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A rationale for long-lived quarks and leptons at the LHC: low energy flavour theory

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ABSTRACT: In the framework of gauged flavour symmetries, new fermions in parity symmetric representations of the standard model are generically needed for the compensation of mixed anomalies. The key point is that their masses are also protected by flavour symmetries and some of them are expected to lie way below the flavour symmetry breaking scale(s), which has to occur many orders of magnitude above the electroweak scale to be compatible with the available data from flavour changing neutral currents and CP violation experiments. We argue that, actually, some of these fermions would plausibly get masses within the LHC range. If they are taken to be heavy quarks and leptons, in (bi)-fundamental representations of the standard model symmetries, their mixings with the light ones are strongly constrained to be very small by electroweak precision data. The alternative chosen here is to exactly forbid such mixings by breaking of flavour symmetries into an exact discrete symmetry, the so-called proton-hexality, primarily suggested to avoid proton decay. As a consequence of the large value needed for the flavour breaking scale, those heavy particles are long-lived and rather appropriate for the current and future searches at the LHC for quasi-stable hadrons and leptons. In fact, the LHC experiments have already started to look for them.

KEYWORDS: Beyond Standard Model, Discrete and Finite Symmetries

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1 Introduction

Experiments at the Large Hadron Collider (LHC) have published their first results on the production of quasi-stable hadrons that are stopped and decay in the detectors [1–5]. Besides special cases of supersymmetric and extra-dimensional models, other models of different sorts also predict heavy particles with delayed decays [6], and many will be checked in turn at the LHC. In this letter, we discuss a new class of metastable particles suitable for those searches that arise in the framework of flavour models.

We argue that, in spite of the generic bound on the symmetry breaking scales of gauge flavour theories from rare transitions [7] being many orders of magnitude above the LHC reach, heavy quarks and/or leptons are plausibly expected with masses in the TeV range. Indeed, their masses can be much below the flavour breaking scale because they would be as protected by the flavour symmetries as the lighter Standard Model (SM) quarks that get masses much below the Fermi scale. As elaborated below, this may happen if the *raison d'être* of those heavy fermions is to cancel the anomalous couplings between the flavour and the SM gauge bosons induced by the light fermion sector.

The heavy fermion metastability is mainly due to an exact discrete symmetry that survives the flavour and electroweak symmetry breakings. This additional symmetry is

needed to prevent proton decay without forbidding neutrino masses. In order to build a consistent effective theory, we then assume a conserved symmetry, the so-called proton hexality [8], which can be written as $(B - 3L)/6$ in terms of the usual baryon and lepton numbers. Since we assume that it results from the breaking of the continuous flavour symmetry, it is local and non-anomalous.¹

Because these new fermions can be consistently endowed with non-standard discrete baryon/lepton numbers they do not mix at all with SM fermions but can decay into three light fermions through dimension-six operators. Therefore their decay lengths are² $c\tau_F \sim 48\pi^3\Lambda^4 m_F^{-5}$, where their masses are $m_F = O(\text{TeV})$ and Λ^{-2} is the coefficient of the corresponding four-fermion interactions. From the analysis of the analogous flavour changing neutral current (FCNC) dimension-six operators for light fermions, flavour physics puts a limit $\Lambda > 5 \cdot 10^4 \text{ TeV}$ [10]. Hence $c\tau_F > 2 \text{ km}$ — actually, much more in the explicit models displayed below.

Of prime importance is the fact that the new states belong to real representations of the SM gauge symmetry, so that their masses mostly arise from their couplings to the flavon(s) rather than to the Higgs boson, and mass mixings with light fermions are forbidden by the discrete symmetry. Thus the severe bounds on those mixings are naturally avoided and their contributions to electroweak precision tests (EWPT) are quite suppressed.

2 A model

For the sake of argument, we take the instance of a single charge abelian flavour symmetry group $U(1)_X$ broken by the vacuum expectation value (*v.e.v.*) of a single complex scalar (ϕ), the so-called flavon, and assign (different) X -charges to the quarks and leptons of either chirality in the three families so to yield an acceptable description of the observed mass hierarchies and mixings. Furthermore, by a suitable combination with the hypercharge, we choose for the Higgs, $X(H) = 0$ and normalize it so that $X(\phi) = -1$. Most important in our analysis are the chiral charge differences of quarks and leptons, $\chi_f^{ij} = X(f_L^i) - X(f_R^j) \equiv f_L^i - f_R^j$, where $f_L = q, l$, $f_R^i = u, d, e$, and $i, j = 1, 2, 3$ are family indices. All the SM model indices are dropped, in a notation that is quite standard for the SM fermions of both chiralities. The chiral charge differences forbid the couplings of all fermions but the top quark to the Higgs at the renormalizable level, providing the needed chiral protection for their masses.

This defines a Froggat-Nielsen effective Lagrangian [11, 12] with a cutoff Λ way above the electroweak scale after integrating out the flavon and the $U(1)_X$ gauge boson, denoted X_μ . We then basically follow the steps in the supersymmetric version in [9], to which we refer for more details, although the presence of scalars there makes the effective theory, as much as its phenomenology, very different from the one discussed below (in particular, the supersymmetric heavy states are short-lived).

¹In the supersymmetric version [9], it replaces the R -parity. We refer to this paper for a more detailed discussion of some issues below.

²The symbol “ \sim ” is used here to denote equality up to an $O(1)$ factor.

3 Charged fermion and neutrino masses

The flavon *v.e.v.*, $\langle \phi \rangle = \epsilon \Lambda$, breaks the flavour symmetry and allows for the fermions to acquire masses from higher dimension operators with a Higgs field and χ_f^{ij} insertions of the field ϕ (ϕ^\dagger) and divided by χ_f^{ij} factors of the cutoff Λ . The mass matrices are then $m_f^{ij} \sim \epsilon^{|\chi_f^{ij}|} v$, where $v = 174$ GeV is the *v.e.v.* of the Higgs field and $f = u, d, l$. From the masses and (left-handed) mixings one can find several solutions for the X -chiralities matrices $|\chi_u^{ij}|$, $|\chi_d^{ij}|$ and $|\chi_l^{ij}|$ that can be found in the literature [13–16]. Here we mostly need their traces: $\text{Tr}|\chi_f| \approx \ln \det(m_F/v)/\ln \epsilon$ and the value of the parameter ϵ that comes out close to the Cabibbo angle, $\epsilon \sim \theta_C \approx 0.2$. Thus one finds: $\text{Tr}|\chi_u| \approx 12$, $\text{Tr}|\chi_d| \approx \text{Tr}|\chi_l| \approx 15$. The fits are better if $\text{Tr}\chi_u$ and $\text{Tr}\chi_d$ have the same sign, which we choose to be positive.

The effective dimension-five operator $(\nu_i \cdot H)(\nu_j \cdot H)/\Lambda$ has charge $X = l_i + l_j$, therefore, the resulting neutrino mass matrix becomes $m_\nu^{ij} \sim \epsilon^{|l_i+l_j|} v^2/\Lambda$. Experiments allow for only a little hierarchy in the eigenvalues and relatively large mixings, so the best one can do in our case is to choose the same charge $l_i = l$ for all families, so that,

$$\epsilon^{|2l|} \sim \frac{\sqrt{\Delta m_{\text{atm}}^2} \Lambda}{v^2} = \frac{\Lambda}{6 \times 10^{11} \text{ TeV}}. \tag{3.1}$$

This fixes $|l|$ for a given cutoff. By imposing $\Lambda \gtrsim 5 \cdot 10^4$ TeV to meet the requirements from flavour physics discussed above, it follows that $|2l| \leq 11$.

4 Exact discrete symmetry

Omitting all Lorentz and SM indices, there are six B and L violating four-fermion operators that can be schematically written as: $lqqq$ (2 operators), $eudu$ (2 operators), $euqq$, and $lqdu$ [17, 18]. They are related to proton decay and preserve $B-L$. Thus, they are forbidden by the discrete symmetry generated by the charge $Z' = B/6 - L/2 \pmod 1$, defining proton hexality, which stabilizes the proton. One can also combine $Z = (Z' - Y/3) \pmod 1$ which defines a Z_6 symmetry as can be seen from the Z' and Z quantum numbers displayed in table 1. Neutrino Majorana masses are allowed and the flavon ϕ is invariant. We embed this symmetry in the flavour one $U(1)_X$, so that it is local and non-anomalous in the most economical way. With this choice, the $U(1)_X$ symmetry is broken by the ϕ *v.e.v.* into the discrete one and Z' coincides with the fractional part of X . Notice that Z' is vector-like to allow for fermion masses, family-independent to allow for family mixings in weak interactions, and anomaly-free with respect to the SM. Since $B-L$ allows for proton decay, one is left with Z' to avoid a fast proton decay.

Concerning the new heavy states to be introduced later on, one must efficiently protect the light fermion mass matrices and efficiently suppress the mixing with the light ones by either giving exotic quantum numbers with respect to the SM gauge group, or by assigning different Z' charges to them. In the first case the mixings will be reduced by factors of v/Λ , but the structure in the mass matrices might be affected in a model-dependent way, so that we choose the second one where the mixings are just forbidden. It seems natural to compensate the anomalies from the SM sector by a choice of heavy fermions with, instead,

f	SU(3)	SU(2)	Y	Z_6	$18Z'$
$q_{(L)}$	3	2	1/6	0	1
$u_{(R)}$	3	1	2/3	-1/6	1
$d_{(R)}$	3	1	-1/3	1/6	1
$l_{(L)}$	1	2	-1/2	-1/3	-9
$e_{(R)}$	1	1	-1	-1/6	-9

F	F'	SU(3)	SU(2)	Y	Z_6	$18Z'$
$Q_{(L)}$	$Q'_{(R)}$	3	2	1/6	1/3	7
$U_{(R)}$	$U'_{(L)}$	3	1	2/3	1/6	7
$D_{(R)}$	$D'_{(L)}$	3	1	-1/3	1/2	7
$L_{(L)}$	$L'_{(R)}$	1	2	-1/2	0	-3
$E_{(R)}$	$E'_{(L)}$	1	1	-1	1/6	-3

Table 1. Notation and quantum numbers of light and heavy fermions with $Z = Z' - \frac{1}{3}Y$. In the text the chiralities (L, R) are omitted.

the same SM quantum numbers, and denote them by capitals as $F = Q, L, U, D, E$, with the same chiralities as the light ones ($Q \equiv Q_L, U \equiv U_R$, etc.). Their parity conjugated states are denoted F' . Furthermore, the notation for the charges is simplified below: $X_F \equiv F$.

5 Heavy quarks and leptons

There are arbitrarinesses in the choice of heavy states needed for the cancellation of anomalies. Let us assume asymptotic freedom so that the total number of heavy quarks is limited to ten. If we further assume whole families of heavy fermions analogous to the SM ones together with their parity conjugated states, there can be at most two of them. For simplicity, here we concentrate on just one heavy family, the generalization to two being straightforward as much as for the case of incomplete families.

The main contributions to the masses of the heavy fermions, m_F , are reduced from their natural scale Λ by their X-chiralities, defined as $\chi_F = F_L - F_R$ so that,

$$m_F \sim \epsilon^{|\chi_F|} \Lambda \sim (3 \times 10^{12}) \epsilon^{|\chi_F|+|2l|} v \tag{5.1}$$

where (3.1) has been used.

In the analysis of the model that follows, there are two bounds one can derive on m_F which imply on limits on the largest $|\chi_F|$. A lower bound that all new fermions have to oblige since they have not been observed yet, i.e., $(m_F)_{\min} \gtrsim v$, and another that selects models with fermions that can be reasonably produced at the LHC, by an upper bound $(m_F)_{\min} \lesssim \epsilon^{-2}v \sim 4 \text{ TeV}$ for the lightest heavy fermion . Therefore we impose the condition:

$$16 - 2|l| \leq |\chi_F|_{\max} \leq 18 - 2|l|. \tag{5.2}$$

Notice that $|\chi_F|_{\max}$ corresponds to the lightest long-lived particle (LLLP), although more than one state could share the same value of $|\chi_F|$ (see one example below), or the next lightest particle could be more easily produced. Finally, one can introduce two whole or incomplete families to compensate the anomalies. Yet, some states could be much lighter than the others, leading to results similar to the one family case in analogy with what happens with the three light families.

Now, since the new states cannot be stable for cosmological reasons, one must allow for their coupling to three light fermions via dimension-six operators. Their SM quantum numbers allow only for the analogous of the proton-decay four-fermion operators described above, see section 4, with any one light state replaced by its corresponding heavy one. This leads to the following patterns of decays:

$$\begin{array}{llll}
 L \rightarrow \bar{q}\bar{q}\bar{q} & L \rightarrow \bar{u}\bar{d}\bar{q} & E \rightarrow \bar{u}\bar{q}\bar{q} & E \rightarrow \bar{u}\bar{u}\bar{d} \\
 Q \rightarrow \bar{q}\bar{q}\bar{l} & Q \rightarrow \bar{d}\bar{u}\bar{l} & Q \rightarrow \bar{q}\bar{u}\bar{e} & D \rightarrow \bar{q}\bar{u}\bar{l} \\
 D \rightarrow \bar{u}\bar{u}\bar{e} & U \rightarrow \bar{q}\bar{u}\bar{l} & U \rightarrow \bar{q}\bar{q}\bar{e} & U \rightarrow \bar{d}\bar{u}\bar{e}
 \end{array} \tag{5.3}$$

where all indices have been skipped.

The heavy fermion Z' are to be chosen *ad hoc*, and their parity conjugated states have the same values to allow for singlet masses. They are displayed in table 1: heavy quarks have $Z' = 7/18$, heavy leptons have $Z' = -1/6$ like antinucleons.

Of most relevance are the flavour charges since they control the suppression of the decay rates: they select the main modes in terms of light fermion flavours. In most of the cases, lifetimes are increased by many orders of magnitude with respect to the lower bound of a few meters (cf. the examples below). The experimental long-life signatures are not very sensitive to the lifetimes as far the decay is always delayed enough, but the family dependence of the decay products are model dependent.

6 Electroweak precision tests

Besides the dominant masses above, the heavy states get mass contributions also from their couplings to the Higgs, analogous to the SM ones. The set of electroweak precision tests (EWPT) are very sensitive to these couplings, being the most important new physics contributions summarized by the oblique parameters [19, 20]. Let us concentrate, for definiteness, into the system $\{Q, U, Q', U'\}$, keeping in mind the definition of their X -chiralities. The order of magnitude of mass terms are given by:

$$\epsilon^{|\chi_Q|} \Lambda \bar{Q}' Q + \epsilon^{|\chi_U|} \Lambda \bar{U} U' + \epsilon^{|\chi_{QU}|} v \bar{U} Q + \epsilon^{|\chi_{Q'U'}|} v \bar{Q}' U' + \text{h.c.} \tag{6.1}$$

where $\chi_{QU} = Q - U$ and $\chi_{Q'U'} = Q' - U'$. The last two ('non-diagonal') masses are reduced by a factor $v/\Lambda < \epsilon^8$ with respect to the dominant ('diagonal') ones, however, their actual ratios depend on model dependent flavour factors. The other mass terms are analogously obtained by the replacements $Q, U \rightarrow Q, D$ and $Q, U \rightarrow L, E$.

The heavy fermion loops contribute to the S and T parameters with, respectively, two and four insertions of Higgs in the new fermion lines, given by the appropriate combinations

of the non-diagonal masses m_{QU} , m_{QD} , m_{LE} , $m_{Q'U'}$, $m_{Q'D'}$, $m_{L'E'}$. Roughly, S and T are proportional to products of two mass ratios, e.g., $m_{QU}m_{Q'U'}/m_Q^2 \sim \epsilon^{|\chi_{QU}|+|\chi_{Q'U'}|-2|\chi_Q|}(v/\Lambda)$, and similar ones. General upper bounds can be obtained from the approximations displayed in ref. [21] for the contributions ΔS and ΔT from the heavy quark sector. Replacing all the non-diagonal masses in the numerators of expressions therein by the largest one, denoted \hat{m} , and the diagonal ones in the denominators by the smallest one, \hat{M} , leads to the upper limits:

$$\Delta T \lesssim \frac{N_c}{2\pi} \frac{\hat{m}^4}{v^2 \hat{M}^2} \quad \Delta S \lesssim \frac{N_c}{8\pi} \frac{\hat{m}^2}{\hat{M}^2} \quad (6.2)$$

where $N_c = 3$ for heavy quarks and 1 for heavy leptons. Since $\hat{m} \lesssim v$ while $\hat{M} \gtrsim v$ from the existent bounds, the contributions of the heavy fermions might be dangerous only when³ \hat{m}/\hat{M} is $O(1)$, namely, when both are $O(v)$. Therefore, the model would be constrained only in a very small portion of its parameter space. Indeed, the discrete symmetry introduced to protect the mass matrix textures also plays a role for the protection of the EWPT parameters by avoiding additional mixed couplings to the Higgs.

7 FCNC and CP violations

The main obstacle to lower the flavour symmetry breaking scale comes from the impressive description of flavour changing and CP violating processes by the SM, leaving in some cases almost no space for new physics contributions. In our approach here, the dangerous operators are $B = L = 0$ four-fermion operators with arbitrary flavour structure, suppressed by Λ^2 , as well as powers of ϵ due to X -charge imbalance, which are relatively model dependent. Other dimension-six operators, $f_R^* H f_L \mathcal{F}$ where \mathcal{F} stands for the photon or the gluon field strength, give rise to flavour changing magnetic moments and electric dipole moments, hence strong limits on Λ .

However the utmost limit on the cutoff Λ turns out to be generic in our framework [7]. Indeed, the exchange of massive flavour gauge bosons give rise to current-current interactions of range $\epsilon\Lambda$, which of course are flavour diagonal in the basis where X is diagonal. Since the fermion charges are different, in going to the mass eigenstate basis, the mixings in flavour currents are given by the mass matrices. Thus for the $s \rightarrow d$ transitions one expects a mixing $O(\theta_C) \sim \epsilon$. Comparison with rare K-physics processes [10] requires $\Lambda > 5 \cdot 10^4$ TeV, as already asserted, puts the lower bound on the cutoff.

8 Anomaly cancellations

Let us now turn to the main ingredient, the cancellation of anomalies introduced by the SM fermion loops. Each one has to be canceled by the heavy sector comprising fermions with SM quantum numbers similar to the light ones. For simplicity we write the expressions for only one regular family of heavy fermions, the generalizations being evident. The anomalies with one flavour boson coupling to the parity conserving gauge bosons of QCD and QED,

³We are assuming only one non-diagonal mass to be dominant. As asserted in ref. [21], (6.2) are overestimates for $\hat{m} \sim \hat{M}$.

depend only on the X-chiralities of the fermion in the loops. They can be combined⁴ into two relations:

$$2\chi_Q + \chi_U + \chi_D = -\text{Tr}(\chi_u + \chi_d) \quad (\approx -27), \quad (8.1)$$

$$\chi_L + \chi_E - \chi_Q - \chi_D = -\text{Tr}(\chi_l - \chi_d) \quad (\geq -30, \leq 0), \quad (8.2)$$

where the values evaluated above are used to obtain the value and limits in brackets, since the signs of χ_l are not fixed.

The cancellation of the anomalous coupling of X_μ to two W 's results in:

$$3\chi_Q + \chi_L = -\text{Tr}(3q + l) \approx -3\text{Tr}(q) + 3l. \quad (8.3)$$

depends on the sum of X -charges that are very model-dependent. Notice that the contribution of the fractionary charges is $9Z'_q - 3Z'_l$, which, of course, vanishes for $B - L$ and is integer (-1) for $Z' = B/6 - L/2$ so that the heavy states can cancel the anomaly. Finally, the anomalous coupling of the photon to two X_μ 's is zero if:

$$\sum_{F=Q,U,D,L,E} \text{Tr} Q_F \chi_F (F_L + F_R) = - \sum_{f=u,d,l} \text{Tr} Q_f \chi_f (f_L + f_R) \quad (8.4)$$

where Q_f, Q_F , are the electric charges and Tr is taken on all SM and flavour indices. This equation also depends explicitly on the X -charges.

Clearly, since there are many more parameters than conditions, there are many solutions to these equations even with a single heavy family. However, the values taken by r.h.s.'s in (8.1) and (8.2) indicate that for many solutions the LLLP would fit into the visibility range (5.2).

Interestingly enough, (8.1) implies that $|\chi_F|_{\max} \geq 7$, hence $|2l| \leq 11$. From (3.1), this is equivalent to the limit on the cutoff from FCNC experiments, i.e., in the models discussed here with only one heavy family, the solutions to the anomaly conditions with $\Lambda < 5 \cdot 10^5$ TeV would require a metastable fermion too light to have escaped observation.

9 Case studies

From (8.3), one sees that the solutions depend quite strongly upon $\text{Tr}q$ and, of course, upon l which also fix the scale Λ , and $\text{Tr}\chi_l$ which is not fixed by the charged lepton masses. For the sake of illustration, we consider only a few simple choices that lead to satisfactory mass matrices. Let us first introduce a notation for the integer part of the X -charges: $x_f = X_f - Z'_f$. The results are shown in table 2. Of course there are many more solutions.

Notice that only relatively high scales, $\Lambda \sim 5 \cdot 10^5 (10^7)$ TeV for $x_l = 5$ or -4 (4) are obtained in these examples. These translate into very long lifetimes. In one case an incomplete family is enough to compensate the anomalies and the metastable lepton is quite light. Heavy quarks decay with a very energetic lepton in the final state, with a characteristic energy distribution.

⁴See, e.g., [9] for details.

x_Q	x_U	x_D	x_L	x_E	LLLP	Mass	Decay Modes	$c\tau$
$x_{Q'}$	$x_{U'}$	$x_{D'}$	$x_{L'}$	$x_{E'}$				
$\chi_l = 15 \quad x_l = 5 \quad x_d = (3 \ 3 \ 3) \quad x_q = x_u = (4 \ 2 \ 0) \quad x_e = (2 \ 0 \ -2)$								
-8	+3	+6	-8	+5	Q or L	$m_Q \sim m_L \sim \epsilon^8 \Lambda$ $\sim 1 \text{ TeV}$	$Q \rightarrow ldu, que$ $L \rightarrow duq$	$\sim 4 \times 10^3 \text{ km}$ $\sim 2 \times 10^4 \text{ km}$
-6	+1	+2	-5	+5	Q or L	$m_Q \sim m_L \sim \epsilon^8 \Lambda$ $\sim 1 \text{ TeV}$	$Q \rightarrow que, ldu$ $L \rightarrow duq$	$\sim 4 \times 10^3 \text{ km}$ $\sim 4 \times 10^3 \text{ km}$
$\chi_l = 15 \quad x_l = 4 \quad x_d = (3 \ 3 \ 3) \quad x_q = x_u = (4 \ 2 \ 0) \quad x_e = (-1 \ 1 \ 3)$								
-6	+3	0	-3	+4	Q or E	$m_Q \sim m_E \sim \epsilon^{10} \Lambda$ $\sim 1 \text{ TeV}$	$Q \rightarrow qql$ $E \rightarrow qqu$	$\sim 10^9 \text{ km}$ $\sim 10^9 \text{ km}$
-6	+3	+5	+5	+6	E	$m_E \sim \epsilon^{11} \Lambda$ $\sim 0.2 \text{ TeV}$	$E \rightarrow qqu$	$\sim 5 \times 10^{12} \text{ km}$
-1	+2	-3	-6	+3	U	$m_U \sim \epsilon^{10} \Lambda$ $\sim 1 \text{ TeV}$	$U \rightarrow ued, uql$	$\sim 10^9 \text{ km}$
+7	-8	-4	-1	-1				
$\chi_l = -15 \quad x_l = -4 \quad x_e = (0 \ -2 \ -4) \quad x_q = x_u + 2 = (5 \ 3 \ 1) \quad x_d = (2 \ 2 \ 2)$								
-13	+4	-2	-6	+3	E	$m_E \sim \epsilon^9 \Lambda$ $\sim 0.3 \text{ TeV}$	$E \rightarrow qqu$	$\sim 10^7 \text{ km}$
-6	-2	-9	-13	+12				

Table 2. Examples of solutions to the anomaly cancellation equations, with $x = X - Z'$, for some choices of the charges in the light sector. In each case, the lightest metastable particle (LLLP), and its mass, decay modes and mean path are displayed.

10 Experimental searches: production and detection

The new heavy quarks and leptons can be pair produced at the LHC with sizeable cross sections as shown in figure 1 for center-of-mass energies of 7 and 14 TeV. As seen above, the decay of the new heavy states takes place through suppressed operators, leading to long lifetimes. The search for new long-lived states is conducted by the ATLAS and CMS collaboration using several experimental techniques, such as dE/dx , time of flight and decays during beam collision intervals [1–4].

The out-of-time decays [5] constrains the heavy particle production cross section times stopping probability to be smaller than $\simeq 0.2 \text{ fb}$. Assuming that the stopping probability for the hadron containing new heavy quarks is similar to the one for stop hadrons [1–4], i.e., of the order of 20%, the presently available data exclude new long-lived quarks with mass $\lesssim 400 \text{ GeV}$. Certainly, this bound should be taken with a pinch of salt since it depends upon the unknown strong interactions of hadrons exhibiting a heavy quark. In the case of new heavy leptons, the out-of-time analysis excludes new leptons with masses $\lesssim 100 \text{ GeV}$, assuming that the stopping probability is equal to the one for staus, i.e. 5%.

More stringent limits can be obtained from the search for slowly moving charged particles [5]. Assuming that the detection of hadrons containing heavy quarks is similar to the one for hadrons possessing stops we obtain that the new heavy quarks must have a mass in excess of $\simeq 800 \text{ GeV}$. Analogously, the search for slow moving staus can be translated to an lower limit on the new lepton mass of $\simeq 400 \text{ GeV}$.

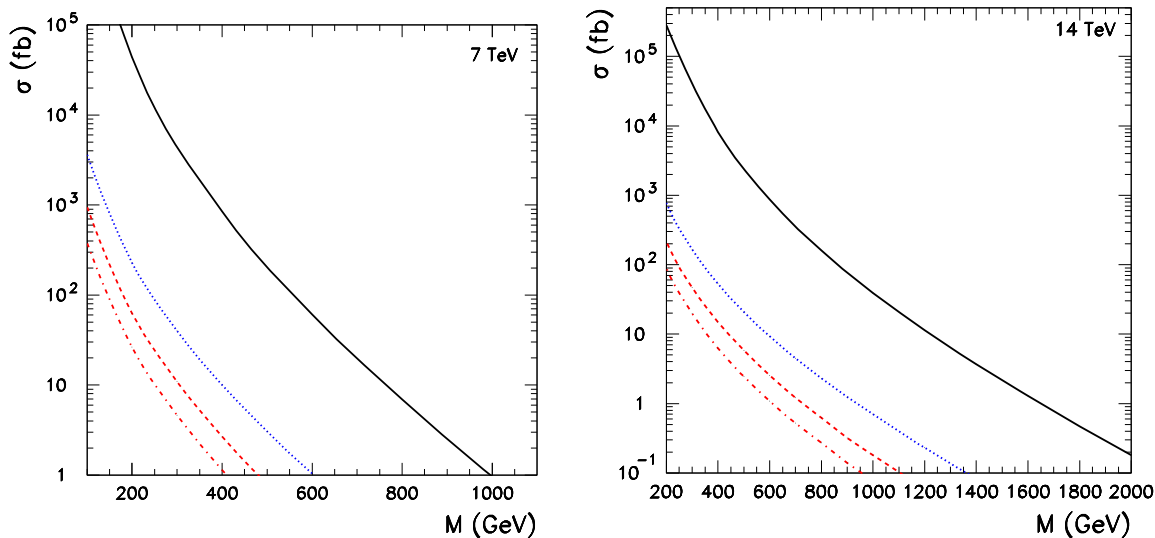


Figure 1. Total production cross section at the LHC as a function of the mass of the LLLP at a center-of-mass energy of 7 TeV (left panel) and 14 TeV (right panel). The black solid line stands for the strong interaction production of new heavy quarks while the blue dotted, red dashed and red dot-dashed lines represent the production of doublet charged lepton pairs, doublet charged and neutral lepton pairs and singlet charged leptons, respectively.

11 Conclusions

We have presented a new class of quasi-stable hadrons and leptons that naturally arise in the framework of flavour models. Those particles can plausibly be expected to show up at the LHC as their masses are protected by the flavour symmetry and can lie in the TeV range. This is much lower than the symmetry breaking scale that is constrained by FCNC and CP violating experimental data to be $\Lambda > 5 \cdot 10^4$ TeV. Limits on the new heavy fermions from EWPT can be avoided by suppressing mixed couplings of the new fermions with the SM ones through the Higgs by the same residual proton-hexality symmetry that prevents rapid proton decay and explains the observed texture of fermion masses and mixings. Presently there are ongoing searches at the LHC that can reveal their existence.

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