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Non-universal Z' from fluxed GUTs

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ABSTRACT

We make a first systematic study of non-universal TeV scale neutral gauge bosons Z' arising naturally from a class of F-theory inspired models broken via $SU(5)$ by flux. The phenomenological models we consider may originate from semi-local F-theory GUTs arising from a single E_8 point of local enhancement, assuming the minimal $\mathscr{Z}_2$ monodromy in order to allow for a renormalisable top quark Yukawa coupling. We classify such non-universal anomaly-free $U(1)'$ models requiring a minimal low energy spectrum and also allowing for a vector-like family. We discuss to what extent such models can account for the anomalous B-decay ratios R_K and R_{K^*}.

1. Introduction

In the Standard Model (SM) of strong and electroweak interactions the gauge couplings to lepton fields are flavour independent. This is a well known property of the SM gauge interactions which is usually called Lepton Flavour Universality (LFU). Among the various processes that can be used to test LFU, are the flavour changing decays involving quark fields such as $b \to s\ell^+\ell^-$ where $\ell^\pm$ stand for $e^\pm, \mu^\pm, \tau^\pm$ leptons. Recent experimental results [1] indicate violations of LFU which may be due to a new neutral boson Z' [2,3] which both changes quark flavour, namely $b \to s$, and couples predominantly to muons rather than electrons, with left-handed quark and lepton couplings preferred [4]. The flavourful Z' models are many and varied. One attractive possibility is to consider a vector-like family with non-universal Z' families which mixes with the three chiral families of quarks and leptons with universal Z' couplings, thereby introducing non-universality by the back door [5].[1]

It is commonly believed that the SM is not the final theory of fundamental interactions, but just an effective low energy limit of some (partially) grand unified theory (GUT), possibly embedded in a string scenario, leading to new physics phenomena and deviations from the SM predictions. For example, new gauge bosons Z' associated with additional abelian symmetries arise from such string embeddings [2]. Within the framework of F-theory [8] there are many such extra gauged abelian symmetries [9], which however are normally assumed to be broken at the high energy scale. Usually these gauge bosons are assumed to have universal couplings to quarks and leptons [10], however there is no reason in principle from the point of view of F-theory why this should be the case. Despite this there has been no systematic study of non-universal Z' gauge bosons at the low energy scale arising from F-theory to our knowledge.

Motivated by the above recent hints for non-universality, we make a first systematic study of anomaly-free non-universal Z' models arising from F-theory. From a phenomenological point of view, we refer to such models as fluxed GUTs. These models assume a set of rules extracted from F-theory and which can form the basis for a phenomenological analysis of models. In particular we shall use these rules to search for a set of anomaly free models[2] in which the low energy effective theory consists of the Minimal Supersymmetric Standard Model (MSSM) augmented by an extra gauged $U(1)'$ group broken at the TeV scale, together allowing also for a possible vector-like family at that scale. We shall discover that none of the models with a minimal low energy matter spectrum is capable of explaining the recent B-physics anomalies, while it is possible to account for these anomalies by allowing an extra vector-like family as in [5].

The layout of the paper is as follows. In section two we introduce the notion of fluxed GUTs which are based on essential features playing instrumental rôle in local F-theory constructions.

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[1] Extensive references to other non-universal Z' models are also included in [5].

[2] For recent anomaly free models with non-universal Z'- couplings see [6,7].

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In particular, we describe the rôle of fluxes on the breaking of the gauge symmetry, the splitting of GUT representations and the fermion chirality. In section three we implement these rules to construct $SU(5)$ GUTs with additional abelian symmetries allowing non-universal couplings to fermions. We focus on specific models where both $SU(5)$ GUT and abelian symmetries are embedded in E_8 and formulate the anomaly cancellation conditions which constrain the spectrum of the emerging effective models. In section four we present the solutions of these conditions and classify them according to the properties of the resulting light particle spectrum. We first present characteristic examples of $SU(5)$ with the minimal supersymmetric spectrum having non-universal couplings with the extra $U(1)$. We also show solutions containing vector-like fermions with non-universal couplings and discuss their implications to low energy physics. In section five we present our conclusions.

2. Fluxed GUTs

In our quest for a possible interpretation of such kind of experimental results with the incorporation of a new gauge boson Z', we will examine GUT models that can in principle be derived in the context of F-theory model building [9]. We stress that this particular string theory framework provides the appropriate tools for a natural inclusion of abelian gauge bosons with non-universal gauge couplings to fermions. Indeed, among other basic ingredients of F-theory constructions, a prominent rôle for the incarnation of the aforementioned scenario is played by the particular structure of the compact manifold and the (abelian) fluxes.

We recall first that the compactification space consists of an elliptically fibred Calabi–Yau manifold (CY_4 fourfold) of four complex dimensions whose geometric singularities display a group structure and, in F-theory, are associated with the gauge symmetries of the effective field theory model [8]. We will assume that the internal manifold has a structure characterised by the maximal exceptional symmetry E_8, where 7-branes wrap an appropriate divisor accommodating the $SU(5)$ gauge group of the effective field theory. In this context, the zero mode spectrum descends from the decomposition of the E_8 adjoint, therefore the following picture emerges in the effective model

$$E_8 \supset SU(5)_{GUT} \times SU(5)_\perp \tag{1}$$

$$248 \to (24,1) + (1,24) + \boxed{(10,5) + (\bar{5},10)} + (\bar{5},\overline{10}) + (5,\overline{10})\,, \tag{2}$$

where in (1), the group factor $SU(5)_{GUT}$, is identified with the grand unified symmetry of the effective model, and in (2) the representations in the box contain the $SU(5)_{GUT}$ fields. As is well known, ordinary matter is accommodated in 10 and $\bar{5}$ representations which, in the present construction, appear with non-trivial transformation properties under the (perpendicular) symmetry $SU(5)_\perp$. For the purposes of the effective field theory description, it is adequate to work in the Higgs bundle picture and express the transformation properties of the $SU(5)_{GUT}$ matter in terms of the Cartan generators of $SU(5)_\perp$ and the five weights $t_i, i = 1, 2, \ldots, 5$ satisfying the $SU(5)$ tracelessness condition

$$t_1 + t_2 + t_3 + t_4 + t_5 = 0\,\cdot \tag{3}$$

In this case, the superpotential can be maximally constrained by four $U(1)_\perp$'s ("perpendicular" to $SU(5)_{GUT}$) according to the breaking pattern

$$E_8 \supset SU(5)_{GUT} \times SU(5)_\perp \supset SU(5)_{GUT} \times U(1)_\perp^4\,\cdot \tag{4}$$

With respect to t_i notation, the non-trivial $SU(5)_{GUT}$ representations given in (2) are designated as $\bar{5}_{t_i+t_j}$, 10_{t_i} and the $SU(5)_{GUT}$ singlets emerging from $SU(5)_\perp$ adjoint decomposition are denoted with $1_{t_i-t_j} \equiv \theta_{ij}$. This way, the $U(1)_\perp^4$ invariance of each superpotential term is ensured as soon as the tracelessness condition $\sum_{i=1}^5 t_i = 0$ is satisfied.[3]

The second important available toolkit in F-theory model building includes the various fluxes which are turned on along the various abelian factors and can be used for the following tasks. Firstly, appropriate fluxes are introduced for the breaking symmetry mechanism and can be regarded as the surrogate of the vacuum expectation value (vev) of the Higgs field, particularly in the case of manifold geometries (such as del Pezzo surfaces) where the latter is absent from the massless spectrum. The $SU(5)_{GUT}$ symmetry in particular, breaks to the Standard Model with a hypercharge flux. Secondly, fluxes are used to generate the observed chirality of the massless spectrum. In order to describe their implications in the present construction we distinguish them into two classes. Initially, a flux is introduced along a $U(1)_\perp$ and its geometric restriction along a specific matter curve Σ_{n_j} is parametrised with an integer M_{n_j}. Then, the chiralities of the $SU(5)$ representations are given by

$$\#5_i - \#\bar{5}_i = M_{5_i} \tag{5}$$

$$\#10_j - \#\overline{10}_j = M_{10_j} \tag{6}$$

The hypercharge flux introduced to break $SU(5)_{GUT}$ is also responsible for the splitting of the $SU(5)$ representations. If some integers $N_{i,j}$ represent hyperfluxes piercing certain matter curves, the 10-plets and 5-plets split according to:

$$10_{t_j} = \begin{cases} n_{(3,2)_{\frac{1}{6}}} - n_{(\bar{3},2)_{-\frac{1}{6}}} &= M_{10_j} \\ n_{(\bar{3},1)_{-\frac{2}{3}}} - n_{(3,1)_{\frac{2}{3}}} &= M_{10_j} - N_j \\ n_{(1,1)_{+1}} - n_{(1,1)_{-1}} &= M_{10_j} + N_j \end{cases}\,, \tag{7}$$

$$5_{t_i} = \begin{cases} n_{(3,1)_{-\frac{1}{3}}} - n_{(\bar{3},1)_{+\frac{1}{3}}} &= M_{5_i} \\ n_{(1,2)_{+\frac{1}{2}}} - n_{(1,2)_{-\frac{1}{2}}} &= M_{5_i} + N_i\,\cdot \end{cases} \tag{8}$$

The splitting of $SU(5)_{GUT}$ representations has important and far reaching implications in model building. With a suitable choice of the flux integers on the Higgs curve(s), $\Sigma_{5_{h_u}}$, $\Sigma_{5_{h_d}}$, one can generate the doublet-triplet splitting of the Higgs 5-plets and suppress dimension-five operators contributing to proton decay. We note in passing that this is a novel feature of F-theory and there is no counterpart in ordinary GUTs. Another important observation is that, by virtue of the hypercharge flux, the SM representations of the same family may no longer be components of the same 5-plet. As a result, the $U(1)_\perp$ charges may be different not only between various families, but even within SM fields of the same fermion generation. It is exactly this property that is expected to allow the existence of non-universal flavour couplings of the corresponding Z' boson to lepton fields. Therefore, we conclude that the hyperflux splitting mechanism provides a natural way to implement the idea of a new neutral Z' gauge boson coupled differently either to the third- or to a vector-like family within the F-theory motivated framework.

[3] *In the geometric language, we say that the 10-plets and 5-plets are found in the intersections of the $SU(5)_{GUT}$ divisor with 7-branes associated with the $U(1)_\perp$'s, usually called matter curves and designated as $\Sigma_{10_{t_i}}$ and $\Sigma_{5_{t_i+t_j}}$.*

3. An $SU(5)$ model

Having discussed the essential features of F-GUT model building, we now proceed in a specific model. We have seen already that, in principle, there are at most four independent roots t_i. On the other hand, the four Cartan generators corresponding to $U(1)_\perp^4$ are expressed as:

$$\begin{aligned} H_1 &= \frac{1}{2}\text{diag}(1,-1,0,0,0), \\ H_2 &= \frac{1}{2\sqrt{3}}\text{diag}(1,1,-2,0,0), \\ H_3 &= \frac{1}{2\sqrt{6}}\text{diag}(1,1,1,-3,0), \\ H_4 &= \frac{1}{2\sqrt{10}}\text{diag}(1,1,1,1,-4). \end{aligned} \tag{9}$$

Throughout this paper we shall assume the minimal $\mathscr{Z}_2$ monodromy[4]

$$\mathscr{Z}_2 \text{ monodromy} : t_1 \leftrightarrow t_2 \tag{10}$$

As a consequence, this action reduces the number of the abelian factors by one, $U(1)_\perp^4 \to U(1)_\perp^3$, and allows a tree-level Yukawa coupling for the top-quark.

For the rest of our analysis, it is of interest to consider the following sequence of flux breaking, which may be associated with different scales

$$E_8 \supset E_6 \times U(1)_\perp^2 \tag{11}$$

$$\supset SO(10) \times U(1)_\psi \times U(1)_\perp^2 \tag{12}$$

$$\supset SU(5)_{GUT} \times U(1)_\chi \times U(1)_\psi \times U(1)_\perp^2. \tag{13}$$

Furthermore, as described above, we assume a $U(1)_Y$ flux which realises the $SU(5)_{GUT}$ breaking and at the same time triggers the doublet-triplet splitting and generates the chiral families of the Standard Model through (7) and (8).

It is convenient to choose a basis for the weight vectors such that the charge generators have the form where $Q_\perp$ is the charge of the $U(1)_\perp$ in the breaking pattern of Eq. (11) that remains after imposing the $t_1 \leftrightarrow t_2$ monodromy. In the conventional basis for the $SU(5)_\perp$ generators in Eq. (9) the unbroken generators are identified as follows:

$$H_2 = Q_\perp, \;\; H_3 = Q_\psi, \;\; H_4 = -Q_\chi. \tag{14}$$

This almost trivial equivalence shows that the $SU(5)_{GUT}$ states in Eq. (2) have well defined E_6 charges Q_χ and Q_ψ. For example $SU(5)$ singlets will in general carry Q_χ and Q_ψ charges which originate from E_6 and which may be unbroken. The equivalence will provide insights into both anomaly cancellation and the origin of R-parity for example, in terms of the underlying E_6 structure, in the explicit models discussed later.

[4] In this semi-local approach, several essential features of the effective field theory model are determined from the topological properties of the associated spectral surface, $\mathscr{C}$, described by a fifth degree polynomial $\mathscr{C} : \sum_{k=0}^{5} b_k t^{5-k} = 0$. The essential ingredients specifying these properties are the coefficients b_k with well defined topological properties. The roots of the polynomial determine the weight vectors $t_{1,\ldots,5}$. In general, however, depending on the properties of the specific compact manifold, the solutions $t_i = t_i(b_k)$'s imply that there is an action on the roots t_i of a non-trivial monodromy group which is a subgroup of the Weyl group $W(SU(5)_\perp) = S_5$. Such subgroups are the alternating groups $\mathscr{A}_n$, the dihedral groups $\mathscr{D}_n$ and cyclic groups $\mathscr{Z}_n$, $n \le 5$ and the Klein four-group $\mathscr{Z}_2 \times \mathscr{Z}_2$.

In order to work consistently, we use the normalised $SU(5)_\perp$ generators

$$Q_\perp = \frac{1}{2\sqrt{3}}\text{diag}(1,1,-2,0,0) \tag{15}$$

$$Q_\psi = \frac{1}{2\sqrt{6}}\text{diag}(1,1,1,-3,0) \tag{16}$$

$$Q_\chi = \frac{1}{2\sqrt{10}}\text{diag}(1,1,1,1,-4) \tag{17}$$

and the low energy $U(1)'$ is generated by

$$Q' = c_1 Q_\perp + c_2 Q_\psi + c_3 Q_\chi \tag{18}$$

which, in order to retain $SU(5)_\perp$ normalisation, are subject to the condition

$$c_1^2 + c_2^2 + c_3^2 = 1 \tag{19}$$

In order to proceed, we need to determine the spectrum and assign the corresponding Q' charges for all fields. We first remark that, in order to obtain three chiral families, the set of integers M_{n_j} associated with the $U(1)_\perp$ flux in formulae (6), (5) must obey

$$\sum_i M_{5_i} = -\sum_j M_{10_j} = -3 \tag{20}$$

This automatically implies also that the total flux vanishes, $\sum M_{5_i} + \sum M_{10_j} = 0$. Furthermore, the $SU(5)$ representations are localised on matter curves possessing certain topological characteristics which determine the various properties of the effective SM theory. Important rôle, in particular, is played by the homologies of these curves which are determined in accordance with the $\mathscr{Z}_2$ factorisation of the spectral cover polynomial. It turns out that the homologies of the whole set of matter curves can be expressed in terms of only three arbitrary parameters $\chi_{6,7,8}$ [11]. If $\mathscr{F}_Y$ represents the hypercharge flux, its 'dot' product with $\chi_{6,7,8}$ determines three hyperflux integers

$$N_7 = \mathscr{F}_Y \cdot \chi_7, \;\; N_8 = \mathscr{F}_Y \cdot \chi_8, \;\; N_9 = \mathscr{F}_Y \cdot \chi_9\,, \tag{21}$$

and we define $\tilde{N} = N_7 + N_8 + N_9$. In conclusion, multiplicities of the SM states are determined by $N_i, i = 7, 8, 9$ associated with hyperflux restrictions and the $U(1)_\perp$ flux integers M_{n_j}. More specifically, for a given choice of these numbers we can determine the exact number of the SM states residing on each matter curve using the formulae (7), (8) and then apply the definition (18) to find their corresponding charges. We refer to previous work for the details [12–14], while here, we only present selected properties of the spectrum in Table 1.

3.1. Anomalies

The presence of an additional $U(1)'$ symmetry in the effective model is associated with contributions to cubic and mixed anomalies with the SM gauge group G_{SM}. On the other hand, we have seen that the multiplicities in the SM spectrum are determined by the flux parameters M_{n_i}, N_j and therefore these numbers are directly involved in the anomaly constraints. These involve the $U(1)'$ anomalies with the G_{SM} currents as well as quartic and trace anomaly conditions. Thus, the flux integers M_{n_i}, N_j are subject to certain restrictions. All anomalies were systematically computed using the SUSYNO package [15], those involving G_{SM} currents in particular are as follows:

Table 1
The first column shows the matter curve accommodating the $SU(5)$ representations, and the $U(1)_\perp$ weights ($\pm$ refer to $10/\overline{10}$ and $\bar{5}/5$ respectively). The second column displays their Q' charges and columns 3 and 4 the $U(1)_{Y,\perp}$ flux integers N_i, M_{n_j}. The last column shows the SM multiplicities in terms of the N_i, M_{n_j} numbers. Singlet fields θ_{ij} are not presented in this Table (see however text of section 2.)

Curve name	Q'	N_Y	M	SM content
$\Sigma_{5_{Hu}(-2t_1)}$	$-\frac{c_1}{\sqrt{3}}-\frac{c_2}{\sqrt{6}}-\frac{c_3}{\sqrt{10}}$	$\tilde{N}$	$M_{5_{Hu}}$	$M_{5_{Hu}}\overline{d^c}+(M_{5_{Hu}}+\tilde{N})\overline{L}$
$\Sigma_{5_{1,\pm(t_1+t_3)}}$	$\frac{5\sqrt{3}c_1-5\sqrt{6}c_2-3\sqrt{10}c_3}{30}$	$-\tilde{N}$	M_{5_1}	$M_{5_1}\overline{d^c}+(M_{5_1}-\tilde{N})\overline{L}$
$\Sigma_{5_{2,\pm(t_1+t_4)}}$	$-\frac{c_1}{2\sqrt{3}}+\frac{c_2}{\sqrt{6}}-\frac{c_3}{\sqrt{10}}$	$-\tilde{N}$	M_{5_2}	$M_{5_2}\overline{d^c}+(M_{5_2}-\tilde{N})\overline{L}$
$\Sigma_{5_{3,\pm(t_1+t_5)}}$	$\frac{-10\sqrt{3}c_1-5\sqrt{6}c_2+9\sqrt{10}c_3}{60}$	$-\tilde{N}$	M_{5_3}	$M_{5_3}\overline{d^c}+(M_{5_3}-\tilde{N})\overline{L}$
$\Sigma_{5_{4,\pm(t_3+t_4)}}$	$\frac{c_1}{\sqrt{3}}+\frac{c_2}{\sqrt{6}}-\frac{c_3}{\sqrt{10}}$	N_7+N_8	M_{5_4}	$M_{5_4}\overline{d^c}+(M_{5_4}+N_7+N_8)\overline{L}$
$\Sigma_{5_{5,\pm(t_3+t_5)}}$	$\frac{20\sqrt{3}c_1-5\sqrt{6}c_2+9\sqrt{10}c_3}{60}$	N_7+N_9	M_{5_5}	$M_{5_5}\overline{d^c}+(M_{5_5}+N_7+N_9)\overline{L}$
$\Sigma_{5_{6,\pm(t_4+t_5)}}$	$\frac{5\sqrt{6}c_2+3\sqrt{10}c_3}{20}$	N_8+N_9	M_{5_6}	$M_{5_6}\overline{d^c}+(M_{5_6}+N_8+N_9)\overline{L}$
$\Sigma_{10_{t,\pm t_1}}$	$\frac{10\sqrt{3}c_1+5\sqrt{6}c_2+3\sqrt{10}c_3}{60}$	$-\tilde{N}$	M_{10_t}	$M_{10_t}Q+(M_{10_t}+\tilde{N})u^c+(M_{10_t}-\tilde{N})e^c$
$\Sigma_{10_{2,\pm t_3}}$	$\frac{-20\sqrt{3}c_1+5\sqrt{6}c_2+3\sqrt{10}c_3}{60}$	N_7	M_{10_2}	$M_{10_2}Q+(M_{10_2}-N_7)u^c+(M_{10_2}+N_7)e^c$
$\Sigma_{10_{3,\pm t_4}}$	$\frac{\sqrt{10}c_3-5\sqrt{6}c_2}{20}$	N_8	M_{10_3}	$M_{10_3}Q+(M_{10_3}-N_8)u^c+(M_{10_3}+N_8)e^c$
$\Sigma_{10_{4,\pm t_5}}$	$-\sqrt{\frac{2}{5}}c_3$	N_9	M_{10_4}	$M_{10_4}Q+(M_{10_4}-N_9)u^c+(M_{10_4}+N_9)e^c$

$$
\begin{aligned}
\mathscr{A}_{G_{SM}^2\times U(1)'} =& \left(-\frac{c_1}{2\sqrt{3}}-\frac{c_2}{2\sqrt{6}}-\frac{c_3}{2\sqrt{10}}\right)M_{5_{Hu}} \\
&+\left(\frac{\sqrt{3}c_1}{4}+\frac{1}{4}\sqrt{\frac{3}{2}}c_2+\frac{3c_3}{4\sqrt{10}}\right)M_{10_t} \\
&+\left(\frac{c_1}{4\sqrt{3}}-\frac{c_2}{2\sqrt{6}}-\frac{c_3}{2\sqrt{10}}\right)M_{5_1} \\
&+\left(-\frac{c_1}{4\sqrt{3}}+\frac{c_2}{2\sqrt{6}}-\frac{c_3}{2\sqrt{10}}\right)M_{5_2} \\
&+\left(-\frac{c_1}{4\sqrt{3}}-\frac{c_2}{4\sqrt{6}}+\frac{3c_3}{4\sqrt{10}}\right)M_{5_3} \\
&+\left(\frac{c_1}{2\sqrt{3}}+\frac{c_2}{2\sqrt{6}}-\frac{c_3}{2\sqrt{10}}\right)M_{5_4} \\
&+\left(\frac{c_1}{2\sqrt{3}}-\frac{c_2}{4\sqrt{6}}+\frac{3c_3}{4\sqrt{10}}\right)M_{5_5} \\
&+\left(\frac{1}{4}\sqrt{\frac{3}{2}}c_2+\frac{3c_3}{4\sqrt{10}}\right)M_{5_6} \\
&+\left(-\frac{1}{2}\sqrt{3}c_1+\frac{1}{4}\sqrt{\frac{3}{2}}c_2+\frac{3c_3}{4\sqrt{10}}\right)M_{10_2} \qquad (22)\\
&+\left(\frac{3c_3}{4\sqrt{10}}-\frac{3}{4}\sqrt{\frac{3}{2}}c_2\right)M_{10_3}-\frac{3c_3}{\sqrt{10}}M_{10_4} \\
&+\frac{1}{4}\sqrt{3}c_1N_7+\left(\frac{c_1}{4\sqrt{3}}+\frac{c_2}{\sqrt{6}}\right)N_8 \\
&+\left(\frac{c_1}{4\sqrt{3}}+\frac{c_2}{4\sqrt{6}}+\frac{1}{4}\sqrt{\frac{5}{2}}c_3\right)N_9 \\
\mathscr{A}_{U(1)_Y\times U(1)'^2} =& \frac{3}{2}c_1^2N_7+\left(\frac{c_1^2}{6}+\frac{2}{3}\sqrt{2}c_2c_1+\frac{4c_2^2}{3}\right)N_8 \\
&+\left(\frac{c_1^2}{6}+\frac{c_2c_1}{3\sqrt{2}}+\sqrt{\frac{5}{6}}c_3c_1+\frac{c_2^2}{12}\right. \\
&\left.+\frac{5c_3^2}{4}+\frac{1}{2}\sqrt{\frac{5}{3}}c_2c_3\right)N_9
\end{aligned}
$$

$$
\mathscr{A}_{G_{SM}^3} = \frac{5}{36}\left(M_{5_{Hu}}+M_{10_t}+M_{5_1}+M_{5_2}+M_{5_3}+M_{5_4}\right. \\
\left.+M_{5_5}+M_{5_6}+M_{10_2}+M_{10_3}+M_{10_4}\right)\,.
$$

In addition, we have to consider the trace anomaly, which is related to the D-flatness constraint and the associated Green–Schwarz cancellation mechanism, and the cubic anomaly that involve contributions from all singlet fields θ_{ij}. To avoid clutter in this short note, these are not included in the above equations, however, they can be computed in a straightforward manner and are taken into account in our subsequent computations. We should point out however, that in order to have a sizable effect in anomalous B-decays, the solution of the anomaly cancellation equations must be compatible with a TeV scale mass of the corresponding Z' boson.

Hence, our strategy is as follows: we start by working out the solutions with the anomalies involving SM currents, $G_{SM}^2\times U(1)'$ and $U(1)_Y\times U(1)'^2$ and solve for all c_i coefficients subject to the normalisation condition $c_1^2+c_2^2+c_3^2=1$. This will provide a solution $c_i=c_i(M_{n_j},N_k)$. First, we notice that the MSSM chirality trivially solves for $\mathscr{A}_{G_{SM}^3}=0$. Then, we check whether there is a solution for $U(1)'^3$, Tr Q' anomalies using the c_i derived before, keeping only the solutions for which the charges can be written in a single radical form.[5] These last two constraints will be solved for multiplicities of singlet fields θ_{ij}. And in our search for anomaly free solutions, we will seek models with the MSSM spectrum, as well as including those with complete vector-like families.

4. Minimal MSSM type spectrum but with a non-universal Z'

We start with models that have the same spectrum of the MSSM with the gauge group extended by a single $U(1)'$ factor. In order to have the MSSM chirality, the flux data must respect the geometric constraints (20) which are used to solve for M_{5_6} and M_{10_4}. Notice that this solves G_{SM}^3 anomaly trivially. Further, the MSSM spectrum was imposed

[5] This means that the charges for each state can all factor out $\sqrt{a}$, where a is a rational number, such that it can be absorbed into the definition of the Z' coupling constant and hence making all charges appear rational numbers themselves.

Table 2
Flux data and c_i coefficients for explicit models 1 to 4 of the classes 1 to 4, respectively.

Model	c_1	c_2	c_3	$M_{5_{H_u}}$	M_{5_1}	M_{5_2}	M_{5_3}	M_{5_4}	M_{5_5}	M_{5_6}	M_{10_t}	M_{10_2}	M_{10_3}	M_{10_4}	N_7	N_8	N_9
1	$-\frac{\sqrt{5}}{3}$	$\frac{1}{6}\sqrt{\frac{5}{2}}$	$-\frac{1}{2}\sqrt{\frac{3}{2}}$	0	0	0	−1	−1	0	−1	1	1	1	0	0	1	0
2	$-\frac{\sqrt{5}}{3}$	$-\frac{1}{3}\sqrt{\frac{5}{2}}$	$\frac{1}{\sqrt{6}}$	0	0	−1	0	0	−1	−1	1	1	0	1	0	0	1
3	$-\sqrt{\frac{5}{6}}$	$-\frac{1}{8}\sqrt{\frac{5}{3}}$	$\frac{3}{8}$	0	0	0	−1	0	−1	−1	1	1	0	1	0	0	1
4	$-\sqrt{\frac{5}{6}}$	$\frac{1}{4}\sqrt{\frac{5}{3}}$	$-\frac{1}{4}$	0	0	0	0	−1	−1	−1	1	0	1	1	0	1	0

Table 3
Explicit Models 1 to 4 using the data in Table 2.

Curve	Weights	Model 1		Model 2		Model 3		Model 4	
		$Q'\sqrt{15}$	SM content	$Q'\sqrt{15}$	SM content	$Q'\sqrt{10}$	SM content	$Q'\sqrt{10}$	SM content
5_{H_u}	$-2t_1$	2	H_u	$\frac{1}{2}$	H_u	$\frac{3}{2}$	H_u	$\frac{3}{2}$	H_u
5_1	t_1+t_3	$\frac{1}{2}$	L	2	d^c+2L	1	L	1	L
5_2	t_1+t_4	−2	H_d	$-\frac{1}{2}$	H_d	$-\frac{1}{4}$	L	$-\frac{3}{2}$	H_d
5_3	t_1+t_5	$\frac{1}{2}$	d^c+2L	$-\frac{7}{4}$	L	$-\frac{3}{2}$	d^c+2L	$-\frac{1}{4}$	L
5_4	t_3+t_4	$\frac{1}{2}$	d^c	2	d^c	$-\frac{9}{4}$	–	1	d^c
5_5	t_3+t_5	−3	–	$-\frac{3}{4}$	–	1	d^c	$\frac{9}{4}$	d^c+L
5_6	t_4+t_5	$\frac{1}{2}$	d^c	$-\frac{7}{4}$	d^c	$-\frac{1}{4}$	d^c	$-\frac{1}{4}$	d^c
10_t	t_1	−1	$Q+2u^c+e^c$	$-\frac{1}{4}$	$Q+2u^c$	$-\frac{3}{4}$	$Q+2u^c$	$-\frac{3}{4}$	$Q+2u^c$
10_2	t_3	$\frac{3}{2}$	$Q+u^c+e^c$	$\frac{9}{4}$	–	$\frac{7}{4}$	$Q+u^c+e^c$	$\frac{7}{4}$	–
10_3	t_4	−1	$Q+2e^c$	$-\frac{1}{4}$	$2Q+u^c+3e^c$	$\frac{1}{2}$	–	$-\frac{3}{4}$	$Q+2e^c$
10_4	t_5	$\frac{3}{2}$	–	$-\frac{3}{2}$	–	$-\frac{3}{4}$	$Q+2e^c$	$\frac{1}{2}$	$Q+u^c+e^c$

$$\#L + \#\overline{L} = 5 \tag{23}$$

$$\#d^c = \#Q = \#u^c = \#e^c = 3 \tag{24}$$

where L represents either the lepton doublets or Higgs doublets and we demanded the existence of a solution for the doublet-triplet splitting problem

$$|N_7| + |N_8| + |N_9| \neq 0\,. \tag{25}$$

Since every state inside the same matter curve shares the same charge Q', a non-universal model will require the families to be supported in different matter curves. In order to still scan models that allow for that to happen either only for the Σ_5 or Σ_{10} matter curves, we allow the flux parameters to vary between

$$N_i,\ M_{n_j} \in [-3, 3]\,. \tag{26}$$

In order to guarantee a renormalisable top quark Yukawa coupling, we identify $H_u \in 5_{-2t_1}$ and $t \in 10_{t_1}$ as these states are involved in the renormalisable Yukawa interaction

$$10_{t_1} 10_{t_1} 5_{-2t_1} \equiv 10_t 10_t 5_{H_u}\,. \tag{27}$$

This accounts for

$$M_{5_{H_u}} + \tilde{N} \geq 1,\ M_{10_t} \geq 1,\ M_{10_t} + \tilde{N} \geq 1\,, \tag{28}$$

and we further assume $M_{10_t} = 1$ to fix the left-handed top t_L.

We also restrict our searches to solutions that have a single and isolated H_u in the 5_{H_u} curve, as we are preventing exotics in the spectrum. This accounts for

$$M_{5_{H_u}} = 0\,,\ \tilde{N} = 1\,, \tag{29}$$

where the last equation can be used to solve for one of the N_i flux

$$N_7 = 1 - N_8 - N_9\,. \tag{30}$$

After scanning the entire parameter space, we find that there are only 48 solutions respecting all the conditions above, which fall in just four different classes of models. We notice that these are the only classes of models for an $SU(5)$ GUT with a $\mathscr{Z}_2$ monodromy with a consistent Z' interaction. The models inside the same class differ as to how the states are distributed amongst different curves and the perpendicular charges they carry. The four distinct classes are, in the $G_{SM} \times U(1)'$ basis

Class 1:

$$(H_u)_{2\frac{1}{\sqrt{15}}} + (H_d)_{-2\frac{1}{\sqrt{15}}} + 3 \times \overline{5}_{\frac{1}{2}\frac{1}{\sqrt{15}}} + 2 \times 10_{-\frac{1}{\sqrt{15}}} + 10_{\frac{3}{2}\frac{1}{\sqrt{15}}} \tag{31}$$

Class 2:

$$(H_u)_{\frac{1}{2}\frac{1}{\sqrt{15}}} + (H_d)_{-\frac{1}{2}\frac{1}{\sqrt{15}}} + 2 \times \overline{5}_{2\frac{1}{\sqrt{15}}} + \overline{5}_{-\frac{7}{4}\frac{1}{\sqrt{15}}} + 3 \times 10_{-\frac{1}{4}\frac{1}{\sqrt{15}}} \tag{32}$$

Class 3:

$$(H_u)_{\frac{3}{2}\frac{1}{\sqrt{10}}} + \overline{5}_{-\frac{1}{4}\frac{1}{\sqrt{10}}} + \overline{5}_{\frac{1}{\sqrt{10}}} + (2L + d^c)_{-\frac{3}{2}\frac{1}{\sqrt{10}}} + 2 \times 10_{-\frac{3}{4}} + 10_{\frac{7}{4}\frac{1}{\sqrt{10}}} \tag{33}$$

Class 4:

$$(H_u)_{\frac{3}{2}\frac{1}{\sqrt{10}}} + (H_d)_{-\frac{3}{2}\frac{1}{\sqrt{10}}} + \overline{5}_{-\frac{1}{4}\frac{1}{\sqrt{10}}} + \overline{5}_{\frac{1}{\sqrt{10}}} + \overline{5}_{\frac{9}{4}\frac{1}{\sqrt{10}}} + 2 \times 10_{-\frac{3}{4}\frac{1}{\sqrt{10}}} + 10_{\frac{1}{2}\frac{1}{\sqrt{10}}} \tag{34}$$

in addition to $Q' \to -Q'$.

In Table 2 we present an explicit element of each class, and the distribution of the states amongst different matter curves is presented in Table 3.

As the left-handed top-quark is always identified as being in 10_t, we notice that in models 1, 3 and 4 the first and second family left-handed quarks have different $U(1)'$ charges, which can be

Table 4
Flux data and c_i coefficients for explicit models 5 to 7 of the classes 5 to 7, respectively.

Model	c_1	c_2	c_3	$M_{5_{H_u}}$	M_{5_1}	M_{5_2}	M_{5_3}	M_{5_4}	M_{5_5}	M_{5_6}	M_{10_t}	M_{10_2}	M_{10_3}	M_{10_4}	N_7	N_8	N_9
5	0	$\frac{1}{2}\sqrt{\frac{15}{34}}$	$\frac{11}{2\sqrt{34}}$	0	0	0	0	−1	−3	1	1	2	−1	1	1	0	0
6	$-\frac{\sqrt{\frac{5}{6}}}{2}$	$-\frac{5\sqrt{\frac{5}{3}}}{8}$	$\frac{3}{8}$	0	1	−1	0	−1	−2	0	2	−1	1	1	0	0	1
7	$\frac{\sqrt{3}}{2}$	$-\frac{1}{4}\sqrt{\frac{3}{2}}$	$\frac{1}{4}\sqrt{\frac{5}{2}}$	0	0	0	1	−3	−1	0	2	1	1	−1	0	1	0

checked for all elements of the respective classes. As a result, the Z' mass for these models is constrained to be $m_{Z'} \gtrsim 10^5$ TeV or otherwise Z' would contribute too strongly at tree-level to $K^0 - \overline{K}^0$ oscillations. In addition, model 2 leaves a certain ambiguity about what L state should be identified as H_d.

5. Models with a vector-like family with non-universal couplings to a Z'

Another interesting possibility is that there are new vector-like states in the spectrum such that the left-handed states have a different charge under $U(1)'$. As shown in [5], this can lead to non-universality in the regular matter through mixing in the new neutral current coupling to Z'.

In order to preserve unification, we need these vector-like states to form complete $SU(5)$ representations. In particular, we look for models that allow for a complete vector-like family. In addition, as we wish to preserve the top-quark assignment, the following simplifying assumptions were taken

$$M_{5_{H_u}} = 0\,,\ \tilde{N} = 1\,, \tag{35}$$

in order to obtain models where the H_u is isolated in its own matter curve, as can be seen from (8). With this simplification, computational time is significantly reduced and we are able to probe the entire region of parameter space of interest.[6]

There are 4067 models with one extra vector-like family, but only 397 classes when considering only $G_{SM} \times U(1)'$ charge assignments. Here we present three of these classes of models

Class 5:

$$(H_u)_{4\frac{1}{\sqrt{85}}} + (H_d)_{-4\frac{1}{\sqrt{85}}} + 3\times\overline{5}_{\frac{7}{2}\frac{1}{\sqrt{85}}} + \overline{5}_{\frac{3}{2}\frac{1}{\sqrt{85}}} + 5_{6\frac{1}{\sqrt{85}}} + 3\times 10_{2\frac{1}{\sqrt{85}}} + 10_{-\frac{11}{2}\frac{1}{\sqrt{85}}} + \overline{10}_{\frac{1}{2}\frac{1}{\sqrt{85}}} \tag{36}$$

Class 6:

$$(H_u)_{\frac{3}{2}\frac{1}{\sqrt{10}}} + (H_d)_{-\frac{3}{2}\frac{1}{\sqrt{10}}} + \overline{5}_{\frac{1}{\sqrt{10}}} + \overline{5}_{-\frac{1}{4}\frac{1}{\sqrt{10}}} + \overline{5}_{-\frac{9}{4}\frac{1}{\sqrt{10}}} + L_{\frac{1}{\sqrt{10}}} + d^c_{-\frac{1}{4}\frac{1}{\sqrt{10}}} + 5_{-\frac{1}{\sqrt{10}}} + 3\times 10_{-\frac{3}{4}\frac{1}{\sqrt{10}}} + 10_{\frac{7}{4}\frac{1}{\sqrt{10}}} + \overline{10}_{-\frac{1}{2}\frac{1}{\sqrt{10}}} \tag{37}$$

Class 7:

$$(H_u)_{-\frac{1}{2}} + (H_d)_{\frac{1}{2}} + 3\times\overline{5}_{-\frac{1}{4}} + \overline{5}_{\frac{3}{4}} + 5_0 + 3\times 10_{\frac{1}{4}} + 10_{-\frac{1}{2}} + \overline{10}_{\frac{1}{4}}\,, \tag{38}$$

and we notice the first and last ones provide models where regular matter is universal under the new interaction, while the non-universality is strictly induced by the extra vector-like states. As $K^0 - \bar{K}^0$ mixing imposes strong constraints on non-universality in the first two families, these two cases, in principle, could be consistent with a light (few TeV) Z' boson mass and could be sufficient to explain the B-decay anomaly. On the contrary, models from the Class 6 will probably require a heavy Z', as the first and second family right-handed down-quarks have different charges.

In Table 4 we present the flux data and solutions for the coefficients c_i for three explicit realisations of models respecting these conditions. The spectrum distribution amongst curves of these three models can be seen in Table 5.

5.1. A brief discussion of model 7

Next, we briefly discuss the main phenomenological properties of the emergent low energy effective theory with the spectrum given by model 7 of Table 5. Because of the constraints imposed by the $K^0 - \bar{K}^0$ oscillations, as discussed above, the left-handed quark doublets of the three chiral fermion families should be accommodated in $10_t, 10_3$ which have equal charges, $Q' = \frac{1}{4}$, with respect to $U(1)'$. Similarly, the three down quarks and lepton doublets are distributed in the fiveplets $\bar{5}_1, \bar{5}_4$ and carry the same charge $Q' = -\frac{1}{4}$, as shown in Table 5. The third family fermions receive masses from the tree-level couplings

$$\lambda_t 10_t 10_t 5_{H_u} + \lambda_b 10_t \bar{5}_4 \bar{5}_{H_d}\,.$$

We observe that there are more than one families on the same matter curves and the present configuration implies a single intersection point for the corresponding Yukawa couplings. In addition, by turning on non-zero fluxes we obtain hierarchical textures for the quark and charged lepton fermion mass matrices as described in [16].

The vector-like family is accommodated as follows: Left-handed doublet and right-handed up quarks together with right-handed electrons, are found in the $\overline{10}_4 + 10_2$ pair. Furthermore, the lepton doublet in 5_3 and the colour triplet in 5_6 form a complete fiveplet, and, together with $\bar{5}_5$ constitute a pair in $\bar{5}_5 + (5_6 + 5_3)$ which completes the vector-like family. These extra states form the following trilinear couplings involving singlet fields

$$\overline{10}_4 10_2 \theta_{53} + \bar{5}_5 5_3 \theta_{13} + \bar{5}_5 5_6 \theta_{43} \tag{39}$$

It is possible that these singlets acquire vevs just below the GUT scale and give masses of this order to vector-like fields which eventually decouple from the spectrum. Indeed, this would happen, for example, in the case of an anomalous $U(1)'$ symmetry where the vevs of the singlets are fixed to be close to the GUT scale by virtue of the Green–Schwarz (GS) mechanism. In this case, the $U(1)'$ symmetry breaking scale is also associated with the GUT scale and the Z' boson becomes superheavy so that any New Physics effects from its non-universal couplings are highly suppressed. Yet, it is possible that the GS mechanism is realised with large vevs not involving all or some of the singlets in (39). This allows the possibility to freely choose low scale vevs for these latter singlet fields and obtain TeV masses to the vector-like fields, giving rise to interesting phenomenological implications. However,

[6] Note that in order for the extra vector-like $(\mathbf{10} + \overline{\mathbf{5}}) + (\overline{\mathbf{10}} + \mathbf{5})$ to induce non-universality from mixing, it has to have different charges than regular matter, therefore we limit the values of flux data to be within $[-3, 3]$.

Table 5
Explicit models 5 to 7 using the data in Table 4.

Curve	Weights	Model 5		Model 6		Model 7	
		$Q'\sqrt{85}$	SM content	$Q'\sqrt{10}$	SM content	Q'	SM content
5_{H_u}	$-2t_1$	-4	H_u	$\frac{3}{2}$	H_u	$-\frac{1}{2}$	H_u
	$2t_1$	–	–	–	–	–	–
5_1	$-t_1-t_3$	–	–	$\frac{1}{4}$	$\overline{d^c}$	–	–
	t_1+t_3	4	H_d	–	–	$-\frac{1}{4}$	L
5_2	$-t_1-t_4$	–	–	–	–	–	–
	t_1+t_4	$\frac{3}{2}$	L	1	d^c+2L	$\frac{1}{2}$	H_d
5_3	$-t_1-t_5$	–	–	–	–	0	$\overline{d^c}$
	t_1+t_5	$-\frac{7}{2}$	L	$-\frac{3}{2}$	H_d	–	–
5_4	$-t_3-t_4$	–	–	–	–	–	–
	t_3+t_4	$\frac{3}{2}$	d^c	$\frac{9}{4}$	d^c+L	$-\frac{1}{4}$	$3d^c+2L$
5_5	$-t_3-t_5$	–	–	–	–	–	–
	t_3+t_5	$-\frac{7}{2}$	$3d^c+2L$	$-\frac{1}{4}$	$2d^c+L$	$-\frac{3}{4}$	d^c+L
5_6	$-t_4-t_5$	6	$\overline{d^c}+\overline{L}$	-1	$\overline{L}$	0	$\overline{L}$
	t_4+t_5	–	–	–	–	–	–
10_t	t_1	2	$Q+2u^c$	$-\frac{3}{4}$	$2Q+3u^c+e^c$	$\frac{1}{4}$	$2Q+3u^c+e^c$
	$-t_1$	–	–	–	–	–	–
10_2	t_3	2	$2Q+u^c+3e^c$	–	–	$-\frac{1}{2}$	$Q+u^c+e^c$
	$-t_3$	–	–	$-\frac{1}{2}$	$\overline{Q}+\overline{u^c}+\overline{e^c}$	–	–
10_3	t_4	–	–	$\frac{7}{4}$	$Q+u^c+e^c$	$\frac{1}{4}$	$Q+2e^c$
	$-t_4$	$\frac{1}{2}$	$\overline{Q}+\overline{u^c}+\overline{e^c}$	–	–	–	–
10_4	t_5	$\frac{11}{2}$	$Q+u^c+e^c$	$-\frac{3}{4}$	$Q+2e^c$	–	–
	$-t_5$	–	–	–	–	$\frac{1}{4}$	$\overline{Q}+\overline{u^c}+\overline{e^c}$

if these vector-like fields are to be relevant for the anomalies in the B-decays, the singlets involved in these terms should acquire non-zero vevs $\langle\theta_{ij}\rangle$ of the order few TeV. Besides, there are also couplings of the vector-like family with the three generations such as $\overline{10}_4 10_3\theta_{54} + \overline{10}_4 10_t\theta_{51}$ etc. Assuming this then, after mixing with the vector-like family, this model can lead to non-universal couplings of the light mixed states which can be selected in such a way as to account for the recent hint for anomalous B-decays, along the lines of [5]. Therefore, we conclude that a scenario such as [5] can, in principle, be realised in this fluxed GUT framework.

The above discussion shows that the masses of the vector-like fermions arise from the VEVs of the singlet fields, and are controlled by these VEVs. The main point we should emphasise is that the masses of the vector-like fermions are protected by the $U(1)'$ and so the spectrum of these masses is controlled by the different singlet VEVs which break the $U(1)'$. In our model we require that the $U(1)'$ is broken at the few TeV scale, hence the vector-like fermion masses must also be around this scale. Clearly the singlets θ_{ij}, with weights $t_i - t_j$, play a crucial role in our model, and are discussed in more detail in Appendix A. In section 3.1, in order to satisfy the anomaly cancellation conditions some of the singlet fields θ_{ij} $(i, j = 1, \ldots, 5)$ charged under the four $U(1)$ factors acquire non zero VEVs. We can choose their VEVs so that the $U(1)'$ combination defined above remains unbroken and consequently the mass of the corresponding gauge boson Z' does not receive a GUT scale mass. Singlets charged under this particular $U(1)'$ combination may acquire appropriate VEVs implying masses for the Z' boson and the vector-like fields of the order of a few TeV.

It is worth to discuss in a bit more detail whether the scheme we have presented can really explain the anomalous B-decay ratios R_K and R_{K^*} consistently with all the other phenomenological constraints. The first requirement is that the Z' couples to left-handed muons more strongly than to electrons. In model 7, the left-handed muons in the lepton doublet L_2 are contained in the $\overline{5}$ representation of the second family which arises from the 5_4 curve while the left-handed electrons in L_1 are contained in the $\overline{5}$ representation of the first family which arises from the 5_1 curve (see Table 5). Both these curves have the same $U(1)'$ charges. However there is an additional vector-like pair of lepton doublets arising from the curves 5_5 and 5_6 which have non-universal $U(1)'$ charges and can mixing differently with the lepton doublets L_1 and L_2, inducing non-universal couplings in the physical left-handed electrons and muons. For example the vector lepton doublet $\bar{L}$ with charge $-3/4$ in 5_5 can mix predominantly with the muon doublet in 5_4 with charge $-1/4$, effectively enhancing the coupling to Z' of the mixed state. This mechanism is discussed in more detail in [5]. As discussed above, such mixing is controlled by the different singlet VEVs θ_{ij}, which we assume to be arbitrary here.

There are important constraints on lepton flavour violating (LFV) processes such as $\mu \to eee$ and $\mu - e$ conversion as recently discussed for example in [17]. However, as discussed there, violations of lepton universality does not always lead to lepton flavour violation: it depends on a mixing angle in the charged lepton Yukawa matrix, and hence is model dependent. Nevertheless, in general it would be expected that such LFV processes should accompany lepton universality violation at some level.

Indeed there are other flavour violating processes such as $B_s - \bar{B}_s$ mixing which can also rule out models, due to the Z' coupling to bs, as also discussed for example in [17] (and references therein). The main requirement here is that the Z' coupling to bs is sufficiently small compared to the Z' coupling to $\mu\mu$ in order to satisfy such constraints. In some models, for example the $SO(10)$ model discussed in [17], it is not possible to explain R_K and R_{K^*} while remaining consistent with $B_s - \bar{B}_s$ mixing. In the $SU(5)$ model 7 here, we see that the three quark doublets Q all have $U(1)'$ charge of $1/4$, while the fourth family quark doublet Q_4 which mixes with Q_3 has charge $-1/2$. However the induced Z' coupling to bs is suppressed by a quark mixing of order $(V_{dL})_{32}$ which is strictly speaking independent of the CKM element V_{ts} so in principle it could be small enough. It then becomes a quantitative question whether the $B_s - \bar{B}_s$ mixing constraint can be satisfied while explaining R_K and R_{K^*} and satisfying the LHC constraint on the Z' mass.

To be a little more quantitative, the relevant terms in the Lagrangian can be written as,

$$\mathscr{L} \supset Z'_\mu \left(g_{bb}\bar{b}_L\gamma^\mu b_L + g_{\mu\mu}\bar{\mu}_L\gamma^\mu \mu_L \right). \tag{40}$$

One possible explanation of the R_K and R_{K^*} measurements in LHCb is that the low-energy Lagrangian below the weak scale contains an additional contribution to the effective 4-fermion operator with left-handed muon, b-quark, and s-quark fields:

$$\Delta\mathscr{L}_{\text{eff}} \supset G_{bs\mu}(\bar{b}_L\gamma^\mu s_L)(\bar{\mu}_L\gamma_\mu \mu_L) + \text{h.c.},$$
$$G_{bs\mu} \approx \frac{1}{(31.5\ \text{TeV})^2}. \tag{41}$$

If this operator arises due to Z' exchange then we can express the coefficient $G_{bs\mu}$ as function of the couplings in Eq. (40),

$$G_{bs\mu} = -\frac{V_{ts}g_{bb}g_{\mu\mu}}{M^2_{Z'}}, \tag{42}$$

where we have assumed that the bs flavour mixing is given by the CKM element $V_{ts} \approx 0.04$. Together, Eqs. (41) and (42) imply the constraint on the parameters g_{bb}, $g_{\mu\mu}$ and $M_{Z'}$:

$$\frac{g_{bb}g_{\mu\mu}}{M^2_{Z'}} \approx \frac{1}{(6.4\ \text{TeV})^2}. \tag{43}$$

The Z' coupling to bs leads to an additional tree-level contribution to $B_s - \overline{B}_s$ mixing due to the effective operator arising from integrating out the Z' at tree level:

$$\Delta\mathscr{L}_{\text{eff}} \supset -\frac{G_{bs}}{2}(\bar{s}_L\gamma^\mu b_L)^2 + \text{h.c}, \qquad G_{bs} = \frac{g^2_{bs}}{M^2_{Z'}} = \frac{g^2_{bb}V^2_{ts}}{M^2_{Z'}}. \tag{44}$$

Such a new contribution is highly constrained by the measurements of the mass difference ΔM_s of neutral B_s mesons as discussed in 2015 leading to the bound of $(G_{bs})^{-1/2} \sim 150$ TeV [18], with a very recent 2017 bound from updated lattice results of $(G_{bs})^{-1/2} \sim 500$ TeV [19]. However the stronger 2017 bound arises from a discrepancy with the Standard Model which could in principle disappear. If we take the milder 2015 bound then this leads to the bound,

$$\frac{g^2_{bb}}{M^2_{Z'}} \lesssim \frac{1}{(6.0\ \text{TeV})^2}. \tag{45}$$

Taking the ratio of Eqs. (45) to (43) we find,

$$\frac{g_{bb}}{g_{\mu\mu}} \lesssim 1, \tag{46}$$

which is certainly satisfied for our model. However this would reduce by an order of magnitude for the stronger 2017 bound $(G_{bs})^{-1/2} \sim 500$ TeV [19]. We conclude that the model here satisfies the 2015 bound but not the stronger 2017 bound (which was presented after this paper was submitted).

6. Conclusions

In this work, we have presented a class of fluxed $SU(5)$ GUTs which naturally incorporate an abelian symmetry, $U(1)'$, where the associated gauge boson Z' displays non-universal gauge couplings to fermion families. This paper represents the first study of non-universal Z' gauge bosons in F-theory phenomenology. In our analysis we have considered an F-theory inspired framework where $SU(5)$ GUT emerges as a subgroup of the maximal exceptional gauge symmetry $E_8 \supset SU(5) \times SU(5)_\perp$. Non-trivial universal fluxes are assumed along the Cartan generators of the custodial group factor, $U(1)^4_\perp \subset SU(5)_\perp$, breaking the symmetry and generating chirality of the $SU(5)$ spectrum. A $\mathscr{Z}_2$ monodromy is imposed which allows rank-one fermion mass matrices and reduces the original symmetry to $SU(5) \times U(1)^3_\perp$, while a hypercharge flux breaks the $SU(5)$ GUT down to the Standard Model and, at the same time, induces the observed fermion chirality to the low energy spectrum. We then assume that the $U(1)^3_\perp$ gauge group is broken to a single non-universal $U(1)'$, by some unspecified Higgsing mechanism.

We have analysed the case where the non-universal $U(1)'$ is an anomaly-free linear combination of the remaining three $U(1)_\perp$ factors, $U(1)' = \sum_{i=1}^{3} c_i U(1)_{\perp,i}$. We chose a basis where two of $U(1)_{\perp,i}$ symmetries are identified with $U(1)_\psi \in E_6$ and $U(1)_\chi \in SO(10)$ and we formulated the gauge anomaly cancellation conditions in terms of a set of integers M_a, N_b introduced to parametrise the universal and hypercharge flux effects. We have scanned for solutions for the three coefficients c_i to determine the multiplicities of the corresponding light spectra in terms of the flux integers M_a, N_b. We found a plethora of models which are classified with respect to the nature of the non-universal couplings to fermions and the emergent spectrum. There are several classes of models with the minimal low energy MSSM spectrum, as well as several classes with additional vector-like fields.

In particular we highlight the phenomenologically promising classes with universal couplings to the three chiral fermion generations with only the single vector-like family having non-universal couplings. When the vector-like family mixes with the chiral families, non-universality involving the mixed light states may result. Such models naturally suppress New Physics contributions and models can be found which are universal in the first two light quark families, compatible with known processes such as the $K^0 - \bar{K}^0$ mixing, while involving non-universality in the third family. In principle these types of models may alter the branching ratios of the B-decays in accordance with the recently observed deficit in the $\mu^+\mu^-$ channel. We discuss the plausibility of the above scenario in relation the anomaly cancellation conditions and the scale of the New Physics effects.

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Appendix A. Singlets

For completeness, in this short appendix we present the singlet fields with their 'charge' assignments.

The superpotential involving only the singlet fields reads (see Table 6)

$$\mathscr{W} = \mu_{ij}\theta_{ij}\theta_{ji} + \lambda_{ijk}\theta_{ij}\theta_{jk}\theta_{ki} \tag{47}$$

where for each term $i \neq j \neq k \neq i$. The F-flatness conditions can be readily computed and obtain the following compact form

$$\frac{\partial\mathscr{W}}{\partial\theta_{ij}} = \mu_{ij}\theta_{ji} + \lambda_{ijk}\theta_{jk}\theta_{ki}$$

Table 6
Master table for singlet fields.

Singlet field	Weights	Q'	Multiplicity
θ_{13}	$t_1 - t_3$	$\frac{\sqrt{3}c_1}{2}$	$M_{\theta_{13}}$
θ_{14}	$t_1 - t_4$	$\frac{c_1+2\sqrt{2}c_2}{2\sqrt{3}}$	$M_{\theta_{14}}$
θ_{15}	$t_1 - t_5$	$\frac{1}{12}\left(2\sqrt{3}c_1 + \sqrt{6}c_2 + 3\sqrt{10}c_3\right)$	$M_{\theta_{15}}$
θ_{34}	$t_3 - t_4$	$\frac{\sqrt{2}c_2-c_1}{\sqrt{3}}$	$M_{\theta_{34}}$
θ_{35}	$t_3 - t_5$	$\frac{1}{12}\left(-4\sqrt{3}c_1 + \sqrt{6}c_2 + 3\sqrt{10}c_3\right)$	$M_{\theta_{35}}$
θ_{45}	$t_4 - t_5$	$\frac{1}{4}\left(\sqrt{10}c_3 - \sqrt{6}c_2\right)$	$M_{\theta_{45}}$
θ_{31}	$t_3 - t_1$	$-\frac{1}{2}\left(\sqrt{3}c_1\right)$	$M_{\theta_{31}}$
θ_{41}	$t_4 - t_1$	$-\frac{c_1+2\sqrt{2}c_2}{2\sqrt{3}}$	$M_{\theta_{41}}$
θ_{51}	$t_5 - t_1$	$\frac{1}{12}\left(-2\sqrt{3}c_1 - \sqrt{6}c_2 - 3\sqrt{10}c_3\right)$	$M_{\theta_{51}}$
θ_{43}	$t_4 - t_3$	$\frac{c_1-\sqrt{2}c_2}{\sqrt{3}}$	$M_{\theta_{43}}$
θ_{53}	$t_5 - t_3$	$\frac{1}{12}\left(4\sqrt{3}c_1 - \sqrt{6}c_2 - 3\sqrt{10}c_3\right)$	$M_{\theta_{53}}$
θ_{54}	$t_5 - t_4$	$\frac{1}{4}\left(\sqrt{6}c_2 - \sqrt{10}c_3\right)$	$M_{\theta_{54}}$

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