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# A $\nu$ supersymmetric anomaly-free atlas

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ABSTRACT: Extensions of the minimal supersymmetric standard model (MSSM) gauge group abound in the literature. Several of these include an additional  $U(1)_X$  gauge group. Chiral fermions' charge assignments under  $U(1)_X$  are constrained to cancel local anomalies in the extension and they determine the structure and phenomenology of it. We provide all anomaly-free charge assignments up to a maximum absolute charge of  $Q_{\text{max}} = 10$ , assuming that the chiral superfield content of the model is that of the MSSM plus up to three Standard Model (SM) singlet superfields. The fermionic components of these SM singlets may play the rôle of right-handed neutrinos, whereas one of the scalar components may play the rôle of the flavon, spontaneously breaking  $U(1)_X$ . Easily scanned lists of the charge assignments are made publicly available on Zenodo. For the case where no restriction is placed upon  $Q_{\text{max}}$ , we also provide an analytic parameterisation of the general solution using simple techniques from algebraic geometry.

**KEYWORDS:** Supersymmetry Phenomenology

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# 1 Introduction

Quantum field theories of vector bosons are notoriously problematic unless they arise from gauge symmetries, whence non-renormalisability and non-unitarity can be tamed. It is thus imperative that the gauge symmetry of the renormalisable ultra-violet completion of any such model should not contain any quantum field theoretic gauge *anomalies*, where quantum corrections spoil the gauge symmetry that was imposed upon the tree-level theory. The Standard Model (SM) itself is anomaly-free and can thus remain a self-consistent theory up to very large renormalisation scales. Despite this, there are good reasons to expect the SM to be an effective field theory resulting from decoupling other fields. Many reasons have been invoked to motivate extending the Lie algebra<sup>1</sup>  $\mathfrak{sm} := \mathfrak{su}(3) \oplus \mathfrak{su}(2) \oplus \mathfrak{u}(1)_Y$  of the

 $<sup>^1\</sup>mathrm{We}$  shall refer to the Lie algebra (as opposed to the Lie group) in  $\mathfrak{mathfrat}$  script.

Standard Model (SM) by a spontaneously broken gauged  $\mathfrak{u}(1)_X$  summand, for example. Such extensions have been used to explain measurements of the anomalous magnetic moment of the muon [1], to provide axions [2] or leptogenesis [3], to provide fermion masses through the Froggatt-Neilsen mechanism [4], or explain measurements of the  $b \to sl^+l^-$  transition which are currently in tension with SM predictions [5–13]. In general, the X charge assignments of the models can be family dependent, resulting in family-dependent couplings of a resulting massive Z' vector boson. Indeed, in several applications (the last two in our aforementioned list) it is a necessary requirement that the X charges are family dependent, since the symmetry and the Z' are respectively used to explain family non-universal effects.

In  $\mathfrak{u}(1)_X$  extensions, the phenomenology of the Z' is often key and is dictated by the integer X charges of the other fields in the model (integer X charges results from an implicit assumption that the extension is compact). The X charges of the chiral fermions in particular dictate the contribution to perturbative local anomalies of such models. There is therefore a non-trivial cross-over between the extensions' phenomenology and anomaly cancellation via the chiral fermions' charge assignments. Unfortunately, in general, with a fixed chiral fermion content, anomaly cancellation conditions (ACCs) are difficult to solve, the number theory state-of-the art being the solution of a single cubic in three unknown integer parameters [14].

Some recent progress has been made in this direction, however. In ref. [15], the gravitational and gauge anomalies of a U(1) gauge symmetry (i.e. with no SM gauge group but with charged chiral fermionic fields) were solved analytically for the charges of a priori fixed numbers of chiral fermions via an ingenious algebraic method;<sup>2</sup> this was soon understood from a geometric perspective [17] by using a theorem due to Mordell [14]. Similar geometric methods were employed to find an analytic solution to the more difficult problem of  $\mathfrak{sm} \oplus \mathfrak{u}(1)$  anomaly-free charge assignments in the specific case of SM fermion content, plus three right-handed (RH) neutrinos (i.e. SM-singlet chiral fermion fields which may carry X charge) [18]. The number of solutions is formally infinite,<sup>3</sup> unlike the case of semi-simple SM extensions with identical fermionic field content, where there is a list of 340 [19]. Unfortunately, the geometric methods employed only solve a small family of similar cases and cannot be deployed on general chiral fermionic contents. Furthermore, the analytic solution, whilst of intrinsic interest in and of itself, comes with a significant drawback for model-builders interested in using it: each charge is parameterised in terms of a fourth-order polynomial of integer parameters. Whilst it is easy to input these parameters and achieve anomaly-free charges, model builders often want to fix a function of them to certain values for phenomenological purposes, but this is a difficult and currently unsolved problem, because it involves solving a system of coupled fourth-order diophantine equations.

Fortunately, when appropriately employed, computers come to the rescue of the reverseengineering model builder. In an  $\mathfrak{sm} \oplus \mathfrak{u}(1)$  'anomaly-free atlas' [20], all solutions of the

<sup>&</sup>lt;sup>2</sup>The algebraic approach was partially extended to  $U(1)^n$  gauge symmetries in ref. [16].

<sup>&</sup>lt;sup>3</sup>One way of seeing this is to set the X charges of the first family of particles to be equal to their hypercharges, the second family to be equal to some integer multiplied by baryon number minus lepton number B - L, and the third family to have zero charge. Any such charge assignment solves the anomaly cancellation conditions. Since there are an infinite number of constants we can multiply the second family by, each of which leads to a distinct chiral solution, there are an infinite number of solutions.

ACCs for integer charges between -10 and 10 for 18 chiral fermion gauge representations in the SM plus three RH neutrinos were found by a scan.<sup>4</sup> Cases which are in a sense equivalent (where the charges differ by a common multiple which can be absorbed into the  $u(1)_X$  gauge coupling, or which differ by a permutation of the family indices within a *species*-fields which have identical SM representations) were only counted once (and aside from some rare cases, only scanned over once). Anomaly-free solutions are scarce: only roughly one in  $10^9$  was anomaly-free from the whole sample. The list of anomaly-free fermionic charge assignments was made publicly available. It is a list of over 21 000 000 solutions that is easy and quick to search through and filter with the aid of a simple computer program. As such, it is user friendly for would-be  $U(1)_X$  gauge extension model builders who can search through the list and filter for charge assignments with various desired properties. The charges are limited in *height* (the maximum absolute value of a charge in any solution), but have the advantage of being easily useable provided one can adapt or write a simple computer program that reads the list in and filters it.

Heretofore, there has been no similar list made for supersymmetric (SUSY) models. SUSY model building has several motivations, the primary one being that it does not suffer from the technical hierarchy problem, where radiative corrections to the Higgs mass tend to drag it up to the largest fundamental energy scale (for example the Planck mass  $\sim 10^{19}$  GeV) multiplied by the square root of a loop factor. There are other motivations for supersymmetry too, for example, in an  $\mathcal{N} = 1$  supersymmetrisation of the SM (the MSSM), the experimental measurements of the gauge couplings agree with the gauge coupling unification condition predicted by SUSY grand unified theories. When one includes an extra multiplicative discrete symmetry such as *R*-parity or matter parity<sup>5</sup> the MSSM possesses a stable particle which, depending upon parameters, has the correct properties to constitute the universe's dark matter and potentially dangerous proton decay processes are suppressed. Particular examples of  $\mathfrak{u}(1)_X$  gauge extensions of the MSSM can combine the aforementioned phenomenological benefits of a Z' with those of SUSY models. Some of these have appeared in the literature, for example see refs. [22–27].

It is our intention here to extend the original non-SUSY anomaly-free atlas to the SUSY case and make a new list (a ' $\nu$  SUSY anomaly-free atlas') available to interested SUSY  $\mathfrak{u}(1)_X$ -extension model builders and others. We shall include the addition of up to three MSSM-singlet chiral superfields: the fermionic components of all or some of these can play the rôle of RH neutrinos, resulting in tiny neutrino masses via the see-saw mechanism (below, we call this model the  $\nu$ MSSM). The scalar component of one of these MSSM-singlet chiral superfields is expected to play the rôle of the flavon, which has a necessarily non-zero X charge and acquires a vacuum expectation value, spontaneously breaking  $U(1)_X$ . One might expect that one of the SM-singlet fields must therefore have a non-zero  $\mathfrak{u}(1)_X$  charge, unlike the non-SUSY case, where the charges of the flavon and all SM-singlet fermions were *a priori* unconstrained. However, we won't impose this condition because the field

<sup>&</sup>lt;sup>4</sup>This strategy has also recently been used for the case of U(1) gauge theory with different numbers of Weyl fermions, in a search for scotogenic models [21].

<sup>&</sup>lt;sup>5</sup>Matter parity is defined as  $(-1)^{3(B-L)}$ , where B is baryon number and L is lepton number, whereas R-parity is defined as  $(-1)^{3(B-L)+2s}$ , where s is spin.

content of the model can easily be extended in a way that does not change the ACCs but which effectively removes the condition, as we shall explain below. A functional difference to the original non-SUSY anomaly-free atlas is the appearance of the Higgsino partners of the two MSSM Higgs doublets, augmenting the number of Weyl fermion SU(2) gauge representations by two. This therefore extends the original list of 18 X charges to 20. In case a height larger than 10 is required, we will also provide a general analytic solution to the anomaly cancellation conditions. This relies on using the same geometric framing in which the SM-plus-3 RH neutrino case was solved [18]; we take the opportunity to demonstrate a new technique to solve such problems, although the technique used in ref. [18] would also have worked.

Since our  $\mathfrak{u}(1)_X$  charge assignments are family dependent, some anomaly-free assignments will result in various Yukawa couplings being banned at the renormalisable and tree level. Further model building can describe how various terms could come about after the  $\mathfrak{u}(1)_X$  is spontaneously broken, for example by inducing non-renomalisable terms (the Froggatt-Nielsen mechanism), or from non-holomorphic terms, or elsewhere. We emphasise that the scope of our paper is to provide the atlas and develop the techniques for doing so; it is *not* our aim to go into any proper model building, but to instead provide a list of charge assignments which satisfy the anomaly-free condition. We shall later illustrate the implementation of some basic model selection rules, with the understanding that these in no way constitute serious models for phenomenological consideration, rather they will form a superset of them. Further selection rules and analysis would be required to identify models of proper interest for phenomenology.

The paper proceeds as follows: in section 2, we describe the anomaly cancellation conditions relevant for the Lie algebra  $\mathfrak{mssm} \oplus \mathfrak{u}(1)_X$ , and a chiral superfield content of the  $\nu$ MSSM. In section 3, we describe the computational scan and how the solutions are listed and ordered, giving the number of solutions found up to a height of 10. We provide an analytic method of solution in section 4, along with a parameterisation of the solution. Various consistency checks of the solutions are described in section 5: some are checks solely of the numerical solutions, some are of the analytic solution and some are checks of the analytic solution versus the numeric solutions. Some initial filters of the numerical solutions (chosen for specific phenomenological reasons) are explored in section 6. We provide a summary of the paper and a discussion in section 7.

## 2 $\mathfrak{u}(1)_X$ extension of the $\nu$ MSSM Lie algebra

We shall now detail the field content of the  $\nu$ MSSM and discuss how they change the ACCs. In table 1, we list our notation for the  $\nu$ MSSM chiral fermionic fields and their representations under the gauged Lie algebra. As previously mentioned, the left-handed fermionic fields contained within the two Higgs chiral superfields provide a new feature as regards the ACCs. We note here that the fermionic components of the chiral superfields  $H_d$  and  $L_i$  have identical representations under the SM gauge Lie algebra, but the fermionic component of  $H_d$  may or may not be discriminated by a different quantum number under an imposed symmetry such as matter parity or R-parity.

Fer	mions			
	$\mathfrak{su}(3)$	$\mathfrak{su}(2)_L$	$\mathfrak{u}(1)_Y$	$\mathfrak{u}(1)_X$
LH quark doublets $Q_i$	3	2	1	$X_{Q_i}$
RH neutrinos $n_i$	1	1	0	$X_{n_i}$
RH charged leptons $e_i$	1	1	-6	$X_{e_i}$
RH up quarks $u_i$	3	1	4	$X_{u_i}$
RH down quarks $d_i$	3	1	-2	$X_{d_i}$
LH lepton doublets $L_i$	1	2	-3	$X_{L_i}$
LH down-type Higgsino $\tilde{H}_d$	1	2	-3	$X_{H_d}$
LH up-type Higgsino $\tilde{H}_d$	1	2	3	$X_{H_u}$
Chiral	superfie	lds		
$\hat{Q}_i$	3	2	1	$X_{Q_i}$
$\hat{N}_{i}^{c}$	1	1	0	$-X_{n_i}$
$\hat{E}^c_i$	1	1	6	$-X_{e_i}$
$\hat{U}_i^c$	$\overline{3}$	1	-4	$-X_{u_i}$
$\hat{D}_i^c$	$\overline{3}$	1	2	$-X_{d_i}$
$\hat{L}_i$	1	2	-3	$X_{L_i}$
$\hat{H}_d$	1	2	-3	$X_{H_d}$
$\hat{H}_d$	1	2	3	$X_{H_u}$

**Table 1**. Conventions for field content with representations under the gauge Lie algebra. RH stands for right-handed and LH stands for left-handed.  $i \in \{1, 2, 3\}$  is a family index. Note that we have re-scaled a more conventional hypercharge assignment by a factor of 6 to make all hypercharges setwise coprime integers. Such a re-scaling can be absorbed into the hypercharge gauge coupling. <sup>c</sup> denotes charge conjugation on the scalar and fermionic components of the chiral superfield.

We have thus augmented the MSSM, as far as the fermionic X charges go, by 20 parameters which we write in a 20-tuple

$$\mathbf{X} := \{ X_{Q_1}, X_{Q_2}, X_{Q_3}, X_{n_1}, X_{n_2}, X_{n_3}, X_{e_1}, X_{e_2}, X_{e_3}, X_{u_1}, X_{u_2}, X_{u_3}, X_{d_1}, X_{d_2}, X_{d_3}, X_{L_1}, X_{L_2}, X_{L_3}, X_{H_d}, X_{H_u} \}.$$
(2.1)

We take it as understood that, for the case where *R*-parity is *not* a symmetry of the theory, we modify (2.1) such that  $X_{H_d}$  is merged with  $X_{L_i}$  to form  $X_{L_{\alpha}}$ , where  $\alpha \in \{1, 2, 3, 4\}$ . For now though, we shall continue the discussion where  $H_d$  is discriminated from  $L_i$  by a discrete symmetry. Since the gauge extension is here assumed to be compact, **X** is a priori valued in  $\mathbb{Z}^{20}$ .

#### 2.1 Anomaly cancellation conditions

The MSSM *per se* is anomaly free. With the addition of  $\mathfrak{u}(1)_X$ , local anomalies persist unless **X** satisfies the ACCs<sup>6</sup>

$$\mathfrak{su}(3)^2 \oplus \mathfrak{u}(1)_X : \sum_{i=1}^3 (2X_{Q_i} - X_{u_i} - X_{d_i}) = 0, \qquad (2.2)$$

$$\mathfrak{su}(2)^2 \oplus \mathfrak{u}(1)_X : \sum_{i=1}^3 (3X_{Q_i} + X_{L_i}) + X_{H_d} + X_{H_u} = 0, \qquad (2.3)$$

$$\mathfrak{u}(1)_X \text{-gravity}: \sum_{i=1}^{3} (6X_{Q_i} - X_{n_i} - X_{e_i} - 3X_{u_i} - 3X_{d_i} + 2X_{L_i}) + 2X_{H_d} + 2X_{H_u} = 0, \quad (2.4)$$

$$\mathfrak{u}(1)_X^3: \sum_{i=1}^3 (6X_{Q_i}^3 - X_{n_i}^3 - X_{e_i}^3 - 3X_{u_i}^3 - 3X_{d_i}^3 + 2X_{L_i}^3) + 2X_{H_d}^3 + 2X_{H_u}^3 = 0, \quad (2.5)$$

$$\mathfrak{u}(1)_X^2 \oplus \mathfrak{u}(1)_Y : \sum_{i=1}^3 (X_{Q_i}^2 - 2X_{u_i}^2 + X_{d_i}^2 + X_{e_i}^2 + X_{L_i}^2) + X_{H_d}^2 - X_{H_u}^2 = 0, \qquad (2.6)$$

$$\mathfrak{u}(1)_Y^2 \oplus \mathfrak{u}(1)_X : \sum_{i=1}^3 (X_{Q_i} - 6X_{e_i} - 8X_{u_i} - 2X_{d_i} + 3X_{L_i}) + 3X_{H_d} + 3X_{H_u} = 0.$$
(2.7)

These ACCs inherit some in-practice physical equivalences between  $\mathfrak{u}(1)_X$  extensions related by the following operations:

- (i) Permutation of family indices within each species, since this is really just a change of basis.
- (ii)  $\mathbf{X} \to a\mathbf{X}$ , where  $a \in \mathbb{Q} \setminus \{0\}$ , when the gauge coupling only appears in the Lagrangian multiplied by a U(1)<sub>X</sub> charge, since the U(1)<sub>X</sub> gauge coupling may be simultaneously re-scaled by 1/a resulting in no substantive change. This is displayed by the fact that the ACCs are homogeneous.
- (iii)  $\mathbf{X} \to \mathbf{X} + y\mathbf{Y}$ , where  $\mathbf{Y}$  is the 20-tuple of fermionic field hypercharges (in the same field ordering as  $\mathbf{X}$ ) and  $y \in \mathbb{Z}$ . Resulting from a group outer automorphism, this change in fermionic representations can be accounted for by a redefinition of gauge fields [16].

Ideally, we wish to record exactly one entry in a list for each physically inequivalent charge assignment.<sup>7</sup> As we shall describe, we only found a practical way to do this for conditions (i) and (ii), but not (iii), implying that there will remain a few physically equivalent charge assignments in any anomaly-free list that we produce. The existence of these will end up providing us with a check of our computer program in section 5. In any case, such equivalent

<sup>&</sup>lt;sup>6</sup>Note that where necessary, we discriminate between the gauge Lie algebra, which is equivalent to  $\mathfrak{sm} \oplus \mathfrak{u}(1)_X$  and the  $\mathrm{MSSM} \times \mathrm{U}(1)_X$  gauge group, which is strictly only determined up to certain quotients, but this does not affect any of our discussion.

<sup>&</sup>lt;sup>7</sup>Note that together, (ii) with (iii) imply that we should regard  $\mathbf{X} \to x\mathbf{X} + y\mathbf{Y}$  as an equivalent theory, where  $x \in \mathbb{Q} \setminus \{0\}$  and  $y \in \mathbb{Q}$ .

charge assignments are rare, and we do not foresee particular problems resulting from their presence in our list. From now on, we refer to 'inequivalent' solutions to implicitly mean inequivalent under conditions (i) and (ii) only.

To incorporate (i), we take the convention that the family indices in **X** are such that, for each species  $S \in \{Q, n, e, u, d, L\}$ ,  $X_{S_1} \leq X_{S_2} \leq X_{S_3}$  (for the case without additional discrete symmetries to distinguish  $H_d$  and S = L,  $X_{S_3} \leq X_{S_4}$  as well). To take (ii) into account, all integers in the tuple must be setwise coprime but note that this still does not implement the equivalence with  $\mathbf{X}' := \{-X_{Q_3}, -X_{Q_2}, -X_{Q_1}, \ldots, -X_{H_d}, -X_{H_u}\}$ . In order to only list one instance of  $\mathbf{X}, \mathbf{X}'$ , we must define a condition that unambiguously picks one of them: here, we use the *lexicographically smaller* tuple.<sup>8</sup> An *n*-tuple  $a = \{a_1, \ldots, a_n\}$  is lexicographically smaller than another *n*-tuple  $b = \{b_1, \ldots, b_n\}$  (written as a < b) if and only if an  $i \in \{1, \cdots, n\}$  exists such that  $a_i < b_i$  and  $a_j = b_j$  for all  $j \in \{1, \cdots, i-1\}$ .

#### 2.2 Symmetry breaking

Since we are not empirically aware of a long-range force that can be attributed to an unbroken  $U(1)_X$  gauge symmetry, we suppose that it must be spontaneously broken. We further assume that it is broken by (at least) one of the scalars  $\tilde{\nu}_{R_i}$  contained in the SM-singlet chiral superfields  $N_i^c$ , so that it does not break the SM gauge symmetry. In order for a  $\tilde{\nu}_{R_i}$  field to play this rôle, by Goldstone's theorem it must possess a non-zero X charge. Typically, such a field is called a *flavon*. Let us denote it for the purposes of the current discussion, as  $\theta$ . Contrary to the non-SUSY case, we obtain a contribution to the ACCs through its fermionic superpartner  $\tilde{\theta}$ , the flavino. However, we will still solve the ACCs as given above assuming three SM-singlet chiral superfields only:  $N_i^c$ , where  $i \in \{1, 2, 3\}$ . The reasons for not explicitly adding to this number (for example by adding one more SM singlet chiral superfield) are twofold: firstly, we find practical barriers with four (or more) SM-singlets; the  $\nu$  SUSY anomaly-free atlas would take too long to compute and would take up too much disk space to store for the desired height of 10. Secondly, by sticking to three SM-singlet chiral superfields, we are able to find an analytic solution to the ACCs.

In principle, the requirement that at least one  $X_{n_i} \neq 0$  would allow us to reduce the domain of X charges considered in our computational search below, although not by much. We choose not to restrict the domain of X charges in this way however, since one could augment our model by *two* additional SM-singlet chiral superfields  $\theta_1$  and  $\theta_2$  such that  $X_{\theta_1} + X_{\theta_2} = 0$ . The contributions from  $X_{\theta_1}$  and  $X_{\theta_2}$  would cancel in the ACCs, leaving the ACCs above unmodified. This type of extension is commonly used, for example, in  $U(1)_{B-L}$  extensions of the MSSM [23]. More generally one can add several SM-singlet chiral superfields which satisfy the pure  $\mathfrak{u}(1)$  anomaly equations and thus cancel out of the ACCs [15, 17]. We also note that to set a superfield's charge to zero has the same effect on the ACCs as would removing the superfield (or at least its fermionic component) entirely from the model.

<sup>&</sup>lt;sup>8</sup>Lexicographical ordering is a much simpler condition than the one used in the original anomaly-free atlas [20].

With the constraints (or lack thereof) listed above, our inequivalent numerical set of solutions will have the following subsets:

- The SM plus three RH neutrinos corresponds to the subset with  $X_{H_u} = X_{H_d} = 0$ (since this is equivalent to the model obtained by removing the  $H_u$  superfield and the  $H_d$  Higgsino from our current set-up).
- The MSSM with up to three  $U(1)_X$ -charged RH neutrino chiral superfields, where the  $U(1)_X$  is broken by (at least) one of the RH sneutrinos, is the subset where at least one  $X_{n_i} \neq 0$ .
- The MSSM with three RH neutrinos and two additional chiral superfields  $\theta_1$  and  $\theta_2$  charged such that  $X_{\theta_1} + X_{\theta_2} = 0$ . This is a possibility for the subset with  $X_{n_1} = X_{n_2} = X_{n_3} = 0$ .

# 3 Numerical solutions up to a height of 10

We produce a list of solutions to the ACCs using a modification of the computer program that was used to produce the original anomaly-free atlas [20]. To search for solutions, we scan over integer values in the domain  $|X_i| \leq Q_{\text{max}}$ . For example, if  $Q_{\text{max}} = 1$ , our solutions will only contain X charges in the set  $\{-1, 0, 1\}$ . Each set of solutions can then be classified by  $Q_{\text{max}}$ , its maximum possible height. Note that in this definition, a list of solutions  $Q_{\text{max}} = N$  also contains all solutions with  $Q_{\text{max}} < N$ .

There are a priori  $(2Q_{\max} + 1)^{20}$  solutions to be checked as solutions to the ACCs. Computing speed is the biggest limiting factor in our search for solutions. The original computer program is described in detail in section 3 of ref. [20], where we direct the curious reader. Here, we find it expedient to not discriminate a priori between  $H_d$  and  $L_i$ : within the program, we therefore remove the explicit  $H_d$  charge, replacing it by  $L_4$ , equivalent to considering the theory without a discrete symmetry to distinguish them. The search for solutions is sped up by removing some equivalent solutions from the scan, and by using the four linear ACCs to directly fix the values of four of the charges. This still leaves us with a large solution space to consider, as compared to the original non-supersymmetric anomaly-free atlas. We further improve the speed by parallelising the three outer loops (i.e. over  $X_{Q_1}, X_{Q_2}$  and  $X_{Q_3}$ ). For  $Q_{\max} \leq 4$  the parallelisation improvement is minimal, but for  $Q_{\max} \geq 5$  we find that this step is necessary to produce solutions in a reasonable amount of time.

#### **3.1** Binary search algorithm

The output of the computer program for  $Q_{\text{max}} = 10$  is a large lexicographically ordered list (the ASCII file is around 125 Gb in size) of inequivalent solutions which solve the ACCs, each one comprised of a line made of the 20 integers which form **X**. We have not assumed a discrete symmetry that distinguishes  $L_i$  from  $H_d$  in the output and so we have four  $X_{L_{\alpha}}$ charges listed. This file forms one of the two most important outputs of the present paper (the other being the analytic solution for any height described in section 4). As mentioned

$Q_{\max}$	$\# \nu {\rm MSSM}$	$\# R_p \nu MSSM$	# Non-SUSY	# No ACC condition
1	111	267	37	$4.5 \times 10^6$
2	2321	6882	357	$2.3  imes 10^{10}$
3	44212	143707	4115	$8.6\times10^{12}$
4	401129	1367991	24551	$8.2\times10^{14}$
5	2582166	9063191	111151	$3.3  imes 10^{16}$
6	13553325	48681027	435304	$7.6\times10^{17}$
7	54699483	199275965	1358387	$1.1 \times 10^{19}$
8	185454955	682827818	3612733	$1.2\times 10^{20}$
9	598267488	2224178673	9587084	$1.0 \times 10^{21}$
10	1628002737	$6\overline{094894134}$	21546919	$6.8 \times 10^{21}$

Table 2. Number of inequivalent anomaly-free charge assignments found for  $U(1)_X$  extensions of the  $\nu$ MSSM where  $L_i$  and  $H_d$  are not discriminated ( $\nu$ MSSM), or where they are ( $R_p\nu$ MSSM) or for the original non-supersymmetric anomaly-free atlas (Non-SUSY). The leftmost columns specifies the height  $Q_{\max}$ , with the corresponding solutions containing all integer charges  $X_i$  in the range  $-Q_{\max} \leq X_i \leq Q_{\max}$ . The rightmost column lists the number of potential inequivalent solutions in the SUSY case before ACCs are applied and where a discrete symmetry distinguishes  $L_i$  from  $H_d$ .

in section 1, we envisage that our output file may be used by supersymmetric model builders by scanning through it with a computer program and filtering the results. Since the file is so large though, we have facilitated the decrease of the complexity of algorithms used to analyse the file in order to speed them up. The fact that our list of solutions is ordered lexicographically means that one can take advantage of the binary search algorithm. This reduces the complexity of finding a solution in the list from  $\mathcal{O}(n)$  to  $\mathcal{O}(\log n)$ . One usually has an intuitive understanding of the binary search algorithm since it is roughly how one usually finds numbers in a phone book or words in a dictionary, as follows. Let us say we have a solution we want to find in our list or show that it does not exist in the list. The binary search algorithm goes half-way down the list and determines if our solution is less then or equal to the solution at the half-way point. If it is present in the first half then we throw away the second half and keep the first, and if it isn't we discard the first half and keep the second. This is then repeated until a list with a single item which will (if it exists) match the one we are trying to find.

## 3.2 Output

We display some basic statistics characterising the number of solutions found in table 2. The number of inequivalent solutions increases rapidly as a function of the maximum height searched over,  $Q_{\text{max}}$ . In all, we find over 1.6 billion solutions for  $Q_{\text{max}} = 10$  in the case where a discrete symmetry does not pick out one of the  $X_{L_{\alpha}}$  charges to be  $X_{H_d}$ , a far larger number than the original anomaly-free atlas (which counts under 22 million inequivalent solutions). As expected, this increases almost four-fold for the case where one does pick an  $X_{L_{\alpha}}$  to be  $H_d$ . Solutions to the ACCs are scarce; their density decreases with



Figure 1. Fraction of otherwise possible solutions that are anomaly free as a function of  $Q_{\text{max}}$ . Specifically, the fraction is equal to the number of inequivalent solutions divided by the number of inequivalent possible assignments before the local anomaly cancellation requirements are imposed.

Model	Q	Q	Q	n	n	n	e	e	e	u	u	u	d	d	d	L	L	L	L	$\tilde{H}_u$
$Y'_3$	-1	-1	1	0	0	0	-6	6	6	-4	-4	4	-2	2	2	-3	3	3	3	-3
$B'_3$	-1	-1	1	-3	3	3	-3	3	3	-1	-1	1	-1	-1	1	-3	3	3	3	-3
ref. [22], table 3	0	0	0	-3	0	0	-1	1	3	-1	-1	-1	1	1	1	-2	0	1	2	-1
ref. [22], table 4	-1	-1	-1	-9	0	0	-9	0	0	1	1	1	1	1	1	0	0	0	9	0
ref. [24]	-1	0	0	-1	0	4	-1	0	4	-1	0	0	-1	0	0	-1	0	0	4	0
ref. [25]	-1	0	0	1	1	1	1	1	1	-1	0	0	-1	0	0	0	1	1	1	0
SUSY $B-L$ [23]	-1	-1	-1	3	3	3	3	3	3	-1	-1	-1	-1	-1	-1	0	3	3	3	0
TFHM [10]	-1	0	0	0	0	0	0	0	6	-4	0	0	0	0	2	0	0	0	3	0

**Table 3.** Some examples of anomaly-free charge assignments found. Here, we list the  $\mathfrak{u}(1)_X$  charges of the (left-handed or right-handed) chiral fermions of each model. These include the solutions  $Y'_3$ ,  $B'_3$  used to derive the analytic solution of section 4 as well as the non-SUSY Third Family Hypercharge (TFHM) solution which we expect to be contained within our list. Note that all solutions have been rescaled and reordered to satisfy the format of our list as detailed in section 3. There is no identification yet of a field with its family; that requires more model building of fermion masses and mixings.

increasing height. For a height of 10, for example, only approximately 1 in  $10^{12}$  possible inequivalent charge assignments are anomaly free. We display this fraction for various different values of  $Q_{\text{max}}$  in figure 1 for the case where no symmetry discriminates between  $H_d$  and  $L_i$  ( $\nu$ MSSM) and the case where it does ( $R_p \nu$ MSSM).

In table 3, we display some solutions that appear in the literature and in our list. All of the solutions shown were found using the binary search algorithm sketched in section 3.1. Their presence in the list is a check of some expected and found solutions. Two solutions  $(Y'_3 \text{ and } B'_3)$  will be useful for our analytic solution, which we turn to now.

#### 4 Analytic solution

In this section, we will first frame our problem in a geometric language that will facilitate our analytic solution of the ACCs. We shall then go on to sketch the geometric method by which the solution is obtained. Then we shall derive the solution in detail algebraically, eventually providing an explicit parameterisation of the 20 integer charges of the  $\nu$ MSSM chiral superfields in terms of some integer parameters. We then provide a right inverse, which, given a solution to the ACCs, returns parameters which will lead to that solution. Such an inverse has the dual purpose of facilitating checks between the numerical and analytic solutions and of providing an additional proof that our solution is generic.

## 4.1 Geometric framing of the problem

The ACCs form a set of polynomial equations in the integers — otherwise called *diophantine* equations. Suppose we take account of only the physical equivalence defined by scaling (point (ii) in section 2). It then does not matter, from a mathematical point of view, whether we use the label  $L_4$  or  $H_d$  for the relevant chiral superfield; here we shall choose the latter. We can view the unknown charges as corresponding to points in the projective space  $\mathbb{PQ}^{19}$ . This is formed by considering the charges as living in the rationals  $\mathbb{Q}^{20}$ , removing the origin and providing an equivalence relation between points in  $\mathbb{Q}^{20}$  differing by rational multiples. The points satisfying the ACCs in  $\mathbb{PQ}^{19}$  are said to form a projective variety.

One might expect our solution to be parameterised by 14 independent integer-valued parameters (starting with 20 and subtracting 6 for the ACCs). However, as we shall see, we shall have to add 9 parameters to cover exceptional cases, making the total number of integer parameters 23. Our solution will then take the form of a map from  $\mathbb{Z}^{23}$  to  $\mathbb{PQ}^{19}$  which satisfies the following properties: its image is completely within the projective variety, it surjects onto the projective variety and its value depends only on the projection onto a  $\mathbb{Z}^{14}$  subspace of  $\mathbb{Z}^{23}$  in all but a few classes of exceptional cases.

Within  $\mathbb{PQ}^{19}$ , the ACCs (2.5) and (2.6) define a cubic and a quadratic hypersurface, respectively (we shall below refer to these as 'the cubic' and 'the quadratic', respectively, for brevity).

To solve systems of diophantine equations, number theorists often use a small set of solutions as a tool for finding all solutions. Given our extensive numerical scan we are in a position to make an attempt in this manner. For a generic set of equations this is not guaranteed to be possible, however we are lucky in that for our particular set of ACCs, at least two distinct methods exist. The first mirrors the method of ref. [18] which exploits a special point of  $PQ^{19}$  that is a 'double point' of both the cubic and quadratic. A second new method is presented here.

# 4.2 Sketch of the method

The linear ACCs are easy to deal with. Their solution defines a projective subspace PL of  $PQ^{19}$ . It is in PL that we must discuss the quadratic and the cubic.

Given a single solution to the quadratic ACC it is possible to find all solutions to the quadratic ACC by constructing all possible lines through this known solution: along each line there must be one further solution to the quadratic, since every rational quadratic in one dimension has either two or zero rational roots. In a similar vein, given a single solution to the cubic, with the special property that all first order partial derivatives vanish at this point, it is possible to find all solutions to the cubic ACC by constructing lines through this point. Such a point is called a *double point* of the cubic.

To solve both the quadratic and the cubic simultaneously it is sufficient to have a *line* on which every point is a solution to the quadratic and every point is a double point of the cubic (although as noted above, other methods do exist). In fact for us there is only one such line (up to permutations of charges within the  $\mathfrak{su}(2)$ -doublet,  $\mathfrak{su}(3)$ -singlet sector, and other species), which is the one between the points  $Y_3$  and  $B_3$  given in table 4. These two points are a reordering of the charges within  $Y'_3$  and  $B'_3$ , respectively, from table 2. The first point,  $Y_3$ , corresponds to hypercharge except for the third family, which has had its charges sign changed. The second point,  $B_3$ , corresponds to B - L where the third family has had its charge's sign changed and the charges  $X_{H_u}$  and  $X_{H_d}$  are modified from their usual values of zero. We will denote the line between them  $Y_3B_3$ .

To see how  $Y_3B_3$  will enable us to find all solutions, let us first define the space PL', defined to be the subspace of PL whose points are orthogonal to  $Y_3$  and  $B_3$  with respect to the standard scalar product on  $\mathbb{Q}^{20}$ . Every point in PL lies on a plane  $Y_3B_3R$  formed by  $Y_3B_3$  and a point  $R \in PL'$ . Thus, we can restrict our attention to looking at such planes, and the points within them which satisfy the ACCs.

Generically (we will look at the few exceptions shortly), the intersection of the quadratic with  $Y_3B_3R$  consists of the union of  $Y_3B_3$  and another line  $L_q$ , as we will see explicitly in the next subsection. In a similar way, the intersection of the cubic with  $Y_3B_3R$  consists of the union of  $Y_3B_3$  and another line  $L_c$ . The intersection of the projective variety defined by the ACCs, and  $Y_3B_3R$  then consists of the line  $Y_3B_3$  and a single point which is the intersection of  $L_q$  and  $L_c$  as shown in the top left-hand panel of figure 2. Finding this point, which is a new solution to the ACCs, is a trivial task, as we shall shortly see.

Let us now look at the exceptional cases, all of which are illustrated in figure 2. They correspond to the following situations: (a) the whole plane lies in the quadratic but not the cubic; (b) the whole plane lies in both the quadratic and the cubic; (c) the whole line  $L_q$  lies in the cubic. The asymmetry between the quadratic and the cubic here is simply a manifestation of the ordering in which we will do our manipulations in the next subsection, and nothing more subtle.

We reiterate that since every point in PL lies in a plane  $Y_3B_3R$ , by finding all solutions in all such planes (considering either the generic case or the exceptional cases) we can find every point in the projective variety. In section 4.4 we will give an explicit parameterisation of the solution formed by such considerations. In section 4.5 a right-inverse to this parametrisation will be given explicitly demonstrating its full generality.

	Q	Q	Q	n	n	n	e	e	e	u	u	u	d	d	d	L	L	L	$\tilde{H}_d$	$\tilde{H}_u$
$Y_3$	1	1	-1	0	0	0	-6	-6	6	4	4	-4	-2	-2	2	-3	-3	3	-3	3
$B_3$	1	1	-1	-3	-3	3	-3	-3	3	1	1	-1	1	1	-1	-3	-3	3	-3	3

**Table 4.** A new ordering for the anomaly-free charge assignments  $Y'_3$  and  $B'_3$  given in table 2, adapted for the analytic solution. Each row lists the  $\mathfrak{u}(1)_X$  charges of the (left-handed or right-handed) chiral fermions of a model.



Figure 2. A schematic of our method. In the generic case (top left), the plane contains one further line in the quadratic and one further line in the cubic. These intercept at a point  $P_0$ , which is our new solution to all ACCs. The exceptional cases are shown in diagrams (a), (b) and (c). In (a),  $L_c$ is a line of solutions to all ACCs. In (b), the  $Y_3B_3R$  is a whole plane of solutions to all ACCs and in (c), the line  $L_q$  is a line of solutions to all ACCs.

#### 4.3 Derivation of the analytic solution

We now give a more detailed description of our solution. To this end we define

$$q(X,Z) := \sum_{i=1}^{3} (X_{Q_i} Z_{Q_i} - X_{L_i} Z_{L_i} - 2X_{u_i} Z_{u_i} + X_{d_i} Z_{d_i} + X_{e_i} Z_{e_i}) + X_{H_u} Z_{H_u} - X_{H_d} Z_{H_d},$$

$$c(W,X,Z) := \sum_{i=1}^{3} (6W_{Q_i} X_{Q_i} Z_{Q_i} + 2W_{L_i} X_{L_i} Z_{L_i} - 3W_{u_i} X_{u_i} Z_{u_i} - 3W_{d_i} X_{d_i} Z_{d_i} - W_{e_i} X_{e_i} Z_{e_i} - W_{n_i} X_{n_i} Z_{n_i}) + 2W_{H_u} X_{H_u} Z_{H_u} + 2W_{H_d} X_{H_d} Z_{H_d},$$

$$(4.1)$$

which are respectively derived from the quadratic and cubic ACCs with e.g.  $X_{Q_1}^3$  replaced with  $W_{Q_1}X_{Q_1}Z_{Q_1}$ . The maps q and c are the unique trilinear forms which return, respectively, the quadratic (2.6) ACC and the cubic ACC (2.5), when all inputs coincide.

A point  $R \in PL'$  can be parameterised by the 12 charges  $R_{S_1}$  for  $S \in \{Q, n, e, u, L, d\}$ ,  $R_{S_2}$  for  $S \in \{e, L, d\}$ ,  $R_{d_3}$ ,  $R_{H_u}$  and  $R_{H_d}$  as well as an extra two parameters  $R_1$  and  $R_2$ . The remaining charges are given by

$$\begin{aligned} R_{Q_2} &= 3(R_1 + R_{H_d} + R_{L_1} + R_{L_2} + R_{Q_1}) + 4R_2 + 2R_{d_3} + R_{e_1} + R_{e_2}, \\ R_{Q_3} &= -4(R_1 + R_2 + R_{Q_1}) - 2(R_{d_3} + R_{e_1} + R_{e_2}) - 3(R_{H_d} + R_{L_1} + R_{L_2}), \\ R_{n_2} &= 4R_1 + R_2 + 3(R_{e_1} + R_{e_2}) - R_{n_1} + R_{Q_1}, \\ R_{n_3} &= -(2R_1 + R_2 + R_{d_1} + R_{d_2} + R_{d_3} + R_{e_1} + R_{e_2} + R_{Q_1}), \\ R_{e_3} &= (R_{d_1} + R_{d_2} + R_{d_3}) + 3(R_{e_1} + R_{e_2}) + 4R_1, \\ R_{u_2} &= -(6R_1 + R_2 + R_{d_1} + R_{d_2} + 2R_{d_3} + 4(R_{e_1} + R_{e_2}) + R_{Q_1} + R_{u_1}), \\ R_{u_3} &= 4R_1 + R_2 + R_{d_3} + 2(R_{e_1} + R_{e_2}) + R_{Q_1}, \\ R_{L_3} &= 3(R_1 + R_{e_1} + R_{e_2}) - (R_{L_1} + R_{L_2}) - (R_{H_u} + R_{H_d}). \end{aligned}$$

$$(4.2)$$

Substituting the generic point,  $\alpha Y_3 + \beta B_3 + \gamma R$ , on the plane  $Y_3 B_3 R$  into the quadratic gives

$$\gamma(2\alpha q(Y_3, R) + 2\beta q(B_3, R) + \gamma q(R, R)) = 0.$$
(4.3)

Putting the exceptional cases to one side for now, this equation generically has two lines of solutions: one specified by  $\gamma = 0$ , namely  $Y_3B_3$ , and a new line  $L_q$ , a general point of which is given by

$$P_{c_1,c_2,c_3}^{(c)} = \{c_2q(R,R) - 2c_3q(B_3,R)\}Y_3 + \{2c_3q(Y_3,R) - c_1q(R,R)\}B_3 + \{c_1q(B_3,R) - c_2q(Y_3,R)\}R,$$
(4.4)

where  $c_1, c_2, c_3 \in \mathbb{Z}$  (over-)parameterise the line.<sup>9</sup>

On making the same substitution into the cubic we would get a similar line. However, since we are only interested in the intersection of these two lines, it is sufficient to substitute  $P_{c_1,c_2,c_3}^{(c)}$  into the cubic. This yields

$$4\{c_1q(B_3, R) - c_2q(Y_3, S)\}^2 \{\{2q(B_3, R)c(R, R, R)3q(R, R)c(B_3, R, R)\}c_1 + \{3q(R, R)c(Y_3, R, R) - 2q(Y_3, R)c(R, R, R)\}c_2 + 6\{q(Y_3, R)c(B_3, R, R) - q(B_3, R)c(Y_3, R, R)\}c_3\} = 0.$$
(4.5)

Solving for  $c_1, c_2, c_3$  generically gives the new solution to the ACCs

$$P_{0} = \{3q(R, R)c(B_{3}, R, R) - 2q(B_{3}, R)c(R, R, R)\}Y_{3} + \{2q(Y_{3}, R)c(R, R, R) - 3q(R, R)c(Y_{3}, R, R)\}B_{3} + 6\{q(B_{3}, R)c(Y_{3}, R, R) - q(Y_{3}, R)c(B_{3}, R, R)\}R.$$

$$(4.6)$$

Let us now return to the exceptional cases.

(a) The plane lies entirely in the quadratic, but not in the cubic: this occurs when  $q(Y_3, R) = 0$ ,  $q(B_3, R) = 0$  and  $q(Y_3, R) = 0$  but at least one of  $c(Y_3, R, R)$ ,  $c(B_3, R, R)$  and c(R, R, R) is non-zero. In this case, we have a line of solutions (over-)parameterised by  $a_1, a_2$ , and  $a_3 \in \mathbb{Z}$  and given by

$$P_{a_1,a_2,a_3}^{(a)} = \{a_2c(R,R,R) - 3a_3c(B_3,R,R)\}Y_3 + \{3a_3c(Y_3,R,R) - a_1c(R,R,R)\}B_3 + 3\{a_1c(B_3,R,R) - a_2c(Y_3,R,R)\}R.$$
(4.7)

<sup>&</sup>lt;sup>9</sup>Our use of projective space allows us to use, by clearing denominators,  $\mathbb{Z}$  for these parameters rather than  $\mathbb{Q}$ .

(b) The plane lies entirely within the quadratic and the cubic: this occurs when q(Y<sub>3</sub>, R) = 0, q(B<sub>3</sub>, R), q(Y<sub>3</sub>, R) = 0, c(Y<sub>3</sub>, R, R) = 0, c(B<sub>3</sub>, R, R) = 0 and c(R, R, R) = 0. In this case, every point on the plane lies in the variety. We then parameterise the plane withb<sub>1</sub>, b<sub>2</sub>, b<sub>3</sub> ∈ Z:

$$P_{b_1,b_2,b_3}^{(b)} = b_1 Y_3 + b_2 B_3 + b_3 R \tag{4.8}$$

(c) The line in the quadratic and the cubic are the same lines: this occurs (excluding the case where the line is just  $\alpha Y_3 + \beta B_3$ ) when

$$2q(B_3, R)c(R, R, R) = 3q(R, R)c(B_3, R, R),$$
  

$$2q(Y_3, R)c(R, R, R) = 3q(R, R)c(Y_3, R, R),$$
  

$$q(B_3, R)c(Y_3, R, R) = q(Y_3, R)c(B_3, R, R).$$
(4.9)

In this case our solution is the line  $P_{c_1,c_2,c_3}^{(c)}$ .

It is possible to combine these exceptional cases and the generic case into one parameterisation of the solution using Kronecker delta functions. This overall parameterisation is given by

$$P = P_{0} + \delta_{q(Y_{3},R),0} \delta_{q(B_{3},R),0} \delta_{q(R,R),0} \{ P_{a_{1},a_{2},a_{3}}^{(a)} + \delta_{c(Y_{3},R,R),0} \delta_{c(B_{3},R,R),0} \delta_{c(R,R,R),0} P_{b_{1},b_{2},b_{3}}^{(b)} \} + \delta_{2q(B_{3},R)c(R,R,R),3q(R,R)c(B_{3},R,R)} \delta_{2q(Y_{3},R)c(R,R,R),3q(R,R)c(Y_{3},R,R)} \times \delta_{q(B_{3},R)c(Y_{3},R,R),q(Y_{3},R)c(B_{3},R,R)} P_{c_{1},c_{2},c_{3}}^{(c)}.$$

$$(4.10)$$

This parameterisation is written in terms of the 12 charges and two extra parameters specifying R, as well as the parameters  $a_1$ ,  $a_2$ ,  $a_3$ ,  $b_1$ ,  $b_2$ ,  $b_3$ ,  $c_1$ ,  $c_2$  and  $c_3$ , which are needed in the exceptional cases. Taking these parameters to be integers returns an integer-valued solution.

## 4.4 Explicit parameterisation

To write the parameterisation more explicitly, we define

$$\begin{split} \Gamma &:= \{ 3q(R,R)c(B_3,R,R) - 2q(B_3,R)c(R,R,R) \} \\ &+ \delta_{q(Y_3,R),0} \delta_{q(B_3,R),0} \delta_{q(R,R),0} \\ &\quad (a_2c(R,R,R) - 3a_3c(B_3,R,R) + \delta_{c(Y_3,R,R),0} \delta_{c(B_3,R,R),0} \delta_{c(R,R,R),0} b_1) \\ &+ \delta_{2q(B_3,R)c(R,R,R),3q(R,R)c(B_3,R,R)} \delta_{2q(Y_3,R)c(R,R,R),3q(R,R)c(Y_3,R,R)} \\ &\quad \delta_{q(B_3,R)c(Y_3,R,R),q(Y_3,R)c(B_3,R,R)} (c_2q(R,R) - 2c_3q(B_3,R)), \\ \Sigma &:= \{ 2q(Y_3,R)c(R,R,R) - 3q(R,R)c(Y_3,R,R) \} \\ &+ \delta_{q(Y_3,R),0} \delta_{q(B_3,R),0} \delta_{q(R,R),0} \\ &\times \{ 3a_3c(Y_3,R,R) - a_1c(R,R,R) + \delta_{c(Y_3,R,R),0} \delta_{c(B_3,R,R),0} \delta_{c(R,R,R),0} b_2 \} \\ &+ \delta_{2q(B_3,R)c(R,R,R),3q(R,R)c(B_3,R,R)} \{ 2c_3q(Y_3,R) - c_1q(R,R) \}, \end{split}$$

$$\Lambda := 6\{q(B_3, R)c(Y_3, R, R) - q(Y_3, R)c(B_3, R, R)\} 
+ \delta_{q(Y_3, R),0}\delta_{q(B_3, R),0}\delta_{q(R, R),0} 
\times \{3\{a_1c(B_3, R, R) - a_2c(Y_3, R, R)\} + \delta_{c(Y_3, R, R),0}\delta_{c(B_3, R, R),0}\delta_{c(R, R, R),0}b_3\} 
+ \delta_{2q(B_3, R)c(R, R, R),3q(R, R)c(B_3, R, R)}\delta_{2q(Y_3, R)c(R, R, R),3q(R, R)c(Y_3, R, R)} 
\delta_{q(B_3, R)c(Y_3, R, R),q(Y_3, R)c(B_3, R, R)}\{2c_1q(B_3, R) - 2c_2q(Y_3, R)\}.$$
(4.11)

Then the charges are given explicitly by fourth order polynomials in the coordinates of R:

$$\begin{split} X_{Q_1} &= \Gamma + \Sigma + \Lambda R_{Q_1}, & X_{Q_2} = \Gamma + \Sigma + \Lambda R_{Q_2}, & X_{Q_3} = -\Gamma - \Sigma + \Lambda R_{Q_3}, \\ X_{n_1} &= -3\Sigma + \Lambda R_{n_1}, & X_{n_2} = -3\Sigma + \Lambda R_{n_2}, & X_{n_3} = 3\Sigma + \Lambda R_{n_3}, \\ X_{e_1} &= -6\Gamma - 3\Sigma + \Lambda R_{e_1}, & X_{e_2} = -6\Gamma - 3\Sigma + \Lambda R_{e_2}, & X_{e_3} = 6\Gamma + 3\Sigma + \Lambda R_{e_3}, \\ X_{u_1} &= 4\Gamma + \Sigma + \Lambda R_{u_1}, & X_{u_2} = 4\Gamma + \Sigma + \Lambda R_{u_2}, & X_{u_3} = -4\Gamma - \Sigma + \Lambda R_{u_3}, \\ X_{L_1} &= -3\Gamma - 3\Sigma + \Lambda R_{L_1}, & X_{L_2} = -3\Gamma - 3\Sigma + \Lambda R_{L_2}, & X_{L_3} = 3\Gamma + 3\Sigma + \Lambda R_{L_3}, \\ X_{d_1} &= -2\Gamma + \Sigma + \Lambda R_{d_1}, & X_{d_2} = -2\Gamma + \Sigma + \Lambda R_{d_2}, & X_{d_3} = 2\Gamma - \Sigma + \Lambda R_{d_3}, \\ X_{H_u} &= 3\Gamma + 3\Sigma + \Lambda R_{H_u}, & X_{H_d} = -3\Gamma - 3\Sigma + \Lambda R_{H_d}. \end{split}$$

$$(4.12)$$

# 4.5 Right inverse

As previously mentioned, this analytic solution has a right inverse, demonstrating its complete generality. Specifically, let T be a known solution and define the point  $G = 108T - (Y_3 \cdot T - B_3 \cdot T)Y_3 - (2B_3 \cdot T - Y_3 \cdot T)B_3$ , where '·' is the usual scalar product. The point G can be thought of as T with its components in the line  $\alpha Y_3 + \beta B_3$  projected out. The parameters  $R_{X_j} = G_{X_j}$  (for  $X_j$  as above), and

$$R_{1} = -\sum_{i=1}^{3} (G_{d_{i}} + G_{L_{i}}) + G_{e_{3}} - G_{H_{u}} - G_{H_{d}},$$

$$R_{2} = \sum_{i=1}^{3} (G_{d_{i}} - G_{e_{i}} + 2G_{L_{i}}) - G_{e_{3}} + 2G_{H_{u}} + 2G_{H_{d}} - G_{Q_{1}} - G_{n_{2}},$$

$$a_{1} = c(B_{3}, T, T), \quad a_{2} = -c(Y_{3}, T, T),$$

$$a_{3} = -c(B_{3}, T, T)(Y_{3} \cdot T - B_{3} \cdot T) + c(Y_{3}, T, T)(2B_{3} \cdot T - Y_{3} \cdot T).$$

$$b_{1} = (Y_{3} \cdot T - B_{3} \cdot T) \quad b_{2} = (2B_{3} \cdot T - Y_{3} \cdot T) \quad b_{3} = 1$$

$$c_{1} = q(B_{3}, T), \quad c_{2} = -q(Y_{3}, T),$$

$$c_{3} = -q(B_{3}, T)(Y_{3} \cdot T - B_{3} \cdot T) + q(Y_{3}, T)(2B_{3} \cdot T - Y_{3} \cdot T), \quad (4.13)$$

return the point T when substituted into the above analytic solution. In fact, they return T up to a multiplicative constant given by

$$6 \times 108^{4} (q(B_{3},T)c(Y_{3},T,T) - q(Y_{3},T)c(B_{3},T,T)) + \delta_{q(Y_{3},T),0}\delta_{q(B_{3},T),0}(3 \times 108^{3}(c(B_{3},T,T)^{2} + c(Y_{3},T,T)^{2}) + 108\delta_{c(Y_{3},T,T),0}\delta_{c(B_{3},T,T),0}) + 2 \times 108^{2}\delta_{q(B_{3},T)c(Y_{3},T,T),q(Y_{3},T)c(B_{3},T,T)}(q(B_{3},T)^{2} + q(Y_{3},T)^{2})$$
(4.14)

but given that our discussion above has been implicitly in projective space, such multiplicative factors are not relevant. In the Zenodo repository [28] we provide a Mathematica<sup>TM</sup> script containing the analytic solution, allowing one to generate solutions at will.

## 5 Checks of the solutions

The material content of section 3 is a list of all inequivalent anomaly-free charge assignments up to a fixed  $Q_{\text{max}}$ . A skeptic could justly ask the question: how does one know this list is complete without redundancies? The algorithm used does guarantee it, but one wishes to mitigate potential errors involved in its computer implementation. A similar level of scrutiny can be applied to the analytic solution of section 4. Although here one might hope the correctness of the solution is mathematically clear-cut, due diligence requires that we should try to ensure that no fallacies have been committed. Happily, several checks can be carried out to satisfy all but the most fastidious skeptic. These checks work in three different modes: consistency checks within the numerical solutions, consistency checks within the analytic solution alone, and cross-checks between the two. The ability to do cross-checks between the two is one of several advantages for providing both. Let us discuss the checks performed for each mode in turn. We note in passing that all checks were carried out successfully.

For any computer program, one useful check is to make a second structurally different program but with the same expected outcome. To this end, we produced a second different program (this one did not use lexicographic ordering, but instead used an ordering similar to that in ref. [10]). The two outputs where then compared and found to agree.

The addition of hypercharge to any solution also leads to a solution, as stated in (iii) of section 2.1. This provides a check of the computer program as follows: each solution for a given  $Q_{\text{max}}$  had multiples of hypercharge added or subtracted from it up to three times. If the resulting charges had a height less than or equal to 10, the binary search method discussed in section 3 was used to confirm that the solution was present in our  $Q_{\text{max}} = 10$  list.

Turning to the analytic solution, the most primitive check is to randomly choose parameters, generate the corresponding charges and confirm that they satisfy the ACCs. This check was carried out on  $10^5$  randomly generated solutions.

The fact that we have a right inverse for our parameterisation means that we can take a solution, apply the inverse and then the parameterisation to return another solution. If our analysis is correct this new solution should agree with the one we started with (up to a scaling). This was carried out on, again,  $10^5$  randomly generated solutions. It was also carried out on all the scanned solutions in our list for  $Q_{\text{max}} = 10$ , thereby providing the first cross check between the numerical and analytic solutions.

The second cross-check between the numerical and analytic solutions was to generate random solutions using the analytic solution, then to identify those of height less than or equal to 10 and confirm that these appear in the numerical solution via the binary search algorithm.

#### 6 Examples of filters

As mentioned in section 3, we envisage that our list of solutions to ACCs (2.2)–(2.7) may be used by supersymmetric model builders by scanning with a computer program, filtering out the charges suited to a particular model. In this section, we provide examples of how the list of solutions might be filtered in order to identify sets of charge assignments with various possible desirable phenomenological properties or uses.<sup>10</sup> Note that in what follows, as in section 4, we will distinguish  $H_d$  from  $L_i$ , and the number of solutions satisfying each constraint is therefore to be compared with the second column of table 2.

We summarise the phenomenological constraints of this section in table 9. We emphasise that the filters used throughout this section provide an initial exploration into the constraints we expect will be most commonly needed by model builders. We expect that the scope of this list is much broader than the phenomenological applications dealt with here, and by making the list of charge assignments publicly available on Zenodo [28] we encourage model builders to search for charge assignments of more specific interest.

#### 6.1 The superpotential

In general, interactions between the chiral supermultiplets of the MSSM are given by the superpotential  $W = W_{R_p} + W_{LV} + W_{BV}$ , where

$$W_{R_{p}} = \mu \hat{H}_{u} \hat{H}_{d} + (y_{u})_{ij} \hat{U}_{i}^{c} \hat{Q}_{j} \hat{H}_{u} + (y_{d})_{ij} \hat{D}_{i}^{c} \hat{Q}_{j} \hat{H}_{d} + (y_{e})_{ij} \hat{E}_{i}^{c} \hat{L}_{j} \hat{H}_{d},$$

$$W_{LV} = \frac{1}{2} \lambda^{ijk} \hat{L}_{i} \hat{L}_{j} \hat{E}_{k}^{c} + \lambda'^{ijk} \hat{L}_{i} \hat{Q}_{j} \hat{D}_{k}^{c} + \mu'^{i} \hat{L}_{i} \hat{H}_{u},$$

$$W_{BV} = \frac{1}{2} \lambda''^{ijk} \hat{U}_{i}^{c} \hat{D}_{j}^{c} \hat{D}_{k}^{c}.$$
(6.1)

 $\hat{U}_i^c, \hat{D}_i^c, \hat{Q}_i, \hat{L}_i, \hat{E}_i^c, \hat{H}_u$  and  $\hat{H}_d$  denote the chiral supermultiplets containing of table 3, and we denote flavour indices by  $i, j, k \in \{1, 2, 3\}$ .  $\lambda^{ijk}, \lambda'^{ijk}, \lambda''^{ijk}, (y_{u,d,e})_{ij}$  are all dimensionless coupling constants and  $\mu, \mu'^i$  each have mass dimension 1. Gauge indices have been suppressed. Note that here we ignore the neutrino chiral supermultiplets  $\hat{N}_i^c$ , postponing their discussion until section 6.3. Here  $W_{R_p}$  denotes terms invariant under *R*-parity, whereas *R*-parity is violated in the *L* and *B*-violating terms  $W_{LV}$  and  $W_{BV}$  respectively.

## 6.1.1 The $\mu$ problem

The MSSM has a fine tuning problem associated with the  $\mu \hat{H}_u \hat{H}_d$  term. Given that this term respects supersymmetry and gauge symmetry, there is no explicitly stated reason for the scale of  $\mu$  to be small. The gauge group can be extended by  $U(1)_X$  to provide a solution to this so-called  $\mu$  problem [29]. This is achieved by charging  $\hat{H}_u$  and  $\hat{H}_d$  under  $U(1)_X$  such that the  $\mu$  term above is forbidden by the  $U(1)_X$  symmetry. Instead, the flavon  $\theta$  is charged, allowing a term of the form (where h is a dimensionless coupling constant)

$$W \supset h\theta \hat{H}_u \hat{H}_d \to h \langle \theta \rangle \hat{H}_u \hat{H}_d, \tag{6.2}$$

<sup>&</sup>lt;sup>10</sup>Computer programs implementing these filters are available on Zenodo [28].

such that when the U(1)<sub>X</sub> symmetry is spontaneously broken, the scalar component of  $\theta$  acquires a vacuum expectation value  $\langle \theta \rangle$  at the TeV scale i.e. the  $\mu$  term is dynamically generated.<sup>11</sup> The  $\mu\nu$ SSM [30–33] also solves the  $\mu$  problem in precisely this manner. Any model with such a dynamically generated  $\mu$  term is often referred to as the next-to-minimal supersymmetric standard model (NMSSM). The NMSSM has received much attention in the literature [34–37].

Remembering that we shall pick one of the  $\hat{N}_i^c$  chiral superfields with a non-zero charge to be the flavon chiral superfield  $\theta^c$ , which has a non-zero X charge out of necessity, we search for such solutions in our list of charges by applying the conditions

$$\exists i \in \{1, 2, 3\}: \quad X_{H_u} + X_{H_d} = X_{n_i} \neq 0, \tag{6.3}$$

where we take  $\theta^c$  to be the  $\hat{N}_i^c$  superfield which satisfies this condition.<sup>12</sup> We find a total of 77 solutions satisfying these constraints with  $Q_{\text{max}} = 1$ , constituting ~ 30% of the full  $Q_{\text{max}} = 1$  list. This percentage reduces to 20% when  $Q_{\text{max}} = 4$ , and 11% when  $Q_{\text{max}} = 10$ , providing in this case a total of 649 831 168 options for a dynamically generated  $\mu$  term. Note that these solutions may not all produce viable models, and that the specifics of further model building will place further constraints on the fermion charges. Here we provide only an initial filter as a preliminary example of how the list of solutions may be queried for the purpose of phenomenological model building.

#### 6.1.2 A renormalisable Yukawa sector

In contrast to the rather weak constraints of (6.3), we may place strong conditions on the Yukawa sector by requiring that all renormalisable Yukawa couplings of charged fermions are allowed in the superpotential  $W_{R_p}$  by being  $U(1)_X$  gauge invariant, i.e. they must satisfy the following equations  $\forall i, j \in \{1, 2, 3\}$ :

$$X_{Q_i} + X_{H_u} - X_{u_j} = 0, \qquad X_{Q_i} + X_{H_d} - X_{d_j} = 0, \qquad X_{L_i} + X_{H_d} - X_{e_j} = 0.$$
(6.4)

(6.4) implies family universality for the species Q, e, u, L and d. For the non-supersymmetric case, it has been shown that anomaly-free charge assignments exist which allow all of the renormalisable Yukawa terms [20]. One can show that in the  $\nu$ MSSM, we obtain one solution for each non-supersymmetric solution of [20], where we must additionally fix  $H_u$  and  $H_d$  to satisfy

$$3X_{H_u} = -3X_{H_d} = -3\sum_{i=1}^3 X_{Q_i} - \sum_{i=1}^3 X_{n_i}.$$
(6.5)

(6.5) means that there cannot be any overlap with the solutions satisfying (6.3), i.e. none of these solutions can simultaneously solve the  $\mu$  problem. By filtering through our list of charges, we find 2 solutions allowing a fully renormalisable Yukawa sector with  $Q_{\text{max}} = 1$  and 5 with  $Q_{\text{max}} = 4$ , as shown in table 5. The full list of  $Q_{\text{max}} = 10$  solutions comprises 38 such solutions.

<sup>&</sup>lt;sup>11</sup>Further detailed model building is required to make sure that  $\langle \theta \rangle \sim \mathcal{O}(\text{TeV})$ , but we shall merely assume here that this is possible.

<sup>&</sup>lt;sup>12</sup>The U $\mu\nu$ SSM [38] uses (6.3) in a certain U(1)' extension of the  $\nu$ MSSM (involving additional quark fields) to solve the  $\mu$  problem, also.

Q	Q	Q	n	n	n	e	e	e	u	u	u	d	d	d	L	L	L	$\tilde{H}_d$	$\tilde{H}_u$
0	0	0	-1	0	1	0	0	0	0	0	0	0	0	0	0	0	0	0	0
0	0	0	-1	-1	-1	1	1	1	-1	-1	-1	1	1	1	0	0	0	1	-1
-1	-1	-1	3	3	3	3	3	3	-1	-1	-1	-1	-1	-1	3	3	3	0	0
-1	-1	-1	2	2	2	4	4	4	-2	-2	-2	0	0	0	3	3	3	1	-1
-1	-1	-1	4	4	4	2	2	2	0	0	0	-2	-2	-2	3	3	3	-1	1

**Table 5.** Anomaly-free charge assignments with  $Q_{\max} = 4$  allowing all Yukawa terms at the renormalisable level. Each row lists the  $\mathfrak{u}(1)_X$  charges of the (left-handed or right-handed) chiral fermions of a model. Note that all listed solutions satisfy  $X_{H_u} + X_{H_d} = 0$ , reducing the ACCs to those of the SM after substitution.

We will now relax the assumption that all Yukawa terms must be present in the Lagrangian at the renormalisable level. We will enforce that the top and bottom quark and the tau lepton tree-level Yukawa terms can be present (since they are closer to order 1 and so more difficult to explain by non-renormalisable or loop interactions, which imply a suppression below order 1) by applying the constraints

$$\exists \sigma_1, \sigma_2, \sigma_3, \sigma_4, \sigma_5 \in S_3 : \quad X_{Q_{\sigma_1(3)}} + X_{H_u} - X_{u_{\sigma_2(3)}} = 0,$$

$$X_{Q_{\sigma_1(3)}} + X_{H_d} - X_{d_{\sigma_3(3)}} = 0,$$

$$X_{L_{\sigma_4(3)}} + X_{H_d} - X_{e_{\sigma_5(3)}} = 0,$$

$$(6.6)$$

where  $S_3$  is the group of permutations of 3 objects. We expect  $Q_{\sigma_1(3)}$ ,  $u_{\sigma_2(3)}$  and  $d_{\sigma_3(3)}$  to be predominantly third generation quarks, and similarly  $L_{\sigma_4(3)}$  and  $e_{\sigma_5(3)}$  to be predominantly composed of third generation leptons. We will further assume that tree-level renormalisable Yukawa terms are not present for the first and second generation fermions by forbidding all other terms in the Yukawa matrices. We can express these constraints by first defining

$$P_{ijklmn} := (X_{Q_i} + X_{H_u} = X_{u_j}) \land (X_{Q_n} + X_{H_d} = X_{d_k}) \land (X_{L_l} + X_{H_d} = X_{e_m}),$$
(6.7)

(where  $\wedge$  means logical 'and') and then imposing

$$(\exists ! i, j, k, l, m, n \in \{1, 2, 3\} : P_{ijklmn}) \land (\forall i, j, k, l, m, n \in \{1, 2, 3\} P_{ijklmn} \Rightarrow n = i), \quad (6.8)$$

where, in standard logic notation,  $\exists!$  means 'there exists a unique'.

This choice is made with the fermion mass problem in mind: it allows larger masses to be generated for the top, bottom and tau through the standard Yukawa terms, but forbids them for the light quarks, producing a mass hierarchy between the light and heavy fermions. In ref. [22] it was shown that the chiral fermions can obtain their masses at loop level through the interactions with their superpartners by including non-holomorphic soft terms in the Lagrangian density. Alternatively, light fermion masses may be acquired through non-renormalisable operators after the flavon  $\theta$  breaks  $U(1)_X$ . Either of these mechanisms require the Lagrangian density to contain terms which will further constrain the charges. We shall assume that all first and second generation fermions acquire their masses through some mechanism such as one of these two, but leave the more model dependent effect of any additional constraints to future investigations. We find that when  $Q_{\text{max}} = 1$ , the list contains 2 solutions satisfying the constraints of (6.8). At  $Q_{\text{max}} = 4$  a total of 15 818 solutions pass these constraints, and at  $Q_{\text{max}} = 10$ this number grows to 34 646 735. This makes clear that by imposing these constraints, not only do we begin to address the fermion mass problem, but we make way for a larger number of options for model-building compared to those of a fully allowed renormalisable Yukawa sector. For example, when  $Q_{\text{max}} = 2$  there are 8 solutions which simultaneously solve the  $\mu$  problem and satisfy (6.8). This overlap grows to 2 954 solutions when  $Q_{\text{max}} = 4$ and 4 088 200 solutions when  $Q_{\text{max}} = 10$ . Furthermore, the constraints (6.8) are inherently flavour non-universal, and thus have the potential to address the *B* anomalies. This overlap will be discussed in more detail in section 6.2.

#### 6.1.3 *R*-parity violation

In contrast to the SM, L and B violating terms are allowed by the field content and gauge symmetries of the MSSM, as shown in (6.1). The simultaneous presence of both B and Lviolating terms will lead to proton decay in contravention to experimental bounds unless one introduces a large degree of fine tuning. Usually, all terms in  $W_{LV}$  and  $W_{BV}$  are forbidden by the imposition of R-parity. In the case that R-parity is not imposed though, we may ask that our  $U(1)_X$  symmetry maintains the stability of the proton instead. We can form three broad sets of solution within this requirement: where *all* R-parity violating terms are banned (this will also maintain the stability of the lightest supersymmetric particle, which may have the properties to constitute cold dark matter), where all terms in  $W_{BV}$ are banned but where at least one term in  $W_{LV}$  is allowed, and those where all terms in  $W_{LV}$  are banned but at least one in  $W_{BV}$  is allowed. Terms such as those in  $W_{LV}$  give a Majorana mass term to left-handed neutrinos (sometimes through loop diagrams) without the need for right-handed neutrinos [39]. Terms in  $W_{BV}$ , on the other hand, can assist in baryogenesis [40].

We may ban terms in  $W_{BV}$  by imposing  $\forall i, j, k \in \{1, 2, 3\}$ 

$$X_{u_i} + X_{d_i} + X_{d_k} \neq 0 \tag{6.9}$$

where  $j \neq k$  since the antisymmetry of  $\lambda''^{ijk}$  in j, k forbids the j = k terms from appearing in the superpotential. Similarly, we may ban all terms in  $W_{LV}$  by imposing the conditions  $\forall i, j, k, l, m, n, p \in \{1, 2, 3\}$ 

$$X_{L_i} + X_{L_j} - X_{e_k} \neq 0, \qquad X_{L_l} + X_{Q_m} - X_{d_n} \neq 0, \qquad X_{L_p} + X_{H_u} \neq 0, \tag{6.10}$$

where  $i \neq j$  because  $\lambda^{ijk}$  is antisymmetric in i, j. At  $Q_{\max} = 1$  we find 8 solutions which ban all *R*-parity violating terms. These solutions are listed in table 6. We find a total of 51 solutions which ban  $W_{BV}$  while allowing terms in  $W_{LV}$ . We find no solutions which ban  $W_{LV}$  while allowing terms in  $W_{BV}$ , i.e. the only solutions which ban  $W_{LV}$  are those which ban all *R*-parity violation.

By increasing the maximum charge  $Q_{\text{max}}$ , we find solutions which ban  $W_{LV}$  while allowing *B*-violation. At  $Q_{\text{max}} = 10$ , we find 444 357 847 solutions which forbid  $W_{LV}$  while allowing terms in  $W_{BV}$ . We find 2 916 984 840 solutions which forbid  $W_{BV}$  while allowing

Q	Q	Q	n	n	n	e	e	e	u	u	u	d	d	d	L	L	L	$\tilde{H}_d$	$\tilde{H}_u$
-1	-1	1	1	1	1	1	1	1	-1	-1	1	-1	-1	1	1	1	1	1	1
-1	-1	1	1	1	1	1	1	1	-1	-1	1	-1	-1	1	1	1	1	0	0
-1	-1	1	1	1	1	1	1	1	-1	-1	1	-1	-1	1	1	1	1	1	0
-1	-1	1	1	1	1	1	1	1	-1	-1	1	-1	-1	1	1	1	1	1	-1
0	0	0	-1	-1	-1	1	1	1	-1	-1	-1	1	1	1	0	0	0	-1	1
0	0	0	-1	-1	-1	1	1	1	-1	-1	-1	1	1	1	0	0	1	1	-1
0	0	0	-1	-1	1	-1	1	1	-1	-1	1	-1	1	1	0	0	0	-1	1
0	0	0	-1	0	0	-1	1	1	-1	-1	1	-1	1	1	0	0	0	-1	1

**Table 6.** At  $Q_{\text{max}} = 1$ , we find 8 anomaly-free charge assignments in our list banning all *R*-parity violating terms in the MSSM superpotential. Each row lists the  $\mathfrak{u}(1)_X$  charges of the (left-handed or right-handed) chiral fermions of a model.

for terms in  $W_{LV}$  at  $Q_{\text{max}} = 10$ , and a total of 885 951 137 solutions which ban all *R*-parity violating solutions, constituting 14% of the list of charge assignments.

#### 6.2 *B* anomalies

Family-dependent charges in the quark and lepton sectors are well-motivated by the recent hints at lepton flavour non-universality associated with  $b \to s\ell^+\ell^-$  transitions [5–13], also known as 'B anomalies'. Global fits incorporating angular distributions and branching fractions point towards new physics contributions to the Wilson coefficients  $C_9$ ,  $C_{10}$  of weak effective theory Hamiltonian operators  $O_9$ ,  $O_{10}$ , respectively, where

$$O_9 = (\bar{s}'_L \gamma_\mu b'_L) (\bar{\mu}' \gamma^\mu \mu') \quad O_{10} = (\bar{s}'_L \gamma_\mu b'_L) (\bar{\mu}' \gamma^\mu \gamma^5 \mu').$$
(6.11)

Here the primes denote that the fermionic fields are in the mass eigenbasis. A vector-like new physics contribution to  $C_9$  with  $C_{10} = 0$ , or a new physics coupling to left-handed muons through the combination  $C_9 = -C_{10}$ , are both favoured by global fits [41] in comparison to the SM.

We will filter through our list in search of solutions potentially capable of explaining the so-called *B* anomalies via the mediation of flavour-changing Z' interactions, resulting from the spontaneously broken U(1)<sub>X</sub> symmetry. We will begin by searching for solutions for which there exists  $i, j \in \{1, 2, 3\}$  with  $Q_i$  and  $L_j$  charged. These will play the role of the left-handed bottom/top quark doublet and muon respectively, contributing to the effective operator  $(\bar{b}_L \gamma^{\mu} b_L)(\mu_L \gamma_{\mu} \mu_L) + \ldots$  once the heavy Z' is integrated out of the effective field theory. We will assume that a rotation to the mass eigenbasis will mix the down-type quarks such that the necessary  $\bar{b}'_L \gamma^{\mu} s'_L$  coupling is produced. As well as this, we will require that the left-handed leptons are not completely flavour universal, i.e.  $\exists k \in \{1, 2, 3\}$  such that  $X_{L_k} \neq X_{L_{\mu}}$ . This will ensure we can have the necessary  $\mu - e$  flavour non-universality to explain the  $b \to s \ell^+ \ell^-$  data.

We find 114 solutions with  $Q_{\text{max}} = 1$  satisfying these conditions, constituting approximately 43% of the total list. When  $Q_{\text{max}} = 10$  this number grows to 1567 142 472, roughly 25% of the full list of charge assignments. Such large numbers indicate that these conditions



Figure 3. We filter through our list to determine the number of solutions capable of solving the  $\mu$  problem and the *B*-anomalies, as well as those allowing only 3rd family Yukawa terms. We find an overlap between these applications, with 2 solutions at  $Q_{\text{max}} = 2$  satisfying all constraints and 1556 at  $Q_{\text{max}} = 4$ .

Q	Q	Q	n	n	n	e	e	e	u	u	u	d	d	d	L	L	L	$\tilde{H}_d$	$\tilde{H}_u$
-1	0	0	-2	1	2	1	2	2	-2	-1	0	-1	0	2	-1	0	2	1	1
-1	0	0	1	2	2	-2	1	2	-1	0	2	-2	-1	0	-1	0	2	1	1

**Table 7.** As depicted in figure 3, at  $Q_{\text{max}} = 2$ , only 2 anomaly-free charge assignments satisfy the constraints required to solve the  $\mu$  problem and the *B*-anomalies while allowing 3rd family Yukawa terms in the Lagrangian. Each row lists the  $\mathfrak{u}(1)_X$  charges of the (left-handed or right-handed) chiral fermions of a model.

leave the charges quite unconstrained, and thus we query the list further for interesting solutions. Firstly, there are solutions within this set which can simultaneously address the  $\mu$  problem and allow only renormalisable tree-level Yukawa terms for the top, bottom and tau. The overlap between each set of constraints is depicted in figure 3. Only 2 solutions can account for all three conditions when  $Q_{\text{max}} = 2$ , and are shown in table 7. This overlap grows when  $Q_{\text{max}} = 4$ , with 1556 solutions solving all three conditions.

Secondly, we will filter through the list for solutions that aren't obviously in danger of violating experimental constraints. Following the motivation of refs. [9, 42], we search for solutions with uniform light quark charges so as to avoid constraints on flavour-violation in the light quark sector. Additionally, we will search for solutions which feature zero coupling of the electron to the associated Z' i.e.  $\exists i, j \in \{1, 2, 3\}$  such that  $X_{L_i} = 0$  and  $X_{e_j} = 0$ . This is motivated by the strong experimental constraints originating from  $e^+e^-$  collisions at LEP. We find 21 such solutions that also allow only third family Yukawa terms and address the  $\mu$  problem in our list with  $Q_{\max} = 10$ . A selection of 10 of these solutions are listed in table 8. Here, in contrast to other tables, the index on each fermion denotes the family number (since these are used in the constraints), and we use  $\tilde{\theta}$  to denote the RH neutrino that plays the role of the (RH) flavino.

	$Q_1$	$Q_2$	$Q_3$	$\tilde{\theta}$	$n_1$	$n_2$	$e_1$	$e_2$	$e_3$	$u_1$	$u_2$	$u_3$	$d_1$	$d_2$	$d_3$	$L_1$	$L_2$	$L_3$	$\tilde{H}_d$	$\tilde{H}_u$
(a)	-3	-3	3	-6	0	10	0	9	5	-3	-3	-2	0	0	2	0	9	6	-1	-5
(b)	0	0	-2	6	-3	4	0	3	2	0	0	-1	-3	-3	3	0	3	-3	5	1
(c)	0	0	-2	10	-1	10	0	0	-7	3	3	5	-8	-8	1	0	6	-10	3	7
(d)	0	0	-1	8	-9	1	0	0	6	-2	-2	0	-2	-2	6	0	-4	-1	7	1
(e)	0	0	-5	6	8	10	0	-1	7	4	4	-4	-7	-7	0	0	7	2	5	1
(f)	0	0	-3	10	-9	3	0	8	6	-5	-5	2	0	0	2	0	-2	1	5	5
(g)	0	0	-3	2	0	3	0	7	6	-1	-1	-5	0	0	1	0	5	2	4	-2
(h)	0	0	-2	6	-6	-3	0	8	7	-5	-5	-1	2	2	3	0	-2	2	5	1
(i)	0	0	-2	6	-6	3	0	4	5	-3	-3	1	0	0	1	0	-2	2	3	3
(j)	0	0	-1	-4	-6	0	0	7	9	-5	-5	-4	7	7	-2	0	-3	10	-1	-3
(k)	0	0	-1	2	2	8	0	-5	-1	4	4	0	-5	-5	0	0	3	-2	1	1

**Table 8.** At  $Q_{\text{max}} = 10$  we find 21 solutions which simultaneously solve the  $\mu$  problem and B anomalies, allow 3rd family Yukawa terms and are well-suited to avoid strong experimental constraints from LEP and quark flavour violation between the first two families. A selection of 10 of these are shown here. Each row lists the  $\mathfrak{u}(1)_X$  charges of the (left-handed or right-handed) chiral fermions of a model.

The 10 solutions shown all feature suppressed couplings of the Z' to the light quarks, either because the light RH down-type quarks have zero charge, as in solution (a), or because the light LH quarks have zero charge as in solutions (b)-(k). In solutions (a) and (b), the muon has equal RH and LH charge i.e.  $L_2 = e_2$ . This results in a purely vector-like coupling with  $C_{10} = 0$ . Similarly, solutions (c), (d) and (e) are particularly interesting in that they all produce negative values of the ratio  $C_9/C_{10}$ , with (c) and (d) giving exactly  $C_9 = -C_{10}$ and solution (e) satisfying  $C_9 = -\frac{3}{4}C_{10}$ . In section 3 we queried the full list of charge assignments in search of known solutions in the literature, listed in table 3. None of these solutions are found in the list of 21 solutions passing our constraints: either because they cannot solve the  $\mu$  problem and address the 3rd family Yukawa terms simultaneously, or because they do not satisfy the constraints we impose to facilitate solving the B anomalies.

#### 6.3 Neutrino masses

Finally, we turn to the neutrinos. The inclusion of RH neutrinos has allowed us the flexibility to solve the ACCs while simultaneously addressing the phenomenological constraints of section 6.1 and section 6.2, as evidenced by the fact that these solutions often have nonzero charges for the RH neutrinos. In particular, this can be seen from table 8 in which all of the solutions feature nonzero charges for at least one of the RH neutrinos. It is then useful to ask what these charge assignments imply for the neutrino masses and mixings.

In order to describe neutrino masses and mixings, we extend the superpotential to include the following terms,

$$W = W_{R_p} + (y_{\nu})_{ij} N_i^c L_j H_u + (M_{\nu_R})_{ij} N_i^c N_j^c, \qquad (6.12)$$

where  $(y_{\nu})_{ij}$  is a 3 by 3 matrix of dimensionless Dirac Yukawa coupling constants and  $(M_{\nu_R})_{ij}$  is a 3 by 3 matrix of Majorana mass terms (of mass dimension 1) for the RH neutrinos. Neutrino masses are then produced through a Type-1 see-saw mechanism. Many alternative mechanisms for producing neutrino masses in the MSSM exist in the literature. Bilinear *R*-parity violating models extend the superpotential to include the *L*-violating  $\mu'_i L_i H_u$  terms which produce neutrino masses through mixing with the neutralinos [43, 44]. In [45], a suppressed Dirac mass term is produced after U(1)<sub>X</sub>-breaking, through the flavon's vacuum expectation value  $\langle \theta \rangle$ . The  $\mu \nu SSM$  extends the MSSM to produce neutrino masses through the inclusion of the trilinear term  $\kappa_{ijk} N_i^c N_i^c N_k^c$  in the superpotential [30–33].

While an investigation into each of these mechanisms and models is beyond the scope of this paper, we will filter through our list in search of solutions which allow all of the terms of (6.12), allowing all possible neutrino masses and mixings via the see-saw mechanism. Note that further model building will further constrain the fermion charges. Here we provide only an initial filter as a preliminary example of how the list of solutions may be queried for the purpose of generating neutrino masses via the see-saw mechanism. The solutions must satisfy the following constraints  $\forall i, j \in \{1, 2, 3\}$ 

$$X_{L_i} + X_{H_u} - X_{n_i} = 0, \qquad X_{n_i} + X_{n_i} = 0, \tag{6.13}$$

implying  $X_{n_i} = 0$  and  $X_{L_i} = -X_{H_u}$ . We find a total of 3 solutions with  $Q_{\max} = 1$  in our list satisfying these constraints. At  $Q_{\max} = 4$  a total of 118 solutions exist, and at  $Q_{\max} = 10$  the list contains 4 878 of these solutions.

# 7 Summary

Specific models incorporating the MSSM with an additional  $U(1)_X$  gauge group can combine the phenomenological advantages of supersymmetry with potential uses of the additional gauge factor and they have received quite some attention in the literature, particularly for the case where the  $U(1)_X$  charges are family dependent. We have found, for the first time, all charge assignments of the MSSM plus three SM-singlet chiral superfields which are free of local anomalies (the SM-singlets can produce neutrino masses as well as spontaneously break the  $U(1)_X$  symmetry). Chiral superfields in real representations can be added to any anomaly-free matter content and result in an anomaly-free solution, since the additional fermionic content will be in a vector-like representation of the gauge group and so its effects cancel in the anomalies. The local anomaly cancellation conditions described in section 2 constitute a system of six homogeneous coupled diophantine equations (2.2)–(2.7), the like of which are notoriously difficult to solve, in general.

Global anomalies are beyond the scope of our work; however, for the case of U(1) extensions of the usual SM gauge group, there are none [46]. One may question whether a quantum field theory absolutely has to be free from anomalies; after all, in an infra-red effective field theory (such as we might expect the  $MSSM \times U(1)_X$  to be) one can in principle add Wess-Zumino terms to the Lagrangian density in order to cancel them. Such terms can result from decoupling a heavy state from the effective field theory. In order to contribute to the anomaly though, the additional heavy state must be a chiral fermion of non-zero U(1)<sub>X</sub>

	Proposition	$\# Q_{\max} = 10$
$\mu$ problem	$\exists i \in \{1, 2, 3\} : X_{H_u} + X_{H_d} = X_{n_i} \land X_{n_i} \neq 0$	649831168
All renormalisable charged	$(\forall i, j, k, l, m, n \in \{1, 2, 3\} P_{ijklmn})$	38
fermion Yukawas		
Only 3rd family renormalisable	$(\exists !  i, j, k, l, m, n \in \{1, 2, 3\}: P_{ijklmn}) \land$	34646735
charged fermion Yukawas	$(\forall i, j, k, l, m, n \in \{1, 2, 3\} P_{ijklmn} \Rightarrow n = i)$	
L-conservation & $B$ -violation	$P_L \wedge \neg P_B$	444357847
B-conservation & $L$ -violation	$P_B \wedge \neg P_L$	2916984840
L & B-conservation	$P_L \wedge P_B$	885951137
B anomalies	$\exists i, j, k \in \{1, 2, 3\} : X_{Q_i} \neq 0 \land X_{L_j} \neq 0 \land$	1567142472
	$X_{L_k} \neq X_{L_j}$	
$B$ anomalies, $\mu$ problem, 3rd	See section 6.2	21
family Yukawa terms & exper-		
imental constraints		
See-saw $\nu$ masses	$\forall i, j \in \{1, 2, 3\} \ X_{L_i} + X_{H_u} = X_{n_j} \land X_{n_i} =$	4878
	$\left -X_{n_{j}}\right $	

**Table 9.** Summary of the phenomenological conditions applied in this paper, along with the number of inequivalent  $Q_{\max} = 10$  solutions which satisfy them. In the above we have used standard logic notation in which  $\forall$  reads as 'for all',  $\land$  as 'and',  $\lor$  as 'or',  $\exists$  as 'there exists',  $\exists$ ! as 'there exists a unique',  $\Rightarrow$  as 'implies', : as 'such that',  $\neg$  as 'not'. For the condition of allowing all renormalisable charged fermion Yukawa terms, we have used the proposition  $P_{ijklm}$  defined as  $P_{ijklmn} := (X_{Q_i} + X_{H_u} = X_{u_j} \land X_{Q_n} + X_{H_d} = X_{d_k} \land X_{L_l} + X_{H_d} = X_{e_m})$ . For the *R*-parity related conditions we have used the propositions  $P_L := \forall i, j, k, l, m, n, p \in \{1, 2, 3\}$   $i = j \lor (X_{L_i} + X_{L_j} - X_{e_k} \neq 0 \land X_{L_l} + X_{Q_m} - X_{d_n} \neq 0 \land X_{L_p} + X_{H_u} \neq 0)$ , and  $P_B := \forall i, j, k \in \{1, 2, 3\}$   $i = j \lor X_{u_i} + X_{d_j} + X_{d_k} \neq 0$ .

charge. It is then not a priori obvious how such a state may acquire a large mass, unless it is linked to the scale of  $U(1)_X$  breaking.<sup>13</sup> One recent non-supersymmetric U(1) gauge extension of the SM [13] has achieved this with some additional fermions that under the SM are in vector-like representations, but which are chiral with respect to  $U(1)_X$ . However, it is far from obvious whether this is possible in general model set-ups, particularly when several mixed anomalies do not cancel. From the model builder's point of view therefore, it is safer to begin with an anomaly-free effective field theory rather than having to worry about how such anomalies are cancelled.

We have provided the general analytic solution for the charges via a new geometric method (a different geometric method was previously employed to solve the anomaly cancellation conditions for non-supersymmetric  $U(1)_X$  extensions of the SM [18]) described in section 4. One inputs 23 integer parameters for each anomaly-free charge assignment.

<sup>&</sup>lt;sup>13</sup>Integrating the top quark out of the SM yields apparent gauge anomalies, but when one includes effective operators resulting from integrating it out, gauge symmetry is restored [47]. This is precisely a case where the heavy mass is linked to the symmetry breaking scale (in this case, of the electroweak symmetry).

A Mathematica<sup>TM</sup> program has been made publicly available which, given the input parameters, produces one such assignment. The general analytic solution passed various internal consistency checks. Whilst the general analytic solution can be difficult for model builders to use, it is useful for (among other things) providing non-trivial checks of any list of numerical solutions.

Anomaly-free charge assignments are scarce: for example, for heights up to 10, as figure 1 shows, only one out of some  $10^{12}$  (or so) inequivalent assignments is anomaly free. Despite their scarcity, the different assignments are still legion (we have identified over 1.6 billion up to a height of 10). The model builder is therefore faced with an enormous haystack in which to find the proverbial needle.

An explicit list of all of these 1.6 billion inequivalent charge assignments up to a maximum absolute value of 10 has been produced via a computer program described in section 3 and made publicly available [28]. Each entry in the list comprises 20 integers, the  $U(1)_X$  charge assignments of 20 chiral superfields of the model. Extensive checks of the list have been made using the analytic solution as well as those of internal consistency. With the aid of a computer, such a list is easily and quickly searched and filtered, looking for charge assignments with various desirable properties. For example, if fewer than three SM-singlets are required for the model, one can filter the list and find all solutions where one of the SM-singlet  $U(1)_X$  charges is zero. As far as anomalies go, having a zero charge for the superfield is equivalent to removing it from the model. We have shown some simple example filters, looking for different desirable properties of the charge assignments in section 6 as a tutorial in their implementation. We hope that the list will be of use for beyond-the-MSSM builders in terms of inspiration and phenomenology.

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