Physics Letters B 748 (2015) 251–254

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Non-generic couplings in supersymmetric standard models

Evgeny I. Buchbinder [a,*], Andrei Constantin [b], Andre Lukas [b]

[a] The University of Western Australia, 35 Stirling Highway, Crawley, WA 6009, Australia
[b] Rudolf Peierls Centre for Theoretical Physics, Oxford University, 1 Keble Road, Oxford, OX1 3NP, UK

ARTICLE INFO

Article history:
Received 25 May 2015
Accepted 7 July 2015
Available online 10 July 2015
Editor: M. Cvetič

ABSTRACT

We study two phases of a heterotic standard model, obtained from a Calabi–Yau compactification of the $E_8 \times E_8$ heterotic string, in the context of the associated four-dimensional effective theories. In the first phase we have a standard model gauge group, an MSSM spectrum, four additional $U(1)$ symmetries and singlet fields. In the second phase, obtained from the first by continuing along the singlet directions, three of the additional $U(1)$ symmetries are spontaneously broken and the remaining one is a B–L symmetry. In this second phase, dimension five operators inducing proton decay are consistent with all symmetries and as such, they are expected to be present. We show that, contrary to this expectation, these operators are forbidden due to the additional $U(1)$ symmetries present in the first phase of the model. We emphasise that such "unexpected" absences of operators, due to symmetry enhancement at specific loci in the moduli space, can be phenomenologically relevant and, in the present case, protect the model from fast proton decay.

1. Introduction

A widely accepted dictum is that all the couplings that are allowed by the symmetries of an effective field theory (EFT) should be present in the Lagrangian. In the present letter, we point out that care has to be taken when this principle is applied to EFTs from string theory. We will present an explicit example, in the context of a standard model derived from heterotic string theory on Calabi–Yau manifolds, where this principle appears to be violated, at least when thinking about the associated EFT in the standard way. One of the relevant key facts is that string theory can lead to symmetry enhancement at particular loci in moduli space. These additional symmetries are not directly visible at a generic locus since the corresponding gauge bosons are massive and removed from the EFT. Yet, these symmetries can still forbid certain operators everywhere in moduli space, thereby leading to "unexpected" absences of operators.

Our example model is based on a heterotic line bundle model on a certain Calabi–Yau manifold which has been constructed in two previous publications [1,2]. Here, we will focus on the associated low-energy theory and explain the effect purely in terms of the four-dimensional $N = 1$ EFT. We will discuss and compare two phases of this EFT, both of which have been obtained from a string construction. The first phase arises at a specific locus in moduli space and corresponds to an MSSM with four additional $U(1)$ symmetries and a number of fields uncharged under the Standard Model gauge group. The additional singlets in this model can be interpreted as bundle moduli and, from a low-energy perspective, they are candidates for right-handed neutrinos. The second phase which arises at a more generic locus in moduli space corresponds to an MSSM with an additional $U_{B-L}(1)$ symmetry. In low-energy terms, it can be obtained from the first phase by continuation along the singlet directions thereby spontaneously breaking three of the four $U(1)$ symmetries while leaving $U_{B-L}(1)$ unbroken.

Our point concerns the allowed operators in the second, generic phase with $U_{B-L}(1)$ symmetry.[1] It is well-known that dimension five operators inducing proton decay are allowed by $U_{B-L}(1)$. Following the general lore, we should, therefore, expect that these operators are present in the generic phase of our model. This would imply a serious phenomenological problem with proton stability. However, it turns out that the enhanced $U(1)^4$ gauge symmetry

[1] The $U_{B-L}(1)$ symmetry is a linear combination of the hypercharge and an additional $U(1)$ symmetry with massive gauge boson. The latter $U(1)$ manifests itself at low energies as a *global* symmetry. This approach is different from the one studied in [3,4]. To prevent the proton from fast decay the authors in [3,4] considered models with *local* $U_{B-L}(1)$ symmetry which then has to be violated by radiative corrections at scales below the string scale but higher than the electroweak scale [5,6].

* Corresponding author.
E-mail addresses: evgeny.buchbinder@uwa.edu.au (E.I. Buchbinder), andrei.const@gmail.com (A. Constantin), lukas@physics.ox.ac.uk (A. Lukas).

http://dx.doi.org/10.1016/j.physletb.2015.07.012
0370-2693/

which arises at the specific locus in moduli space comes to the rescue. Not only does this enhanced symmetry forbid the dimension five operators in question, it also forbids all such operators with additional singlet insertions. This means that these operators remain forbidden even when we turn on singlet vacuum expectation values and move away from the enhanced symmetry locus.

To explain this in detail, we define the two relevant effective field theories in Sections 2 and 3, respectively, and present their field content and their allowed superpotential couplings. We discuss the implications of the $U(1)$ symmetries at the enhanced symmetry locus and throughout the moduli space, in particular the resulting absence of dimension five proton decay operators. We conclude in Section 4.

2. The theory with enhanced symmetry

We begin by describing the four-dimensional $N = 1$ theory at the locus with enhanced symmetry, starting with the particle spectrum and followed by the key features of the effective action.

2.1. Spectrum

We consider an effective field theory with the standard model gauge group $G_{\rm SM} = SU(3) \times SU(2) \times U(1)$ and with an additional $U(1)^4$ symmetry group which significantly constrains the theory. Such models with extra $U(1)$ symmetries arise from compactifications of the $E_8 \times E_8$ heterotic string theory at specific loci in the moduli space where the structure group of the vector bundle degenerates [7–9]. The gauge bosons of these extra $U(1)$ groups can be massive or massless depending on the details of the model. If the gauge boson is massive the corresponding $U(1)$ group appears at low energies as a global symmetry. It is convenient to describe these additional $U(1)$ symmetries by the group $S(U(1)^5)$ whose factors we label by indices $a, b, \ldots = 1, \ldots 5$. Its representations are denoted by five-dimensional integral vectors

$$\mathbf{q} = (q_1, \ldots, q_5) \tag{1}$$

with the understanding that two charge vectors $\mathbf{q}$ and $\mathbf{q}'$ are identified, if $\mathbf{q} - \mathbf{q}' \in \mathbb{Z}\mathbf{n}$, where $\mathbf{n} = (1, 1, 1, 1, 1)$.

The gravitational spectrum of the model consists of the dilaton, S, four Kähler moduli $T^i = t^i + i\chi^i$ (where t^i are the geometrical fields, measuring the size of Calabi–Yau two-cycles, and χ^i are the associated axions) plus complex structure moduli which will not play an essential role in our discussion. The axions χ^i transform non-linearly under the $S(U(1)^5)$ symmetry[2] as

$$\delta_a \chi^i = -k_a{}^i \, . \tag{2}$$

In the rest of the paper we will concentrate on the specific model constructed in [1,2]. For our model, the integers $k_a{}^i$ are explicitly given by

$$\left(k_a{}^i\right) = \begin{pmatrix} -1 & -1 & 0 & 1 & 1 \\ 0 & -3 & 1 & 1 & 1 \\ 0 & 2 & -1 & -1 & 0 \\ 1 & 2 & 0 & -1 & -2 \end{pmatrix} \tag{3}$$

Let us review the properties of the resulting low-energy theory (see [1,2] for details). As was discussed above, the symmetry group of the low-energy effective theory is $G_{\rm SM} \times S(U(1)^5)$. In this case, three out of the four $U(1)$ gauge bosons receive string scale Stückelberg masses and the remaining one is massless.

The matter spectrum consists of the following multiplets

$$\begin{array}{llllll} 2\, Q_{\mathbf{e}_2} & 2\, u_{\mathbf{e}_2} & 2\, e_{\mathbf{e}_2} & Q_{\mathbf{e}_4} & u_{\mathbf{e}_4} & e_{\mathbf{e}_4} \\ & 2\, L_{\mathbf{e}_4+\mathbf{e}_5} & 2\, d_{\mathbf{e}_4+\mathbf{e}_5} & L_{\mathbf{e}_2+\mathbf{e}_5} & d_{\mathbf{e}_2+\mathbf{e}_5} & \\ & H_{\mathbf{e}_2+\mathbf{e}_4} & \bar{H}_{-\mathbf{e}_2-\mathbf{e}_4} & & & \\ 3\, S_{\mathbf{e}_2-\mathbf{e}_1} & 3\, S_{\mathbf{e}_4-\mathbf{e}_1} & 5\, S_{\mathbf{e}_2-\mathbf{e}_3} & 3\, S_{\mathbf{e}_2-\mathbf{e}_5} & S_{\mathbf{e}_4-\mathbf{e}_3}\, , & \end{array} \tag{4}$$

where the subscripts indicate the $S(U(1)^5)$ charges and $\mathbf{e}_a$ denote the standard unit vectors in five dimensions. The first three lines represent a perfect MSSM spectrum, however with specific $S(U(1)^5)$ charges for each multiplet. In addition, we also have a spectrum of singlets, S, which are neutral under the standard model group but charged under $S(U(1)^5)$. Note that the $S(U(1)^5)$ charge of the standard model multiplets only depends on the $SU(5)$ GUT multiplet they reside in, so that, for each family, the multiplets in $\mathbf{10} = [Q, u, e]$ have the same $S(U(1)^5)$ charge, as do the multiplets in $\bar{\mathbf{5}} = [d, L]$. This fact is related to the underlying group structure of the model, which originates from an $SU(5)$ GUT broken by a Wilson line.

The above spectrum is apparently anomalous. Indeed, one can compute the mixed $U(1)$–$G_{\rm SM}^2$ anomaly to find

$$\mathcal{A}_{U(1)-G_{\rm SM}^2} = \sum_{\text{all families}} \left(3\,\mathbf{q}(\mathbf{10}) + \mathbf{q}(\bar{\mathbf{5}})\right) = (0, 7, 0, 5, 3)\, . \tag{5}$$

However, these anomalies (as well as the cubic and mixed gravitational anomalies) are cancelled by the Green–Schwarz mechanism, facilitated by the axionic shifts (2).

If we describe linear combinations of the $U(1)$ symmetries by vectors $\mathbf{v} = (v^a)$ (demanding that $\mathbf{v} \cdot \mathbf{n} = 0$ to remove the overall $U(1)$), then massless vector bosons are characterised by the equation $k_a{}^i v^a = 0$. Applying this to Eq. (3) shows that, for our model, three of the four $U(1)$ symmetries are Stückelberg massive, while the linear combination $\mathbf{v} = (-4, 1, 6, -4, 1)$ remains massless.

2.2. Effective action

The Kähler potential has the standard form

$$K = -\log(S + \bar{S}) - \log(\kappa) + K_{\rm cs} + G_{IJ} C^I \bar{C}^J\, , \tag{6}$$

where $K_{\rm cs}$ is the complex structure Kähler potential which will not be needed explicitly and C^I collectively denote all matter fields listed previously. The specific form of the matter field Kähler metric G_{IJ} is not relevant to our discussion and it will be sufficient to know that it is positive definite. The pre-potential, κ, for the Kähler moduli is explicitly given by[3]

$$\kappa = d_{ijk} t^i t^j t^k = 12(t_1 t_2 t_3 + t_1 t_2 t_4 + t_1 t_3 t_4 + t_2 t_3 t_4)\, , \tag{7}$$

and this equation defines the topological numbers d_{ijk} for our model. We also note that the allowed range of the moduli t^i (the Kähler cone of the underlying manifold) is $t^i > 0$, for $i = 1, 2, 3, 4$.

From this Kähler potential and the $S(U(1)^5)$ symmetry transformations given earlier, we can compute the $S(U(1)^5)$ D-terms D_a. Their general form is [8][4]

$$D_a = \frac{3}{\kappa} k_a{}^i d_{ijk} t^j t^k + \sum_{I,J} q_a(C_I) C^I \bar{C}^J \tag{8}$$

where $q_a(C_I)$ denotes the $S(U(1)^5)$ charges of the matter fields. Due to the special unitary nature of the group these D-terms

[2] The dilatonic axion also receives a non-trivial transformation at one-loop order. However, this does not affect our discussion.

[3] For ease of notation, we will write explicit indices of the fields t^i as subscripts.

[4] In addition, there is also a one-loop correction to this D-term, resulting from the transformation of the dilatonic axion, which we omit. This correction does not affect our discussion.