note that the absolute value of the corrections to the $Z'_N$ width due to $cE_6SSM$ particles is somewhat larger here, growing by about 13 GeV.). Hence, cross section, line shape, and forward-backward asymmetry studies are feasible for BM4 too and should enable the accurate profiling of the $Z'_N$ state.

*Exotics.*—The exotic $D$-fermions are much too heavy for this benchmark and their detectability, even at 14 TeV, will be challenging, though not impossible. This is illustrated in Fig. 3 (bottom frame), where their inclusive cross section is shown to be at the fb level (including all three generations) and to require very high invariant masses for the final state, where the control of the SM background is more uncertain. There is however scope at large luminosities (the situation here is not dissimilar from the case of $D_3$ at the 14 TeV LHC for BMC).

An interesting feature of BM4 is that it contains relatively light exotic scalars ($\tilde{D}_1$), for all generations, unlike the case of BMC. Because these exotic scalars have masses about 312 GeV they are expected to be leptoquarks. Such light leptoquarks will be efficiently produced at the LHC, both at 7 and particularly 14 TeV, where the cross section is approximately 0.53 and 4.9 pb, respectively. They decay into quark-lepton final states mainly through $Z_2^H$ violating operators involving quarks and leptons of the

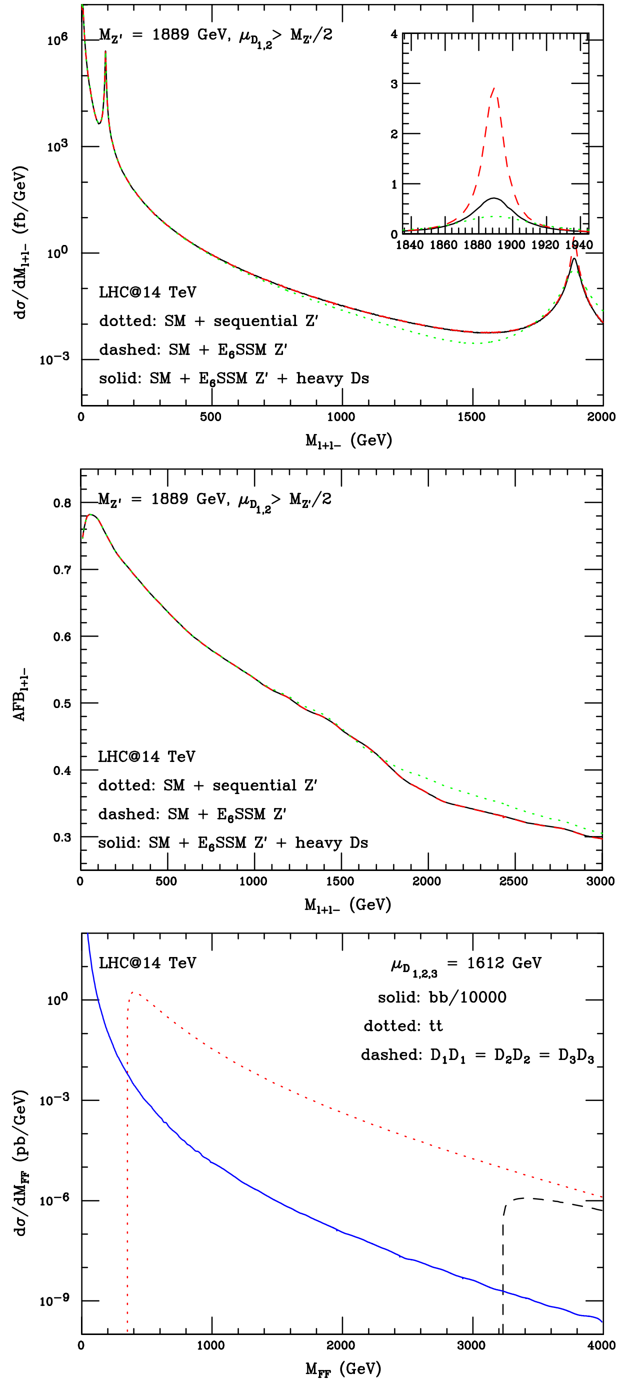


FIG. 3 (color online). Results for benchmark 4 at the 14 TeV LHC. *Top:* Differential cross sections for Drell-Yan production, with respect to the lepton pair invariant mass. *Middle:* Forward-backward asymmetries. *Bottom:* Production cross sections of exotic $D$-quark pairs, in comparison to bottom- and top-quark pair production. The total production rates are $\sigma(D_1D_1) = \sigma(D_2D_2) = \sigma(D_3D_3) = 0.9$ fb.

third generation, i.e., $\tilde{D} \to t\tau$. This leads to an enhancement of $pp \to t\bar{t}\tau\bar{\tau}$ (without missing energy) at the LHC.

## IV. CONCLUSIONS

We have previously proposed a constrained version of the exceptional supersymmetric standard model, the cE$_6$SSM, based on a universal high-energy soft scalar mass $m_0$, soft trilinear mass $A_0$ and soft gaugino mass $M_{1/2}$. The cE$_6$SSM predicts a characteristic SUSY spectrum containing a light gluino, a light winolike neutralino and chargino pair, and a light binolike neutralino, with other sparticle masses except the lighter stop being much heavier. In addition, cE$_6$SSM allows the possibility of light exotic color triplet charge $1/3$ $D$ fermions and scalars, and predicts an observable $Z'_N$ spin-1 gauge boson.

In this paper, motivated by the fact that the cE$_6$SSM allows the spectrum above to be quite light and observable with the first data from the LHC, we have focused on two of the most characteristic and striking LHC signatures of the cE$_6$SSM, namely, the prediction of a $Z'_N$ gauge boson and exotic $D$-fermions, and the interplay between these two predictions. In particular we have shown how the $Z'_N$ gauge boson may decay into exotic $D$-fermions, increasing its width and modifying its line shape. For example, we find that the width may increase by a factor of 2, which effectively reduces the Drell-Yan cross section into charged lepton pairs also by a factor of 2, relaxing the current Tevatron limits from 892 GeV down to about 820 GeV. In addition we have calculated the LHC production cross section of the $D$-fermions and discussed their decay patterns.

The added value of the cE$_6$SSM, compared to previous studies, is that it provides a predictive framework for the experimental study of such signatures via the use of benchmark points. We illustrated this by considering two of the benchmark points previously proposed in some detail. The first benchmark point C, which has low values of $(m_0, M_{1/2})$ around (700, 400) GeV and a $Z'_N$ gauge boson with mass around 1 TeV, gave rise to signatures corresponding to an "early LHC discovery" using "first data."[11] We also examined benchmark point 4 with higher values of $(m_0, M_{1/2})$ around (1100, 700) GeV and a $Z'$ gauge boson with mass around 2 TeV, providing a more challenging scenario corresponding to late discovery using all accumulated data at the CERN collider. Further, in both scenarios, the singlinos are always light (this is a generic feature of the model in fact) and thus contribute a very sizable amount to the $Z'_N$ width (also thanks to their strong couplings to the latter), so that they could possibly be accessed at the LHC from the study of the $Z'_N$ line shape, whichever the $Z'_N$ mass. For both benchmark scenarios we find rather copious production of exotics, in the case of point C primarily of $D$-fermions while in the case of point 4 primarily of $\tilde{D}$-scalars, both yielding peculiar signatures involving third-generation SM fermions (i.e., $b$ and/or $t$-(anti)quarks plus $\nu_\tau$ and/or $\tau$-(anti)leptons). Finally, note that the above values of $(m_0, M_{1/2})$ in the cE$_6$SSM yield a squark and gluino spectrum roughly equivalent to that in the cMSSM with $m_0$ about $3/2$ times larger and $M_{1/2}$ about 4 times smaller than the corresponding cE$_6$SSM values.

If a $Z'_N$ gauge boson and/or $D$-exotics (fermions or scalars, depending on the model configuration) were discovered at the LHC, identified by measurements of their mass, cross section and decay signatures as discussed here, this would not only represent a revolution in particle physics, but would also point towards a possible underlying high-energy E$_6$ gauge structure, providing the first glimpse into superstring theory.

[11] After we submitted our paper for publication ATLAS collaboration published new results of the search for SUSY particles [35]. These results set a more stringent lower limit on the gluino mass and, as a consequence, a more stringent upper bound on the gluino production rate at the LHC. It seems that the new ATLAS exclusion limits rule out the BMC type spectrum. Despite the fact that BMC might be already excluded it still us allows to demonstrate at least two crucial features of the cE$_6$SSM. Indeed, within the cE$_6$SSM one can always find a solution with relatively small $\kappa_{1,2}$ so that the corresponding $D$ fermion states are light giving rise to remarkable signature that might be observed at the LHC in the near future. BMC also demonstrates that $D$-fermions and other exotic states give a substantial contribution to the $Z'$ width if these states are light.


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## APPENDIX: FEYNMAN RULES

In this appendix the $Z'_N$ Feynman rules of the E$_6$SSM for the considered benchmarks are presented.

The couplings shown in Fig. 4 are determined as follows. For the scalar partners of fermions with substantial mixing the couplings are given by

$$\begin{aligned} f_{11} &= (\tilde{Q}^N_{f_L}\cos^2\theta_{\tilde{f}} - \tilde{Q}^N_{f^c}\sin^2\theta_{\tilde{f}}), \\ f_{22} &= (\tilde{Q}^N_{f_L}\sin^2\theta_{\tilde{f}} - \tilde{Q}^N_{f^c}\cos^2\theta_{\tilde{f}}), \\ f_{12} &= f_{21} = -(\tilde{Q}^N_{f^c} + \tilde{Q}^N_L)\sin\theta_{\tilde{f}}\cos\theta_{\tilde{f}}, \end{aligned} \tag{A1}$$

where $\theta_{\tilde{f}}$ is the mixing

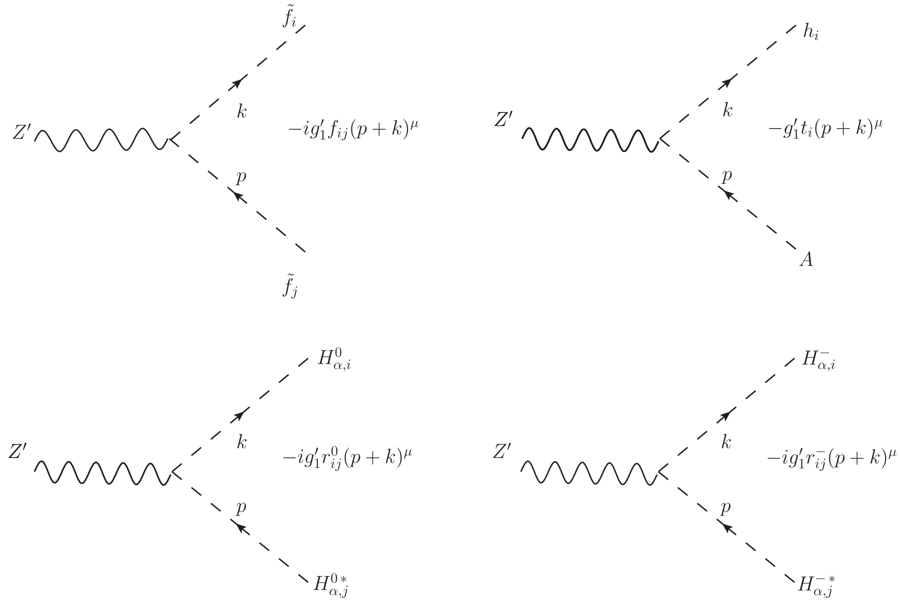


FIG. 4. Feynman rules: $Z'_N$ coupling to scalars.

$$\begin{pmatrix} \tilde{f}_1 \\ \tilde{f}_2 \end{pmatrix} = \begin{pmatrix} \cos\theta_{\tilde{f}} & \sin\theta_{\tilde{f}} \\ -\sin\theta_{\tilde{f}} & \cos\theta_{\tilde{f}} \end{pmatrix} \begin{pmatrix} \tilde{f}_L \\ \tilde{f}_R \end{pmatrix} \tag{A2}$$

The relation between the Higgs gauge and mass eigenstates is $H_i^0 = U_{ji}^{-1} H_j + i V_{ji}^{-1} A_j$, where $H_i^0 = \{H_u^0, H_d^0, S\}$ and $A_j = \{A, G', G^0\}$, where the form of $V$ can be read off from Eqs. (58)–(59) of Ref. [1] and $U$ is found when the $CP$-even Higgs mass matrix is diagonalized. The Higgs $Z'_N$ Feynman rules shown in Fig. 4 then take the form $t_j = Q'_i U_{ji}^{-1} V_{1i}^{-1} (p+k)^\mu$.

The inert Higgs come from two generations of "up" and "down" type doublets and each generation has 8 degrees of freedom. The charged and neutral components are almost degenerate, but are split by $D$-term contributions. The physical states are formed by mixing the up and down type Higgs as follows:

$$H^0_{\alpha,1} = \cos\theta^0_\alpha H^{d,0}_\alpha + \sin\theta^0_\alpha H^{u,0}_\alpha, \tag{A3}$$

$$H^0_{\alpha,2} = \cos\theta^0_\alpha H^{u,0}_\alpha - \sin\theta^0_\alpha H^{d,0}_\alpha, \tag{A4}$$

$$H^-_{\alpha,1} = \cos\theta^-_\alpha H^{d,-}_\alpha + \sin\theta^-_\alpha H^{u,+*}_\alpha, \tag{A5}$$

$$H^-_{\alpha,2} = \cos\theta^-_\alpha H^{u,+*}_\alpha - \sin\theta^-_\alpha H^{d,-}_\alpha. \tag{A6}$$

and the couplings shown in Fig. 4 are then of the form

$$\begin{aligned} r^0_{11} &= (\tilde{Q}^N_{H_1}\cos^2\theta^0 - \tilde{Q}^N_{H_2}\sin^2\theta^0), \\ r^-_{11} &= (\tilde{Q}^N_{H_1}\cos^2\theta^- - \tilde{Q}^N_{H_2}\sin^2\theta^-) \end{aligned} \tag{A7}$$

TABLE IV. Scalar couplings to $Z'_N$.

| | BMC | BM4 |
|---|---|---|
| Stops $g'_1 f_{11}$ | −0.0.04537 | −0.05128 |
| Stops $g'_1 f_{22}$ | 0.05827 | −0.06423 |
| Stops $g'_1 f_{12}$ | 0.04518 | −0.04682 |
| Sbottoms $g'_1 f_{11}$ | −0.07910 | −0.07624 |
| Sbottoms $g'_1 f_{22}$ | −0.1592 | −0.1571 |
| Sbottoms $g'_1 f_{12}$ | −0.01574 | −0.03302 |
| Staus $g'_1 f_{11}$ | 0.09303 | 0.09377 |
| Staus $g'_1 f_{22}$ | 0.1474 | −0.1487 |
| Sups $g'_1 f_{11}$ | −0.06540 | −0.06787 |
| Sups $g'_1 f_{22}$ | 0.07812 | 0.080821 |
| Sdowns $g'_1 f_{11}$ | 0.07812 | 0.0821 |
| Sdowns $g'_1 f_{22}$ | −0.1562 | −0.1616 |
| Selectron $g'_1 f_{11}$ | 0.09306 | 0.09377 |
| Selectron $g'_1 f_{22}$ | −0.14737 | −0.1487 |
| Scalar exotic D's 3rd Gen $g'_1 f_{11}$ | 0.06668 | 0.04380 |
| Scalar exotic D's 3rd Gen $g'_1 f_{22}$ | 0.001274 | 0.02407 |
| Scalar exotic D's 3rd Gen $g'_1 f_{12}$ | 0.1859 | 0.1953 |
| Scalar exotic D's 1st/2nd Gen $g'_1 f_{11}$ | 0.04114 | 0.04380 |
| Scalar exotic D's 1st/2nd Gen $g'_1 f_{22}$ | 0.002426 | 0.02407 |
| Scalar exotic D's 1st/2nd Gen $g'_1 f_{12}$ | 0.1888 | 0.1953 |
| Neutral inert Higgs $g'_1 r^0_{11}$ | −0.1581 | −0.1042 |
| Neutral inert Higgs $g'_1 r^0_{22}$ | 0.06511 | 0.01042 |
| Neutral inert Higgs $g'_1 r^0_{12}$ | 0.1585 | 0.1870 |
| Charged inert Higgs $g'_1 r^-_{11}$ | −0.1443 | −0.1013 |
| Charged inert Higgs $g'_1 r^-_{22}$ | 0.05129 | 0.007607 |
| Charged inert Higgs $g'_1 r^-_{12}$ | 0.1674 | 0.1878 |
| Higgs $g'_1 t_1$ | −0.01964 | −0.006459 |
| Higgs $g'_1 t_2$ | 0.1183 | 0.1210 |
| Higgs $g'_1 t_3$ | 0.002314 | 0.0001778 |

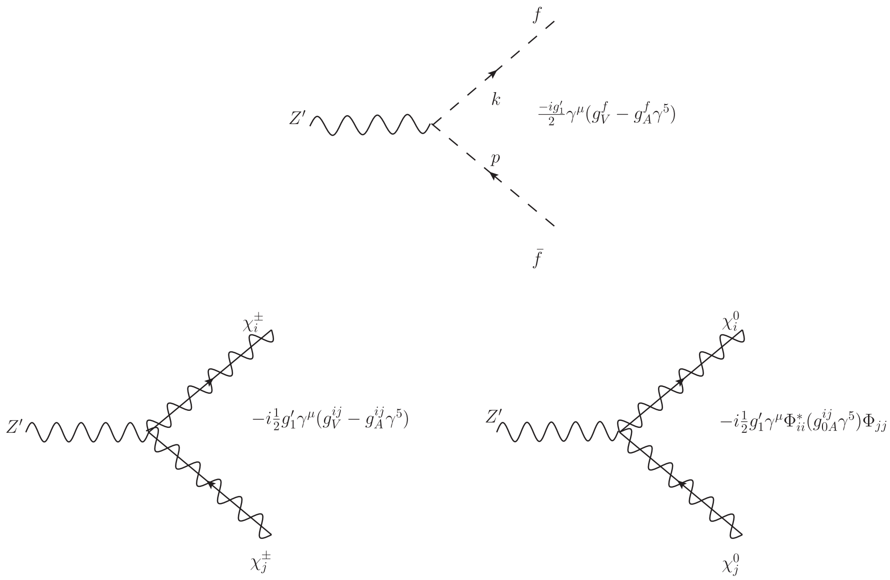


FIG. 5. Feynman rules: $Z'$ couplings to fermions.

TABLE V. Vector and axial fermion couplings to $Z'_N$.

| | BMC Vector $g_V$ | BMC Axial $g_A$ | BM4 Vector $g_V$ | BM4 Axial $g_A$ |
|---|---|---|---|---|
| $Z'l\bar{l}$ | 0.1108 | 0.4901 | 0.1110 | 0.4900 |
| $Z'\nu_l\bar{\nu}_l$ | 0.3004 | 0.3004 | 0.3005 | 0.3005 |
| $Z'u\bar{u}$ | 0.02630 | 0.3004 | 0.02618 | 0.3005 |
| $Z'd\bar{d}$ | −0.1634 | 0.4901 | −0.1633 | 0.4900 |
| $Z'D_i\bar{D}_i$ | 0.1371 | −0.7906 | 0.1372 | −0.7906 |
| $Z'\tilde{H}_\alpha\bar{\tilde{H}}_\alpha$ | −0.1897 | −0.7906 | −0.1895 | −0.7906 |
| $Z'\tilde{s}_\alpha\bar{\tilde{s}}_\alpha$ | 0.7906 | 0.7906 | 0.7906 | 0.7906 |
| $Z'\chi_1^+\bar{\chi}_1^+$ | −0.013565 | −0.01372 | 0.003494 | −0.003558 |
| $Z'\chi_2^+\bar{\chi}_2^+$ | −0.1760 | −0.7768 | 0.1861 | −0.7870 |
| $Z'\chi_1^+\bar{\chi}_2^+$ | −0.07755 | −0.08396 | 0.04450 | −0.03834 |
| $Z'\chi_1^0\bar{\chi}_1^0$ | 0 | −0.002220 | 0 | −0.0004995 |
| $Z'\chi_1^0\bar{\chi}_2^0$ | 0 | −0.003660 | 0 | 0.0008260 |
| $Z'\chi_1^0\bar{\chi}_3^0$ | 0 | −0.03365 | 0 | −0.01560 |
| $Z'\chi_1^0\bar{\chi}_4^0$ | 0 | −0.003298 | 0 | 0.01565 |
| $Z'\chi_1^0\bar{\chi}_5^0$ | 0 | −0.03511 | 0 | 0 |
| $Z'\chi_1^0\bar{\chi}_6^0$ | 0 | −0.001813 | 0 | 0.0007258 |
| $Z'\chi_2^0\bar{\chi}_2^0$ | 0 | −0.006037 | 0 | −0.001366 |
| $Z'\chi_2^0\bar{\chi}_3^0$ | 0 | −0.05440 | 0 | 0.02579 |
| $Z'\chi_2^0\bar{\chi}_4^0$ | 0 | −0.005385 | 0 | −0.02588 |
| $Z'\chi_2^0\bar{\chi}_5^0$ | 0 | −0.0005868 | 0 | 0 |
| $Z'\chi_2^0\bar{\chi}_6^0$ | 0 | −0.002947 | 0 | −0.001222 |
| $Z'\chi_3^0\bar{\chi}_3^0$ | 0 | −0.04902 | 0 | −0.4866 |
| $Z'\chi_3^0\bar{\chi}_4^0$ | 0 | −0.4852 | 0 | 0.4890 |
| $Z'\chi_3^0\bar{\chi}_5^0$ | 0 | −0.007955 | 0 | 0.03565 |
| $Z'\chi_3^0\bar{\chi}_6^0$ | 0 | −0.02591 | 0 | 0.03806 |
| $Z'\chi_4^0\bar{\chi}_4^0$ | 0 | −0.4791 | 0 | −0.4901 |
| $Z'\chi_4^0\bar{\chi}_5^0$ | 0 | −0.05635 | 0 | 0.009647 |
| $Z'\chi_4^0\bar{\chi}_6^0$ | 0 | −0.05165 | 0 | −0.01211 |
| $Z'\chi_5^0\bar{\chi}_5^0$ | 0 | −0.4569 | 0 | 0.4559 |
| $Z'\chi_5^0\bar{\chi}_6^0$ | 0 | 1.08906 | 0 | 1.08910 |
| $Z'\chi_6^0\bar{\chi}_6^0$ | 0 | 0.52064 | 0 | 0.52270 |

$$r^0_{22} = (\tilde{Q}^N_{H_1}\sin^2\theta^0 - \tilde{Q}^N_{H_2}\cos^2\theta^0),$$
$$r^-_{22} = (\tilde{Q}^N_{H_1}\sin^2\theta^- - \tilde{Q}^N_{H_2}\cos^2\theta^-) \tag{A8}$$

$$r^0_{12} = (\tilde{Q}^N_{H_2} + \tilde{Q}^N_{H_1})\sin\theta^0\cos\theta^0,$$
$$r^-_{12} = (\tilde{Q}^N_{H_2} + \tilde{Q}^N_{H_1})\sin\theta^-\cos\theta^-. \tag{A9}$$

The numerical values of the scalar couplings for our benchmarks are given in Table IV.

The fermion Feynman rules are shown in Fig. 5. The chargino masses are found by a bi-unitary diagonalization of the chargino mass matrix

$$U^*_- X U^{-1}_+ = M_{ch} = \begin{pmatrix} m_{\chi^\pm_2} & 0 \\ 0 & m_{\chi^\pm_1} \end{pmatrix}$$

$$\text{where } U_\pm = \begin{pmatrix} \cos\theta_\pm & \sin\theta_\pm \\ -\sin\theta_\pm & \cos\theta_\pm \end{pmatrix} \quad \text{if } \det(X) > 0$$

and for $\det(X) < 0$ $U_+ \to \sigma_3 U_+$ gives us the correct matrix to diagonalize $X$ such that all masses are positive.

This leads to the chargino couplings taking the form

$$g^{ij}_V = \tilde{Q}_{H_2}U_{+i2}U_{+j2} - \tilde{Q}_{H_1}U_{-i2}U_{-j2}, \tag{A10}$$

$$g^{ij}_A = \tilde{Q}_{H_2}U_{+i2}U_{+j2} + \tilde{Q}_{H_1}U_{-i2}U_{-j2}. \tag{A11}$$

For neutralino couplings the situation is similar, but in this case we have neutral Majorana fermions so the vector couplings vanish. In addition, when we diagonalize the mass matrix numerically, we find the mixing matrix $N$ which diagonalizes the neutralino mass matrix though $N^* M_{\rm neut} N^{-1}$ to give diagonal masses, $m(i)$, which can be negative or positive. To obtain positive masses we can then also perform a phase rotation $\Phi^*(N^* M_{\rm neut} N^{-1})\Phi^{-1}$ where $(\Phi)_{jk} = (i)^{\theta(j)}\delta_{jk}$, where $\theta(j) = 0(1)$ if $m(j)$ is positive (negative). Neutralino couplings take then the form

$$g^{ij}_A = \sum_k 2Q_k N_{ik} N^*_{jk} \quad \text{where } Q_k = (0, 0, \tilde{Q}_{H_1}, \tilde{Q}_{H_2}, \tilde{Q}_S, 0). \tag{A12}$$

The form of the couplings $g^f_V$ and $g^f_A$ were given in and are only reproduced here for convenience. For the fermions of ordinary matter one has $g^f_V = \tilde{Q}_{f_L} - \tilde{Q}_{f^c}$ and $g^f_A = \tilde{Q}_{f_L} + \tilde{Q}_{f^c}$. For the exotic colored objects we have similarly: $g^D_V = \tilde{Q}_D - \tilde{Q}_{\bar{D}}$ and $g^D_A = \tilde{Q}_D + \tilde{Q}_{\bar{D}}$. For the inert Higgsinos one gets: $g^{\tilde{H}_\alpha}_V = \tilde{Q}_{H1} - \tilde{Q}_{H2}$ and $g^{\tilde{H}_\alpha}_A = \tilde{Q}_{H1} + \tilde{Q}_{H2}$.

The numerical values of these couplings for the benchmarks studied in this paper are shown in Table V.


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