Surprising phenomenology in the non-Universal U(1) gauge extended $\mu\nu$SSM

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Abstract

So far the most sophisticated experiments have shown no trace of new physics in the TeV scale. Consequently, new models with unexplored parameter regions are necessary to explain current results, analyse further the existing data, and propose new experiments. In this letter, we present a modified version of the $\mu\nu$SSM supersymmetric model where a non-Universal extra U(1) gauge symmetry is added in order to restore an effective R-parity that ensures proton stability. We show that anomalies equations cancel without having to add any exotic matter, and find that it is the viability of the model through anomalies cancellation what defines the conditions in which fermions interact with dark matter candidates via the exchange of $Z'$ bosons. The strict condition of universality violation across families in the charge under the extra gauge symmetry means that LHC constraints for a $Z'$ mass do not apply directly to our model, allowing for a yet undiscovered light $Z'$, as we discuss. Moreover, we explore the possibility of isospin violating interactions of dark matter with nuclei for in model; we observe that this interaction depends, surprisingly, on the Higgs charges under the new symmetry, both limiting the number of possible models and allowing to analyse indirect dark matter searches in the light of well defined, particular scenarios.
1 Introduction

Barring the Higgs discovery \[1,2\], no signs of new physics beyond the Standard Model (SM) have been seen so far after run 1 of the LHC. In particular, regarding Supersymmetry (SUSY), there are no signals of the coloured states \[3\], namely squarks and gluinos, that were predicted to be abundant in the TeV scale. These data impose severe constraints on SUSY models, pushing the coloured states to masses beyond 1 TeV. However, several recast analyses showed that –despite the not very satisfactory run 1 LHC data– there still is much room for light SUSY states \[4–6\].

The minimal supersymmetric standard model (MSSM) is the most simple realisation of a \(N = 1\) SUSY model. In the MSSM construction however, the mass term responsible of the electroweak symmetry breaking (EWSB) is added \(\text{ad hoc}\), not specifying its origin \[7–10\]. Another important issue is that the MSSM is unable to explain the fact that neutrinos do have mass \[11, 12\]. An elegant proposal to solve both problems at once comes from the so-called “\(\mu\) from \(\nu\)” supersymmetric model (\(\mu\nu\)SSM) \[1\], which adding right-handed neutrinos efficiently solves the \(\mu\)-problem and clarifies the origin of the neutrinos masses, achieving it only with two new terms in the superpotential, as can be seen in Eq. (1.1). The first one, \(\lambda_i \tilde{\nu}^c_i \tilde{H}_d \tilde{H}_u, i = 1, 2, 3\), is a trilinear coupling between the Higgs and the three families of right-handed neutrinos. When EWSB takes place the supersymmetric partners of the right-handed neutrinos, the right-handed sneutrinos \(\tilde{\nu}^c_i\), develop vacuum expectation values (VEVs), giving rise to an effective \(\mu\) term, \(\mu \sim \lambda v^c\). The second new term in the superpotential, \(\kappa_{ijk} \tilde{\nu}^c_i \tilde{\nu}^c_j \tilde{\nu}^c_k\), provides Majorana masses to the right-handed neutrinos after EWSB takes place. As the right neutrinos couple to the leptons, a low-scale see-saw mechanism is induced and the light left-handed neutrinos become massive \[14,15\].

\[
W = \epsilon_{ab} \left( Y^{ij}_u \tilde{H}^b_2 \tilde{Q}^a_i \tilde{\nu}^c_j + Y^{ij}_d \tilde{H}^a_1 \tilde{Q}^b_i \tilde{\nu}^c_j + Y^{ij}_e \tilde{H}^b_1 \tilde{L}^a_i \tilde{\nu}^c_j + Y^{ij}_\nu \tilde{H}^b_2 \tilde{L}^a_i \tilde{\nu}^c_j \right) - \epsilon_{ab} \lambda^i \tilde{\nu}^c_i \tilde{H}^a_2 \tilde{H}^b_1 + \frac{1}{3} \kappa_{ijk} \tilde{\nu}^c_i \tilde{\nu}^c_j \tilde{\nu}^c_k
\]

As a result of the heavy mixing occurring within the neutral and charged sectors, the \(\mu\nu\)SSM presents a very rich phenomenology, markedly different from the usual collider scenarios \[16–21\]. This means not only that new parameter regions open up for SUSY searches but also that the \(\mu\nu\)SSM model predictions could have escaped unnoticed so far. Nonetheless the \(\mu\nu\)SSM has issues as well; both new terms added to the superpotential explicitly break \(R\)-parity (\(R_p\)). Such breaking comes directly from the right-handed neutrinos and it is governed by the value of the neutrino Yukawa coupling, \(Y_\nu\). As \(R_p\) is no longer a symmetry of the model, dangerous lepton and baryon number violating terms are allowed in the superpotential. Moreover, the stability of the proton is no longer guaranteed. In order to forbid such terms one can either impose a discrete symmetry or a \(U(1)\) gauge symmetry. The former introduces a cosmic domain wall problem \[22,22\], while the latter is usually present in stringy realisations of the SM (see for example \[23–26\]).

The idea of adding an extra \(U(1)\) symmetry to the gauge group of the SM is by no means new \[27\], and adding it to the MSSM is also a well trodden path, of which \[28,35\] are only a few examples. In fact for the \(\mu\nu\)SSM it has already been tentatively explored \[36\],

\[^{1}\text{See \[13\] and references therein for an extensive review.}\]
leading to promising results. The price to pay however is having to recalculate the anomalies cancellation conditions, which for the SM matter content and gauge groups are known to be “miraculously” fulfilled. In Ref. \[36\] it was found that for the $\mu\nuSSM$ to be consistent with an extra U(1) gauge symmetry, the matter content of the model has to be enlarged by several extra colour triplets, left doublets and singlet fields. Furthermore, it is a general rule for all gauge extended models that exotic fields are needed to cancel all anomalies. While the possibility of exotic fields cannot be excluded its presence is problematic not only due to lack of evidence but because it disrupts the unification of coupling constants at the GUT scale. Hence a minimalist solution is always desired.

In this letter we present a solution of the U(1) enlarged $\mu\nuSSM$ which is anomaly free by means of having non-universal charges of the superfields under the extra symmetry. In our model, neither meaningless singlet fields, nor exotic fields are needed, which to our knowledge is a first on gauge extended models. We solve the anomalies equations by assuming that each of the families can have a different $U(1)_X$ charge, finding several groups of solutions depending on few unconstrained extra charges. In addition, we explore some possible phenomenological consequences of having non-Universality in our model. Concretely, we study the extra charge dependence of the $Z'$ interaction with fermions as a result of the mixing between the extra U(1) boson and the usual Z boson \[37,38\]. Such dependence implies that the production limits of a $Z'$ at the LHC are no longer valid and have to be recalculated for the specific models allowed by the anomalies cancellation, leading to scenarios where a light $Z'$ is possible, a common feature of stringy constructions \[37,39,40\]. In addition, we study possible scenarios of $Z'$ mediated spin independent dark matter (DM) interactions \[41\], finding a family of anomalies equations solutions with isospin violating scalar and vector DM-quark interactions. For the non-Universal U(1) gauge enlarged $\mu\nuSSM$ the possible isospin violating scenarios are parametrised by the Higgs fields charges under the extra symmetry, rendering a series of finite, discrete values that could be discriminated in experiments. Hence experimental detection of DM is not only conditioned by the specific model realisation but could also be used to provide with clear, testable predictions to discern among DM and string compactification scenarios.

2 The non-Universal U(1)$\mu\nuSSM$

In this section we present the necessary conditions for anomalies cancellations and the implications for model building. In addition, we explore the possible phenomenology characteristic of the model, which comes imposed by the extra charges assignment of the fields, which themselves are constrained by the anomalies cancellation conditions.

2.1 Non-universal anomaly cancellation in the U(1)$\mu\nuSSM$

The gauge group of the U(1)$\mu\nuSSM$ is $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X$, where each of the superfields composing the model’s spectrum has now an extra ($Q_X$) charge. Consequently, all anomalies equations involving the new symmetry have to be recalculated if we want the model to be anomaly free. The analysis is two-fold: On the one hand, the terms appearing in the superpotential must have vanishing total charge, and on the other hand, anomalies must also cancel. Hence, restrictive bounds are imposed on each superfield charge $Q_X$. Furthermore, we can use certain constraints to not only allow terms in the superpotential, but also to explicitly banish undesirable terms.
We will work under the assumption that no exotic matter is needed. To that end, we consider the same matter content as for the original $\mu\nu$SSM. The anomalies equations that must be fulfilled with this matter content are

\[ \sum_i (2Q_{Q_i} + Q_{u_i} + Q_{d_i}) = 0 \]
\[ \sum_i (3Q_{Q_i} + Q_{L_i}) + Q_{H_1} + Q_{H_2} = 0 \]
\[ \sum_i (\frac{1}{6}Q_{Q_i} + \frac{1}{3}Q_{d_i} + \frac{4}{3}Q_{u_i} + \frac{1}{2}Q_{L_i} + Q_{e_i}) + \frac{1}{2}(Q_{H_1} + Q_{H_2}) = 0 \]
\[ \sum_i (Q_{Q_i}^2 + Q_{d_i}^2 - 2Q_{u_i}^2 - Q_{L_i}^2 + Q_{e_i}^2) - Q_{H_1}^2 + Q_{H_2}^2 = 0 \]
\[ \sum_i (6Q_{Q_i}^3 + 3Q_{d_i}^3 + 3Q_{u_i}^3 + 2Q_{L_i}^3 + Q_{e_i}^3 + Q_{\nu_i}^3) + 2Q_{H_1}^3 + 2Q_{H_2}^3 = 0 \]
\[ \sum_i (6Q_{Q_i} + 3Q_{u_i} + 3Q_{d_i} + 2Q_{L_i} + Q_{e_i} + Q_{\nu_i}) + 2Q_{H_1} + 2Q_{H_2} = 0. \]

To solve equations 2.2 we need a set of constraints, which we in addition use to ensure that our model has certain desired properties arising naturally. Prime among them is forbidding a $\mu$ term from appearing in the superpotential, as its absence is not otherwise automatically guaranteed. Thus, we impose $Q_{H_1} \neq -Q_{H_2}$. Furthermore, the Higgs mass term is obtained from the right sneutrinos $\nu^c$ singlet fields acquiring VEVs at the EWSB scale. Therefore a term coupling the right neutrinos and the Higgs fields must also be allowed in the superpotential, which imposes $Q_{H_1} + Q_{H_2} + Q_{\nu^c} = 0$ for at least one of the three families of right neutrinos. Moreover, as the $\mu\nu$SSM was born to answer the neutrino mass problem, we impose that tree-level mass terms below the soft breaking scale must appear for the right neutrinos, such that a see-saw mechanism is implemented in our model. It is therefore a condition that $Q_{L_i} + Q_{\nu^c} + Q_{H_2} = 0$. In addition, we would like to directly forbid certain operators –such as those violating baryon number– from the superpotential, which means $Q_u \neq -2Q_d$. The remaining mass terms are $a\ priori$ not imposed in the superpotential, permitting thus the different fields to acquire mass either at tree-level order with Yukawa couplings or at first order in perturbation theory, via non-holomorphic 1-loop mass terms. The choice of either is to be fixed accordingly with the anomalies equations.

Giving mass to certain fields via non-holomorphic couplings means that such mass must be provided by SUSY-breaking operators introduced via radiative, first loop corrections, and which appear naturally in gravity mediated SUSY-breaking scenarios [42]. It has been demonstrated that either mechanism is in principle indistinguishable in experiments but for the heaviest particles [43], namely the top quark and the $\tau$ lepton, for which tree level Yukawa terms need to be imposed. Thus $Q_{Q_3} + Q_{u_3} + Q_{H_2} = 0$ and $Q_{L_3} + Q_{e_3} + Q_{H_1} = 0$. With this conditions, the first, second, and sixth equations in 2.2 become

\[ 2Q_{Q_1} + 2Q_{Q_2} + Q_{Q_3} + Q_{d_1} + Q_{d_2} + Q_{d_3} + Q_{u_1} + Q_{u_2} - Q_{H_2} = 0 \]
\[ Q_{H_1} + Q_{H_2} + Q_{L_1} + Q_{L_2} + Q_{L_3} + 3Q_{Q_1} + 3Q_{Q_2} + 3Q_{Q_3} = 0 \]
\[ 3Q_{d_1} + 3Q_{d_2} + 3Q_{d_3} + 3Q_{e_1} + 3Q_{e_2} + Q_{H_1} + Q_{L_1} + Q_{L_2} \]
\[ + 6Q_{Q_1} + 6Q_{Q_2} + 3Q_{Q_3} + 3Q_{u_1} + 3Q_{u_2} - 4Q_{H_2} = 0 \]
which we can use to fix the conditions for quark charges. Should we try for all the up quarks to have tree-level mass terms, then all down quarks necessarily acquire mass through non-holomorphic terms. But from the third equation this imposes that just one lepton has tree-level Yukawa mass term. Putting everything back into the second equation it would lead to $Q_{H_1} = -Q_{H_2}$, reintroducing the $\mu$ term in the superpotential. And the same happens if we try to have both first and second family of up quarks with non-holomorphic mass terms. Consequently, the only possibility is for either the first or the second family of up quarks to have tree-level Yukawa coupling, while the other acquires its mass through a non-holomorphic term. On the contrary, the necessary condition for the down-type quarks is to have two of the families having non-holomorphic mass terms and one a tree-level Yukawa. There is however freedom in choosing which family acquires its mass via which mechanism, a fact that will be of importance for the phenomenology of the model as we shall see in what follows. The left leptons (L) mimic the behaviour of the up-quarks.

To conclude the analysis, we must choose either of the groups of solutions, the rest being symmetric. In particular, fixing the first and third families of up-quarks with superpotential tree-level Yukawa couplings, and establishing the remaining quarks and leptons accordingly, immediately fulfils equations 1 and 6 from 2.2 and leaves the second and third the same and equal to

$$Q_{H_1} + Q_{H_2} + Q_{L_1} + Q_{L_2} + Q_{L_3} + 3Q_{Q_1} + 3Q_{Q_2} + 3Q_{Q_3} = 0.$$  \hfill (2.4)

Clearing $Q_{L_1}$ and replacing it in the non-linear anomalies we obtain for equation four in 2.2

$$(Q_{H_1} + Q_{H_2})(Q_{H_1} + Q_{L_2} + 3(Q_{Q_1} + Q_{Q_3})) = 0.$$  \hfill (2.5)

As the first brackets cannot be equal to zero, or the $\mu$ term would be back into the superpotential, the second brackets must cancel, which fixes $Q_{L_2} = -Q_{H_1} - 3Q_{Q_1} - 3Q_{Q_3}$, with which the fifth equation in 2.2 is simplified to

$$\begin{align*}
(Q_{H_1} + Q_{H_2}) & \left(-2Q_{L_3}(Q_{H_2} + Q_{L_3}) + Q_{H_1}(Q_{H_2} - 3Q_{Q_2}) - 3(Q_{H_2} + 2Q_{L_3})Q_{Q_2} - 9Q_{Q_2}^2\right) = 0.
\end{align*}$$  \hfill (2.6)

We thus have solutions depending on $Q_{H_1}, Q_{H_2}, Q_{L_3}, Q_{Q_1}, Q_{Q_2}, Q_{Q_3}$, a repeating characteristic of the model independently of which fields have non-holomorphic mass term, and where any combination is valid as long as Eq. 2.6 has a real solution and a term of the form $Q_{H_1} + Q_{H_2} + Q_{\nu_i} = 0$ is allowed. For this to happen, the corresponding $Q_{L_i}$ must be equal to $Q_{H_1}$, which means that not all right neutrinos will have a tree-level coupling with the Higgs fields, being nonetheless guaranteed that some will, thus providing a natural mass term for the Higgs particle. In addition, note that by having the quark families opposite one another, it is guaranteed that no baryon number violating operator is allowed in the superpotential as long as Higgs and quarks have different charges. We thus have shown that within the framework of the U(1) extended $\mu\nu$SSM, anomalies are cancelled without the need to add exotic matter, at the price of having non-universal charges, and with the gain of forbidding most troublesome operators in the superpotential. In the appendix A an specific $Q_X$ charge distribution for the above described family of solutions can be found, together with and an altogether different scenario, with important DM phenomenological consequences. In the remaining part of the article we elaborate on the novel, particular, and potentially relevant phenomenological characteristics of the non-Universal U(1)$\mu\nu$SSM.
2.2 Phenomenological consequences of Universality violation across fermion families.

The phenomenological manifestation of an extra U(1) gauge symmetry comes, mainly, from the mixing of the massive neutral components of the vector bosons from the gauge sector, namely the $Z$ and the $Z'$ gauge bosons. For the case of the enlarged $\mu\nu$SSM, and contrary to other models where similar mixing occurs (see for example [50] and references therein), the mixing happens naturally within the neutralino mixing matrix as part of the right sneutrinos acquiring vacuum expectation value. Such fact enriches greatly the phenomenology of the model and, as we will describe, imposes conditions for both collider and dark matter interactions, similar to what happens within the $\mu\nu$SSM alone [13]. The presence of the non-Universal extra gauge symmetry however introduces a new dependence on the specific charge of each field under the extra symmetry which will condition the possible interactions of fermions with the $Z'$, introducing a very particular phenomenology specific of the non-Universal U(1)$_{\mu\nu}$SSM.

In that sense the phenomenology can change depending on the mixing of the new sector. To parametrise the influence that the extra U(1) has, we define the mixing factor $R$ [36],

$$R = \frac{(M_Z^2 Z')^2}{M_Z^2 M_{Z'}^2}, \quad (2.7)$$

where the entries $M_{ij}$ correspond to the terms of the mixing matrix between $Z$ and $Z'$. In principle $R$ should be smaller than $10^{-3}$ given the experimental constraints available [44], with the consequence that $M_{ZZ'}$ has to be smaller than $M_{ZZ'} \lesssim 56 \text{ GeV}$ when $M_{Z'}=1 \text{ TeV}$, with such constraint becoming weaker as the mass of the $Z'$ gets heavier. Thus only heavier masses for the $Z'$ would fulfil such condition together with the ones coming from accelerator searches, and would require a somewhat large fine-tuning. Nonetheless, these limits are calculated for when the extra charges are Universal, which does not occur in our model. Hence the bounds presented have to be taken carefully, as the couplings of the physical states are now dependent on the $Q_X$ charges (as well as on the vacuum expectation value of $\nu^c$). However, in the rest of the paper we will consider that the mixing in the $Z - Z'$, i.e. $R \rightarrow 0$.

In our model the physical coupling of the $Z'$ to the fermionic sector is

$$g'^{Z'}_\alpha = g' Q_X^\alpha, \quad (2.8)$$

where $\alpha$ corresponds to the matter field $\psi_\alpha$ and $Q_X^\alpha$ is the charge of this field under the U(1)$_X$. $g'$ corresponds to the coupling constant of the U(1)$_X$ gauge symmetry. Thus, once the value of $g'$ is fixed, the way the $Z'$ couples to the different fermions depends strictly on the charges $Q_X$. The values these charges can have are fixed by the anomalies cancellation conditions, with only certain discrete values allowed. Moreover, since those charges break universality among fermions (see for example the models presented in the Appendix A), each fermion family will hold in general a different value of the charge and hence, will couple with different strength to the $Z'$, having deep phenomenological consequences.

$^2$A complete description of the entries $M_{ij}$ and their dependencies with the parameter of the model can be found in Ref. [36].

$^3$However the new configuration of charges and the relation among the different vacuum expectation values can induce a sizeable mixing in the $Z - Z'$ system.
The physical couplings of the SM particle content to the $Z'$ is described as follows. According to Eq. (2.8) the left and right handed components of the SM fermions do not have to share the same couplings to the $Z'$, as they depend on the charge assignment. Usually, the couplings of a vector boson can be expressed in its vector and axial forms. The vectorial coupling is defined as the sum of the left and right components, for example the vectorial coupling of the quarks is,

$$C^V_{u_i} = g_{Z' u_i L} + g_{Z' u_i R} = g'(Q_{u_i L} + Q_{u_i R}), \quad (2.9)$$

$$C^V_{d_i} = g_{Z' d_i L} + g_{Z' d_i R} = g'(Q_{d_i L} + Q_{d_i R}), \quad (2.10)$$

where $i = 1, 2, 3$ stands for the three families of both up and down quarks since they could have different charges. On the contrary, the axial coupling is defined as the difference of the components,

$$C^A_{u_i} = g_{Z' u_i L} - g_{Z' u_i R} = g'(Q_{u_i L} - Q_{u_i R}), \quad (2.11)$$

$$C^A_{d_i} = g_{Z' d_i L} - g_{Z' d_i R} = g'(Q_{d_i L} - Q_{d_i R}), \quad (2.12)$$

As we see the vector and axial couplings are different, this is not strange as within the SM the $Z$ boson behaves similarly. As the charges can be different in the up and down sectors the result is that $Z'$ does not couple in the same way to up and down quarks. The consequences are twofold. On the one hand, the production of a $Z'$ in collider experiments has to be recalculated taking into account that each of the fermion pairs that could produce it has a different coupling value, which makes the current constraints and limits invalid in this model. And on the other hand, direct dark matter searches, which are heavily dependent on the DM particle interaction with protons and neutrons, are affected by the fact that now the coupling is quark-family dependent, and will not interact the same with protons (which have more up-quark content) and with neutrons, affecting again to current experimental searches and imposing different limits.

In the remaining part of the section we will describe briefly the interesting phenomenological consequences that we just commented focusing on the collider and DM aspects.\footnote{A more detailed and involved description of the phenomenology of these models will be described in a forthcoming work.}

### Collider Phenomenology of the $Z'$

The $Z'$ could be eventually produced in the LHC. Both ATLAS and CMS have searches on high mass resonances decaying into a pair of leptons or hadronically (see for example Ref. \[46–48\]). As no signal of a $Z'$ is found, bounds can be set on the production and subsequent decay of a $Z'$, $pp \rightarrow Z' \rightarrow \psi \bar{\psi}$ for a defined mass. However, in the set of non-Universal models one can avoid such strong limits given the fact that different fermion families couple differently to the $Z'$. It could be the case that the up and down quark families have charges such that the effective coupling to the $Z'$ gets suppressed together with its production.

The general expression for the $Z'$ production and subsequent decay into fermions at the LHC is \[49,50\],

$$\sigma_{ff} \simeq \left( \frac{1}{3} \sum_q \frac{dL_{q\bar{q}}}{dm^2_{Z'}} \times \hat{\sigma}(q\bar{q} \rightarrow Z') \right) \times BR(Z' \rightarrow f\bar{f}), \quad (2.13)$$
where $dL_{qq}/dm^2_\bar{q}$ stands for the parton luminosities, $\hat{\sigma}(q\bar{q} \rightarrow Z')$ is the peak cross section for the $Z'$ boson, and $BR(Z' \rightarrow f\bar{f})$ is the branching ratio for the $Z'$ decaying into a fermion pair. As it was pointed out in Ref. [49], one can describe those parameters as a function of the sum of the different production rates for each quark and its $Z'$ coupling,

$$\sigma_{ff}^{\text{LO}} = \sum_{i=1}^{3} [c_{u_i}\tilde{\omega}_{u_i}(s,m^2_{Z'}) + c_{d_i}\tilde{\omega}_{d_i}(s,m^2_{Z'})] \times BR(Z' \rightarrow f\bar{f}).$$

(2.14)

Here, $c_q$ are defined as $c_q = (C_q^V)^2 + (C_q^A)^2$ and the functions $\tilde{\omega}_q(s,m^2_{Z'})$ contain all the information related with the parton distribution function, NLO corrections, etc.

The most important part of the model in the $Z'$ collider phenomenology is the fact that all type of fermions, no matter the family or the flavour, couple differently to the $Z'$. This would weaken the experimental searches of this kind of particles that ATLAS and CMS perform. The are several ways in which the $Z'$ production is diminished. One can have small quark couplings giving a tiny production cross section in such a way that the $Z'$ is barely produced in the LHC even for light masses of the $Z'$. Together with this effect the charges to the leptonic sector could be small as well reducing the total amount of observable events. However, in this model the charge assignment is not free since it is fixed by the cancellation of anomalies. This situation induces that one cannot arbitrarily make the couplings as small as it would be required to directly avoid collider searches. The charge assignment determines totally the $Z'$ phenomenology, so clear predictions for the LHC can be established, while on the other hand, a $Z'$ discovery would severely constraint the possible models, thus hinting towards the specific realisation in nature of the non-Universal U(1)$_{\mu\nu}$SSM. In that sense a deeply study should be done in the future [45] to determine the consequences of such charge assignment.

**Dark Matter Phenomenology**

The $\mu\nu$SSM has several DM candidates [13,20,51], including the neutralinos, the gravitino or even new particles charged accordingly to be good DM candidates. As a consequence of the presence of right neutrinos, neutralinos mix with neutrinos, and the Higgs with the sneutrinos, leading to several possible benchmark scenarios with distinctive signals [52]. In the case of the $U(1)$ enlarged $\mu\nu$SSM, such signal comes from the $Z'$ mediated spin independent interaction with a dark matter particle $\psi$. We can parametrise the effective Lagrangian of DM particle interaction with protons $p$, and neutrons, $n$, mediated by a vector boson as,

$$\mathcal{L}_{\psi}^{\text{SI}} = f_p(\bar{\psi}\gamma_\mu\psi)(\bar{p}\gamma^\mu p) + f_n(\bar{\psi}\gamma_\mu\psi)(\bar{n}\gamma^\mu n),$$

(2.15)

where the vector couplings $f_p$ and $f_n$ are defined through their nucleon valence quark content as [41]

$$f_p = 2b_u + b_d; \; f_n = b_u + 2b_d,$$

(2.16)

with $b_{u,d}$ the effective $Z'$ mediated vector couplings

$$b_{(u,d)} = \frac{gdmC_{(u,d)}^V}{2m^2_{Z'}}.$$

(2.17)

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5For more information one can see Ref. [49] or the Appendix of Ref. [50]
Using the definitions for the vector coupling obtained in Eqs. 2.10 we have that

\[ b_{(u)} = \frac{g_d m'_g}{2m^2_{Z'}} (Q_{u_L} + Q_{u_R}), \]  
(2.18)

\[ b_{(d)} = \frac{g_d m'_g}{2m^2_{Z'}} (Q_{d_L} + Q_{d_R}), \]  
(2.19)

such that the effective coupling of the DM particle to protons and neutrons is,

\[ f_p = \frac{g_d m'_g}{2m^2_{Z'}} (2Q_{u_L} + 2Q_{u_R} + Q_{d_L} + Q_{d_R}), \]  
(2.20)

\[ f_n = \frac{g_d m'_g}{2m^2_{Z'}} (Q_{u_L} + Q_{u_R} + 2Q_{d_L} + 2Q_{d_R}). \]  
(2.21)

We can define the amount of the isospin violation as the ratio \( f_n/f_p \), that in our case is given by

\[ f_n/f_p = \frac{Q_{u_L} + Q_{u_R} + 2Q_{d_L} + 2Q_{d_R}}{2Q_{u_L} + 2Q_{u_R} + Q_{d_L} + Q_{d_R}}. \]  
(2.22)

As we can see, the ratio \( f_n/f_p \), depends exclusively on the charges of the corresponding quarks under the extra \( U(1)_X \) symmetry. Having a non-Universal extra gauge symmetry means that in general these charges will not be the same, and therefore the amount of isospin violation will in general be different than the usual ±1 of universal models, thus providing a distinctive, particular signal in the cross section of DM-nucleus elastic scattering experiments. Moreover, notice that the vector coupling ratio is independent of the value of the gauge coupling \( g' \).

The striking feature of the non-Universal \( U(1)_\mu \nu \) SSM is that the amount of isospin violation can only acquire a discrete set of values. In the non-Universal \( U(1)_\mu \nu \) SSM there is a class of solutions for which both the up and down quarks have the same kind of mass terms (i.e, either tree-level or non-holomorphic). This means in particular that it is the opposite Higgs, namely \( H_1 \) or \( H_2 \) which gives the mass to each of them. As each Higgs has a different extra charge under the new gauge symmetry, for this class of models the amount of isospin violation is parametrised by the Higgs charges as follows

\[ \frac{f_n}{f_p} = \frac{Q_{H_2} + 2Q_{H_1}}{2Q_{H_2} + Q_{H_1}}. \]  
(2.23)

Hence, as long as the Higgs extra charges are different, a condition necessary in order to forbid the \( \mu \) term from the superpotential, there will be isospin violation. It is important to note that the ratio \( f_n/f_p \) depends only on the Higgs charges and not on the up and down ones. For example, the model presented in the Appendix has isospin violation \( \approx -1.75 \). Of utmost importance is to stress that not just any value of isospin violation is allowed but, as the extra charges must fulfil certain conditions, the number of possible models is constrained, allowing experiments to discriminate among realisations of the supersymmetric model, which could bear direct relation with the kind of low energy scale string realisation.
3 Conclusions

In this letter, we have presented a new supersymmetric model in which, by adding a non-Universal U(1) gauge symmetry to the already explored $\mu\nu$SSM model, not only both the $\mu$ term problem and the neutrino masses problem are solved, but the stability of the proton is ensured by recovering an effective R-parity, forbidding at the same time baryon number violating operators and avoiding a possible domain wall problem. By allowing non-Universal charges in all the fields, we demonstrate that there exist families of solutions which require of no exotic matter whatsoever to cancel anomalies. Moreover, in solving the anomalies equations we find that there only is a discrete set of possible extra charges allowed, a fact that has deep implications in the possible phenomenology of the model.

Universality breaking implies, in the U(1)$_{\mu\nu}$SSM, that the SM fermions will have in general a different value of the charge $Q_X$ both for each family and also between up and down doublet components. This means that, while stringent bounds are already imposed in the possible production of a $Z'$ at accelerators, these bounds do not apply directly to our model. As the production rates are calculated assuming that all fermions will couple with the same strength to the $Z'$ boson, and such coupling depends on the specific fermion extra charge, they have to be recalculated for each specific realisation of anomalies cancellation in the non-Universal U(1)$_{\mu\nu}$SSM, meaning that new scenarios open up in which even a light $Z'$ boson could have escaped detection at LHC. A forthcoming work will address in detail the specifics of such scenarios [45].

The extra charges dependence of fermionic coupling to a $Z'$ also modifies interaction rates with a DM particle, which is dominated by the lightest $Z'$ mass eigenstate. The conditions imposed by anomalies cancellation lead to a family of scenarios where isospin violation is realised and depends exclusively on the Higgs extra charges, a particle in principle completely unrelated with DM interactions. The implications for DM detection are profound, as the specific model realisation implies a very specific interaction rate with protons and/or neutrons, rendering particular experiments more or less suitable, and modifying the conditions for existing ones. Benchmark scenarios with concrete realisations of extra charge distributions will be analysed in [45].

Summarising, the non-Universal U(1)$_{\mu\nu}$SSM is an attractive model, which could be easily related with specific intersecting brane constructions and which, through a very particular phenomenology consequence of the discrete and constrained extra charge values, could when and if SUSY is discovered, be related directly with the kind of low energy stringy realisation.

Acknowledgements

V.M.L. acknowledges support of the Consolider MULTIDARK project CSD2009-00064, the SPLE ERC project and the BMBF under project 05H15PDCAA. S.O.C. acknowledges support from MINECO FEDER funds FIS2015-69512-R and Fundación Séneca (Murcia, Spain) Project No. ENE2016-79282-C5-5-R.
A Example of charge assignation

We present here two examples of solutions for the anomalies equations, representative of the two main families of models presented in the main text.

Example 1

A minimalist charge assignment that fulfils the anomalies equations and permits the effective $\mu$ term, having a quarks fields hierarchy, and with tree-level Yukawas and non-holomorphic terms appearing at opposing families in quarks. Particularly, in the example presented is the second family of up-quarks and the first, and third families of down quarks, the ones which acquire mass via non-holomorphic terms.

\[
\begin{array}{|c|c|c|c|c|c|c|}
\hline
& Q_Q & Q_u & Q_d & Q_L & Q_e & Q_{\nu^c} \\
\hline
1^{st} \text{ Family} & -\frac{1}{3} & \frac{1}{3} & \frac{1}{3} & 0 & -2 & 0 \\
\hline
2^{nd} \text{ Family} & -\frac{2}{3} & \frac{2}{3} & -\frac{1}{3} & 7 & -7 & -7 \\
\hline
3^{rd} \text{ Family} & -\frac{3}{3} & \frac{3}{3} & \frac{3}{3} & 2 & -4 & -2 \\
\hline
\end{array}
\]

With this charges assignment, the superpotential below the soft breaking scale reads

\[
-\mathcal{L}_{\text{eff}} = Y_u(u_L)^c q_u H_2 + \tilde{Y}_e(e_L)^c q_e H_1 + Y_t(t_L)^c q_t H_2 \\
+ \tilde{Y}_d(d_L)^c q_d H_2 + Y_s(s_L)^c q_s H_1 + \tilde{Y}_b(b_L)^c q_b H_2 \\
+ Y_e(e_L)^c L_e H_1 + \tilde{Y}_\mu(\mu L)^c L_\mu H_2 + Y_\tau(\tau L)^c L_\tau H_1 \\
+ \epsilon_{ab} Y_{\nu^i}^j H^b_i L^a_i \nu^c_j + \nu^c \tilde{H}_1 \tilde{H}_2 + +\nu^c \tilde{H}_1 \tilde{H}_2,
\]

where the tilded $Y$ represent the Yukawas generated by non-holomorphic interactions.

Example 2

Now we present an example of solution for the family of models which have non-trivial violation of isospin, with consequences for the phenomenology of dark matter. In this case, the condition is that both first families of up and down quarks have the same kind of mass term. In this case, having both with a Yukawa at tree-level, a possible anomalies cancellation charges distribution is

\[
\begin{array}{|c|c|c|c|c|c|c|}
\hline
& Q_Q & Q_u & Q_d & Q_L & Q_e & Q_{\nu^c} \\
\hline
1^{st} \text{ Family} & \frac{1}{3} & \frac{1}{3} & -\frac{2}{3} & -10 & 25 & 25 \\
\hline
2^{nd} \text{ Family} & -\frac{1}{3} & \frac{1}{3} & 0 & \frac{8}{3} & -13 & -1 \\
\hline
3^{rd} \text{ Family} & \frac{2}{3} & -\frac{1}{3} & -\frac{14}{3} & 2 & -4 & -\frac{1}{3} \\
\hline
\end{array}
\]

\[
Q_{H_1} = 2, \quad Q_{H_2} = -\frac{2}{3}
\]

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