Baryon number violation
at the TeV scale

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A la veille de ma défense privée, de nervosité, je fais les cent pas dans ma chambre d'enfant réinvestie entre deux déménagements. Est-ce sa couverture rouge qui capte mon attention ? J'ouvre ce guide pratique d'origami et redécouvre certaines figures reproduites à l'époque. Je m'arrête ensuite sur ces quelques pages d'introduction, probablement peu ou pas lues. "L'origami est poésie." Il ne nécessite ni adresse, ni matériel particulier ; un peu de papier, de concentration et de patience lui suffisent. Le rapprochement avec une certaine physique est de circonstance, celle des idées qui m'est si savoureuse.

"L'art rapproche les Hommes puisqu'il nécessite la paix pour se développer," l'introduction se poursuit, comme mon parallèle avec les sciences. A l'heure de clôturer ces années de doctorat, les Hommes que je souhaite remercier ne sont pas seulement Hommes de sciences. Une œuvre d'art n'est pas le produit de quelques minutes, heures, semaines, mois de travail mais le résultat de toute une vie, aurait dit Picasso.

Il y a les membres de mon jury, Giacomo Bruno, Christophe Grojean et Jernej Kamenik qui ont donné de leur temps pour lire, écouter, critiquer et proposer des améliorations.

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avons passées à discuter, surtout de physique, lors de notre première rencontre à Lyon ont été pour moi fantastiques ; ton expérience partagée de la vie de physicien dans ses aspects extra-scientifiques, un enseignement de grande valeur.

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Il y a les amis et colocataires qui auront eu, entre autres mérites, celui probablement inconscient mais non moins apprécié, de m'offrir une fenêtre portant au delà de ce microcosme de l'université, de la recherche, de la physique et de rappeler l'existence d'une autre normalité, du concept d'heures de bureau et de week-ends !

Il y a tous ceux dont j'ai croisé la route, avec qui j'ai partagé un bout de chemin, échangé quelques mots ou regardé s'écouler un peu de notre temps. A votre contact se façonne ce que je suis.

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Foreword

The elementary constituents of matter and their interactions are the object of particle physics. The established description of all known fundamental forces but gravity in a quantum framework is named *standard model*. Completed in the seventies, it encompasses in a common framework the electromagnetic force, *i.e.* a unified description of electricity, magnetism and optics, as well as the weak and strong nuclear forces, respectively responsible for radioactivity and the binding of atomic nuclei. A fifth force thanks to which the fundamental particles acquire their masses and recently confirmed experimentally [1] is also included. At the Large Hadron Collider (LHC) which accelerates protons against each others, a huge experimental effort has indeed led to the discovery of a new boson with properties very much alike the ones predicted by the Brout-Englert-Higgs mechanism [2] in its minimal implementation.

The standard model is however known to be incomplete. Many different extensions have been suggested over the years to solve important issues raised by experimental observations or theoretical considerations. The neutrino masses and mixings are notably not accounted for; too numerous apparently free parameters call for an elucidation; the large hierarchies between some of them look unnatural; what is called *dark matter* remains unknown; there is no fully satisfactory explanation for the matter–anti-matter asymmetry of the universe; etc. Standard-model extensions should not only be self-consistent from a theoretical point of view, they also have to be compatible with presently available experimental data and new predictions should make them testable experimentally. The measurement indeed provides the final criterion to confirm or exclude any theoretical hypothesis.

The strong link required between pure theory and experiment is however not straightforward to establish as the technicality of both approaches requires a high level of specialization. So-called *phenomenology* is a branch of particle physics which is concerned with bridging possible gaps between these two sub-fields by developing theoretical models with interesting observational consequences, establishing strategies for testing efficiently theoretical assumptions, and interpreting experimental measurements in those frameworks. Phenomenologists also provide experimentalists with the theoretical input, the strategies

and tools needed to analyse the rough information obtained from their impressive detectors.

The data already collected and still to be gathered by the LHC on the most energetic collisions ever produced in controlled conditions provides physicists with invaluable information on what *new physics* should complement the standard model. No clear evidence has been found so far and many proposed scenarios are increasingly constrained by more and more precise observations. In the present quest for further experimental evidences, generic approaches relying on few assumptions about what exactly is to be discovered are therefore interesting to pursue. Those *model-independent* or *effective* techniques readily complement the efforts devoted to the construction of explicit models.

Exploiting what has already been firmly established is essential for targeting the most realistic directions of extension. Symmetries are paramount guiding principles in this endeavour. A symmetry is an operation against which a physical law remains unchanged. Translations are good examples: fundamental laws of nature do obviously not change as they are observed in one laboratory or the other. Before the twentieth century, symmetries were basically seen as interesting consequences of known dynamical laws. Physicists then noticed that they are actually primary features of our understanding of nature that, to the least, constrain the allowed dynamical laws, or even dictate their very form. Symmetries acting globally or locally, on space-time or on the internal spaces of particle properties therefore play a fundamental role in modern physics.

Under the hypothesis of their conservation, only a manageable variety of extensions of our standard description are allowable. In this way, we are therefore provided with very indicative information on where to search for new phenomena. The detection of those would then help us very much in determining what new dynamics—if any—is to complement the known one, what new law should replace the so-far cherished rules.

I. Introduction

This dissertation describes some contributions [3–5][1] to the current new-physics search program of particle physics. The spirit of effective approaches played an important guiding role and will therefore be the subject of this first introductory chapter. An effective-field-theory view on the standard-model will then be provided, through its construction, step by step, using symmetries and the mathematical objects called *fields* that describe observed particles. This will naturally lead us to consider the interactions, in the form of *higher-dimensional operators*, that extend the set constituting our standard model. They are constructed using exactly the same fields, *gauge* (local, internal) and Lorentz (global, space-time) symmetry requirements. A special focus will be devoted to interactions involving the heaviest known elementary particle, the *top quark*.

A first framework for describing new physics is therefore the effective-field-theory one. The subset of operators violating a questionable symmetry called *baryon number* will be examined with particular attention in the second chapter. An approximate conservation of this symmetry could imply that the protons and neutrons composing the nuclei of atoms are actually unstable. Not a single of their decays has been observed so far but baryon number violation is a generic feature of almost all theoretical models and a necessary ingredient for explaining why a so little amount of anti-matter is seen in the universe. Some interesting processes these operators could trigger at the LHC will be described. They involve the top quark, which could possibly have behaviours distinct from its more common and lighter homologues.

The construction of a generic effective theory is only possible with known particles. The indirect effects of hypothetical heavier states are described effectively but this framework cannot properly account for new light particles directly producible in the experiment considered. A third chapter will therefore address the possibility that new physics could be directly produced at the LHC, in *resonant* transitions between known particles. Resonant baryon-number-violating processes could occur without conflicting with existing experimental data if, on top of the gauge symmetries of the standard model, new physics also preserves the breaking pattern of its *flavour* (global, internal) sym-

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[1]Conference proceedings on those topics were also produced [6, 7].

metries. This assumption is also very well supported by precise measurements. The heaviest families or *generations*—including the top quark—are then preferentially involved in baryon-number-violating interactions, along with their lighter counterparts. A generic but qualitative description of some promising LHC signatures is then provided.

A fourth chapter first discusses how this standard-model flavour symmetry breaking pattern can systematically and precisely be imposed on new-physics interactions through the so-called *minimal-flavour-violation* prescription. A quantitative study of a specific *supersymmetric* scenario including, explicitly, new particles that could mediate resonant transitions at the LHC is then carried out. In that particular framework, the baryon-number-violating signature involving top quarks and generically identified in the previous chapter is shown to be generically expected and especially well suited to probe for this type of model.

I.1 Describing nature effectively

Let us here introduce the effective-approach spirit exploited in what follows.[2]

In all areas of natural sciences, widely different scales and parameters can potentially enter the modelling of processes. In a specified regime, a good description providing insight in what is actually happening and allowing the production of the most accurate results with a minimum of technical efforts inevitably includes first the most relevant mechanisms and treats (very) small parameters as vanishing. The comparison of model predictions with precise experimental data may sometimes require considering next-to-leading terms in the small parameter perturbative expansion. A systematic method for doing so had therefore better be provided together with the effective model definition.

A more fundamental description may be available and could potentially be used to determine what those corrections are. If such a model depending on fewer parameters is not available or if useful results cannot be derived in the situation of interest, we may fix the parameters of our effective description using (a fraction of) the experimental data available and confront the predictions that can then be made with other measurements. A classical example from particle physics is provided by the strong force whose strength is so important at low energies that our usual perturbative description is intractable. Perturbative computations can only be carried out in an effective theory that actually features different degrees of freedom than those of the high energy description.

This picture is very familiar to all natural scientists: general relativity is unnecessary for describing the fall of an apple off a tree; quantum electrodynamics is of little relevance for studying biological processes; etc. In describing the elementary building blocks of matter and their interactions, one may be tempted to extrapolate models to arbitrarily small distances or arbitrarily high energies. Requiring this extrapolation to make sense as a criterion for a model to be acceptable is very demanding. Does such an ultimate description just exists? Does quantum field theory—the framework we currently use to describe particles and interactions—remain valid up to arbitrary high energies?

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[2]Useful reviews on effective field theories include Refs. [8].

There are no definite answers to those questions. Humbler paths are however available for improving our description of phenomena.

Effective approaches characterized by a range of validity established as limited from the onset are pragmatic methods making the relevant physics transparent and calculations easier, if not just possible at all. Limiting our description to a specific range does, on the other hand, not restrict our ability of comparison with experimental data. An effective field theory can be built to describe particle physics in this phenomenological way, within a field theory. Fortunately enough, we are not embarked on this construction without any guide and believe this method to be very general.

Symmetries

Our guides are symmetries.[3] They are defined as transformations like translations in space or time—just to mention the simplest ones—that leave our description of phenomena unchanged.

The fundamental role they play was only recognized during the twentieth century. However, it is actually difficult to imagine what science would be about without the symmetries of nature. No generality could for instance be extracted from the sum of particular events if experiments carried out at different places and times were to yield distinct conclusions. Our understanding of nature is indeed all about regularities: independently of initial conditions, physicists would like to encode in laws the way systems evolve. Specific quantities like the energy or charges are sometimes conserved during this evolution. This remarkable feature can be traced back to the presence of (continuous) symmetries in the theory. That is the content of Emmy Noether's theorem [11], dating from 1918.

The major shift in the importance given to symmetries occurred with the advent, first, of the theories of special and general relativity, then of quantum mechanics and ultimately of the standard model of particle physics built from local gauge symmetries, within a quantum field theory. As more and more sorts of symmetries were progressively singled out during the twentieth century, they were no longer seen as mere consequences of dynamical laws but rather considered as primary ingredients of our understanding of nature. The variety of dynamical laws appeared to be very much constrained—if not fully determined—by the assumption of their conservation.

Physicists were therefore no longer primarily concerned with abstracting the regularity of dynamical laws from particular events produced with different initial conditions. The challenge became to unravel the symmetries underlying and determining dynamical laws. In this endeavour, difficulties arise from the fact that symmetries are often slightly broken or hidden in ordinary conditions. Making them manifest may require going to so-far inaccessible energies or temperatures, probing phenomena at smaller and smaller distance scales.

In describing new physics, the effective-field-theory approach makes an extensive use of the tight constraints imposed by symmetries. It assumes the

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[3]The ideas presented here about *The role of symmetries in fundamental physics* are inspired from Refs. [9, 10].

ones that are established at known energy scales are still present just beyond that regime. Within a field-theory framework, an effective field theory actually includes nothing less than all interactions between known particles that are compatible with those assumed symmetries.

One could hardly imagine a more general set-up. It is built upon the elementary principles of quantum field theory together with what we think are the fundamental ingredients of our physical description of nature: symmetries. As Steven Weinberg [12], physicists indeed tend to believe that *although individual quantum field theories have of course a good deal of content, quantum field theory itself has no content beyond analyticity, unitarity, cluster decomposition, and symmetry.* Therefore, *if one writes down the most general possible Lagrangian, including* all *terms consistent with assumed symmetry principles, and then calculate matrix elements with this Lagrangian to any given order in perturbation theory, the result will simply be the most general possible S-matrix consistent with analyticity, perturbative unitarity, cluster decomposition and the assumed symmetry.*

Low energies

The assumed symmetries drastically reduce the range of possibilities but still in principle allow for an infinite number of local interactions between particles. If they are all relevant, not much can be said. However, we do not aim at describing physics at arbitrary high energies. Modestly and pragmatically, we would just like to go slightly beyond what we know already.

In this range, the symmetries observed at low energies should still be present. Well beyond that point they may be embedded in a larger set. New states may also appear at higher energies. However, the effects of heavy particles and the higher-energy modes of light ones can be modelled by local interactions between light particles in the low-energy limit. This is ultimately a consequence of the uncertainty principle of quantum mechanics and the basis of *renormalization*.

Let us imagine some state of energy (or momentum) Λ much larger that the energy E at which we aim to describe phenomena. This may be a new particle with mass of order Λ or the high-energy excitation of a known and much lighter particle. Quantum mechanics allows the classically forbidden production of that state for a time lapse (or propagating over a distance) of order $1/\Lambda$. Such space-time distances are much smaller than the ones of order $1/E$ probed in processes of energy E. So, all the information about energy scales much higher than the one we are practically interested in can be encoded in local interactions between light particles.

Even if an infinite number of local interactions is allowed, only a finite and definite set is required to be taken into account in order to obtain predictions of a given accuracy. Elementary dimensional analysis is sufficient to understand this fact.

Power counting

Interactions described by operators O_i involving more and more fields and their derivatives will have higher and higher dimensions d_i. As the Lagrangian

defining our effective field theory has dimension four (the dimension of space-time), the introduction of a dimensionful quantity is needed. This quantity is in principle arbitrary but let us choose it to be the scale Λ (much higher than the one at which we wish to describe phenomena) that defines the threshold between light states explicitly considered and heavy states whose effects are effectively taken into account through local interactions between light fields. We can then expect the remaining dimensionless coefficients C_i to be of order one since each operator can in principle receive contributions from physics of characteristic scale Λ. So we write:

$$\mathcal{L}(\phi) = \sum_i C_i \, \Lambda^{4-d_i} \, O_i(\phi) \tag{1}$$

where ϕ collectively represents the light fields of the theory.

Operators of dimensionality d_i much larger than four will therefore be associated with a high negative power of Λ. We can then expect their contributions at energy E to a dimensionless observable to scale roughly as $(E/\Lambda)^{d_i-4}$. For $E \ll \Lambda$ and $d_i > 4$, such contributions are thus small.

This power counting procedure establishes a hierarchy between operators with the lowest-dimensional ones being in principle the most relevant. Therefore, in many cases, we need not worry about an infinite number of interactions and can often rather restrict ourselves to the lowest-dimensional ones. If Λ is known, the dimension d_i of operators that should be taken into account to achieve a given accuracy for the effective-field-theory predictions can actually be determined.

Renormalizability

The above power suppression does not affect operators of dimension smaller or equal to four. Theories restricted to such operators have thus enjoyed a preferred status during the early days of quantum field theory. The above argument shows that they provide a first-order approximation in the low-energy limit, $E \ll \Lambda$.

These theories were qualified as *renormalizable*. Let us explain the origin of this term. As we already mentioned, quantum effects allow classically forbidden states of energy higher than the energy of the process considered to be produced for limited durations. Quantum corrections therefore in principle include contributions from states of arbitrary high energies. Those make many of them formally infinite. The standard means of making sense of those infinite quantum corrections is to introduce what is called a *regulator*. This is an artificial modification of the high-energy behaviour of the theory that makes the answer of quantum calculations finite. The simplest way of regularizing a theory is not to take into account states of energies higher than some fixed cut-off scale which is most naturally chosen to be Λ. This regularization scheme is not necessarily the most convenient one in practice but it makes the conceptual discussion more transparent.

As explained before, the effects of this modification of the high-energy behaviour of our theory can be modelled by local operators. In case those operators are already present in the original theory, a modification of their coefficients

could fully compensate for the introduction of the regulator. So, in practice, we can just calculate observables in the modified theory and correct for the error that we have made in doing so by fixing the value of operator coefficients so as to reproduce the physically observed results at a given scale. The latter procedure was called renormalization and the running of the theory parameters when Λ is continuously decreased is called *renormalization group* flow and was formulated in these terms by Kenneth Wilson [13].

In theories without operators of dimension larger than four, only a finite set of coefficients needs to be renormalized for making predictions including an arbitrarily high level of quantum corrections (at each order, a new renormalization should however be carried out). This is the reason they were called renormalizable. Because they considered regulators as spurious, pioneers of field theory wanted the cut-off scale Λ to be pushed to infinity. This was leading to very puzzling situations were the renormalization of operator coefficients reproducing physical results was infinite. However we do not know all the states that appear in nature at arbitrarily high energies and we do not even know if a field theory can possibly be used to describe phenomena of arbitrarily high energies. So, it is much more physically sensible to keep the cut-off finite: it truly represents the limits of our physical knowledge.

On the other hand, the computation of quantum corrections involving interactions of higher and higher dimensions would require the introduction of more and more renormalized coefficients. An infinite number of measurements would therefore be needed to perform this infinite number of renormalizations and any predictivity would seem to be lost. Theories containing operators of dimension larger than four were therefore qualified as *non-renormalizable* and considered as pathological. However, we need only make predictions for observable with a finite accuracy. Therefore, the introduction of an arbitrary number of interactions of dimension higher than four is never required in practice. The power counting argument of the previous section showed only a limited number of insertions of a limited number of operators is necessary. Theories involving operators of dimension higher than four are actually renormalizable order by order in $1/\Lambda$.

Gravity is a remarkable example of effective field theory where quantum corrections make perfect sense (see for instance Ref. [14]). The relevant degrees of freedom, perturbations of the metric (the flat Minkowski one $\eta_{\mu\nu}$ if the cosmological constant is taken vanishing), are encoded in a spin-two tensor $h_{\mu\nu}$ whose interactions compatible with general covariance are of the form

$$(\partial_\mu h_{\alpha\beta})(\partial_\nu h^{\beta\alpha})h^{\mu\nu}, \qquad \text{or} \qquad (\partial_\alpha h_{\mu\nu})(\partial_\beta h^{\mu\alpha})h^{\nu\beta}, \qquad \text{or...}$$

and have at least dimension five. The characteristic scale associated with those effects is the Planck mass $M_P \simeq 10^{19}$ GeV related to the Newton gravitational constant by $\sqrt{\hbar c / G_N}$. The immeasurable size of this scale compared with laboratory energies makes gravitational effects entirely negligible in particle physics experiments.

Naturalness

We have argued above that operators of dimension higher than four cause no problem in a quantum field theory. Let us here briefly comment on the operators of dimensions smaller than four, qualified as *super-renormalizable*.

As mentioned before, if a hard cut-off is used as regulator, all local operators in principle receive contributions from (known or unknown) physics at scale Λ and beyond. We therefore expect all dimensionful quantities appearing in the low-energy Lagrangian to be of order Λ. This would in particular hold true for mass terms. Though, this is a kind of contradiction because particles with masses of order of the cut-off should not be included in the low-energy theory since their effects had better be modelled by local interactions amongst lighter fields.

This line of reasoning indicates it should not be possible to write explicit mass terms directly in Lagrangians. There should be symmetries forbidding them. As those symmetries could be somehow broken, this does not mean that no particle much lighter than the threshold of validity of theory can be present. The chiral symmetry forbidding fermion masses is explicitly broken by fermion mass terms themselves and the gauge symmetry forbidding vector boson masses can be *spontaneously* broken in the vacuum.

Scalar masses, on the other hand, are more troublesome. In the standard model of particles physics, no symmetry forbids them. The Higgs boson mass is said to be *unnatural*: one would not expect it to be far from a scale Λ at which new physics appears. Yet, this particle has been observed with a mass of about 125 GeV and nothing else has so far been found around that scale.

Another operator of dimension lower than four is allowed by all symmetries of the standard model: the unit operator. When coupling matter with gravity, this operator gives rise to a cosmological constant. From the argument above, its coefficient is naturally expected to be of order of Λ^4 that we would at least expect to be around $(1 \text{ TeV})^4$. Though, cosmology provides us with a much tinier observed value of order $(2.2 \text{ meV})^4$ [15].

Those two naturalness puzzles are not *sensu stricto* inconsistencies of the standard model but still worry theorists.

Dimensional regularization

The introduction of a cut-off was very useful in the previous conceptual discussion. For performing actual computations however, it has some disadvantages. Gauge and Lorentz invariances are for instance no longer manifest at each step of the computation and are only recovered in the final result once all allowed contributions have been taken into account.

Dimensional regularization is another widely used regulator. Although less intuitive, it makes calculations easier and keeps gauge and Lorentz invariances explicit. It proceeds by considering the number of space-time dimensions as a continuous variable D. An expansion around $D = 4$ can then be performed. While not obvious at first sight, changing the number of space-time dimensions actually modifies the high-energy behaviour of a theory. This modification can then be absorbed in local operators, as it is with a cut-off regulator. The

reasonings made above therefore remain equally applicable with dimensional regularization.

Howard Georgi [8] put forward the following argument showing how a variation of the space-time dimensions modifies a quantum field theory at high energies. The typical integrals over high-energy states that appear in quantum corrections are of the form (in Euclidean space):

$$\int \frac{d^4l}{(2\pi)^4} \frac{1}{(l^2 + A^2)^\alpha}$$

with α an integer, l the energy-momentum of a virtual state, and A a function of external particles momenta and masses. Let us now change the number of dimensions to $D = 4 + \delta$,

$$\int \frac{d^4l}{(2\pi)^4} \int \frac{d^\delta l_\delta}{(2\pi\mu)^\delta} \frac{1}{(l_\delta^2 + l^2 + A^2)^\alpha}$$

and introduce a dimensionful quantity μ in order to preserve the dimensionality of the integral. This so-called *renormalization scale* is to be taken close to the energy of the process considered for the perturbation theory to be well behaved. Usually the integrals over the D dimensions are performed all at once but let us here carry out first the integral over the extra δ dimensions. The usual techniques then yield

$$\int \frac{d^4l}{(2\pi)^4} \frac{1}{(l^2 + A^2)^\alpha} \quad \frac{\Gamma(\alpha - \delta/2)}{\Gamma(\alpha)} \quad \left(\frac{l^2 + A^2}{4\pi\mu^2}\right)^{\delta/2}. \tag{2}$$

This expression shows that for δ small and l^2 of the order of A^2 and μ^2, the physics has not changed. In that regime, indeed, the ratio of gamma functions is close to one, as is the last factor. However, for l^2 much larger than A^2 and μ^2 (more precisely, for $\frac{1}{2} \ln \frac{l^2 + A^2}{4\pi\mu^2} \gg \frac{1}{|\delta|}$), the integrand is very different from the $D = 4$ one. A small variation of the dimension of space-time actually changed the high-energy behaviour of the theory.

Another advantage of dimensional regularization is that it does not introduce any unphysical dimensionful quantity but μ which only appears inside logarithms (this can be seen by expanding Eq. (2) for δ tending to zero). This regularization scheme, by itself, does therefore not mix operators of different dimensions. For coefficients C_i approximately of order one, Λ is then no longer an artificially chosen cut-off but a combination of truly physical scales that therefore contains some information about the characteristic energies of the physics lying above the range directly probed.

Effective field theory for new physics

Particle physics up to scales of the order of the TeV is well described by a renormalizable quantum field theory, the standard model. As we have argued in this section, the addition of extra non-renormalizable operators would make perfect sense, even when considering quantum corrections. We would like to use

this effective-field-theory approach to parametrize new physics or to quantify precisely its absence.

The non-renormalizable interactions may model the low-energy effects of an extended renormalized theory featuring new heavy particles, but not necessarily. It could also account for some non-perturbative effects due to a new strongly coupled dynamics or intrinsically non-perturbative physics like for instance *instantons* [16] that arise from the topological structure of the vacuum. Like for gravity at energies lower than the Planck scale, it may also be that the only way of modelling yet unknown physics in a quantum field theory is through higher-dimensional operators. So, describing new phenomena effectively may possibly be rather inescapable. The prejudice that the new physics we are looking for can only take the form of a renormalizable quantum field theory may not be completely justified.

The use of an effective theory to describe new physics however rests on the basic assumption that its characteristic scale is somewhat higher than the energies probed at the experiment considered. Otherwise would the order by order treatment in E/Λ be inappropriate. When considering extensions of the standard model, the fact that no new phenomena has so far been conclusively observed may possibly indicate they are suppressed by powers of E/Λ and suitably modelled by higher-dimensional operators.

The way to proceed for constructing an effective field theory able to model new-physics effects is then the following:

1. establish first what fields and what symmetries will be used,
2. out of these fields, construct all possible operators preserving the assumed symmetries,
3. eliminate all operators that have the same physical consequences as combinations of other ones and form an independent basis,
4. for a specific application, determine what are the lowest-dimensional operators contributing when a given order of quantum corrections is considered.

If we indeed face a clear separation between the electroweak and the new-physics scales, it is then also justified to build our effective theory extending the renormalizable standard model, out of the fields we know and to impose the symmetries currently well established.

Experimentally, the above developments lead to a distinction between two options available to discover beyond-the-standard-model phenomena:

– either work at low energies and make very high precision measurements in order to overcome the significant power suppression $(E/\Lambda)^n$,
– or go to higher energies, where the ratio E/Λ is larger and putative new physics effects therefore enhanced.

The LHC has already probed and will further explore this high energy frontier.

I.2 Standard-model particles and symmetries

Before addressing the construction of standard-model operators (of dimensions lower as well as higher than four), let us describe both the standard-model matter content and the symmetries established over the time by experimental tests. A special attention will be devoted to the top quark.

The elementary particles of the standard model and their interactions are described in the context of a quantum field theory based on several symmetry principles of different natures. The first one is a symmetry of space-time under translations (in space or time), rotations, and *boosts* which are simultaneous transformations of both space and time. The matter fields transform in spinorial representations of the Lorentz group $SO(3,1) \supset SU(2) \times SU(2)$. With *spin* $1/2$, they are named *fermions* and can have either *left-* or *right-handed chiralities* depending on the $SU(2)$ factor under which they are transforming as doublets.

Next, gauge symmetries are of prime relevance. When imposed on the dynamics of matter fermions, they require the presence of new particles whose own transformation properties compensate for the fermions ones. Those particles are *vector bosons* of the Lorentz group, with spin 1. The standard-model gauge symmetries are unitary rotations of matter fields that vary from space-time point to space-time point in an internal space. While the matter fermions transform in the fundamental representations of the corresponding unitary groups, the gauge vector bosons are in the adjoint. Three gauge groups have been isolated. A local $SU(3)_c$ symmetry gives rise to the strong interaction (*quantum chromodynamics*, or QCD) between coloured fermions called *quarks*. *Leptons* are on the other hand colourless. Eight vector bosons called *gluons* mediate this interaction. A $SU(2)_L$ symmetry group is associated with the weak interaction carried by a triplet of W bosons. Remarkably, only left-handed fermions transform, non-trivially, as $SU(2)_L$ doublets (hence the $_L$ subscript). A last, a $U(1)_Y$ gauge symmetry shifts all standard-model fermions by a space-time dependent phase proportional to their respective *hyper*charge. A B boson is the corresponding mediator. The standard-model gauge symmetry is therefore the direct product: $SU(3)_c \times SU(2)_L \times U(1)_Y$.

Some global internal symmetries can also be present. They are not imposed from the onset in the construction of interactions but rather arise as fortuitous consequences of the other symmetry requirements and specific matter content of the standard model. They are thus sometimes said to be *accidental* and often conserved in but particular sets of interactions. As will be detailed below, two $U(1)_B$ and $U(1)_L$ groups are global (and only classical) symmetries of all standard-model operators of dimension four, at most. The first one shifts all quarks by a common global phase while the second one does so for leptons. According to Noether's theorem, they are associated with the conservation of two charges, the *baryon* and *lepton* numbers. *Flavour* symmetries that leave invariant the quadratic interactions between fermions will also be of particular importance in what follows.

As a consequence of the chiral character of fermions and of the gauge symmetries, all standard-model particles should be absolutely massless. The $SU(2)_L \times U(1)_Y$ electroweak gauge symmetry is however not realized in the

vacuum. There, only its well-known electromagnetic $U(1)_{EM}$ subgroup actually survives. The minimal mechanism leading to this specific *spontaneous symmetry breaking* requires the presence of a Lorentz scalar, weak doublet, ϕ whose electromagnetically neutral component acquires, in the vacuum, a non-vanishing expectation value $v/\sqrt{2} \simeq 174$ GeV that breaks the electroweak symmetry. This is the Brout-Englert-Higgs mechanism [2]. Through their interactions with the Higgs field ϕ, three combinations of the original W and B bosons then obtain a mass. They are the $W^{\pm}$ and Z bosons. The fourth degree of freedom corresponds to the photon γ that remains massless and mediates the electromagnetic interaction by coupling to particles with a strength proportional to their respective electric charge. Three of the four original degrees of freedom of ϕ have become the longitudinal polarizations of the $W^{\pm}$ and Z bosons. The last electrically neutral one is the physical Higgs particle h which was presumably discovered at the LHC, as announced on the fourth of July 2012 [1]. The exact nature of the resonance observed with a mass of 125 GeV still remains to be determined precisely.

All fermionic degrees of freedom also become massive through their interactions with ϕ. As will be seen below, the operator relevant in the case of *neutrinos*, the electrically neutral components of the lepton doublets, is of dimension five. The resulting massive states that actually propagate in the vacuum do however not correspond to the ones having definite chiral and $SU(2)_L \times U(1)_Y$ transformation properties. Two bases are therefore distinguished: the gauge and mass (or physical) eigenstate bases. Three generations of physical fermions differing only by their masses are observed. Up-type quarks have an electric charge of $+2/3$ (in units of the absolute electron one), down-type quarks have charge $-1/3$, neutrinos are neutral and charged leptons have the electromagnetic charge of the electron. These mass eigenstates receive special *flavour* names:

	First,	second,	third	generations.
Up-type quarks	*up* $u_L + u_R$	*charm* $c_L + c_R$	*top* $t_L + t_R$	
Down-type quarks	*down* $d_L + d_R$	*strange* $s_L + s_R$	*bottom* $b_L + b_R$	
Neutral leptons or *neutrinos*	ν_{eL}	$\nu_{\mu L}$	$\nu_{\tau L}$	
Charged leptons	*electron* $e_L + e_R$	*muon* $\mu_L + \mu_R$	*tau* $\tau_L + \tau_R$	

A $_L$ or $_R$ subscript distinguishes their left- or right-handed components which can be obtained using the $P_{L,R}$ projectors. They are admixtures of the three left- and right-handed gauge eigenstates, respectively. Those will be denoted as:

		$Y/2$
q	for the three left-handed quark doublets, with hypercharge	$+1/6$
u	for the up-type right-handed quark singlets,	$+2/3$
d	for the down-type right-handed quark singlets,	$-1/3$
l	for the three left-handed lepton doublets,	$-1/2$
e	for the right-handed charged lepton singlets,	-1

For clarity, generation indices a, b, c, d,... $\in \{1, 2, 3 = N_g\}$, weak $SU(2)_L$ indices i, j, k, l,... $\in \{1, 2\}$, colour $SU(3)_c$ indices α, β, γ,... $\in \{1, 2, 3 = N_c\}$ and spinorial ones p, q, r, s,... will be dropped when unnecessary.

The top quark   As the heaviest known fundamental particle, the top quark enjoys a special status. Awaited since the postulate of three generations of quarks by Kobayashi and Maskawa in 1973 [17], it was finally discovered in 1995 at the Tevatron experiment [18] with a mass of about 173 GeV, much higher than first expected. The previous discovery of a fermion, the b quark which is about 40 times lighter, had occurred already 18 years beforehand [19].

Unlike all other quarks, the top high mass makes its partial weak decay width to a b and an on-shell W larger than the QCD hadronization scale. It therefore decays before being able to form mesons or baryons.

This very high mass of the top indicates a strong coupling to the Higgs scalar field. The top quark is therefore of primary relevance to the study of the electroweak symmetry breaking mechanism. It is actually the only fermion with a coupling ($m_t \sqrt{2}/v \simeq 1$) that is not strikingly small.

Experimentally, by producing millions of top quark pairs, the LHC is improving the measurements of its properties at astonishing rates. The latest combination of Tevatron and LHC measurements [20] provides a top quark mass value of:

$$m_t = 173.34 \pm 0.76 \text{ GeV}.$$

Its pair production cross section has been computed with a remarkable accuracy of about 5% [21] and measured at the 7 and 8 TeV LHC to be approximately 162 and 240 pb with 7% and 5% uncertainties [22, 23]. The top quark width is predicted with an accuracy (finite m_b, NLO EW and NNLO QCD [24]) exceeding by far the experimental one. The best direct measurement [25],

$$1.10 < \Gamma_t < 4.05 \text{ GeV} \qquad \text{(or } \Gamma_t < 6.38 \text{ GeV at the 95\% CL)},$$

is compatible with a prediction of approximately 1.4 GeV and corresponds to a lifetime of about 10^{-25} s. Angular distributions in the $t \to bW^+ \to b\ell^+\nu_\ell$ decay are also predicted with a good accuracy [26]. The W helicity fraction measurement [27] are notably still an order of magnitude less precise than its theory prediction. Various searches for exotic top interactions have also been performed (see [28, 29]).

I.3 Standard-model interactions

With the fields and symmetries just described, the operators responsible for the interactions between standard-model particles can now be constructed. As stressed before, both *renormalizable* and *non-renormalizable* operators are built from the same requirements. We will concentrate on interactions involving fermions, and the top quark in particular, that will be relevant for the specific applications considered later (a complete basis of dimension-six operators can be found in Ref. [30]). In the process of this construction, some aspects of the standard model will be discussed, flavour structures in particular.

The dimensions of operators set the hierarchy amongst them but, first, symmetries constrain the whole edifice.

Lorentz invariance   In the construction of Lorentz-invariant operators, angular momentum conservation restricts the number of fermions involved to be even. Each of them, the top quark in particular, should therefore appear paired. As all standard-model fermions are either quarks, carrying a baryon number B of $\pm 1/3$, or leptons with $L = \pm 1$. All operators should satisfy the selection rule $\Delta(3B + L) \in 2\mathbb{Z}$ which trivially holds true for $\Delta B = 0 = \Delta L$ operators.

The Lorentz transformation properties of a pair of chiral fermions are determined by group theory. Under $SU(2) \times SU(2) \subset SO(3,1)$, left-handed fermions transform as $(\mathbf{2}, \mathbf{1})$ and right-handed ones as $(\mathbf{1}, \mathbf{2})$. Therefore, when combining two fermions of identical chiralities,

$$(\mathbf{2}, \mathbf{1}) \otimes (\mathbf{2}, \mathbf{1}) = (\mathbf{1}, \mathbf{1}) \oplus (\mathbf{3}, \mathbf{1}),$$
$$(\mathbf{1}, \mathbf{2}) \otimes (\mathbf{1}, \mathbf{2}) = (\mathbf{1}, \mathbf{1}) \oplus (\mathbf{1}, \mathbf{3}),$$

we can form a scalar $(\mathbf{1}, \mathbf{1})$ or the chiral half of an antisymmetric tensor $(\mathbf{1}, \mathbf{3}) \oplus (\mathbf{3}, \mathbf{1})$. On the other hand, combining two fermions of opposite chiralities,

$$(\mathbf{1}, \mathbf{2}) \otimes (\mathbf{2}, \mathbf{1}) = (\mathbf{2}, \mathbf{2}),$$

leads to a Lorentz vector. Furthermore, the combination of two vectors:

$$(\mathbf{2}, \mathbf{2}) \otimes (\mathbf{2}, \mathbf{2}) = (\mathbf{1}, \mathbf{1}) \oplus (\mathbf{3}, \mathbf{1}) \oplus (\mathbf{1}, \mathbf{3}),$$

or the combination of anti-symmetric tensors of identical chirality:

$$(\mathbf{3}, \mathbf{1}) \otimes (\mathbf{3}, \mathbf{1}) = (\mathbf{1}, \mathbf{1}) \oplus (\mathbf{3}, \mathbf{1}) \oplus (\mathbf{5}, \mathbf{1}),$$
$$(\mathbf{1}, \mathbf{3}) \otimes (\mathbf{1}, \mathbf{3}) = (\mathbf{1}, \mathbf{1}) \oplus (\mathbf{1}, \mathbf{3}) \oplus (\mathbf{1}, \mathbf{5}),$$

each contain a scalar while antisymmetric tensors of opposite chirality do not:

$$(\mathbf{3}, \mathbf{1}) \otimes (\mathbf{1}, \mathbf{3}) = (\mathbf{3}, \mathbf{3}).$$

The basis of Dirac matrices $\Gamma \in \{P_{L,R}, \gamma^\mu P_{L,R}, \sigma^{\mu\nu} P_{L,R}\}$[4] allows to form fermions bilinears of definite Lorentz transformation properties: scalars, vectors and tensors. Introducing the charge conjugation matrix C defined to satisfy

$$C^\dagger = C^{-1}, \qquad C^T = -C, \qquad C\gamma^{\mu T} C^\dagger = -\gamma^\mu,$$

the charge (CP, actually) conjugate of a fermion can be written $\psi^c \equiv C\bar{\psi}^T$, so that $(\psi^c)^c = \psi$ and $\overline{\psi^c} = -\psi^T C^\dagger$. From the definition of $\sigma^{\mu\nu} \equiv \frac{i}{2}[\gamma^\mu, \gamma^\nu]$ and of $P_{L,R} \equiv (\mathbb{I} \mp \gamma^5)/2$ in terms of $\gamma^5 \equiv i\gamma^0\gamma^1\gamma^2\gamma^3$, the transposes of Γ matrices are obtained to be

$$(C^\dagger P_{L,R})^T = -C^\dagger P_{L,R},$$
$$(C^\dagger \gamma^\mu P_{L,R})^T = +C^\dagger \gamma^\mu P_{L,R}, \tag{3}$$
$$(C^\dagger \sigma^{\mu\nu} P_{L,R})^T = +C^\dagger \sigma^{\mu\nu} P_{L,R}.$$

---

[4]As $\sigma^{\mu\nu}\gamma^5 = \frac{-i}{2}\epsilon^{\mu\nu\rho\sigma}\sigma_{\rho\sigma}$, one single $\sigma^{\mu\nu}$ tensor would actually suffice, but the two chiral ones are independent.

In Lorentz invariant combinations of two bilinears, the Fierz[5] identities can be used to exchange two fields from different bilinears:

$$
\begin{aligned}
4\,[P_R]^p{}_q[P_R]^r{}_s &= 2\,[P_R]^p{}_s[P_R]^r{}_q + 1/2\,[\sigma^{\mu\nu}P_R]^p{}_s[\sigma_{\mu\nu}P_R]^r{}_q, \\
[\sigma^{\mu\nu}P_R]^p{}_q[\sigma_{\mu\nu}P_R]^r{}_s &= 6\,[P_R]^p{}_s[P_R]^r{}_q - 1/2\,[\sigma^{\mu\nu}P_R]^p{}_s[\sigma_{\mu\nu}P_R]^r{}_q, \\
[\sigma^{\mu\nu}P_R]^p{}_q[\sigma_{\mu\nu}P_L]^r{}_s &= 0, \\
2\,[P_R]^p{}_q[P_L]^r{}_s &= [\gamma^\mu P_L]^p{}_s[\gamma_\mu P_R]^r{}_q, \\
[\gamma^\mu P_R]^p{}_q[\gamma_\mu P_R]^r{}_s &= -[\gamma^\mu P_R]^p{}_s[\gamma_\mu P_R]^r{}_q.
\end{aligned}
\tag{4}
$$

The following identities can also be used to reshuffle fields in scalar and tensor quadrilinears:

$$
\begin{aligned}
[C^\dagger P_R]_{pq}[C^\dagger P_R]_{rs} - [C^\dagger P_R]_{pr}[C^\dagger P_R]_{qs} + [C^\dagger P_R]_{ps}[C^\dagger P_R]_{qr} &= 0, \\
[C^\dagger \sigma^{\mu\nu} P_R]_{pq}[C^\dagger \sigma_{\mu\nu} P_R]_{rs} + [C^\dagger \sigma^{\mu\nu} P_R]_{pr}[C^\dagger \sigma_{\mu\nu} P_R]_{qs} & \\
+ [C^\dagger \sigma^{\mu\nu} P_R]_{ps}[C^\dagger \sigma_{\mu\nu} P_R]_{qr} &= 0.
\end{aligned}
\tag{5}
$$

They are named after Jan Arnoldus Schouten. As Fierz ones, they are obviously also valid for P_L and P_R interchanged.

Colour conservation The elementary invariant tensors of $SU(3)$ being δ^α_β and $\epsilon_{\alpha\beta\gamma}$ ($\equiv \tilde{\epsilon}$), quarks like the top are required to appear as quark–anti-quark pairs $\delta^\alpha_\beta\,\bar{Q}_\alpha Q^\beta$ or as antisymmetric combinations of three quarks (or anti-quarks) $\epsilon_{\alpha\beta\gamma}\,Q^\alpha Q^\beta Q^\gamma$ ($\equiv \tilde{\epsilon}\,QQQ$).

An important distinction between those two elementary colour-invariant structures is that the former carries no net baryon number while the latter has $B = \pm 1$. Moreover, every odd number of quark triads should always come associated with an odd number of (anti-)leptons and reciprocally, since the total number of fermions should remain even.

The minimal combination of fields violating the lepton number only is therefore built upon two (anti-)leptons and leads to $|\Delta L| = 2$ while operators violating baryon number contain at least three (anti-)quarks and one (anti-)lepton and generate $|\Delta B| = 1 = |\Delta L|$ processes. Lepton- and baryon-number-violating operators should thus at least have dimension three and six, respectively.

Out of two triplets, the generators of $SU(3)$, the Gell-Mann matrices $[\lambda^A]^\alpha{}_\beta$ can make an octet. They moreover satisfy

$$
[\lambda^A]^\alpha{}_\beta[\lambda^A]^\gamma{}_\delta = T_c\,(\delta^\alpha_\delta \delta^\gamma_\beta - \frac{1}{N_c}\delta^\alpha_\beta \delta^\gamma_\delta)
\tag{6}
$$

with $T_c = 2$ fixing their normalisation $\mathrm{Tr}\{\lambda^A \lambda^B\} = T_c\,\delta^{AB}$ and $N_c = 3$ the number of colours. By analogy with the Lorentz case presented above, this equality is usually also named after Markus Fierz.

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[5]Markus Fierz is credited with the first occurrence those equalities, although they could be due to Wolfgang Pauli who Fierz himself refers to [31].

$SU(2)_L$ conservation   In $SU(2)$ an anti-doublet is changed into a doublet using the antisymmetric tensor $\epsilon \equiv \epsilon^{ij} = \epsilon_{ij} = \left(\begin{smallmatrix} 0 & 1 \\ -1 & 0 \end{smallmatrix}\right)$ and *vice versa*. A δ^i_j tensor can then be used to form singlets out of doublet–anti-doublet pairs. Similarly to the $SU(3)$ case, the generators of $SU(2)$, the Pauli matrices,

$$[\tau^I]^i{}_j = \left(\left(\begin{array}{cc} 0 & 1 \\ 1 & 0 \end{array} \right), \left(\begin{array}{cc} 0 & -i \\ i & 0 \end{array} \right), \left(\begin{array}{cc} 1 & 0 \\ 0 & -1 \end{array} \right) \right),$$

can also serve to form triplets out of pairs of doublets (note $\epsilon = i\tau^2$). Two triplets can then be combined into a singlet which is actually the one obtained from the following combinations of δ's:

$$[\tau^I]^i{}_j [\tau^I]^k{}_l = T_L \left(\delta^i_l \delta^k_j - \frac{1}{N_L} \delta^i_j \delta^k_l \right) \tag{7}$$

with $N_L = 2$ the span of i, j, k, l indices, and $T_L = 2$ defined through $\mathrm{Tr}\left\{ \tau^I \tau^J \right\} = T_L \delta^{IJ}$.

The $SU(2)_L$ equivalent of Schouten identities will also be useful below:

$$\epsilon_{ij}\epsilon_{kl} - \epsilon_{ik}\epsilon_{jl} + \epsilon_{il}\epsilon_{jk} = 0. \tag{8}$$

It arises as the combination on the left-hand side is a completely antisymmetric tensor with four indices taking only two different values and must therefore vanish.[6]

Hypercharge conservation   is the last symmetry requirement but there is no mystery about it as it is simply additive.

Order by order construction

Let us now proceed with the explicit and systematic construction of standard-model interactions. In order to impose the full standard-model gauge symmetry, the q, u, d, l and e fermionic gauge eigenstates will be used. Since only physical fields are however observed, the interactions constructed here would require to be rotated to the mass eigenstate basis for all practical purposes. As noted above, operators containing a top quark involve at least a fermion pair and therefore have at least dimension three.

Dimension three

Leaving baryon- and lepton-number-violating operators aside for the moment, we consider first lepton–anti-lepton or quark–anti-quark pairs. They have the following Lorentz and gauge transformation properties:

---

[6]A simple proof quoted in Ref. [32] as originating from Burt Ovrut.

	Lorentz	$SU(3)_c$	$SU(2)_L$	$U(1)_{Y/2}$
$q^c q$	$(\mathbf{2},\mathbf{2})$	$\mathbf{1}$ or $\mathbf{8}$	$\mathbf{1}$ or $\mathbf{3}$	0
$u^c u$	$(\mathbf{2},\mathbf{2})$	$\mathbf{1}$ or $\mathbf{8}$	$\mathbf{1}$	0
$d^c d$	$(\mathbf{2},\mathbf{2})$	$\mathbf{1}$ or $\mathbf{8}$	$\mathbf{1}$	0
$q^c u$	$(\mathbf{1},\mathbf{1})$ or $(\mathbf{1},\mathbf{3})$	$\mathbf{1}$ or $\mathbf{8}$	$\bar{\mathbf{2}}$	$+1/2$
$q^c d$	$(\mathbf{1},\mathbf{1})$ or $(\mathbf{1},\mathbf{3})$	$\mathbf{1}$ or $\mathbf{8}$	$\bar{\mathbf{2}}$	$-1/2$
$u^c d$	$(\mathbf{2},\mathbf{2})$	$\mathbf{1}$ or $\mathbf{8}$	$\mathbf{1}$	-1
$l^c l$	$(\mathbf{2},\mathbf{2})$	$\mathbf{1}$	$\mathbf{1}$ or $\mathbf{3}$	0
$e^c e$	$(\mathbf{2},\mathbf{2})$	$\mathbf{1}$	$\mathbf{1}$	0
$l^c e$	$(\mathbf{1},\mathbf{1})$ or $(\mathbf{1},\mathbf{3})$	$\mathbf{1}$	$\bar{\mathbf{2}}$	$-1/2$

Charge-conjugated bilinears should also be kept in mind—for those that are not self-conjugate.

None of above fermionic gauge eigenstate bilinears can form a Lorentz and gauge singlet. In particular, this implies that no explicit fermion mass term is allowed in the standard model.

At dimension three, two (anti-)leptons could also possibly form a $|\Delta L| = 2$ operator. No Lorentz and gauge singlet is however found amongst the three possible combinations:

	Lorentz	$SU(3)_c$	$SU(2)_L$	$U(1)_{Y/2}$
ll	$(\mathbf{1},\mathbf{1})$ or $(\mathbf{3},\mathbf{1})$	$\mathbf{1}$	$\mathbf{1}$ or $\mathbf{3}$	-1
ee	$(\mathbf{1},\mathbf{1})$ or $(\mathbf{1},\mathbf{3})$	$\mathbf{1}$	$\mathbf{1}$	-2
le	$(\mathbf{2},\mathbf{2})$	$\mathbf{1}$	$\mathbf{2}$	$-3/2$

Dimension four

Increasing the dimension by a minimal amount, the smallest addition to a fermion pair is a covariant derivative or a scalar boson. By construction, the covariant derivative acting on a field does not change its gauge transformation properties but is a Lorentz vector $(\mathbf{2},\mathbf{2})$. On the other hand, the standard model scalar—or Higgs—boson transforms as:

	Lorentz	$SU(3)_c$	$SU(2)_L$	$U(1)_{Y/2}$
ϕ	$(\mathbf{1},\mathbf{1})$	$\mathbf{1}$	$\mathbf{2}$	$+1/2$

With those two extra objects, at dimension four, we can form all the two-fermion operators of the standard-model Lagrangian. Those include fermion kinetic terms: $K^a{}_b\,\bar{\psi}_a i\slashed{D}\psi^b$, and Yukawa interactions: $Y^a{}_b\,\bar{\psi}_a\psi'^b\,\phi$. The operator coefficients K and Y carrying two generation indices are *a priori* arbitrary complex matrices. K is only required to be Hermitian in order for the Lagrangian to be real. The kinetic term governs the interactions of fermions with gauge fields as well as their propagation. The construction of other operators (including the Yukawa terms) is therefore best performed with the fields actually propagating in the $SU(3)_c \times SU(2)_L \times U(1)_Y$-invariant theory. They are obtained from a unitary rotation in generation space of the original fermions:

$\psi' \equiv U\psi$, for U diagonalizing $K = U^\dagger L U$. The kinetic terms then write $L_b^a \ \bar\psi'_a i\slashed{D}\psi'^b$ and are moreover customarily set in *canonical* forms $\delta_b^a \ \bar\psi''_a i\slashed{D}\psi''^b$ by rescaling each field to $\bar\psi''^a \equiv \sqrt{L_a^a}\psi'^a$ (no sum). Dropping all primes, we will use these canonically normalized fields in the following and redefine them to be the gauge eigenstates.

Any further global unitary transformation of the fermionic fields in generation space will leave invariant the obtained kinetic Lagrangian,

$$\bar q \, i\slashed{D} \, q, \qquad \bar u \, i\slashed{D} \, u, \qquad \bar d \, i\slashed{D} \, d, \qquad \bar l \, i\slashed{D} \, l, \qquad \bar e \, i\slashed{D} \, e,$$

which therefore possesses a large global *flavour* symmetry [33]:

$$U(N_g)^5 \equiv U(N_g)_q \times U(N_g)_u \times U(N_g)_d \times U(N_g)_l \times U(N_g)_e,$$

for $N_g = 3$, the number of fermionic generations. On the other hand, the Yukawa interactions:

$$Y_u \ \bar u q \ \epsilon\phi, \qquad Y_d \ \bar d q \ \phi^*, \qquad Y_e \ \bar e l \ \phi^*, \qquad +\text{h.c.}$$

that involve simultaneously two different types of fermions explicitly break this symmetry. Out of the five $U(N_g)$'s, all non-Abelian $SU(N_g)$ parts and three Abelian $U(1)$'s are broken. The two $U(1)$'s left (classically) invariant are the baryon-number $U(1)_B$ and the lepton-number $U(1)_L$ under which the quarks and leptons are respectively given a common global phase.

Those interactions were constructed with the canonically normalized fields that propagate in the $SU(3)_c \times SU(2)_L \times U(1)_Y$-invariant theory, the gauge eigenstates. However, the electroweak $SU(2)_L \times U(1)_Y$ is broken down in the vacuum to the electromagnetic $U(1)_{EM}$. The non-vanishing expectation value of the Higgs scalar doublet $(0, v/\sqrt{2})^T$ yields a second type of quadratic terms for fermions, masses:

$$\frac{v}{\sqrt{2}} \ Y_u \ \bar u q^{i=1}, \qquad \frac{v}{\sqrt{2}} \ Y_d \ \bar d q^{i=2}, \qquad \frac{v}{\sqrt{2}} \ Y_e \ \bar e l^{i=2}, \qquad +\text{h.c.}$$

In the vacuum, isolating the massive fields that propagate requires the singular values decomposition of the Yukawas $Y_{u,d,e}$ by bi-unitary transformations:

$$Y_u \equiv V_{u_R}^\dagger y_u V_{u_L}, \qquad Y_d \equiv V_{d_R}^\dagger y_d V_{d_L}, \qquad Y_e \equiv V_{e_R}^\dagger y_e V_{e_L},$$

The left- and right-handed parts of the physical eigenstates are then given by:

$$\begin{aligned} u_R &\equiv V_{u_R} u, & d_R &\equiv V_{d_R} d, & e_R &\equiv V_{e_R} e, \\ u_L &\equiv V_{u_L} q^{i=1}, & d_L &\equiv V_{d_L} q^{i=2}, & e_L &\equiv V_{e_L} l^{i=2}, \end{aligned}$$

and have definite masses proportional to the singular values of the Yukawa matrices $m_{u,d,e} \equiv y_{u,d,e} \, v/\sqrt{2}$.

In the leptonic sector, all effects of this transformation on the $SU(2)_L$ part of the kinetic terms can be absorbed in the definition of a $\nu_L \equiv V_{e_L} l^{i=1}$. So, the two components of the lepton doublet get rotated by the same unitary matrix. If we were to limit ourselves to operator of dimension four constructed

with standard-model fields, the absence of neutrino masses would eliminate any distinction between gauge and physical leptons and allow to make all interactions involving leptons diagonal in flavour. Extending the matter content with a completely neutral gauge eigenstate ν would permit the construction of a second Yukawa interaction involving the left-handed lepton doublet: $\tilde{Y}_\nu\, \bar{\nu} l\, \epsilon\phi$.[7] It would provide neutrinos with Dirac masses and the singular value decomposition of $\tilde{Y}_\nu = V^\dagger_{\nu_R} y_\nu \tilde{V}_{\nu_L}$ would fix the unitary rotation defining the mass eigenstates $\nu_L \equiv \tilde{V}_{\nu_L} l^{i=1}$. As $\tilde{V}_{\nu_L}$ is *a priori* independent of V_{e_L}, the weak interaction between leptonic physical eigenstates would then in general acquire a non-trivial flavour structure: $\overline{l^{i=1}}W^+ l^{i=2} = \overline{\nu_L}W^+ (\tilde{V}_{\nu_L} V^\dagger_{e_L}) e_L$. The resulting mixing matrix $V_{\text{PMNS}} \equiv V_{e_L} \tilde{V}^\dagger_{\nu_L}$ is named after Bruno Pontecorvo, Ziro Maki, Masami Nakagawa and Shoichi Sakata [34]. The flavour diagonality of the interactions amongst leptons is also lost, in general, with operators of dimension higher than four built from standard-model fields only (see below).

In the quark sector, on the contrary, the unitary rotations of the two components of q are fixed independently at dimension four already by the singular values decompositions of Y_u and Y_d. In the physical basis, the $SU(2)_L$ gauge interactions between quarks are therefore non-diagonal: $\overline{q^{i=1}}W^+ q^{i=2} = \overline{u_L}W^+ (V_{u_L} V^\dagger_{d_L}) d_L$. The mixing matrix $V_{\text{CKM}} \equiv V_{u_L} V^\dagger_{d_L}$ is named after Nicola Cabibbo, Makoto Kobayashi and Toshihide Maskawa [17, 35]. It contains the only observable manifestation of the rotations from gauge to mass eigenstates: the $V_{u_{L,R}}$ and $V_{d_{L,R}}$ matrices are not separately measurable.

In both the leptonic and quark sectors, only the *charged-current* interactions involving physical $W^\pm$ are *flavour changing*. In the standard-model, loop-level processes only can generate flavour-changing neutral currents. Because of the unitarity of mixing matrices, those are moreover proportional to potentially small fermion mass differences. This is the so-called GIM mechanism [36], named after Sheldon Lee Glashow, John Iliopoulos and Luciano Maiani.

Dimension five

In addition to a fermion pair, at dimension five, one could include either

1. two covariant derivatives contracted to form a Lorentz scalar, or anti-symmetrized to a field strength tensor,
2. one scalar boson and one covariant derivative,
3. two scalar bosons.

In the list of fermion bilinears presented on p.18, no scalar or tensor entry is a gauge singlet. The possibility 1. is therefore excluded. Similarly, since no vector bilinear has the transformation properties of a scalar boson, option 2. is also to be rejected. Finally, although no baryon- and lepton-number-conserving bilinear has the transformation properties of a pair of scalar bosons,

---

[7]A Majorana mass term $\overline{\nu^c}\nu$, of dimension three, could also be constructed but, for simplicity, will not be considered in the following discussion.

	Lorentz	$SU(3)_c$	$SU(2)_L$	$U(1)_{Y/2}$
$\phi\phi$	$(\mathbf{1}, \mathbf{1})$	$\mathbf{1}$	$\mathbf{3}$[8]	$+1$
$\phi^*\phi$	$(\mathbf{1}, \mathbf{1})$	$\mathbf{1}$	$\mathbf{1}$ or $\mathbf{3}$	0

the ll bilinear in its Lorentz singlet and $SU(2)_L$ triplet version does transform as $\phi\phi$. So we can form one single operator of dimension five: $(\overline{l^c}\epsilon\tau^I l)\,(\phi\epsilon\tau^I\phi)$ which violates the lepton number by two units. Using the Fierz identity of $SU(2)$ given in Eq. (7), this operator can also be rewritten in a more common form:

$$Y_\nu\;(\overline{l^c}\epsilon\phi)\,(l\epsilon\phi) \qquad + \text{h.c.}$$

It is symmetric under the exchange of the two leptons (fermions anti-commute and the Lorentz structure is anti-symmetric). The Y_ν matrix in generation space can therefore be chosen symmetric.

Interestingly, in the vacuum, this operator first constructed by Weinberg in 1979 [37] generates Majorana masses for (left-handed) neutrinos. As in the previous section, the unitary rotation from gauge $l^{i=1}$ to physical ν_L eigenstates is then determined by the singular values decomposition $Y_\nu \equiv V_{\nu_L}^T y_\nu V_{\nu_L}$. A non-trivial mixing matrix $V_{\text{PMNS}} \equiv V_{e_L} V_{\nu_L}^\dagger$ is therefore also generated by this dimension-five operator.

Dimension-six, four-fermion operators

At dimension six, many more possibilities arise. We can form four-fermion operators as well as two-fermion operators including three scalar bosons and covariant derivatives.

Let us first consider systematically the construction of four-fermion operators conserving both lepton and baryon numbers. A simple inspection of the fermion bilinears listed before (on p.18) informs us that the hypercharge is of primary relevance to determine what bilinears are to be combined together. Proper Lorentz, colour and $SU(2)_L$ transformation properties are then to be fixed. Doing so, we obtain the four-fermion operators of Table 1. The Hermitian conjugates of scalar and tensor operators are also tacitly added to this list. For vector operators, having the same general form as their Hermitian conjugate, the reality condition translates into a constraint on the operator coefficient:

$$[C^*]^a{}_b{}^c{}_d = [C]^b{}_a{}^d{}_c, \tag{9}$$

where, as before, a, b, c, d are flavour indices. For operators composed of two identical bilinears (modulo flavour), we obviously also have

$$[C]^a{}_b{}^c{}_d = [C]^c{}_d{}^a{}_b. \tag{10}$$

---

[8]The antisymmetric singlet combination vanishes as the two bosonic fields commute and are identical.

! $\bar{q}\gamma^\mu q$ $\bar{q}\gamma_\mu q$	$.\gamma^\mu\lambda^A.$ $.\gamma_\mu\lambda^A.$	$.\gamma^\mu\tau^I.$ $.\gamma_\mu\tau^I.$	$.\gamma^\mu\lambda^A\tau^I.$ $.\gamma_\mu\lambda^A\tau^I.$	
$\bar{q}\gamma^\mu q$ $\bar{u}\gamma_\mu u$	$.\gamma^\mu\lambda^A.$ $.\gamma_\mu\lambda^A.$			
$\bar{q}\gamma^\mu q$ $\bar{d}\gamma_\mu d$	$.\gamma^\mu\lambda^A.$ $.\gamma_\mu\lambda^A.$			
! $\bar{u}\gamma^\mu u$ $\bar{u}\gamma_\mu u$	$.\gamma^\mu\lambda^A.$ $.\gamma_\mu\lambda^A.$			
$\bar{u}\gamma^\mu u$ $\bar{d}\gamma_\mu d$	$.\gamma^\mu\lambda^A.$ $.\gamma_\mu\lambda^A.$			
! $\bar{d}\gamma^\mu d$ $\bar{d}\gamma_\mu d$	$.\gamma^\mu\lambda^A.$ $.\gamma_\mu\lambda^A.$			
! $\bar{q}u$ ϵ $\bar{q}d$	$.\lambda^A.$ ϵ $.\lambda^A.$	$.\sigma^{\mu\nu}.$ ϵ $.\sigma_{\mu\nu}.$	$.\sigma^{\mu\nu}\lambda^A.$ ϵ $.\sigma_{\mu\nu}\lambda^A.$	

$\bar{q}\gamma^\mu q$ $\bar{l}\gamma_\mu l$	$.\tau^I.$ $.\gamma_\mu\tau^I.$
$\bar{q}\gamma^\mu q$ $\bar{e}\gamma_\mu e$	
$\bar{u}\gamma^\mu u$ $\bar{l}\gamma_\mu l$	
$\bar{u}\gamma^\mu u$ $\bar{e}\gamma_\mu e$	
$\bar{d}\gamma^\mu d$ $\bar{l}\gamma_\mu l$	
$\bar{d}\gamma^\mu d$ $\bar{e}\gamma_\mu e$	
$\bar{q}u$ ϵ $\bar{l}e$	$.\sigma^{\mu\nu}.$ ϵ $.\sigma_{\mu\nu}.$
$\bar{q}d$ $\bar{e}l$	

! $\bar{l}\gamma^\mu l$ $\bar{l}\gamma_\mu l$	$.\gamma^\mu\tau^I.$ $.\gamma_\mu\tau^I.$
$\bar{l}\gamma^\mu l$ $\bar{e}\gamma_\mu e$	
! $\bar{e}\gamma^\mu e$ $\bar{e}\gamma_\mu e$	

Table 1: Dimension-six four-fermion operators conserving the baryon and lepton numbers. The anti-doublets have their $SU(2)_L$ index implicitly contracted with the nearest doublet (preferentially within the same Lorentz bilinear) and $\epsilon \equiv \epsilon^{ij} = \epsilon_{ij}$ is used to contract the $SU(2)_L$ indices of two (anti-)doublets.

Other forms It is possible to use Fierz identities in spinorial, colour or $SU(2)_L$ space to reshuffle fermions and transform all vector and tensor operators into scalar ones, or remove all Pauli and Gell-Mann matrices. The above forms however have the advantage of separating the four-fermion operators into two bilinears that have definite colour and $SU(2)_L$ transformation properties and do not mix quarks and leptons. Each bilinear therefore conserves the lepton, baryon and, *a fortiori*, fermion numbers.

Redundancies All those operators are however not independent due to the presence, in some of them, of identical fields (modulo their flavour, yet unassigned). This is not transparent with the operators expressed in their present form since they may not have apparent transformation properties under the exchange of identical fields. The fields contents for which care must be taken are identified with an exclamation mark '!' in Table 1.

Let us first consider the four operators deriving from the $q^c q \; q^c q$ field con-

tent. Each vector bilinear can form a colour singlet or octet, and a $SU(2)_L$ singlet or triplet. Let us therefore temporarily denote them O_{L1}^{c1}, O_{L1}^{c8}, O_{L3}^{c1} and O_{L3}^{c8}. On the other hand, using the Fierz transformations of Eq. (6) and Eq. (7), colour and $SU(2)_L$ structures that are either symmetric or antisymmetric under the exchange of the two (anti-)quarks can be formed:

$$[\lambda^A]^\alpha_{\ \beta}\,[\lambda^A]^\gamma_{\ \delta} + k_{c\pm}\,\delta^\alpha_\beta\,\delta^\gamma_\delta = T_c\,(\delta^\alpha_\delta\,\delta^\gamma_\beta \pm \delta^\alpha_\beta\,\delta^\gamma_\delta),$$
$$[\tau^I]^i_{\ j}\,[\tau^I]^k_{\ l} + k_{L\pm}\,\delta^i_j\,\delta^k_l = T_L\,(\delta^i_l\,\delta^k_j \pm \delta^i_j\,\delta^k_l),$$

where $k_{c\pm} \equiv T_c\,(1/N_c \pm 1)$ and $k_{L\pm} \equiv T_L\,(1/N_L \pm 1)$.

With those structures, linear combinations of the four $q^c q\, q^c q$ operators that are either symmetric or antisymmetric under the exchange of the two quarks or anti-quarks can be constructed:

$$O_{L\pm}^{c\pm} \equiv \left[O_{L3}^{c8} + k_{c\pm}O_{L3}^{c1}\right] + k_{L\pm}\left[O_{L3}^{c8} + k_{c\pm}O_{L1}^{c1}\right].$$

As fermions anticommute and as the vector Lorentz structure is already antisymmetric (see the last equality of Eq. (4)), O_{L+}^{c+} and O_{L-}^{c-} are symmetric while O_{L+}^{c-} and O_{L-}^{c+} are antisymmetric under the said permutation. Equivalently, introducing generation indices, we have:

$$[O_{L\pm}^{c\pm}]^a_{\ b}{}^c_{\ d} + k_{L\pm}^{c\pm}\,[O_{L\pm}^{c\pm}]^a_{\ d}{}^c_{\ b} = 0, \tag{11}$$

for $k_{L+}^{c+} = -1 = k_{L-}^{c-}$ and $k_{L+}^{c-} = +1 = k_{L-}^{c+}$.

Two choices are then possible for constructing a basis of independent operators. In the $O_{L\pm}^{c\pm}$ basis, the condition (11) is diagonal and can be enforced by constraining the operator coefficients to be either symmetric or antisymmetric under the permutations of generation indices. The other option is to use the $O_{L1/3}^{c1/8}$ basis, without imposing any new condition on operators coefficients. Equations (11) however allow us to retain only two operators out of the four initial ones. Solving for instance for O_{L3}^{c1} and O_{L3}^{c8}, we get

$$\begin{pmatrix} [O_{L3}^{c8}]_{abcd} \\ [O_{L3}^{c8}]_{cbad} \\ [O_{L3}^{c1}]_{abcd} \\ [O_{L3}^{c1}]_{cbad} \end{pmatrix} = \begin{pmatrix} 1 & 1 & 0 & -3/4 \\ 1 & 1 & -3/4 & 0 \\ 0 & -4 & 1 & -1 \\ -4 & 0 & -1 & 1 \end{pmatrix} \begin{pmatrix} [O_{L1}^{c8}]_{abcd} \\ [O_{L1}^{c8}]_{cbad} \\ [O_{L1}^{c1}]_{abcd} \\ [O_{L1}^{c1}]_{cbad} \end{pmatrix}$$

that explicitly demonstrate O_{L1}^{c1} and O_{L1}^{c8} alone form a complete set of $q^c q\, q^c q$ operators.

A similar but simpler reasoning applies in the case of vector operators containing identical fields and for which the construction of colour octets ($u^c u\, u^c u$ and $d^c d\, d^c d$) or $SU(2)_L$ triplets ($l^c l\, l^c l$) is allowed. Either one of the two operators can be eliminated and the flavour structure of its coefficient kept unconstrained or can the basis be symmetrized and the effective coefficients given the corresponding symmetry in generation space.

An additional ingredient is however needed to treat the case of the operator based on the $q^c u\, q^c d$ field content. Their $SU(2)_L$ structure is already antisymmetric, and its colour structures can be symmetrized as before. Though,

out of the scalar and tensor, we need to form Lorentz structures with definite transformation properties under the exchange of the two q^c fields. The relevant Fierz transformations are obtained from the first equality of Eq. (4):

$$[\sigma^{\mu\nu}P_R]^p{}_q\,[\sigma_{\mu\nu}P_R]^r{}_s + k_{T\pm}[P_R]^p{}_q\,[P_R]^r{}_s = T_T([P_R]^p{}_s[P_R]^r{}_q \pm [P_R]^p{}_q[P_R]^r{}_s),$$

where $k_{T\pm} \equiv T_T(1/N_T \pm 1)$ with $T_T \equiv 8$ and $N_T \equiv 2$. With those structures, we can go through steps identical to the ones we made for the $q^c q\ q^c q$ operators and either suppress two structures in the unsymmetrized basis or define four symmetrized operators and restrict their coefficients to be either symmetric or antisymmetric under the exchange of the two q^c generation indices.

Finally, the vector operator based on the $e^c e\ e^c e$ field content is already symmetric under the exchange of the two pairs of (anti-)leptons. The vector Lorentz structure is indeed antisymmetric and fermions anticommute. On top of equalities (9) and (10) its coefficient should therefore satisfy

$$[C]^a{}_b{}^c{}_d = [C]^a{}_d{}^c{}_b.$$

So, for all operators containing identical fields, care must be taken not to include in the basis elements that are not independent of the others. Alternatively, it is possible to construct a basis in which all those redundancies can be accounted for by requiring the operator coefficients to have definite symmetry properties in generation space.

Dimension six, baryon- and lepton-number-violating operators

As already mentioned, the first baryon-number-violating operators also appear at dimension six and contain three (anti-)quarks and one (anti-)lepton.

$B - L$ **selection rule**   For definiteness, let us start with three quarks instead of anti-quarks. They can be right- or left-handed, $SU(2)_L$ singlets or doublets. Lorentz invariance requires operators to be constructed from an even number of left- or right-handed fields while $SU(2)_L$ invariance asks for an even number of doublets.

However, in the standard model, all right-handed particles are $SU(2)_L$ singlets and all left-handed ones are doublets. Corresponding anti-particles have opposite chiralities but anti-doublets are equivalent to doublets (the fundamental representation of $SU(2)$ is *pseudo-real*).

So, if we pick up an even number of left-handed quark doublets (two or none), the leptonic field they should be combined with is a right-handed $SU(2)_L$ singlet: e. Conversely, with an odd number of quark doublets (one or three), the associated leptonic field should be a left-handed $SU(2)_L$ doublet: l. In both cases, the quark triad should thus be complemented with a lepton, not an anti-lepton. The four-fermion operators violating B and L, that can be constructed out of standard-model fields, therefore conserve $B - L$.

Field content   The only missing restriction arises from hypercharge conservation. Amongst all possible combinations of three quarks and a lepton satisfying the chirality and $SU(2)_L$ counting requirements: *qqql* (of total hypercharge

$Y/2 = 0$), $qque$ (0), $qqde$ (-1), $uuql$ (1), $duql$ (0), $ddql$ (-1), $uuue$ (1), $duue$ (0), $ddue$ (-1) and $ddde$ (-2), only four are permitted.

Explicit construction   In baryon- and lepton-number-violating four-fermion operators, there is no clear rationale for forming one fermion bilinear instead of the other. All possible bilinears will anyway violate baryon, lepton and fermion numbers. If an arbitrary pairing is fixed, the operator construction proceeds as in the B and L conserving case. Some equalities however allow to relate all possible pairings.

For a chosen pairing, two independent singlets can be formed out of four $SU(2)_L$ doublets using the $\epsilon_{ij}\,\epsilon_{kl}$ and $[\epsilon\tau^I]_{ij}[\epsilon\tau^I]_{kl}$ tensor structures, similarly as before. Equivalently, we could also choose two $\epsilon_{..}\epsilon_{..}$ structures at will amongst the three possible ones. They are related by the equality (8).

Likewise, out of four fields of identical chirality, the two possible Lorentz singlets have $[C^\dagger P_{L,R}]_{pq}[C^\dagger P_{L,R}]_{rs}$ and $[C^\dagger \sigma^{\mu\nu} P_{L,R}]_{pq}[C^\dagger \sigma_{\mu\nu} P_{L,R}]_{rs}$ structures if the pairing is fixed. With a liberal fermion pairing, two scalar (or tensor) Lorentz structures could be chosen out of the three possible ones satisfying the first (or second) Schouten identity in Eq. (5).

With two left- and two right-handed fields, on the other hand, only one Lorentz-invariant structure can be built since $[(\mathbf{1},\mathbf{2})\otimes(\mathbf{1},\mathbf{2})]\otimes[(\mathbf{2},\mathbf{1})\otimes(\mathbf{2},\mathbf{1})] = (\mathbf{1},\mathbf{1}) \oplus (\mathbf{1},\mathbf{3}) \oplus (\mathbf{3},\mathbf{1}) \oplus (\mathbf{3},\mathbf{3})$ contains one single singlet. Depending on the pairing chosen, vector or scalar bilinears are to be used. These structures are then related by the two last equalities of Eq. (4). As stressed before, no Lorentz-invariant structure is obtained by contracting a left- and a right-handed tensor bilinear since $(\mathbf{1},\mathbf{3}) \otimes (\mathbf{3},\mathbf{1})$ does not contain any Lorentz scalar (see also the third equality of Eq. (4)).

Let us arbitrarily fix a pairing for each of the four allowed field contents as doing so will make the discussion about redundancies an exact replica of the one carried out for baryon- and lepton-number-conserving operators. The naive list of possible combinations of $SU(2)_L$ and Lorentz structures is then:

$$
\begin{array}{|llllllll|}
\hline
! \; \tilde{\epsilon}\,\overline{\boldsymbol{q^c}}\,\epsilon\,\boldsymbol{q} & \overline{\boldsymbol{q^c}}\,\epsilon\,\boldsymbol{l} & \tilde{\epsilon}\,.\,\epsilon\tau^I\,. & .\,\epsilon\tau^I\,. & \tilde{\epsilon}\,.\,\sigma^{\mu\nu}\epsilon\,. & .\,\sigma_{\mu\nu}\epsilon\,. & \tilde{\epsilon}\,.\,\sigma^{\mu\nu}\epsilon\tau^I\,. & .\,\sigma_{\mu\nu}\epsilon\tau^I\,. \\
! \; \tilde{\epsilon}\,\overline{\boldsymbol{q^c}}\,\epsilon\,\boldsymbol{q} & \overline{\boldsymbol{u^c}}\,\boldsymbol{e} & & & & & & \\
\; \tilde{\epsilon}\,\overline{\boldsymbol{d^c}}\,\boldsymbol{u} & \overline{\boldsymbol{q^c}}\,\epsilon\,\boldsymbol{l} & & & & & & \\
! \; \tilde{\epsilon}\,\overline{\boldsymbol{d^c}}\,\boldsymbol{u} & \overline{\boldsymbol{u^c}}\,\boldsymbol{e} & & & \tilde{\epsilon}\,.\,\sigma^{\mu\nu}\,. & .\,\sigma_{\mu\nu}\,. & & \\
\hline
\end{array}
$$

where the $\tilde{\epsilon} \equiv \epsilon_{\alpha\beta\gamma}$ tensor serves to antisymmetrize the colour indices of the three quarks; nearest $SU(2)_L$ indices contractions are implicit; so is also the inclusion charge-conjugate operators; and, as before, an exclamation mark '!' stresses the presence of identical fields leading to redundancies.

Redundancies   First, fermionic fields in a bilinear can be permuted using the Lorentz bilinear transposition properties of Eq. (3) while, obviously, in $SU(2)$, $[\epsilon]^T = -[\epsilon]$ and $[\epsilon\tau^I]^T = +[\epsilon\tau^I]$. Taking also Fermi-Dirac statistics and colour antisymmetrization into account, the $qque$ operator is seen to be symmetric

under the permutation of its two quark doublets. The two corresponding generation indices of its coefficient should therefore have the same property.

From the experience gained in discussing baryon- and lepton-number-conserving four-fermion operators, it is already clear that two of the *qqql* operators could be removed. In doing so, the identical character (modulo flavour) of the first and third quarks has been used. The remaining two operators still exhibit definite symmetry properties under the exchange of the first and second q's (the four initial operators are respectively symmetric, antisymmetric, antisymmetric and symmetric). It is therefore possible to further reduce them to one single operator having no such definite transformation property as, for instance, $\tilde{O}^{(4)} \equiv \tilde{\epsilon} \ \overline{q^c}^i q^j \ \overline{q^c}^k l^l \ \epsilon_{il}\epsilon_{jk}$ [32]. Finally, due to the first Schouten identity of Eq. (5) and its equivalent in $SU(2)$, Eq. (8), permutations of the flavour indices of this operator still satisfies [32]:[9]

$$\tilde{O}^{(4)}_{abcd} + \tilde{O}^{(4)}_{bacd} - \tilde{O}^{(4)}_{cbad} - \tilde{O}^{(4)}_{cabd} = 0. \tag{12}$$

It is equally clear that one of the two *udue* operators is redundant (or that two independent operators with definite symmetry properties under the exchange of the two u fields could be constructed).

Basis The standard basis of baryon- and lepton-number-violating four-fermion operators was established in Refs. [32, 37, 39]. It was chosen to contain only scalar operators:

$$
\begin{aligned}
O^{(1)} &\equiv \ \tilde{\epsilon} \ \overline{d^c} \, u \qquad \overline{q^c} \, \epsilon \, l \\
O^{(2)} &\equiv \ \tilde{\epsilon} \ \overline{q^c} \, \epsilon \, q \qquad \overline{u^c} \, e \\
O^{(3)} &\equiv \ \tilde{\epsilon} \ \overline{q^c} \, \epsilon \, q \qquad \overline{q^c} \, \epsilon \, l \qquad\qquad + \text{h.c.} \\
O^{(4)} &\equiv \tilde{\epsilon} \ \overline{q^c} \, \epsilon\tau^I \, q \quad \overline{q^c} \, \epsilon\tau^I \, l \\
O^{(5)} &\equiv \ \tilde{\epsilon} \ \overline{d^c} \, u \qquad \overline{u^c} \, e
\end{aligned}
\tag{13}
$$

For the *qqql* field content, the two $O^{(3)}$ and $O^{(4)}$ operators with definite symmetry properties under the exchange of the two first quarks have been kept. As they are respectively the symmetric and anti-symmetric part of $\tilde{O}^{(4)}$ under this permutation,

$$-O^{(3)}_{abcd} = \tilde{O}^{(4)}_{abcd} + \tilde{O}^{(4)}_{bacd},$$
$$-O^{(4)}_{abcd} = \tilde{O}^{(4)}_{abcd} - \tilde{O}^{(4)}_{bacd},$$

the equality (12) translates into

$$O^{(3)}_{abcd} - \frac{1}{2}(O^{(3)}_{cabd} + O^{(4)}_{cabd}) - \frac{1}{2}(O^{(3)}_{cbad} + O^{(4)}_{cbad}) = 0.$$

---

[9]The $\tilde{O}^{(4)}$ operator could actually be decomposed according to the irreducible representations of the permutation group of its three quarks into: a completely symmetric tensor (▢▢▢ with $N_g(N_g + 1)(N_g + 2)/6 = 10$ components), two tensors of mixed symmetry properties (▱ with $N_g(N_g^2 - 1)/3 = 8$ components) and a completely antisymmetric tensor (▯ with $N_g(N_g - 1)(N_g - 2) = 1$ component). The equality (12) indicates one mixed tensor actually vanishes [38].

Dimension-six, two fermions operators

Beside four-fermion operators, two fermions associated with three scalar bosons or covariant derivatives (potentially combined in a field strength tensor) also form dimension-six operators. Let us treat the four possible cases one after the other:

- three covariant derivatives,
- two covariant derivatives and a scalar boson,
- one covariant derivative and two scalar bosons,
- three scalar bosons.

When possible, covariant derivatives will be traded for fields using the classical equations of motion to lowest order in $1/\Lambda$:

$$i\slashed{D}q = Y_u^\dagger \, u \, \epsilon\phi^* + Y_d^\dagger \, d \, \phi,$$

$$i\slashed{D}u = Y_u \, q \, \epsilon\phi,$$

$$i\slashed{D}d = Y_d \, q \, \phi^*,$$

$$i\slashed{D}l = Y_e^\dagger \, e \, \phi,$$

$$i\slashed{D}e = Y_e \, l \, \phi^*,$$

$$D_\mu D^\mu \phi = -Y_u^\dagger \, \bar{q}\epsilon u - Y_d \, \bar{d}q - Y_e \, \bar{e}l + \mu^2 \phi - \lambda \left(\phi^*\phi\right)\phi,$$

$$\frac{1}{g_s}(D_\nu G^{\nu\mu})^A = \bar{q}\gamma^\mu \frac{\lambda^A}{2} q + \bar{u}\gamma^\mu \frac{\lambda^A}{2} u + \bar{d}\gamma^\mu \frac{\lambda^A}{2} d,$$

$$\frac{1}{g}(D_\nu W^{\nu\mu})^I = \phi^* \frac{i\vec{D}^{\mu I}}{2}\phi + \bar{q}\gamma^\mu \frac{\tau^I}{2} q + \bar{l}\gamma^\mu \frac{\tau^I}{2} l,$$

$$\frac{1}{g'}D_\nu B^{\nu\mu} = \frac{1}{2}\phi^* i\vec{D}^\mu \phi + \frac{1}{6}\bar{q}\gamma^\mu q + \frac{2}{3}\bar{u}\gamma^\mu u - \frac{1}{3}\bar{d}\gamma^\mu d - \frac{1}{2}\bar{l}\gamma^\mu l - \bar{e}\gamma^\mu e,$$

with $\vec{D}^\mu \equiv \vec{D}^\mu - \overleftarrow{D}^\mu$ and $\vec{D}^{\mu I} \equiv \tau^I \vec{D}^\mu - \overleftarrow{D}^\mu \tau^I$. The replacements they allow are equivalent to non-linear field redefinitions that do not modify the S-matrix elements [40, and references therein]. Only will the replacement based on the equation of motion of ϕ mix operators of different dimensions. It involves the only dimensionful coupling of the standard model, μ.

Three covariant derivatives   A fermion bilinear to be combined with three covariant derivatives should form a gauge-invariant vector. All three derivatives can be chosen to act on one of the two fermionic fields,

$$\bar{\psi} \, D^\mu D^\nu D^\rho \gamma^\sigma \, \psi,$$

as other combinations are equivalent to this one up to a total derivative. The four Lorentz indices can be contracted in four independent ways: they are three different pairing that can be formed using the metric and a totally antisymmetric combination based on $\epsilon_{\mu\nu\rho\sigma}$:

$$\bar{\psi} \, D^\mu D_\mu \slashed{D} \, \psi, \qquad \bar{\psi} \, \slashed{D}D^\mu D_\mu \, \psi, \qquad \bar{\psi} \, D^\mu \slashed{D} D_\mu \, \psi, \qquad \bar{\psi} \, D^\mu D^\nu D^\rho \gamma^\sigma \, \psi \, \epsilon_{\mu\nu\rho\sigma}.$$

The use of the equation of motion for ψ allows to trade the first contraction for operators with two derivatives and a scalar boson. This is also true in the second contraction, where the equation of motion for $\bar{\psi}$ is to be employed and a total derivative is left over. The third contraction reduces to the first two ones, modulo a term containing the covariant derivative of a field strength tensor:

$$2\, D^\mu D_\sigma D_\mu = D_\sigma D^\mu D_\mu + D^\mu D_\mu D_\sigma + (D^\mu X_{\sigma\mu}).$$

Here, $X_{\mu\nu}\psi \equiv [D_\mu, D_\nu]\psi$ is the sum of the field strength tensors of the gauge groups under which ψ transforms non trivially (each multiplied by $-i$ times the corresponding gauge coupling). The use of the equation of motion for $X^{\mu\nu}$ therefore renders this term equivalent to four-fermion operators, two-fermion operators with a covariant derivative and two scalar bosons. In the fourth case, the equality:

$$\gamma^\mu \gamma^\nu \gamma^\rho = -i\epsilon^{\mu\nu\rho\sigma}\gamma_\sigma\gamma^5 + g^{\mu\nu}\gamma^\rho - g^{\mu\rho}\gamma^\nu + g^{\nu\rho}\gamma^\mu$$

can be used to exchange $\epsilon_{\mu\nu\rho\sigma}\gamma^\sigma$ for products of the metric and Dirac matrices. All of the generated terms contain at least a $D\!\!\!/$ suitably placed for the equations of motion of ψ or $\bar{\psi}$ to be applied (potentially also generating a total derivative).

So, by using the classical equations of motion, a fermion bilinear associated with all possible combinations of three covariant derivatives can always be traded for operators with less derivatives and more scalar bosons or with four fermions.

Two covariant derivatives and a scalar boson With two fermions, a scalar boson and two covariant derivatives, the field content is the one of a Yukawa interaction: $u^c q\phi$, $d^c q\phi^*$ or $e^c l\phi^*$ (and charge conjugates) but the fermion bilinear is allowed to have a scalar or tensor Lorentz structure.

Up to a surface term in the action, the two covariant derivatives can be chosen to act on the fermions only. Denoting the generic field content by $\psi^c \psi' \phi$, there are therefore three possible structures:

$$(D^\mu D^\nu \psi^c)\psi'\phi, \qquad \psi^c(D^\mu D^\nu \psi')\phi, \qquad (D^\mu \psi^c)(D^\nu \psi')\phi.$$

The Lorentz indices are then in principle to be contracted with either $g_{\mu\nu}$ or $\sigma_{\mu\nu}$. However, acting with γ^μ or $D\!\!\!/$ on the left of the classical equation of motion for ψ, schematically written as $iD\!\!\!/\psi - Y_\psi \psi'\phi = 0$, and using $\gamma^\mu\gamma^\nu = g^{\mu\nu} - i\sigma^{\mu\nu}$ yields the equalities:

$$ig^{\mu\nu}D_\nu\psi + \sigma^{\mu\nu}D_\nu\psi - Y_\psi\gamma^\mu\psi'\phi = 0,$$

$$iD^\mu D_\mu\psi + \frac{1}{2}\sigma^{\mu\nu}X_{\mu\nu}\psi - Y_\psi D\!\!\!/(\psi'\phi) = 0,$$

Similar relations of course also hold for ψ^c and ψ'. The first one shows that the two possible Lorentz structures for fermion bilinears are related by the equations of motion.

For the $(D^\mu D^\nu \psi^c)\psi'\phi$ and $\psi^c(D^\mu D^\nu \psi')\phi$ structures, we chose to use tensor bilinears. They give rise to operators with a field strength tensor, of generic

form $\bar{\psi}\sigma_{\mu\nu}X^{\mu\nu}\psi'\phi$. In the $(D^\mu\psi^c)(D^\nu\psi')\phi$ case, let us use a scalar structure. Because of the relation,

$$2\,(D^\mu\psi^c)(D_\mu\psi')\phi = -\,(D^\mu D_\mu\psi^c)\psi'\phi - \psi^c(D^\mu D_\mu\psi')\phi + \psi^c\psi'(D_\mu D^\mu\phi),$$
$$+ D^\mu D_\mu(\psi^c\psi'\phi) - 2D^\mu(\psi^c\psi'D_\mu\phi),$$

it also reduces, by the equations of motion for ψ, ψ' and ϕ, to operators with a field strength tensor, a second scalar boson, four fermions or to Yukawa operators.

Therefore, all operators that are to be considered in this category are:

$\bar{u}\,\sigma_{\mu\nu}\,B^{\mu\nu}\,q\;$	$\epsilon\phi,$	$\bar{d}\,\sigma_{\mu\nu}\,B^{\mu\nu}\,q\;$	$\phi^*,$	$\bar{e}\,\sigma_{\mu\nu}\,B^{\mu\nu}\,l\;$	$\phi^*,$
$\bar{u}\,\sigma_{\mu\nu}\,W^{\mu\nu}\,q\;$	$\epsilon\phi,$	$\bar{d}\,\sigma_{\mu\nu}\,W^{\mu\nu}\,q\;$	$\phi^*,$	$\bar{e}\,\sigma_{\mu\nu}\,W^{\mu\nu}\,l\;$	$\phi^*,$
$\bar{u}\,\sigma_{\mu\nu}\,G^{\mu\nu}\,q\;$	$\epsilon\phi,$	$\bar{d}\,\sigma_{\mu\nu}\,G^{\mu\nu}\,q\;$	$\phi^*,$		

together with their Hermitian conjugates.

A covariant derivative and two scalar bosons Given the transformation properties of fermion bilinears (listed on) and of pairs of scalars (on), the only allowed operators containing a covariant derivative, two scalar bosons and a fermion bilinear are of the form:

$$\bar{\psi}\,i\overset{\leftrightarrow}{\slashed{D}}{}^{(I)}\psi \quad \phi^*(\tau^I)\phi, \qquad\qquad \bar{\psi}\gamma^\mu(\tau^I)\psi \quad \phi^*i\overset{\leftrightarrow}{D}{}^{(I)}_\mu\phi.$$

Both singlet and triplet bilinears are allowed only if ψ is a $SU(2)_L$ doublet. The derivative $\overset{\leftrightarrow}{D}{}^{(I)}$ serves to make those operators Hermitian (provided their coefficients that depend on the flavours of the two ψ's are themselves Hermitian). Using the equation of motion for ψ, the first structure can be eliminated in favour of operators with three scalar bosons. We are therefore left with the following operators to be added in our dimension-six basis:

$\bar{q}\gamma^\mu\tau^I q$	$\phi^*i\overset{\leftrightarrow}{D}{}^I_\mu\phi,$	$\bar{l}\gamma^\mu\tau^I l$	$\phi^*i\overset{\leftrightarrow}{D}{}^I_\mu\phi,$
$\bar{q}\,\gamma^\mu\,q$	$\phi^*i\overset{\leftrightarrow}{D}_\mu\phi,$	$\bar{l}\,\gamma^\mu\,l$	$\phi^*i\overset{\leftrightarrow}{D}_\mu\phi,$
$\bar{u}\,\gamma^\mu\,u$	$\phi^*i\overset{\leftrightarrow}{D}_\mu\phi,$	$\bar{e}\,\gamma^\mu\,e$	$\phi^*i\overset{\leftrightarrow}{D}_\mu\phi,$
$\bar{d}\,\gamma^\mu\,d$	$\phi^*i\overset{\leftrightarrow}{D}_\mu\phi.$		

Three scalar bosons As no fermion bilinear listed on has a total hypercharge of $\pm 3/2$, operators composed of a fermion bilinear and three scalar bosons necessarily involve a $\phi^*\phi$ pair and a Yukawa-like $\psi^c\psi'\phi$ structure. Those two pieces should have scalar Lorentz transformation properties but could in principle form both singlet and triplet $SU(2)_L$ combinations ($\bar{\psi}\psi'\,\phi\;\;\phi^*\phi$ and $\bar{\psi}\psi'\,\epsilon\tau^I\phi\;\;\phi^*\tau^I\phi$). Bose-Einstein statistics will however select the symmetric combination of the two ϕ's ($\phi^*\tau^I\;\bar{\psi}\psi'\;\phi\epsilon\tau^I\phi$). There is therefore only one independent structure. If the $SU(2)_L$ singlet one is chosen, we get the following

operators in this category:

$$\begin{array}{llll} \bar{u}\,q & \epsilon\phi & \phi^*\phi, \\ \bar{d}\,q & \phi^* & \phi^*\phi, \\ \bar{e}\,l & \phi^* & \phi^*\phi, \end{array}$$

as well as their Hermitian conjugates.

I.4 Conclusions

In this introductory chapter, the spirit and methods of effective field theories have been presented. After having described the standard-model symmetries and field content, a complete basis of operators involving fermions has been constructed with dimension up to six. In the process of this construction, the presence of certain global flavour symmetries has been highlighted.

Amongst those operators, only the ones with dimension equal or less to five have been firmly established experimentally. Whether new physics contributes to dimension-six operators or modifies the standard-model field content remain open questions. The next chapters explore three different paths that could be followed to try answering them. The first one uses the effective theory just constructed, the second chapter generically identifies how new resonances could manifest themselves, and the third one studies quantitatively the consequences of a specific model. Examples of applications involving the top quark, baryon number violation and standard-model flavour symmetries are developed in each case.

II. Effective field theory

With the relevant effective field theory constructed, let us now turn to its phenomenology and concentrate our attention on the top sector. Above dimension four, one single dimension-five operator violates the lepton number by two units and provides standard-model neutrinos with masses and mixings. At dimension six, a wide variety of operators possibly triggers very distinctive processes: the structure of the *tbW* vertex can be modified, flavour-changing neutral current unsuppressed by the GIM mechanism are possibly generated, etc. The most striking consequence of those six-dimensional operators is however baryon number violation.

II.1 Baryon number violation

The good apparent conservation of the baryon number remains amongst the greatest puzzles of particle physics. Violations occur naturally in most theoretical models while stringent experimental bounds have been set, most notably on nucleon decay processes otherwise forbidden.

The first step towards the introduction of a conservation laws for the baryon number is presumably due to Weyl [41].[10] As, in 1929, the positron was still to be discovered, he considered the electron and the proton on the same footing and imagined the conservation of both of these positive and negative quanta of charge to be associated with two separate gauge symmetries:

> *It is plausible to anticipate that, of the two pairs of components of the Dirac quantity, one belongs to the electron, the other to the proton. Further, two conservation laws of electricity will have to appear, which state (after quantization) that the number of electrons as well as the number of protons remains constant. To these conservation laws must correspond a twofold gauge invariance, involving two arbitrary functions.*
> [41, p.332] (translation by A. Pais quoted in Ref. [42])

In a systematic classification of known particles Stueckelberg introduced in 1936 [45][11] a heavy charge, or *schwere Ladung*, carried by protons and neutrons

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[10]Our historical comments follow Refs. [42–44].
[11]Selected papers of Stueckelberg are reproduced in Ref. [46].

(discovered in 1932) as opposed to a light charge, or *leichte Ladung*, for electrons and neutrinos (postulated in 1930). Aiming to explain the lack of experimental evidences for transitions from the heavy sector to the light one, Stueckelberg moreover suggested explicitly the conservation of this heavy charge in a twofold paper dating from 1938:

> *Besides the* conservation law of electric charge*, which follows from Maxwell's theory, there apparently* [offenbar] *exists a further conservation law: For all observed transformations of matter, no transformations of heavy particles (neutron and proton) into light particles (electron and neutrino) have yet been observed. We therefore wish to require* [fordern] *a* conservation law of the heavy charge.
> <div align="right">[47, 48, p.317] (translation adapted from Ref. [42])</div>

In passing, the second paper of this series remarkably introduced an original mechanism for the construction of a massive unbroken Abelian gauge theory.

The baryon number conservation law was reinvented more than ten years later, in 1949, by Wigner in order to explain more specifically the stability of the proton against the otherwise allowed $p \to \gamma e^+$ decay suggested by analogy with the $e^- \to \gamma \nu_e$ one forbidden by electric charge conservation:

> *It is conceivable, for instance, that a conservation law for the number of heavy particles (protons and neutrons) is responsible for the stability of the protons in the same way as the conservation law for charges is responsible for the stability of the electron. Without the conservation law in question, the proton could disintegrate, under emission of a light quantum, into a positron, just as the electron could disintegrate, were it not for the conservation law for the electric charge, into a light quantum and a neutrino.* [9, p.525]

As far as terminology is concerned, the word *baryon* seems to have been introduced by Pais in 1953 [49]. Nucleons and hyperons were so gathered in this newly defined family.

The following year, a quarter of century after Weyl's hint, a first experimental test was proposed by Goldhaber (unpublished, see Ref. [50]) to set a quantitative limit on the proton lifetime using the *observed spontaneous fission rate of* ^{232}Th. His method led to a limit of the order of 10^{20} years for bound nucleons lifetime. Goldhaber can maybe also be credited with the following extraction of a bound on nucleon lifetime:

> *Why did these three learned gentlemen, Weyl, Stueckelberg, and Wigner, feel so sure that baryons are conserved? Well, you might say that it's very simple, they felt it in their bones. Had their bones been irradiate by the decays of nucleons, they would have noticed effects considerably exceeding "permissible radiological limits" if the nucleon lifetime were $< 10^{16}$ years and if at least 10% of the nucleon rest mass were to appear as radiation absorbable in the body. That is a fairly sensitive measurement, but one can do much better by deliberate experiments.* [51]

The sixties and seventies, on the other hand, provided theorists with many reasons for doubting of the absolute character of the baryon number conservation law. A first model assuming it is only approximate was put forward in 1959 by Yamaguchi [52] that postulated a *superweak* interaction violating the baryon number amongst other symmetries in processes like $p \to e^+ e^+ e^-$. The underlying motivation was the observation that, amongst the known interactions, the stronger ones possess the larger number of symmetries and *vice versa*. An important step was then made in 1967 by Sakharov who stated baryon number violation as one of the three necessary ingredients for a dynamical generation of the baryon–anti–baryon asymmetry observed in the universe, from a symmetric initial condition [53]. In the early seventies, it was also suggested that black holes would lead to the non-conservation of baryon number [54]. A few years later, the advent of Grand Unified Theories [55, 56] and the further theoretical support they brought to baryon number violation sparked off detailed studies and experimental searches.

Violations occur almost automatically beyond, but also within, the standard model. Baryon and lepton numbers are only conserved by by operators of dimension up-to-four constituted of standard-model fields and preserving gauge invariances. Even if we restrict ourselves to operators of dimension four at most, at the quantum level, those two symmetries are nonetheless separately anomalous [16]. In other words, $U(1)_L$ and $U(1)_B$ are fairly accidental global symmetries of the standard-model Lagrangian. Only is the $B - L$ combination anomaly-free (so it can be gauged without extra matter content requirement). Transitions between topologically distinct vacua of the electroweak theory violate $B + L$. At zero temperature, those tunnelling processes are however suppressed by an extremely small non-perturbative factor of the order of $e^{-2\pi/\alpha_W}$. At finite temperatures, a thermal fluctuation could acquire enough energy to overcome the potential barrier separating two adjacent vacua [57]. Called *sphalerons* [58], such processes could have occurred in the universe history either generating a net $B + L$ asymmetry from a symmetric initial condition, or washing out the $B + L$ component of a pre-existing asymmetry. The latter mechanism could for instance generate a baryon asymmetry from a lepton one.

Experimentally, long-standing searches for proton or bound neutron decays ($|\Delta B| = 1$), neutron oscillations or di-nucleon decays ($|\Delta B| = 2$) have reached impressive heights. Those limits, directly constraining processes involving the first generations of fermions only, have been extended to higher generations by studying τ lepton as well as D and B mesons decays to a proton or a Λ meson. At a somewhat higher energy scale, Z decays to a proton and an electron or muon have also been bound. A few of the most stringent limits are displayed in Table 2. Only more recently were baryon-number-violating processes involving the top quark studied. Morrissey, Tait and Wagner [62] considered the $bb \to \bar{t}\bar{\nu}_\tau$ process induced, between third-generation fermions, by the instantons of an extended gauge group. Hou, Nagashima and Soddu [63] also mentioned the top baryon-number-violating decay $t \to \bar{c}\,\bar{b}\,\ell^+$ triggered by dimension-six operators.

$$
\begin{aligned}
\tau_{n\bar{n}} &> & 0.86 \times 10^8 \text{ years} & \quad (90\% \text{ CL})~[59] \\
\tau/\mathrm{Br}(pp \to K^+ K^+) &> & 1.7 \times 10^{32} \text{ years} & \quad (90\% \text{ CL})~[60] \\
\tau_p/\mathrm{Br}(p \to e^+ \pi^0) &> & 1.2 \times 10^{34} \text{ years} & \quad (90\% \text{ CL})~[61] \\
\tau_p/\mathrm{Br}(p \to K^+ \bar{\nu}) &> & 3.9 \times 10^{33} \text{ years} & \quad (90\% \text{ CL})~[61] \\
\tau_p/\mathrm{Br}(p \to K^0 \mu^+) &> & 1.6 \times 10^{33} \text{ years} & \quad (90\% \text{ CL}) \\
\mathrm{Br}(\tau^- \to \Lambda \pi^-) &< & 7.2 \times 10^{-8} & \quad (90\% \text{ CL}) \\
\mathrm{Br}(D^0 \to pe^-) &< & 1.0 \times 10^{-5} & \quad (90\% \text{ CL}) \\
\mathrm{Br}(B^+ \to \Lambda e^+) &< & 3.2 \times 10^{-8} & \quad (90\% \text{ CL}) \\
\mathrm{Br}(Z \to pe^-) &< & 1.8 \times 10^{-6} & \quad (95\% \text{ CL})
\end{aligned}
$$

Table 2: Limits on representative baryon- (and lepton-) number-violating processes. They are taken from Ref. [15] unless otherwise specified.

II.2 LHC phenomenology

With its specific signatures, the top quark provides a clear means of testing directly baryon number violation at much higher energies. Processes involving all three generations can therefore be placed under direct scrutiny. Interestingly also, thanks to its lifetime shorter than the hadronization time, the top quark is the only system in which baryon number violation is probed at the quark level. Finally it is the quark whose charge and hence baryon number can the most easily be identified in detectors.

Operators

The baryon- and lepton-number-violating operators of dimension six, constructed out of standard-model fields and preserving Lorentz and gauge symmetries,

$$
\begin{aligned}
O^{(1)} &\equiv & \tilde{\epsilon}~\overline{d^c}~u & \qquad \overline{q^c}~\epsilon~l \\
O^{(2)} &\equiv & \tilde{\epsilon}~\overline{q^c}~\epsilon~q & \qquad \overline{u^c}~e \\
O^{(3)} &\equiv & \tilde{\epsilon}~\overline{q^c}~\epsilon~q & \qquad \overline{q^c}~\epsilon~l \\
O^{(4)} &\equiv & \tilde{\epsilon}~\overline{q^c}~\epsilon\tau^I~q & \qquad \overline{q^c}~\epsilon\tau^I~l \\
O^{(5)} &\equiv & \tilde{\epsilon}~\overline{d^c}~u & \qquad \overline{u^c}~e
\end{aligned}
\tag{14}
$$

were constructed in Section I.3 and are reminded here for convenience. At tree level, they could arise from the exchange of heavy mediators with the gauge and Lorentz transformation properties listed in Table 3. Only $O^{(1)}$ and $O^{(2)}$ could arise from the tree-level exchange of a heavy vector and only $O^{(3)}$, $O^{(4)}$ and $O^{(5)}$ through the mediation of a tensor. The X_S scalar has the quantum numbers of a down squark but a (renormalizable) R-parity violating supersymmetric theory would actually only generate operator $O^{(1)}$ due to the holomorphicity of the superpotential. A minimal grand unified theory based on $SU(5)$ would include X_S in the fundamental representation embedding also the standard-model scalar boson as well as X_V in the adjoint representation together with standard-model gauge bosons. In this case, all operators but $O^{(4)}$ would therefore be generated.

	$X_S \sim (\mathbf{3,1},-1/3)^S$	$X_T \sim (\mathbf{3,1},-1/3)^T$	$X'_S \sim (\mathbf{3,1},-4/3)^S$	$X'_T \sim (\mathbf{3,1},-4/3)^T$	$X''_S \sim (\mathbf{3,3},-1/3)^S$	$X''_T \sim (\mathbf{3,3},-1/3)^T$	$X_V \sim (\mathbf{3,2},-5/6)^V$	$X'_V \sim (\mathbf{3,2},-1/6)^V$
$O^{(1)}$	•						•	•
$O^{(2)}$	•						•	
$O^{(3)}$	•	•						
$O^{(4)}$					•	•		
$O^{(5)}$	•	•	•	•				

Table 3: $(SU(3)_c,\ SU(2)_L,\ U(1)_{Y/2})_{\text{Lorentz}}$ transformation properties of the heavy tree-level mediator that could possibly give rise to the four-fermion operators of Eq. (14).

The construction of operators based on symmetries could only make sense with fields in their gauge eigenstates. However, below the Fermi scale at which the electroweak symmetry is spontaneously broken, only physical eigenstates are experimentally observed. The basis of gauge-invariant baryon-number-violating operators (14), should therefore be rotated to a physical one. The interactions obtained are of the form

$$UDUE \qquad \text{and} \qquad DUDN$$

where U, D, E, N are flavour-generic up- and down-type quark, charged lepton and neutrino mass eigenstates. In Ref. [3], we performed an effective-theory study of baryon-number-violating processes involving a single top quark and a charged lepton. They are the most relevant at the LHC. Same-sign top pairs requiring a higher centre-of-mass energy could also be produced at hadron $(gD \to \bar{t}\bar{t}E^+)$ or future ep colliders $(eD \to \bar{t}\bar{t})$ if physical operators containing two tops have significant coefficients at the scales probed. Processes involving a neutrino (like $DD \to \bar{t}\bar{N}$ and $t \to \bar{D}\bar{D}\bar{N}$) were discarded as their signatures could be mimicked by flavour-changing neutral currents (like $gU \to t\bar{N}N$ and $t \to gU\bar{N}N$). Due to the $SU(2)_L$ invariance of the original basis, the operators generating $DUDN$ interactions anyway also contribute to the $UDUE$ ones. On the contrary, at hadron colliders, a single charged lepton in the final state, without any missing energy, is in principle an unambiguous evidence for baryon number violation. Because of the selection rule $\Delta(B+L) \in 2\,\mathbb{Z}$ deriving from angular momentum conservation (see Section I.3, on p.15), this single lepton arising from a hadronic initial state—a $|\Delta L| = 1$ process—actually forces ΔB to be odd and, hence, non-vanishing. A similar reasoning would also hold for ep colliders where $eD \to \bar{U}\bar{t}$ without neutrino or charged lepton in the final state departs from flavour-changing charged $(e\bar{D} \to \nu_e \bar{t})$ or neutral $(eU \to et)$ current processes and clearly points at baryon number violation.

Requiring operators to contain a physical top quark and charged lepton

leads to two distinct structures only:

$$O^{(s)} \equiv \tilde{\epsilon} \quad \overline{t^c}(aP_L + bP_R)D \quad \overline{U^c}(cP_L + dP_R)E,$$
$$O^{(t)} \equiv \tilde{\epsilon} \quad \overline{t^c}(a'P_L + b'P_R)E \quad \overline{U^c}(c'P_L + d'P_R)D,$$

where a, a', b,... are dimensionless coefficients whose products $[ac]$, $[ad]$,... in general depend on the flavour of the U, D and E fermions. Note that a structure like $O^{(u)} \equiv \tilde{\epsilon} \quad \overline{t^c}P_L U \quad \overline{D^c}P_L E$ would be redundant because of the Schouten identity of Eq. (5). In case U is also a top quark, the two $O^{(s)}$ and $O^{(t)}$ operators are obviously not longer independent and considering only one of them suffices. The contributions that each product of coefficients receives from the original gauge basis of Eq. (14) are:

$$[ac]_{3b'c'd'} \ [V_{u_L}]^3{}_a[V_{d_L}]^{b'}{}_b[V_{u_L}]^{c'}{}_c[V_{e_L}]^{d'}{}_d = 2\,C^{(3)}_{abcd} + 2\,C^{(4)}_{abcd} - 4\,C^{(4)}_{acbd},$$
$$[ad]_{3b'c'd'} \ [V_{u_L}]^3{}_a[V_{d_L}]^{b'}{}_b[V_{u_R}]^{c'}{}_c[V_{e_R}]^{d'}{}_d = 2\,C^{(2)}_{abcd},$$
$$[bc]_{3b'c'd'} \ [V_{u_R}]^3{}_a[V_{d_R}]^{b'}{}_b[V_{u_L}]^{c'}{}_c[V_{e_L}]^{d'}{}_d = -\,C^{(1)}_{bacd},$$
$$[bd]_{3b'c'd'} \ [V_{u_R}]^3{}_a[V_{d_R}]^{b'}{}_b[V_{u_R}]^{c'}{}_c[V_{e_R}]^{d'}{}_d = -\,C^{(5)}_{bacd},$$
$$[a'c', \ a'd', \ b'c', \ b'd']_{3bcd} = [ac, \ bc, \ ad, \ bd]_{cd3b}.$$

where the V_{u_L}, V_{d_L},... unitary matrices arise from the rotation of gauge eigenstates to physical ones (see Section I.3 on p.20). All these rotation matrices are unknown. Without theoretical prejudice, we have *a priori* no information neither on the flavour structure of operator coefficients in the gauge basis. We are therefore practically constrained to consider only the coefficients of the physical basis operators and to carry out independent measurements of each flavour variant, at some energy scale. The standard-model running and mixing of operator coefficients [32, 38] can also be used to translate the bounds obtained to different scales.

However, the fact that the physical basis derives from a gauge one implies some correlations. All coefficients of $DUDN$ operators,

$$\tilde{O}^{(s)} \equiv \tilde{\epsilon} \quad \overline{D^c}(\tilde{a}P_L + \tilde{b}P_R)U \quad \overline{D^c}(\tilde{c}P_L + \tilde{d}P_R)N,$$

can in particular be expressed in terms of the above $UDUE$ ones and of the CKM as well as PMNS mixing matrices:

$$[\tilde{a}\tilde{c}]_{abcd} = [ac]_{a'b'c'd'} \ [V_{\text{CKM}}]^{a'}{}_a[V^\dagger_{\text{CKM}}]^{b'}{}_b[V_{\text{CKM}}]^{c'}{}_c[V_{\text{PMNS}}]^{d'}{}_d,$$
$$[\tilde{b}\tilde{c}]_{abcd} = [bc]_{a'b'c'd'} \ [V_{\text{CKM}}]^{a'}{}_a[V^\dagger_{\text{CKM}}]^{b'}{}_b[V_{\text{CKM}}]^{c'}{}_c[V_{\text{PMNS}}]^{d'}{}_d,$$
$$[\tilde{a}\tilde{d}] = 0 = [\tilde{c}\tilde{d}].$$

Remarkably, those relations only involve measurable flavour structures and are therefore, in principle, experimentally accessible.

As the same combinations of $C^{(3)}$ and $C^{(4)}$ appear in both $UDUE$ and $DUDN$ physical operators, the effects of $O^{(3)}$ and $O^{(4)}$ seem to be experimentally indistinguishable. The symmetry of $O^{(2)}$ in flavour space also gets exported to $[ad]$.

Processes

The $O^{(s)}$ and $O^{(t)}$ operators give rise to single top production ($UD \to \bar{t}E^+$) or decay ($t \to \bar{U}\bar{D}E^+$) at hadron colliders.

The spin- and colour-summed squared amplitudes for those processes can be obtained using the prescription of Ref. [64] for the Feynman rules of fermion-number-violating interactions. When all fermion masses are neglected with respect to the top one, one single expression is obtained:

$$\sum_{\text{spins, colours}} |M|^2 = \frac{24}{\Lambda^4} \left[(p_t \cdot p_D)(p_U \cdot p_E)\,(A + C) \right. \tag{15}$$
$$-(p_t \cdot p_U)(p_D \cdot p_E)\,C$$
$$\left. +(p_t \cdot p_E)(p_D \cdot p_U)\,(B + C) \right]$$

where Λ is the effective scale of the Lagrangian (1) and the dimensionless parameters,

$$A \equiv |ac|^2 + |ad|^2 + |bc|^2 + |bd|^2$$
$$B \equiv |a'c'|^2 + |a'd'|^2 + |b'c'|^2 + |b'd'|^2$$
$$C \equiv \operatorname{Re}\{(ac)^*(a'c') + (bd)^*(b'd')\}$$

have been introduced for convenience. They arise respectively from the square of $O^{(s)}$, of $O^{(t)}$, and from their interference and inherit the flavour dependence of the a, b, b',... coefficients.

Decay The top baryon-number-violating decay rate is then given by

$$\Gamma_t^{\text{BNV}} = \int\limits_0^{m_t/2} \mathrm{d}E_E\, \frac{m_t^2 E_E^2}{32\pi^3 \Lambda^4} \left[\left(\frac{A}{3} + B + C \right) \left(1 - \frac{2E_E}{m_t} \right) + \frac{A}{6} \right]$$
$$= \frac{m_t^5}{192\pi^3} \frac{1}{16\Lambda^4} [A + B + C],$$

for E_E the energy of the charged lepton in the top rest frame (see Fig. 1). Note the lepton energy spectrum is made harder by $O^{(s)}$ while the contribution due to $O^{(t)}$ vanishes at the endpoint of the spectrum. The hardest spectrum is obtained with a negative interference, for $A = 4B = -2C$.

Neglecting the baryon-number-violating contribution to the total width of the top fixed at 1.4 GeV, a branching fraction of

$$\text{Br}_t^{\text{BNV}} = 1.2 \times 10^{-6} \quad [A + B + C] \quad \left(\frac{m_t}{173 \text{ GeV}} \right)^5 \left(\frac{1 \text{ TeV}}{\Lambda} \right)^4$$

is obtained. At the LHC, assuming $t\bar{t}$ production rates of 170, 250 and 810 pb at 7, 8 and 13 TeV centre-of-mass energies, there would be approximately 0.40, 0.58 and 1.9 baryon-number-violating top decay(s) per inverse femtobarn times the sum over flavours of $A + B + C$.

Those flavours of the decay products are of marginal importance for the rate but heavy flavours could be tagged at colliders. The finite b-quark mass
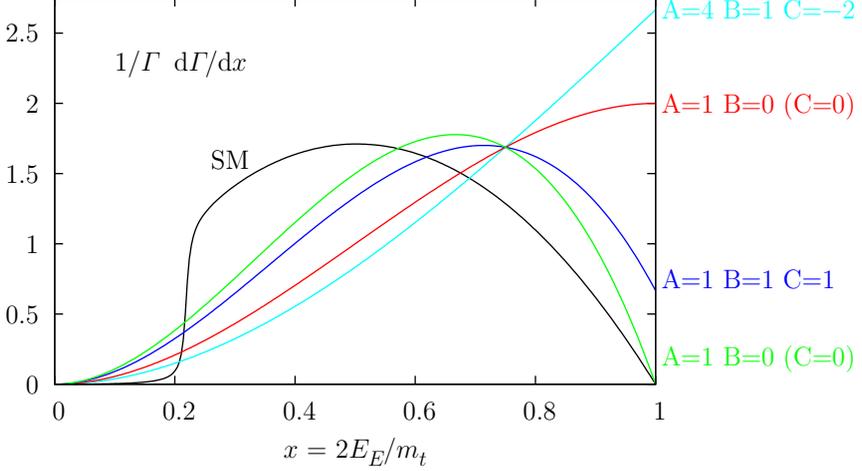
Fig. 1: Comparison of the charged lepton energy spectrum, in the top rest frame, for baryon-number-violating $t \to \bar{U}\bar{D}E^+$ and standard-model $t \to bE^+\nu$ decays. All fermion masses have been neglected with respect to the top one.

effect on the partial baryon-number-violating width could be of the order of 10%. Retaining the full m_D dependence, the partial width actually becomes

$$\Gamma_t^{\mathrm{BNV}}(m_D) = \frac{m_t^5}{192\pi^3} \frac{1}{16\Lambda^4} \left\{ [A + B + C] \, (1 + \delta_1) + [2D + E] \, \delta_2 \right\}$$

$$\text{for} \quad D \equiv \mathrm{Re}\{(ac)(bc)^* + (ad)(bd)^*\},$$
$$E \equiv \mathrm{Re}\{(ad)(b'd')^* + (bc)(a'c')^*\},$$

$$\text{and} \quad \delta_1 \equiv -x_D^2(9 - 9x_D^2 + x_D^4) - 24x_D^4 \ln x_D,$$
$$\delta_2 \equiv 4x_D \left[(1 - x_D^2)(1 + 10x_D^2 + x_D^4) + 12x_D^2(1 + x_D^2) \ln x_D \right],$$
$$\text{with} \quad x_D \equiv m_D/m_t.$$

Taking $m_D = 4.7$ GeV, the corrections amount to $\delta_1 \simeq -0.0066$ and $\delta_2 \simeq 0.11$.

Production On the other hand, the single top baryon-number-violating production in association with a charged lepton has a partonic cross section of

$$\hat{\sigma}_t^{\mathrm{BNV}} = \frac{1}{96\pi\Lambda^4} \int_{-\hat{s}+m_t^2}^{0} d\hat{t} \left[A\frac{\hat{t}(\hat{t} - m_t^2)}{\hat{s}^2} + B\frac{\hat{s} - m_t^2}{\hat{s}} + 2C\frac{\hat{t}}{\hat{s}} \right]$$

$$= \frac{\hat{s}}{96\pi\Lambda^4} \left(1 - \frac{m_t^2}{\hat{s}} \right)^2 \left[\left(\frac{A}{3} + B + C \right) + \frac{m_t^2}{\hat{s}} \frac{A}{6} \right],$$

with the Mandelstam variables $\hat{s} \equiv (p_U + p_D)^2$ and $\hat{t} \equiv (p_U - p_E)^2$, again when all masses are neglected with respect to the top one. The limited range of applicability of the effective theory manifests itself in a $\hat{s}$ growth. For a

partonic centre-of-mass energy larger than Λ, this result cannot any longer be trusted. To avoid overestimating the LHC rates, we will therefore impose a cut at $\sqrt{\hat{s}} < \Lambda$ in what follows. This simple and model-independent means of ensuring unitarity has actually little impact for processes initiated by sea quarks whose parton distribution functions inside the proton (PDFs) quickly decrease at large momentum fractions.

The U and D flavour dependence of the total cross section is strong. The ud initial state would be the most PDF-favoured while the cb one would be the most suppressed. Cross sections at the 8 and 13 TeV LHC for those two extreme cases as well as for the intermediate ub one are displayed in Table 4. The rates of the corresponding charge-conjugate processes are also given for completeness.

	σ [fb]		$ud \to \bar{t}E^+$		$ub \to \bar{t}E^+$		$cb \to \bar{t}E^+$	
A	B	C	$\bar{u}\bar{d} \to tE^-$		$\bar{u}\bar{b} \to tE^-$		$\bar{c}\bar{b} \to tE^-$	
1	0	0	330	(610)	58	(160)	2.7	(12)
			20	(62)	5.7	(22)	2.7	(12)
0	1	0	940	(1 700)	160	(460)	7.3	(32)
			54	(170)	16	(61)	7.3	(32)
1	1	1	2 200	(4 100)	380	(1 100)	17	(76)
			130	(410)	37	(140)	17	(76)

Table 4: Cross section [fb] for baryon-number-violating single top production in association with a charged lepton at the 8 (13) TeV LHC. The CTEQ6L1 PDF [65] has been used. A cut on the partonic centre-of-mass energy has been set at $\sqrt{\hat{s}} < \Lambda = 1$ TeV to avoid overestimating the cross sections with contributions from a regime where the effective field theory predictions can potentially not be trusted. No other cuts were applied.

Signatures and backgrounds

As only baryon number violation could be responsible for the production of a single charged lepton without any missing energy at the LHC. The processes we are considering do not have any irreducible background.

Decay   For definiteness, in baryon-number-violating top decay, let us fix the flavours of final-state fermions as in $t \to \bar{b}\bar{c}\mu^+$. This combination features two pairs of the same generation. After standard-model top pair production, one baryon-number-violating and one fully-hadronic decay are demanded. A single lepton with five jets and no true missing energy is therefore the sought signature.

The required hadronic decay of the second top avoids the presence of true missing energy is signal events. This feature could be used for backgrounds rejection. A more refined analysis could also include semi-leptonic decays. The
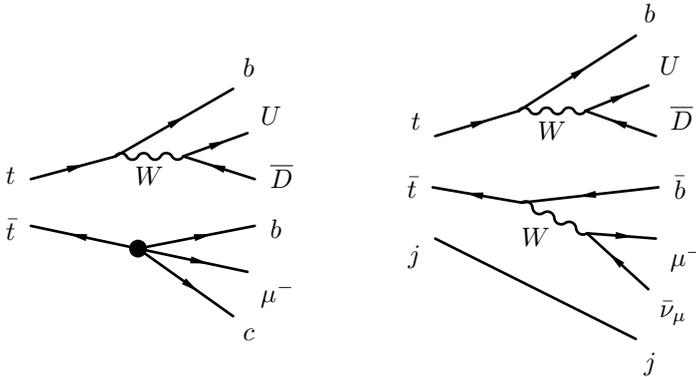
Fig. 2: The baryon-number-violating top decay signature and main semi-leptonic standard-model background. The signal features no true missing energy and two same-sign b quarks while a neutrino and opposite-sign b quarks characterize the background.

dominant background would be top pair production plus an extra jet with one fully-hadronic and one semi-leptonic top decay (see Fig. 2). The presence of a neutrino carrying away some true missing energy and the sign of the b quark in the standard-model semi-leptonic top decay differentiate it from the signal. Semi-leptonic decays of the b quark could provide a handle on this feature. The top decaying in a baryon-number-violating fashion would then preferentially lead to a same-sign lepton pair with a soft lepton on average harder than in a standard-model decay. Experimental collaborations use this feature in top-quark charge measurements [66]. The standard-model production of a W in association with five jets could also mimic the signal but would be suppressed by b tagging and the reconstruction of two tops.
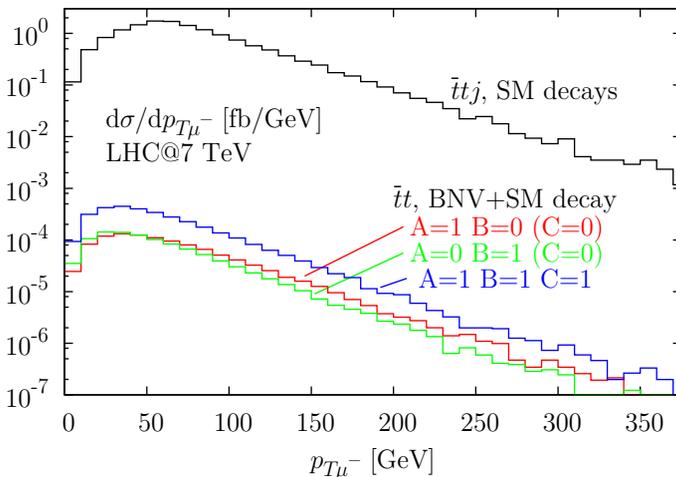
Fig. 3: Transverse momentum of the charged lepton in top baryon-number-violating decay with respect to standard-model semi-leptonic $\bar{t}tj$ background, at the 7 TeV LHC. Leading-order and parton-level results obtained with MADGRAPH5 [67]. $\Lambda = 1$ TeV is chosen. Five central jets are required ($p_{Tj} > 40$ GeV, $|\eta_j| < 2.5$, $\Delta R_{jj} > 0.5$), a central isolated lepton ($|\eta_\mu| < 2.5$, $\Delta R_{j\mu} > 0.5$) and low missing transverse energy ($\not{E}_T < 30$ GeV). The reconstruction of two tops from the visible particles has not been imposed.

The baryon-number-violating interaction has been implemented in a FEYN-RULES [68]–MADGRAPH5 [67] model in order to allow for signal simulation. The lepton transverse momentum for signal and background obtained from this framework is for instance displayed in Fig. 3. The standard-model background distribution appears distorted towards higher transverse momenta by a missing transverse energy cut. It would otherwise be softer than the signal one, as expected from Fig. 1.

Production For definiteness, in baryon-number-violating top production, let us fix the fermion flavours as in $cb \to \bar{t}\mu^+$. Other flavours in the initial state would lead to a significant enhancement of the signal (see Table 4) so that our quantitative conclusions will be conservative. As in the decay case, the anti-top is chosen to decay fully hadronically in order to avoid true missing energy in the final state. The signal signature is therefore a single lepton with three jets reconstructing a top and no true missing energy (see Fig. 4).
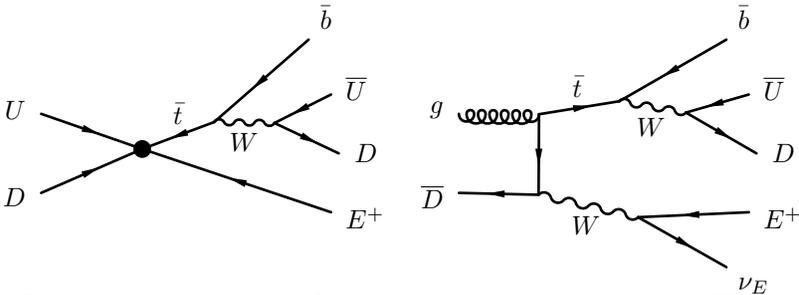
Fig. 4: Baryon-number-violating single top production signature and main tW associated production background. The missing transverse energy carried away by the neutrino notably distinguishes the two processes.

The dominant backgrounds are tW associated production and $Wjjj$ with a fully hadronic decay of the top and leptonic decays of the Ws. The tagging of a b and the reconstruction of a top quark from the three jets significantly suppress $Wjjj$. The absence of true missing energy in the signal also provides a crucial handle to disentangle it from the backgrounds. In the signal, the lepton directly produced from the initial state has a much harder spectrum than the one originating from the W decay in backgrounds (see Fig. 5). The use of a high transverse momentum selection cut would best be combined with a boosted top reconstruction technique but we will not use such algorithm in this simplified analysis. As stressed before, the effective theory can only be trusted in the low-energy regime. So we impose $\sqrt{\hat{s}} < \Lambda$ with $\Lambda = 1$ TeV to avoid overestimating the baryon-number-violating top production signal.

Requiring a 150 GeV cut on the lepton transverse momentum in the $cb \to \bar{t}\mu^+$ signal and $\bar{t}W^+$ and W^+jjj backgrounds, a sensitivity $S/\sqrt{S+B}$ of at least 5 is obtained for:

$$(A+B+C)^{1/4}/\Lambda \quad \geq \quad \begin{cases} 2.6 \ (1.7) \ \text{TeV}^{-1}, & \text{if } B = 0 = C, \\ 2.0 \ (1.3) \ \text{TeV}^{-1}, & \text{if } A = 0 = C, \\ 1.6 \ (1.2) \ \text{TeV}^{-1}, & \text{if } A = B = C, \end{cases}$$

with 20 (100) fb^{-1} of pp collisions collected at 8 (13) TeV.
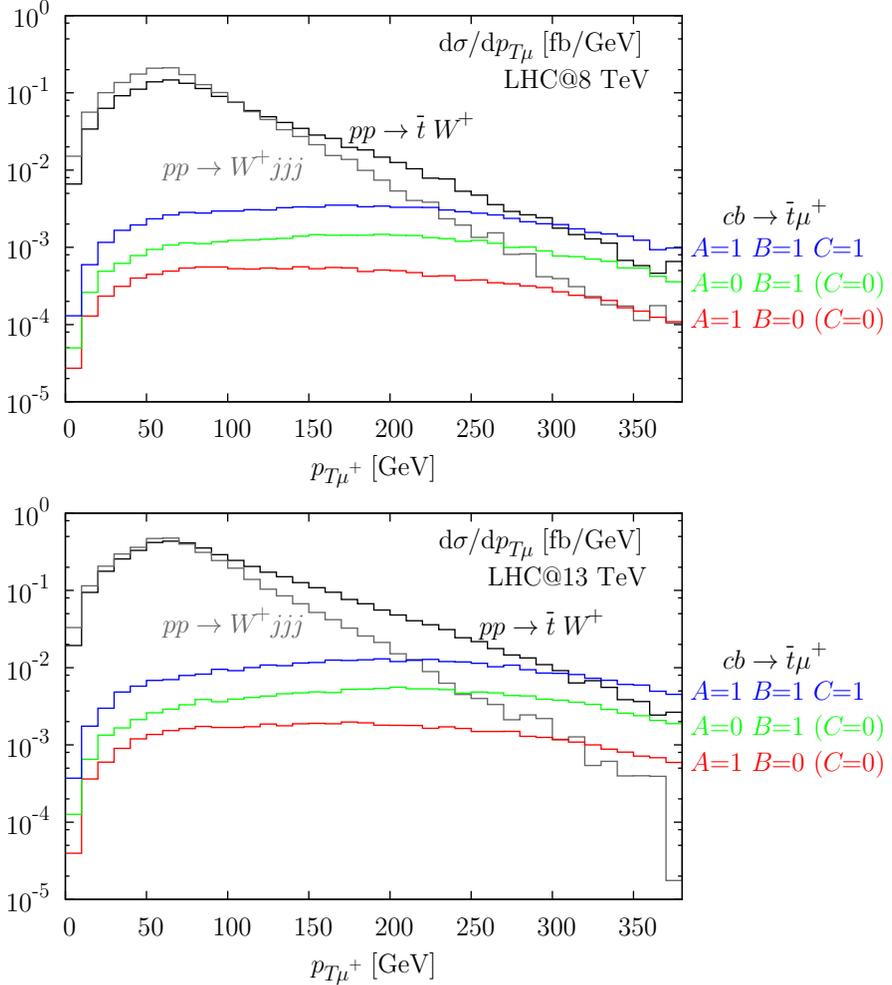
Fig. 5: Lepton transverse momentum in baryon-number-violating top plus lepton production compared with standard-model backgrounds. The top quarks and Ws are hadronically and leptonically decayed, respectively. $\sqrt{\hat{s}} < \Lambda = 1$ TeV has been imposed to avoid signal overestimation. Selected events have three central jets ($p_{Tj} > 40$ GeV, $|\eta_j| < 2.5$, $\Delta R_{jj} > 0.5$) with an invariant mass close to the top one ($|m_{jjj} - m_t| < 40$ GeV) and at least one jet tagged as a b-jet (with a flat tagging efficiency of 70, 10 and 1% for b-, c- and light-jets). We also required a central isolated lepton ($|\eta| < 2.5$, $\Delta R_{j\mu} > 0.5$) and little missing transverse energy ($\not{E}_T < 30$ GeV). These leading-order and parton-level results have been obtained with MADGRAPH5 [67].

II.3 Direct constraints

The CMS collaboration searched for the baryon-number-violating top decay described here. An analysis of 19.5 fb^{-1} at 8 TeV [69] improved the limit set with 5 fb^{-1} of 7 TeV collisions [70]. The branching fraction for the decay of a top in a muon (an electron) and two jets was constrained to be smaller than 1.6 (1.7) $\times 10^{-3}$ at the 95% confidence level. This is advertised as the first direct limit obtained on a baryon-number-violating process involving the top quark. It translates into a $\mathcal{O}(6 \text{ TeV}^{-1})$ bound on $(A + B + C)^{1/4}/\Lambda$.

In this search, a tight and a basic event selection are defined. The tight selection is contained in the basic one but is expected to feature a much larger fraction of signal events, if present. This subdivision allows for a significant reduction of uncertainties as simulation is only required to estimate the fraction of events from the basic selection that is also contained in the tight one. This expectation, for a given baryon-number-violating branching fraction, can then be compared to the measured one and a limit (or, estimate) computed.

Events in the basic selection contain a central and isolated electron or muon, no extra lepton and five jets amongst which at least one is *b* tagged. The tight selection also requires little missing transverse energy and the reconstruction of two top quarks.

II.4 Indirect constraints

The two references that considered the possibility of baryon-number-violating interactions involving the top quark [62, 63] also studied the indirect bounds that would derive from the fantastic limits on nucleon decays. *What do the bounds on nucleon-decay operators imply on processes involving a top?* or, reciprocally: *What would be the consequences, on nucleon stability, of a significant baryon-number-violating operator involving top?* are the interrogations raised.

They dramatically question the flavour structure of baryon-number-violating operators in the physical basis (the one we have direct experimental access to). They ask about the relation between processes bound to the lightest fermion generations, on the one hand, and free to involve all three generations, on the other hand. The large separation between the energies they involve sources this difference in their nature. Our knowledge about flavour is however limited to the fermion masses, CKM and PMNS matrices. Our ignorance about flavour beyond these parameters forbids us to give general and unambiguous answers to those questions. The following discussion of indirect limits will thus always rely the assumption that a single flavour variant of baryon-number-violating operator dominates all the others at the TeV scale. Were it not the case, many different and unknown contributions of the same magnitude would be expected at the GeV scale, from renormalization-group running and mixing of coefficients or from fixed order corrections. No firm conclusion on the rate of a given process could then be drawn. This assumption that a single coefficient dominates makes our discussion but indicative.

A dominant baryon-number-violating interaction between four physical fermions of the first generations at the TeV scale is required to be tiny. Its running

down to the nucleon mass scale is expected to be roughly of order one. At this scale, it could then contribute to matter instability through diagrams like:
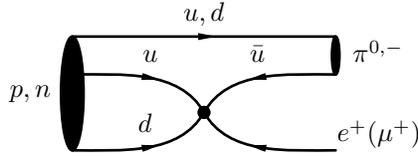
The corresponding nucleon partial decay width would scale as $\Gamma_N \simeq C_{udu\ell}^2 m_p^5 / \Lambda^4$. Given that the proton partial lifetime in that channel is bound to be larger than about 10^{34} years, an approximate limit of 10^{-13} TeV^{-1} would apply on $\sqrt{C_{udd\ell}}/\Lambda$.

The initial idea to constrain operators coefficients involving higher generations was put forward by Maciano [71]. He suggested that the baryon-number-violating decay of a tau lepton into a proton would indirectly be constrained through diagrams featuring a W emission, like:
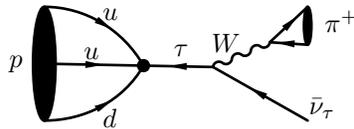
It would naively give a contribution to the proton width of the order of $C_{udu\tau}^2 G_F^2 m_p^{11} / m_\tau^2 \Lambda^4$. A constraint of the order of 10^{-11} TeV^{-1} on a dominant $\sqrt{C_{udu\tau}}/\Lambda$ coefficient would then derive from a limit on the partial lifetime of the proton of the order of 10^{32} years in that channel [15].

Hou, Nagashima and Soddu [63] noted that the same reasoning could be applied for baryon-number-violating interactions involving the top quark. So,
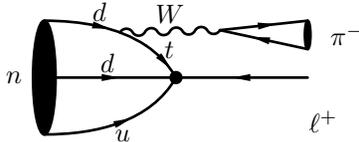
would provide an indirect constraint of the order of 10^{-9} TeV^{-1} on a dominant $\sqrt{C_{tdu\ell}}/\Lambda$ coefficient at the TeV scale (substituting the tau mass for the top one and appending a $|V_{td}|^2$ CKM factor to the previous partial width estimate). Clearly, in this case, there would be little hope to observe a $tdu\ell$ interaction at colliders. A baryon-number-violating top decay to first generations fermions $t \to \bar{u}\bar{d}\ell^+$ would indirectly be constrained to have a branching fraction smaller than 10^{-44}!

Baryon-number-violating interactions featuring more higher-generation fermions could be constrained in the same way, with more virtual W emissions giving rise to final-state pions. However, with three heavy quarks in the core baryon-number-violating interaction, say a $tbc\mu$ vertex, the bound obtained would become rather loose. The contribution to the nucleon decay width would scale as $C_{tbc\mu}^2 G_F^6 m_p^{23} |V_{td}V_{ub}V_{cd}|^2 / m_t^2 m_b^2 m_c^2 \Lambda^4$ multiplied by a four-body

phase space volume[12] factor of the order of 10^{-7} and the channel-independent limit on the nucleon lifetime would be of the order of 10^{30} years. Therefore, $\sqrt{C_{tbc\mu}}/\Lambda$ would only be indirectly constrained in that way to be smaller than 10^{-2} TeV^{-1}. A significantly looser bound of about $10^{+3.5}$ TeV^{-1} is even obtained in Ref. [63] by taking into account the amplitudes to form a nucleon and three pions out of quarks.

In Ref. [62], W loops instead of emissions were used to relate baryon-number-violating interactions involving higher generations and nucleon decay:
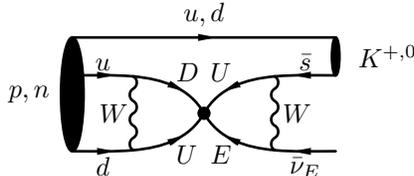
Diagrams of the form:
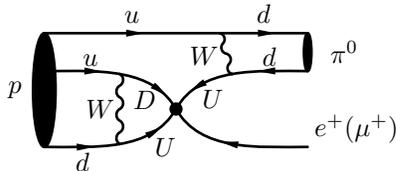
could have been imagined too. In the former case, fixing the inner fermions flavours to $tbc\mu$, we could expect the partial decay width to scale as $C_{tbc\mu}^2 \, m_p^{17}$ $(G_F^2/16\pi^2)^2 \, |V_{ub}V_{td}V_{cs}|^2 \, / \, m_t^2 m_b^2 m_c^2 \Lambda^4$. An indirect upper bound of the order of 10^{-7} TeV^{-1} would then apply on a dominant $\sqrt{C_{tbc\mu}}/\Lambda$ at the TeV scale. This rough estimate is in approximate agreement with the more refined one of Ref. [63]. However, up to two-loop order, many diagrams could in principle contribute to this decay if several flavour variants of baryon-number-violating coefficients were to have comparable magnitudes. In particular, there would be contributions from the above two-loop topology with internal fermion flavours summed over. Some contributions might efficiently cancel each other in some GIM-like mechanism. The individual coefficients could then still have non-negligible magnitudes without conflicting with the nucleon stability constraints. No strict and general statement can be made without more insight into the flavour structure of baryon-number-violating interactions. Only direct measurements at collider can lift those ambiguities and determine whether sizeable baryon-number-violating interactions involve a top quark.

II.5 Conclusions

An effective field theory has been used to model baryon-number-violating interactions involving the top quark. Operators of lowest dimension have been considered and no assumption on their flavour structure has been made. The LHC phenomenology of single top production and decay has been described as well as the indirect bounds arising from matter stability constraints. The

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[12]The phase space volume for n massless particles with a centre of mass energy of $\sqrt{s}$ is given by $2\pi(4\pi)^{2-2n}s^{n-2} \, / \, (n-1)!(n-2)!$ [72].

latter only apply if a single operator coefficient is sizeable at the TeV. In the absence of established flavour theory beyond the CKM matrix, the direct search for baryon-number-violating top decays carried out by CMS presently sets the only unambiguous bounds on the four-fermion operators considered.

While the validity of the effective field theory applied to top decay only assumes a new physics scale larger than the top mass, an accurate effective-field-theory description of production processes at the LHC (especially through valence quarks) is only possible for characteristic scales beyond the TeV range. Below that regime, the effective theory still efficiently encodes the constraints provided by imposed symmetries but fails to provide reliable quantitative estimates. The hierarchy established by power counting breaks down and the lowest dimensional operators are no longer guaranteed to involve the most relevant sets of fields. They nevertheless still single out the simplest cases.

III. Resonances

As illustrated in the previous chapter, the flavour structures of new-physics can be of uttermost relevance for its phenomenological consequences. Precise flavour measurements however clearly exclude fully generic structures. In meson decays and mixings, the important cancellations of flavour-changing neutral currents arising from the standard-model GIM mechanism were observed not to be upset by new-physics contributions. Significant hierarchies between flavour-changing charged currents have also been measured in the quark sector. In the leptons sector, the tiny neutrino masses drastically tame the effects of large mixings.

Guided by the standard-model flavour structures, we will abandon the effective-field-theory parametrization of new physics. A qualitative description of the collider signatures expected in the presence of new resonances will be addressed. The overall conservation of symmetries will serve to establish what are the simplest sets of fields that could dominantly be involved in baryon- and lepton-number-violating interactions. The discussion of low-energy bounds will be postponed to Chapter IV. since a satisfactory quantitative treatment would require the introduction (in Section IV.1) of the *minimal-flavour-violation* framework.

III.1 Extrapolating standard-model flavour structures

As no significant deviation from standard-model flavour structures has so far been established and because fermion masses, CKM and PMNS matrices are the only known flavour parameters, examining the consequences of an extrapolation of those structures to the new-physics sector appears justified.

As discussed in Section I.3, the standard-model gauge sector possesses a large $U(N_g)^5$ global symmetry associated with the independent unitary rotations, in generation space, of each of the five fermion species: q, u, d, l and e. All the five $U(1)$ Abelian factors are however separately *anomalous*, *i.e.* broken at the quantum level [16]. Only the non-Abelian $SU(N_g)_q \times SU(N_g)_u \times SU(N_g)_l \times SU(N_g)_e$ and an anomaly-free $U(1)_{B-L}$ combination survive. The Yukawa interactions that give their masses to the fermions in the vacuum and

cause physical eigenstates to mix through gauge interactions also break explicitly the majority of the initial $U(N_g)^5$ flavour group. All non-Abelian factors are broken and the Abelian ones separately suffer the same fate: only the combinations forming the baryon and lepton numbers are conserved. Eventually, because of the anomaly just mentioned, the sole $U(1)_{B-L}$ factor of the initial $U(N_g)^5$ flavour symmetry survives. We will however not assume it is also preserved by new physics.

In the spirit of the minimal-flavour-violation ansatz, we would like to regard the Yukawas as perturbations of a flavour-symmetric theory and consider no other significant flavour breaking source is introduced by the new-physics sector. New-physics flavour structures will primarily be assumed invariant under the non-Abelian part of the standard-model flavour group:

$$SU(N_g)^5 = SU(N_g)_q \times SU(N_g)_u \times SU(N_g)_d \times SU(N_g)_l \times SU(N_g)_e.$$

The global $U(1)$ symmetries are not imposed because of their anomalous character. This flavour symmetry is only applicable to interactions between gauge eigenstates that also preserve the full $SU(3)_c \times SU(2)_L \times U(1)_Y$ standard-model gauge group.

The primary effect of the Yukawa sector is to make the fermions massive in the vacuum where the standard-model gauge group gets broken down to $SU(3)_c \times U(1)_{EM}$ associated with colour and electric charge conservations. Let us therefore first introduce universal masses by considering Y_u, Y_d and Y_e Yukawas proportional to the identity matrix. Physical and gauge eigenstates are still indistinguishable and no flavour mixing can be generated. Left- and right-handed fields can then not any longer be treated separately so that, at this stage, one non-Abelian factor for quarks and an other one for leptons:

$$SU(N_g)_B \times SU(N_g)_L.$$

are still conserved out of the initial five.

Secondly should the non-universal character of fermion masses be taken into account. Misalignments between gauge and physical eigenstates can then occur and are actually measured. In the standard model, all the information about mixings in encoded in a non-trivial CKM matrix. The unitarity of the latter implies that flavour-changing neutral currents are highly suppressed. Its highly hierarchical character also makes flavour-changing charged currents small. Without neutrino masses, there are no observable misalignments in the leptonic sector.

III.2 The example of baryon number violation

Let us come back to the example of baryon and lepton number violation to apply those ideas about new-physics flavour structures. The underlying assumption is that the dominant sources of flavour symmetry breaking are of standard-model nature. As stressed before, precision flavour measurements are so far in complete agreement with this hypothesis.

First, imposing the full $SU(N_g)^5$ invariance dramatically restricts the gauge eigenstate combinations possibly involved in baryon- and lepton-number-violating processes. The δ_b^a flavour invariant tensors of each of the five factors pair a fermion with its conjugate do not lead to any violation. Flavour antisymmetric structures built upon ϵ_{abc} tensors are therefore to be used. Full $SU(3)_c \times SU(2)_L \times U(1)_Y$ standard-model gauge invariance then requires at least twelve fermions to be involved in baryon- and lepton-number-violating interactions [73]:

$$(duql)^{N_g}, \qquad (qque)^{N_g}, \qquad (qqql)^{N_g}, \qquad (duue)^{N_g}.$$

Those minimal combinations basically amount to three—N_g—copies of the four-fermion operators of Eq. (13), anti-symmetrized in generation space with ϵ_{abc} tensors. They therefore also conserve $B - L$ but with $\Delta B = \Delta L = \pm N_g$ which generalizes to $N_c \Delta B \in N_g \mathbb{Z} \ni \Delta L$ with a non-minimal field content [73].

Introducing universal masses explicitly breaks the flavour symmetry group to $SU(N_g)_B \times SU(N_g)_L$. The minimal field content of a baryon- or lepton-number-violating interaction is then reduced to six fermions ($6 = 2N_g = |N_c\Delta B| + |\Delta L|$) and allows for the four selection rules of Table 5. Two triads of quarks or leptons are anti-symmetrized in generation space using ϵ_{abc} tensors in each *fermionic core*. Extra Higgs fields, or equivalently, $q^c u$, $d^c q$ or $e^c l$ fermion combinations, are required to preserve the full standard-model gauge symmetry in the $(\Delta B; \Delta L) = (0; \pm 6)$ and $(\pm 1; \mp 3)$ cases.

$(\Delta B; \Delta L) =$	$(0; \pm 6),$	$(\pm 1; \pm 3),$	$(\pm 1; \mp 3),$	$(\pm 2; 0)$
	$lll\ lll\ \phi^6$	$uuu\ ell$	$ddd\ \ l^c l^c l^c\ \phi^*$	$udd\ udd$
		$quu\ lll$	$ddq\ \ e^c l^c l^c\ \phi^{*3}$	$udd\ dqq$
		$qqq\ lll\ \phi^2$	$dqq\ \ l^c l^c l^c\ \phi^{*3}$	$dqq\ dqq$
		$uqq\ ell\ \phi^2$	$udd\ \ l^c l^c l^c\ \phi^{*3}$	
		$udq\ lll\ \phi^2$		
		$uqq\ lee\ \phi^2$		
		$duu\ ell\ \phi^2$		

Table 5: Fermionic cores violating baryon or lepton numbers and compatible with $SU(N_g)_B \times SU(N_g)_L$ flavour and $SU(3)_c \times SU(2)_L \times U(1)_Y$ gauge symmetries. Antisymmetric tensors (not displayed) ensure the flavour symmetry requirement.

Next should the non-universal character of fermion masses and standard-model flavour mixings to be taken into account. Fermionic cores should now be written in terms of the U, D, E and N physical eigenstates (see Table 6). As flavour changes are costly in the standard model, the fully antisymmetric structures built upon the minimal six-fermion content would still be the dominant ones. The more a flavour variant departs from the fully antisymmetric one, the more it would be suppressed. Baryon and lepton number violating effects therefore preferentially involve all three generations.

At the GeV scale, the six-fermion structures will form operators of dimension no lower than nine. In particular, the dimension-six operators discussed in the previous chapter are not generated. Low-energy processes will thus at least

$(\Delta B; \Delta L) =$	$(0; \pm 6),$	$(\pm 1; \pm 3),$	$(\pm 1; \mp 3),$	$(\pm 2; 0)$
	$NNN\ NNN$	$UUU\ EEN$	$DDD\ E^c N^c N^c$	$UDD\ UDD$
		$UUD\ ENN$	$UDD\ N^c N^c N^c$	
		$UDD\ NNN$		

Table 6: Baryon- and lepton-number-violating physical-eigenstate fermionic cores obtained after the introduction of non-universal fermion masses and flavour mixings. As flavour changes are small in the standard model, the anti-symmetric flavour combinations are expected to be the dominant ones.

suffer a $(E/\Lambda)^{10}$ suppression. Next is the presence of third-generation fermions and charm quarks kinematically forbidden. The dominant operators involving all three generations do therefore not directly contribute to (di-)nucleon decays or neutron–anti-neutron oscillations. In low-energy processes, an additional suppression therefore arises from the kinematical constraints on the operator flavour structure.

A faithful quantitative description of the expected suppression would require a precise prescription for quantifying the effect of the flavour symmetry breaking. The minimal-flavour-violation one will be discussed in Chapter IV. In that framework, the small Yukawa couplings of light quarks involved in low-energy processes are actually responsible for the major part of their flavour suppression. In the lepton sector, flavour-changing processes are also greatly moderated by the tiny neutrino masses. Anticipating the result of this discussion, let us just mention that with the minimal-flavour-violation prescription, the (di-)nucleon and neutron–anti-neutron oscillation translate into bounds on the baryon- and lepton-number-violation characteristic scale of the order of TeV only.

III.3 LHC phenomenology

If new particles mediating baryon and lepton number violation were to be present at such a low scale, the non-local nature of their interactions could not generally be neglected at the LHC. An effective-field-theory description would then cease to be accurate.

General conclusions can nevertheless be drawn regardless of the non-local process actually taking place. If all new particles decay promptly enough, initial and final states will only contain standard-model fields. Their composition can thus be studied within the flavour framework presented at the beginning of this chapter. The flavour constraints apply globally and are therefore still expected to hold. As all fermionic generations can easily be produced at TeV colliders, the least suppressed fully anti-symmetric structures are expected to be the most relevant ones. Three-generation signatures of baryon- and lepton-number-violating interactions are therefore expected [4].

The organizing principle provided, in the effective field theory, by operator dimensions is however lost in the resonant regime. The minimal field contents listed in Table 6 would give rise to the effective operators of lowest dimension

satisfying our symmetry requirements. Out of the effective-field-theory regime, they are *a priori* no longer associated with the processes of largest rates. Resonances could enhance the cross sections of transitions between non-minimal sets of particles.

Fermionic cores classification All allowed sets of initial- and final-state particles involved in baryon- and lepton-number-violating transitions however reduce to the fermionic cores listed above provided they satisfy the minimal selection rules deriving from the imposed standard-model flavour structure: $|N_c\Delta B| + |\Delta L| = 2N_g$. All such such sets of fields would conversely be obtained by dressing the fermionic cores of Table 5 with flavour-neutral and gauge-invariant combinations of fields conserving the baryon and lepton numbers. Any number of neutral bosons or fermion–anti-fermion pairs could be added to the fermionic cores: D^μ, $\phi^\dagger\phi$, $q^c q$, $u^c u$, $d^c d$, $l^c l$, $e^c e$, etc. After the introduction of flavour mixings, the number of possibly relevant physical-eigenstate combinations is therefore considerably extended beyond the fermionic cores of Table 6. With X^0 symbolizing all neutral combinations of h, γ, Z, g, UU^c, DD^c, EE^c or NN^c and X^+ the charged $W^+ X^0$, $UD^c X^0$ or $NE^c X^0$ ones, Table 7 lists the structures requiring at most two charged $X^\pm$.

Same-sign signatures The simplest six-fermion cases may however first deserve our attention. Remarkably, they feature same-sign quarks and leptons only, *i.e.* either quarks or anti-quarks and either leptons or anti-leptons. This distinctive characteristic is in general lost with non-minimal field contents and may provide an interesting means for signal over background discrimination.

In the initial state, valence quarks have the highest parton luminosities.[13] Departing slightly from the minimal fermionic contents, initial-state gluons may interestingly lead to enhanced rates thanks to their large PDFs. In the final state, the signs of top quarks and charged leptons are the most unambiguously identified. Their lepton and baryon numbers are therefore asserted and characteristic same-sign signatures identifiable. Transitions involving the smallest number of particles, valence quarks or gluons in the initial state as well as tops and charged leptons in the final state are thus of primary phenomenological interest.

With just six fields, the $\epsilon_{abc} U^a D^b D^c \quad \epsilon_{def} U^d D^e D^f$ fermionic core may lead to the $(\Delta B; \Delta L) = (\pm 2; 0)$ production of two same-sign tops while a $(\Delta B; \Delta L) = (\pm 1; \pm 3)$ transition involving the $\epsilon_{abc} U^a U^b U^c \quad \epsilon_{def} E^d E^e N^f$ core may produce one top in association with two charged leptons. The following

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[13]The ij parton luminosity is defined from the parton distribution functions $f_{i,j}(x)$ at momentum fraction x as $\mathcal{L}_{ij}(\hat{s}) \equiv \int_0^1 \mathrm{d}x \int_0^1 \mathrm{d}y\, f_i(x) f_j(y)\, \delta(sxy - \hat{s})$, so that the total cross section writes $\sigma_\alpha(s) = \int_{\hat{s}_0}^s \mathrm{d}\hat{s} \quad \sum_{ij} \mathcal{L}_{ij}(\hat{s})\, \hat{\sigma}_{ij\to\alpha}(\hat{s})$, with $\sqrt{\hat{s}_0}$ the partonic centre-of-mass energy at production threshold.

$(\Delta B; \Delta L) =$	$(0; \pm 6),$	$(\pm 1; \pm 3),$	$(\pm 1; \mp 3),$	$(\pm 2; 0)$
$X^- X^-$:		$UUU\ NNN$	$UUU\ N^c N^c N^c$ $UUD\ E^c N^c N^c$ $UDD\ E^c E^c N^c$ $DDD\ E^c E^c E^c$	$UUD\ UUD$
X^- :		$UUU\ ENN$ $UUD\ NNN$	$UUD\ N^c N^c N^c$ $UDD\ E^c N^c N^c$ $DDD\ E^c E^c N^c$	$UUU\ DDD$
X^0 :	$NNN\ NNN$	$UUU\ EEN$ $UUD\ ENN$ $UDD\ NNN$	$UDD\ N^c N^c N^c$ $DDD\ E^c N^c N^c$	$UDD\ UDD$
X^+ :	$ENN\ NNN$	$UUU\ EEE$ $UUD\ EEN$ $UDD\ ENN$ $DDD\ NNN$	$DDD\ N^c N^c N^c$	$UDD\ DDD$
$X^+ X^+$:	$ENN\ ENN$	$UUD\ EEE$ $UDD\ EEN$ $DDD\ ENN$		$DDD\ DDD$

Table 7: Fermionic eigenstate combinations violating baryon and lepton numbers according to the minimal selection rule $|N_c \Delta B| + |\Delta L| = 2N_g$ imposed by standard-model flavour symmetries. Adding baryon- and lepton-number-conserving combinations of standard-model fields $X^{0,\pm}$ that are flavour- and colour-neutral may possibly lead to enhanced rates in transitions involving resonant new-physics states. Fermionic combinations requiring more than two charged combinations $X^\pm$ are not displayed.

processes are therefore the ones we will focus on:

$$
\begin{aligned}
(\Delta B; \Delta L) = \quad & (\pm 1; \pm 3) & & (\pm 2; 0) \\
& uc \to e^+ \mu^+ \bar{\nu}_\tau\ \bar{t}, & & dd \to \bar{t}\bar{s}\ \bar{t}\bar{s}, \\
& ug \to e^+ \mu^+ \bar{\nu}_\tau\ \bar{t}\bar{c}, & & gd \to \bar{t}\bar{s}\ \bar{t}\bar{s}\bar{d}, \\
& gg \to e^+ \mu^+ \bar{\nu}_\tau\ \bar{t}\bar{c}\bar{u}. & & gg \to \bar{t}\bar{s}\bar{d}\ \bar{t}\bar{s}\bar{d},
\end{aligned}
\tag{16}
$$

Other processes are probably more difficult to reach experimentally.

Charge asymmetries   Corresponding charge-conjugated processes may also be considered. At *pp* colliders like the LHC, the asymmetry of the initial state (in electric charge and baryon number) can however get exported to the final state. The parton luminosities of valence quarks are indeed much higher than the ones of their conjugates. The *dd*-initiated process would for instance occur

much more frequently than its $\bar{d}\bar{d}$-initiated conjugate. Much more $\bar{t}\bar{t}$ than tt pairs are expected in this case. The resulting asymmetry in production rates $\sigma_{\bar{t}\bar{t}} - \sigma_{tt}/\sigma_{\bar{t}\bar{t}} - \sigma_{tt}$ can be estimated from partons luminosities only, in resonant transitions. The partonic cross section is then peaked at centre-of-mass energies close to the production threshold so that only a narrow window of the partons luminosities actually dominates the total cross section integral. The asymmetries between dd (dg, uc, and ug) and $\bar{d}\bar{d}$ ($\bar{d}g$, $\bar{u}\bar{c}$, and $\bar{u}g$) parton luminosities at the 8 and 13 TeV LHC are displayed in Fig. 6. For a resonant intermediate state of mass above the TeV, the asymmetries often exceed 80%.
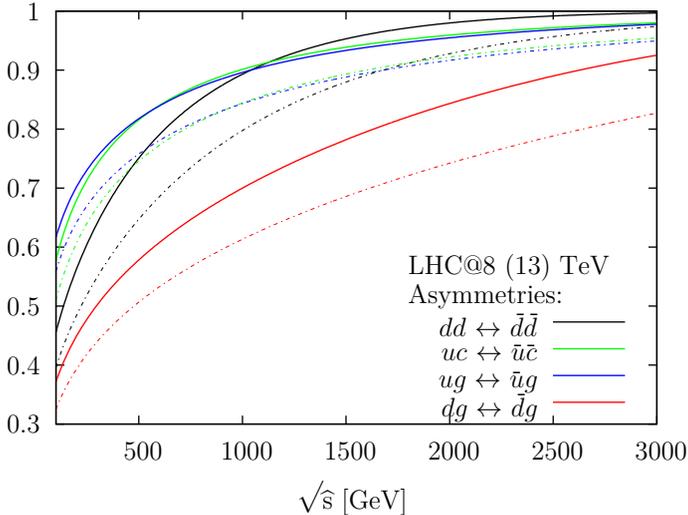
Fig. 6: Asymmetries between ij and $\bar{i}\bar{j}$ partons luminosities $\mathcal{L}_{ij}(\hat{s}) - \mathcal{L}_{\bar{i}\bar{j}}(\hat{s}) \, / \, \mathcal{L}_{ij}(\hat{s}) + \mathcal{L}_{\bar{i}\bar{j}}(\hat{s})$ for $ij = dd$, dg, uc and ug, at the 8 (solid lines) and 13 (dashed lines) TeV LHC. CTEQ6L1 parton distribution functions [65] have been used.

In the $(\Delta B; \Delta L) = (\pm 1; \pm 3)$ case, the anti-top is most easily reconstructed when decaying hadronically. Some missing energy carried away by a neutrino in the production process would render difficult the reconstruction of a semi-leptonically decaying anti-top. On the other hand, the number of hard jets is limited so that combinatorial background should not be too important even with a hadronically decaying anti-top. In that case, the clear signature provided by a single same-sign lepton pair is moreover preserved. A predominance of positively charged same-sign leptons is sourced by LHC parton luminosities while the flavour symmetry favour the pair to involve two different generations.

In the $(\Delta B; \Delta L) = (\pm 2; 0)$ case, the semi-leptonic decays of the two (anti-)tops look preferable. The ambiguous reconstruction of the tops arising from the two invisible neutrinos is to be suffered in compensation for a low-background signature featuring two same-sign leptons, two b's and at least two additional light jets. Remarkably, at the LHC, the parton luminosities favour the production of negatively charged leptons pairs. This is at odd with the naive expectation that a positively-charged pp initial state would tend to produce more positively charged leptons than negatively charged ones.

At the LHC, the dominant standard-model source of same-sign isolated leptons produced in association with b jets is $t\bar{t}W$ production [74–77] which actually features a predominance of positively charged leptons. At 8 TeV, the resulting asymmetry almost amounts to 40% [78] but is significantly diluted (down to about 10%) by a symmetric component arising from events with misidentified lepton or hadrons [74–77]. Beyond the standard model, few realistic scenarios beside baryon number violation lead, at pp colliders, to more negatively charged isolated leptons than positively charged ones. We could actually only think about a model featuring a heavy down-type quark b' having neutral flavour-changing couplings to the first generation (see Fig. 7).
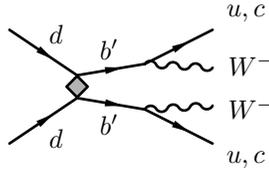
Fig. 7: Possible scenario leading to a higher production rate for negatively-charged same-sign lepton pairs than for positively-charged ones. The flavour-changing neutral current of the new heavy down-type quark b' with the first generation is significantly constrained by precision flavour measurements. Other collider measurements also restrict the existence of such a new b' state.

III.4 Simplified specific models

The standard-model flavour symmetries have been proved useful to determine in what direction to look for resonant new physics violating the lepton and baryon numbers. Beyond the effective-field-theory regime, those principles are not sufficient for making quantitative predictions in terms of theoretical parameters. This section will therefore present specific simplified models in which the signatures discussed are produced by new resonant states in the TeV range. Rates and asymmetries can then be computed for given couplings and masses.

Our intent is not to provide complete and fully realistic models but only to illustrate what sort of scenario could give rise to the signatures identified generically and to get some more quantitative estimates of rates. Though, Chapter IV. will be devoted to the further study of the second scenario sketched here.

Leptoquarks

The $(\Delta B; \Delta L) = (\pm 1; \pm 3)$ processes of Eq. (16) could be mediated by vector and scalar leptoquarks. The list of all possible tree-level bosonic mediators giving rise to the *uuu ell* and *quu lll* interactions is provided in Table 8. With full standard-model gauge invariance imposed, the two possible trilinear leptoquark interactions actually involve two vectors and one scalar or tensor. This motivated our choice of Lorentz transformation properties, even if only colour and charge conservations are imposed on combinations of physical eigenstates.

The couplings of those leptoquarks with fermions are bound to a single generation and taken chiral. The scalar (vector) leptoquarks are given an electric

	$X_S \sim (\mathbf{3,1},-1/3)_S$	$X_T \sim (\mathbf{3,1},-1/3)_T$	$X'_S \sim (\mathbf{3,1},-4/3)_S$	$X'_T \sim (\mathbf{3,1},-4/3)_T$	$X''_S \sim (\mathbf{3,3},-1/3)_S$	$X''_T \sim (\mathbf{3,3},-1/3)_T$	$X'''_S \sim (\mathbf{1,1},-1)_S$	$X'''_T \sim (\mathbf{1,1},-1)_T$	$X''''_S \sim (\mathbf{1,3},-1)_S$	$X''''_T \sim (\mathbf{1,3},-1)_T$	$X_V \sim (\mathbf{3,2},-5/6)_V$	$X'_V \sim (\mathbf{3,2},-1/6)_V$	$X''_V \sim (\mathbf{1,2},-3/2)_V$
$ue-uu-ll$	•	(•)	•	(•)			•	(•)					
$uu-ul-el$			•	(•)							•		•
$ue-ul-ul$	•	(•)									•		
$ql-uu-ll$	•	(•)	•	(•)			•	(•)					
			•	(•)	•	(•)			•	(•)			
$ll-uq-ul$							•	(•)			•		•
									•	(•)	•		•
$ql-ul-ul$	•	(•)									•		
					•	(•)					•		

Table 8: Possible combinations of tree-level bosonic mediators giving rise to the two six-fermion interactions, *uuu ell* and *quu lll*, satisfying the selection rule $(\Delta B; \Delta L) = (\pm 1; \pm 3)$ and preserving the full standard-model gauge symmetry. The possible pairing of two fermions determine the mediators gauge and Lorentz transformation properties. Bosonic mediators produce ⧩⧨ topologies while new fermions would also lead to ⧩⧨ graphs.

charge of $-2/3$ $(1/3)$ so that they can couple a neutrino (charged lepton) with a up-type quark. The trilinear vector-vector-scalar vertex has a completely anti-symmetric structure in generation space imposed by the standard-model flavour symmetry requirement. Gluons couple to the vectors through Yang-Mills–type interactions: $-\frac{1}{2}|D_\mu V_\nu - D_\nu V_\mu|^2 - ig_s V_\mu^\dagger G^{\mu\nu} V_\nu$ [79]. Same-sign lepton production in association with an anti-top is then achieved through diagrams like the ones displayed in Fig. 8. This simplified model has been implemented for simulation, in a MadGraph5 [67] model through FeynRules [68].
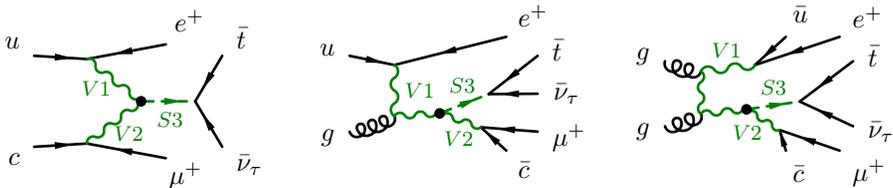
Fig. 8: Some processes leading to the production of a pair of same-sign leptons of different flavours in association with an anti-top, in the simplified leptoquark model described in the text.

Constraints   The masses of leptoquarks are constrained by searching for their QCD pair production at hadron colliders and subsequent decays to leptons (or neutrinos) and jets, with a given branching fraction (see Fig. 9a). In our

scenario, $V1$, $V2$ and $S3$ would have branching fractions of 100% to ue, $c\mu$ and $t\nu_\tau$ respectively. The limits obtained by the ATLAS and CMS collaborations for different types of leptoquarks are summarized in Table 9. None of them directly constrains the leptoquarks involved in the processes considered here. Setting the scalar (vector) masses to 500 GeV (1 TeV) is therefore a conservative choice.

Fig. 9: (a) QCD pair production constraining leptoquark masses for a given branching fraction to leptons (neutrinos) and jets. (b) Contribution of leptoquark t-channel exchange to dileptons. For a fixed mass, it bounds their couplings to fermions.

Leptoquark		Mass limit [GeV] Br = 1	Br = 1/2	Experiment			
ej (νj)	scalar	830	640	CMS	7 TeV	5/fb	[80]
		660	607	ATLAS	7 TeV	1.03/fb	[81]
μj (νj)	scalar	1070	785	CMS	8 TeV	19.6/fb	[82]
		840	650	CMS	7 TeV	5/fb	[80]
		685	594	ATLAS	7 TeV	1.03/fb	[83]
		632	523	CMS	7 TeV	2/fb	[84]
$b\nu$	scalar	450		CMS	7 TeV	4.7/fb	[85]
$b\tau$	scalar	740		CMS	8 TeV	19.7/fb	[86]
		534		ATLAS	7 TeV	4.7/fb	[87]
		525		CMS	7 TeV	4.8/fb	[88]
$b\tau$	vector	760		CMS	7 TeV	4.8/fb	[88]
$t\tau$	scalar	550		CMS	8 TeV	19.5/fb	[89]

Table 9: Lower limits on the masses of leptoquarks obtained at the LHC by searching for their QCD pair production and subsequent decays to the final states listed in the first column, with branching fraction Br. When Br $\neq$ 1, the leptoquarks are assumed to decay to a neutrino and a jet with branching fraction $1-$ Br.

For fixed masses, the couplings of the leptoquarks to first-generations fermions are constrained by the e^+e^- and $\mu^+\mu^-$ spectra measured at hadron colliders [90, 91]. The t-channel exchange of V_1 and V_2 indeed contributes to those final states (see Fig. 9b). A universal coupling to fermions of both vector and scalar leptoquarks, with magnitude fixed at 0.4, is compatible with the present uncertainties on the measured invariant mass spectra (see Fig. 10). Due to its PDF-suppressed character (especially for high momentum fractions), the $\mu\mu$ channel provides a much weaker constraint than the ee one.

The trilinear vector-vector-scalar coupling between leptoquarks remains to be bound by searching for the baryon- and lepton-number-violating processes
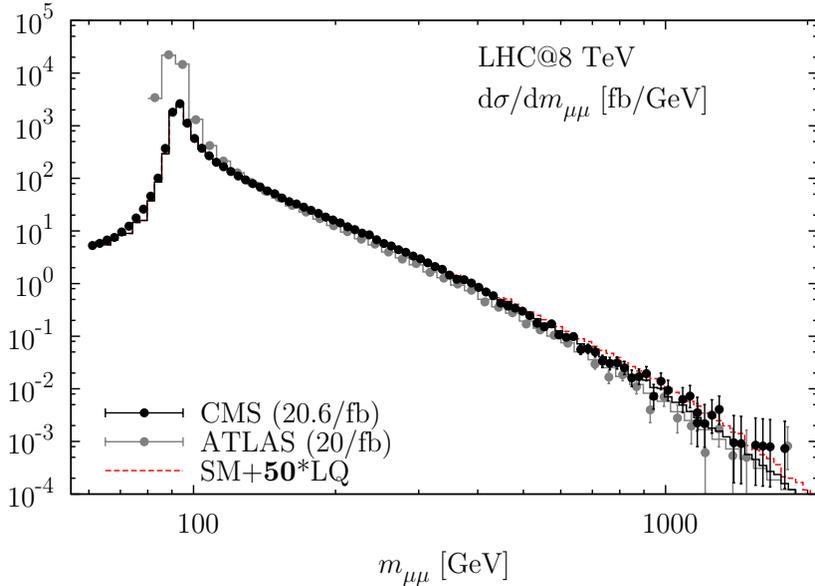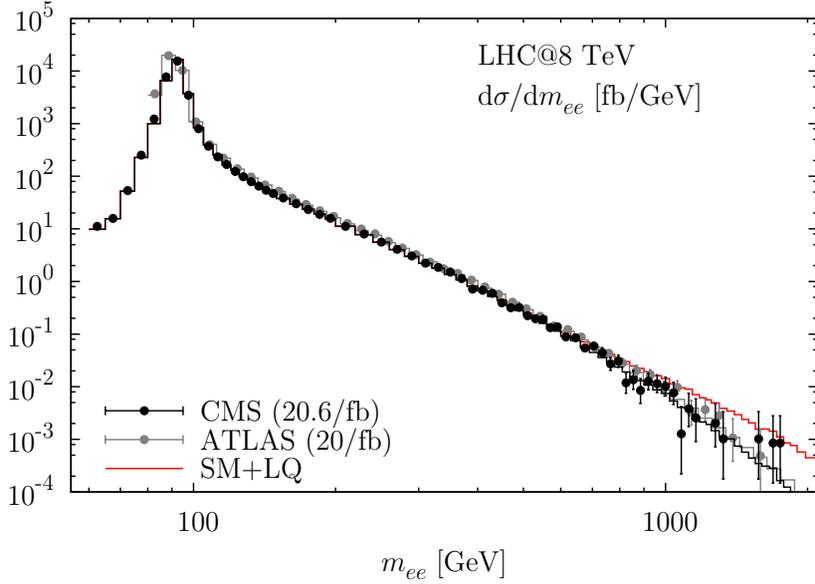
Fig. 10: The contribution of leptoquark t-channel exchange to dilepton spectra compared with CMS (Fig. 2 of Ref. [90]) and ATLAS (Fig. 1 and 2 of Ref. [91]) data and standard-model expectations. The simulated leptoquark contribution in CMS acceptance ($p_{Te} > 35$ GeV, $p_{T\mu} > 45$ GeV, $|\eta_\ell| \lesssim 2.5$) is added to CMS standard-model expectation. The ATLAS acceptance ($p_{Te} > 40$, 30 GeV, $p_{T\mu} > 25$ GeV, $|\eta_\ell| \lesssim 2.5$) is slightly larger than the CMS one in the $\mu\mu$ channel. Finite efficiencies that would reduce somewhat the leptoquark signal actually seen in detectors are neglected. In the $\mu\mu$ channel the signal is magnified by a factor of 50 to become visible on top of the standard-model expectation.

of Fig. 8. The associated (dimensionful) coupling constant is taken to be 1 TeV here.

Rates   At the 8 (13) TeV LHC, the leading order rates and the asymmetries,

$$A_{\ell_a \ell_b}^{\mathrm{NP}} \equiv \frac{\sigma^{\mathrm{NP}}(pp \to \ell_a^+ \ell_b^+ X) - \sigma^{\mathrm{NP}}(pp \to \ell_a^- \ell_b^- \bar{X})}{\sigma^{\mathrm{NP}}(pp \to \ell_a^+ \ell_b^+ X) - \sigma^{\mathrm{NP}}(pp \to \ell_a^- \ell_b^- \bar{X})}, \tag{17}$$

for the production of same-sign leptons through resonant leptoquarks are displayed in Table 10.

Processes	σ^{LQ} [fb]	$A_{e\mu}^{\mathrm{LQ}}$
$uc \to e^+ \mu^+ \bar{\nu}_\tau \ \bar{t}$	0.0011 (0.0076)	+93% (+91%)
$ug \to e^+ \mu^+ \bar{\nu}_\tau \ \bar{t}\bar{c}$	0.014 (0.24)	+96% (+95%)
$gg \to e^+ \mu^+ \bar{\nu}_\tau \ \bar{t}\bar{c}\bar{u}$	0.0018 (0.14)	0%
$pp \to e^+ \mu^+ \bar{\nu}_\tau \ \bar{t}X$	0.017 (0.39)	+78% (+45%)

Table 10: Rates for same-sign leptons production in association with an anti-top at the 8 (13) TeV LHC. Asymmetries in the production rates of positive and negative pairs. The simplified leptoquark model described in the text was implemented in MADGRAPH5 [67] through FEYNRULES [68] for producing these leading-order and parton-level results.

R-parity violation

The second scenario able to produce, through resonances, the $(\Delta B; \Delta L) = (\pm 2; 0)$ signature identified generically is a simplified R-parity violating model. The supersymmetric sector is restricted to super-QCD; all squarks are taken degenerate; one single sizeable R-parity-violating coupling is assumed between a top, a down and a strange (s)quark, with magnitude fixed to 0.1. We refer to Chapter IV. for a more detailed description.

The most relevant processes contributing to same-sign top production are displayed in Fig. 11. Their relative rates depend crucially on the mass hierarchy between squarks and gluinos. Two benchmarks have been chosen: $M_{\tilde{Q}} = 600$, $M_{\tilde{g}} = 750$ GeV and $M_{\tilde{Q}} = 800$, $M_{\tilde{g}} = 650$ GeV. Corresponding rates and charged leptons asymmetries are displayed in Table 11 for the 8 (13) TeV LHC. The values of charge asymmetries are in good agreement with the estimates provided by the parton luminosities of Fig. 6 that assumes threshold productions of the squarks and gluinos.

As stressed before, a remarkable predominance of same-sign lepton pairs of negative charges is expected in processes initiated by valence quarks, after semi-leptonic tops decays. Though, with the gluino lighter than squarks, the same-sign top production is dominated by gg-initiated processes and no total asymmetry is observable. This charge asymmetry could therefore be used to discriminate between the two families of hierarchies.
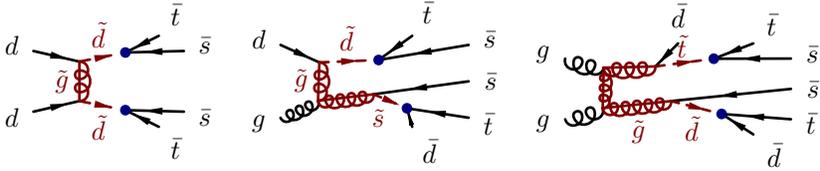
Fig. 11: Some diagrams contributing to the $(\Delta B; \Delta L) = (\pm 2; 0)$ same-sign tops plus jets signature in a simplified supersymmetric model restricted to super-QCD and featuring one single sizeable R-parity-violating coupling between a top, a down and a strange (s)quark.

Processes	$M_{\tilde{Q}} = 600,\ M_{\tilde{g}} = 750$ GeV		$M_{\tilde{Q}} = 800,\ M_{\tilde{g}} = 650$ GeV	
	σ^{RPV} [fb]	A^{RPV}	σ^{RPV} [fb]	A^{RPV}
$dd \to \bar{t}\bar{s}\,\bar{t}\bar{s}$	28 (91)	$-95\%\ (-89\%)$	0.012 (0.056)	$-98\%\ (-93\%)$
$dg \to \bar{t}\bar{s}\,\bar{t}d\bar{s}$	19 (140)	$-81\%\ (-72\%)$	1.7 (14)	$-81\%\ (-72\%)$
$gg \to \bar{t}d\bar{s}\,\bar{t}d\bar{s}$	1.4 (22)	0%	43 (490)	0%
$pp \to \bar{t}\bar{s}\,\bar{t}\bar{s}\,X$	48 (250)	$-84\%\ (-67\%)$	44 (510)	$-1.8\%\ (-1.2\%)$

Table 11: The 8 (13) TeV LHC rates and charge asymmetries (defined in Eq. (17)) for same-sign tops plus jets production through resonant squarks and gluinos, in a simplified supersymmetric scenario restricted to super-QCD and featuring one single significant R-parity-violating coupling between a top, a down and a strange (s)quark, with strength fixed to 0.1. Leading-order results obtained through the FEYNRULES-MADGRAPH5 [67, 68] software chain.

The bounds on such a scenario obtained from same-sign leptons plus b jets searches at hadron colliders will be discussed extensively in the next chapter. S-channel stop production would also lead to dijet resonances:
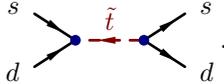
Due to a large QCD background, CMS and ATLAS are however not very sensitive to new-physics contributions in the dijet invariant mass region below 1 TeV. Limits available in this range and the estimated dijet rates for the two benchmark mass hierarchies considered here are displayed in Fig. 12. The leading-order and parton-level rates computed with MADGRAPH5 [67] are orders of magnitude below the limits.

III.5 Conclusions

The standard-model flavour symmetries constrain significantly baryon- and lepton-number-violating interactions between physical eigenstates. At least six fermions should be involved and preferentially belong to all three generations. At the nucleon mass scale, the four-fermion operators discussed in Chapter II. are in particular not generated. Neutron-anti-neutron oscillation and (di-)nucleon decays then constrain the allowed six-fermions operators to be associated with a Λ of the order of the TeV only (see next chapter).

With such a low characteristic scale, baryon- and lepton-number-violating processes would be resonant at the LHC. The global nature of the imposed
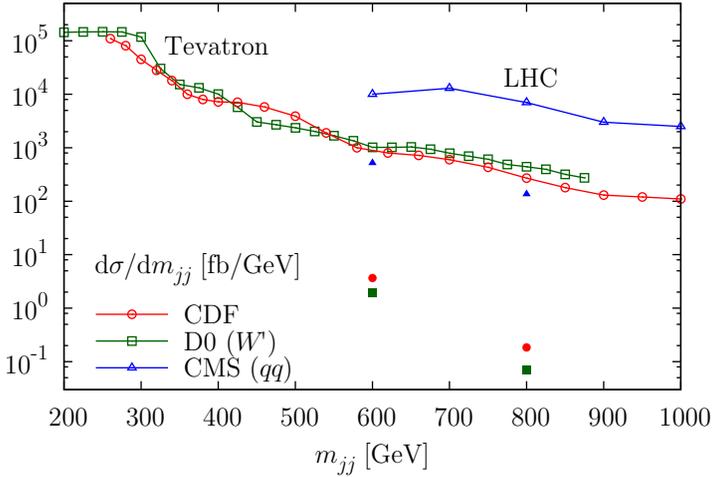
Fig. 12: Limits (empty markers) on non-standard-model dijet invariant mass spectrum below
1 TeV from – CDF [92]: 1.13/fb of $p\bar{p}$ collisions at 1.96 TeV, $|y_j| < 1.0$,
  – D0 [93]: 0.109/fb of $p\bar{p}$ collisions at 1.8 TeV, $|\eta_j| < 1.0$, $|\Delta\eta_{jj}| < 1.6$,
  – CMS [94]: 0.13/fb of pp collisions at 7 TeV, $|\eta_j| < 2.5$, $|\Delta\eta_{jj}| < 1.3$,
Leading-order and parton-level dijet rates (solid markers) in the corresponding acceptances
for the two R-parity violating benchmark scenarios featuring $M_{\tilde{t}} = 600$ and 800 GeV, respec-
tively. The limits relevant for quark jets have been selected when available.

flavour-symmetry requirement nevertheless allowed to classify the possible re-
sulting signatures beyond the effective-field-theory regime. Same-sign (anti-)
tops plus jets and same-sign leptons produced in association with a (anti-)top
were identified to be particularly promising signatures. At pp colliders the par-
ton distribution functions would favour processes initiated by valence quarks
with respect to their conjugates. A predominance of same-sign anti-tops decay-
ing to negatively charged leptons and of positively charged lepton pairs could
then respectively be observed.

   Quantitative estimates for rates and asymmetries were obtained in simpli-
fied scenarios producing the two signatures generically identified. The asymme-
tries in production rates of positively and negatively charged same-sign leptons
were shown to discriminate between new-physics scenarios. In particular, a pre-
dominance of negatively charged lepton pairs does not occur in the standard
model and is a clear indication for baryon number violation.

IV. Specific model

In the previous chapter, an extrapolation of the standard-model flavour structures to new-physics interactions provided valuable indications as to where to search for baryon- and lepton-number violation at the LHC. It rests on the experimental fact that no large deviation from the standard-model flavour structures has been observed in precision measurements. The global imposition of the flavour symmetries allowed to identify particularly interesting new-physics signatures below the regime of applicability of the effective field theory. After having studied baryon-number-violating top production and decay through four-fermion operators, we were therefore able to reach model-independent conclusions in the presence of resonant processes.

This general approach has however limitations. No quantitative estimate could be made without relying on specific models. Studying some of them can also provide useful clues as to how realistic new-physics scenarios could be realized. In this chapter, we will therefore discuss in more details the particular R-parity-violating supersymmetric model already introduced in the previous chapter. Its most generic and promising LHC signatures will then be identified and approximate bounds derived on the relevant parameters.

We will nonetheless not abandon the guidance offered by standard-model flavour-breaking patterns. Extending them to the new-physics sector was done in the previous chapter in a qualitative way. Giving quantitative estimates of the flavour suppressions suffered by operators requires a specific prescription. The minimal-flavour-violation hypothesis will be used here. Let us thus begin by introducing this method.

IV.1 Minimal flavour violation

Minimal flavour violation [95] bases itself on the flavour symmetry of the standard-model gauge Lagrangian whose explicit breakings, originating from Yukawa interactions, are treated as perturbations. The singular values of the Yukawa matrices are indeed all no larger than one. No other significant source of breaking of the flavour symmetry is introduced: flavour structures are aligned with standard-model ones.

The specific minimal-flavour-violation prescription requires the full theory to be made formally invariant under the flavour group. The Yukawas are promoted to spurious fields with the non-trivial flavour transformation properties needed to make all standard-model interactions flavour invariant. New-physics couplings are then to be written as products of those *spurions* whose flavour transformation properties compensate for the ones of the fermions involved in the corresponding interactions. Several combinations of spurions are in general possible. If the Yukawas are the only sources of breaking of the flavour symmetry, all these possible products could possibly appear, potentially combined with coefficients of order one at most.

In the vacuum, as fields are rotated to their mass eigenstates, spurions are frozen to their physical values in terms of Yukawa couplings and mixing matrix elements. The dominant term in the minimal-flavour-violation expansion is considered to provide an estimate of natural magnitude for the flavoured coupling. This estimate is minimal in the sense that new-physics flavour structures unaligned with the standard-model ones could possibly enhance it. Such enhancements are however well constrained experimentally. In this sense are the minimal-flavour-violating couplings actually maximal. Any departure from the standard-model flavour structures would in general spoil the GIM mechanism and yield flavour-changing neutral currents much larger than the ones predicted within the standard model and tightly bound by precision measurements.

Spurions The spurion transformation properties are fixed so as to make the Yukawa interactions,

$$Y_u \, \bar{u}q \, \epsilon\phi, \qquad Y_d \, \bar{d}q \, \phi^*, \qquad Y_e \, \bar{e}l \, \phi^*, \qquad +h.c.$$

invariant under the flavour symmetry of the standard-model gauge sector, $U(N_g)^5$. When no interaction violating neither the baryon nor the lepton number is to be considered, the formal invariance under the five $U(1)$'s can be restored on top of the non-Abelian $SU(N_g)$ factors. The suitable transformations of the Yukawa spurions are given in Table 12.

	$U(N_g)_\psi = SU(N_g)_\psi \times U(1)_\psi$				
	$\psi = q,$	$u,$	$d,$	$l,$	e
Y_u	$(\bar{\mathbf{3}}, -1)$	$(\mathbf{3}, +1)$			
Y_d	$(\bar{\mathbf{3}}, -1)$		$(\mathbf{3}, +1)$		
Y_e				$(\bar{\mathbf{3}}, -1)$	$(\mathbf{3}, +1)$

Table 12: Minimal transformation properties of the spurions required to render the Yukawa interactions formally invariant under the full $U(N_g)^5$ flavour group. For each $U(N_g)_\psi = SU(N_g)_\psi \times U(1)_\psi$ the representation of the non-Abelian part and Abelian charge are given.

Baryon and lepton number violation

As stressed earlier, there are however serious doubts about the fundamental character of the baryon and lepton number symmetries. The interactions be-

yond the (classical) standard model that violate these quantum numbers would inevitably involve fermions and be flavoured. Indications about the magnitude and hierarchies between the different flavour variants of the corresponding couplings can be obtained using the minimal-flavour-violation prescription [73, 96].

The conservation of all five $U(1)_\psi$ flavour groups can then obviously not be assumed. At least one of them should be broken when considering baryon number violation alone and a second one, at least, should be violated in a lepton non-conserving framework. Let us consider there is no preferred Abelian factor and take them all to be broken. The Yukawa spurions are then only given $SU(N_g)_\psi$ transformation properties, without $U(1)_\psi$ charges.

The use of the completely antisymmetric tensors associated with all five $SU(N_g)_\psi$ is then allowed. As in the previous chapter, no dimension-six operator can be constructed and a minimal content of six fermions is required. With the full standard-model gauge invariance imposed, the minimal allowed fermionic combinations (see Table 5, on p.49) are [73]:

$(\Delta B; \Delta L) =$	$(\pm 1; \pm 3)$,	$(\pm 2; 0)$
	uuu ell	*udd udd*
	qqu lll	*udd dqq*
		dqq dqq

The formal invariance under the $SU(N_g)^5$ flavour group can be recovered using products of spurions like:

$$\epsilon_{\mathsf{u_1 u_2 u_3}} \, u^{\mathsf{u_1}} u^{\mathsf{u_2}} u^{\mathsf{u_3}} \quad \epsilon_{\mathsf{l_1 l_2 l_3}} \, (Y_e^\dagger e)^{\mathsf{l_1}} l^{\mathsf{l_2}} l^{\mathsf{l_3}}, \quad \text{or,}$$

$$\epsilon_{\mathsf{q_1 q_2 q_3}} \, (Y_u^\dagger u)^{\mathsf{q_1}} (Y_d^\dagger d)^{\mathsf{q_2}} (Y_d^\dagger d)^{\mathsf{q_3}} \quad \epsilon_{\mathsf{d_1 d_2 d_3}} \, (Y_d Y_u^\dagger u)^{\mathsf{d_1}} d^{\mathsf{d_2}} d^{\mathsf{d_3}}.$$

Those simplest structures are not sufficient to produce all possible fermion flavour assignments. In the first case, the three u's for instance necessarily belong to all three generations. Different flavour variants can be obtained by inserting more spurions. A convenient means of doing so is to introduce [97]:

$$O \equiv \mathbb{I} \oplus X_d \oplus X_u \oplus X_u^2 \oplus X_d^2 \oplus \{X_u, X_d\} \oplus i[X_u, X_d] + ...$$

with $X_{u,d} \equiv Y_{u,d}^\dagger Y_{u,d}$ Hermitian. The $\oplus$ sign indicates an arbitrary linear combination should be taken with order one (possibly complex) coefficients. This matrix O has the transformation properties of the $\mathbf{3} \times \bar{\mathbf{3}} = \mathbf{1} + \mathbf{8}$ representations of $SU(N_g)_q$ and can be inserted in the simplest structures to construct the minimal-flavour-violation expansion. The replacements,

$$q \to Oq, \qquad u \to (\mathbb{I} \oplus Y_u O Y_u^\dagger) u, \qquad d \to (\mathbb{I} \oplus Y_d O Y_d^\dagger) d,$$
$$Y_u^\dagger u \to O Y_u^\dagger u. \qquad\qquad\qquad Y_d^\dagger d \to O Y_d^\dagger d,$$

would then generate all possible $SU(N_g)^5$ invariants.

The *qqu* and *uuu* combinations would for instance contribute to interactions between various species of physical up-type quark through:

$$\epsilon_{\mathsf{q_1q_2q_3}} \ (Oq)^{\mathsf{q_1}}(Oq)^{\mathsf{q_2}}(OY_u^\dagger u)^{\mathsf{q_3}}$$

$$
\begin{aligned}
&\ni \epsilon_{\mathsf{q_1q_2q_3}} \ q^{\mathsf{q_1}} q^{\mathsf{q_2}} (Y_u^\dagger u)^{\mathsf{q_3}} &\ni&& u_L c_L (y_t t_R) &\simeq& 10^0 \ u_L c_L t_R, \\
& &\ni&& u_L t_L (y_c c_R) &\simeq& 10^{-2} \ u_L t_L c_R, \\
&\ni \epsilon_{\mathsf{q_1q_2q_3}} \ q^{\mathsf{q_1}} (X_d q)^{\mathsf{q_2}} (Y_u^\dagger u)^{\mathsf{q_3}} \ni u_L (V_{tb} y_b^2 V_{ub}^* u_L)(y_c c_R) &\simeq& 10^{-8} \ u_L u_L c_R, \\
&\ni \epsilon_{\mathsf{q_1q_2q_3}} \ q^{\mathsf{q_1}} q^{\mathsf{q_2}} (X_d Y_u^\dagger u)^{\mathsf{q_3}} \ \ni \ u_L c_L (V_{tb} y_b^2 V_{ub}^* y_u u_R) &\simeq& 10^{-11} \ u_L c_L u_R.
\end{aligned}
$$

Note the phases arising, notably, from the determinant of unitary matrices

$$\epsilon_{ijk} \ U^i{}_{i'} U^j{}_{j'} U^k{}_{k'} = \det(U) \ \epsilon_{i'j'k'}$$

were dropped. As can be seen in the above example, right-handed light quarks are often associated with significant suppression arising from the corresponding Yukawa couplings. Confining all fields in the same generation leads to especially tiny prefactors. Due to the above identity, structures like

$$\epsilon_{\mathsf{q_1q_2q_3}} \ (X_d q)^{\mathsf{q_1}}(X_d q)^{\mathsf{q_2}}(Y_u^\dagger u)^{\mathsf{q_3}} = \det(X_d) \ \epsilon_{\mathsf{q_1q_2q_3}} \ q^{\mathsf{q_1}} q^{\mathsf{q_2}} (X_d^{-1} Y_u^\dagger u)^{\mathsf{q_3}}$$

are not sufficient for that purpose. Even more spurion insertions are required as for instance in:

$$
\begin{aligned}
\epsilon_{\mathsf{q_1q_2q_3}} \ q^{\mathsf{q_1}} (X_d q)^{\mathsf{q_2}} (X_u X_d Y_u^\dagger u)^{\mathsf{q_3}} &\ni \det(V_{u_L}^\dagger) \ u_L (V_{cb} y_b^2 V_{ub}^* u_L)(y_t^2 V_{tb} y_b^2 V_{ub}^* y_u u_R) \\
&\simeq 10^{-18} u_L u_L u_R.
\end{aligned}
$$

Leptonic sector

We have so far omitted discussing the leptonic sector. If only standard-model Yukawa couplings serve as spurions, there is no flavour mixing between leptons. Neutrino masses and mixings have however been firmly established experimentally.

The introduction of right-handed neutrinos and Dirac mass terms generates mixings in the leptonic sector. No new type of baryon- and lepton-number-violating structure is however constructible with the Yukawa spurion giving rise to Dirac neutrino masses. On the contrary, a Majorana mass term interestingly provides a new spurion able to render formally $SU(N_g)^5$-invariant the baryon- and lepton-number-violating four-fermion operators. Such a spurion has indeed radically different transformation properties under the flavour group. In the **6** representation of $SU(N_g)_{l(\nu)}$ for a left- (right-)handed neutrino Majorana mass term, its two indices of the same nature can be contracted with the one of the single lepton involved in dimension-six operators through a ϵ tensor [73, 96].

Left-handed Majorana mass term Let us introduce the dimension-five Weinberg operator giving rise to left-handed neutrinos Majorana masses (see Section I.3, on p.20):

$$\mathcal{Y}_\nu \ (\overline{l^c}\epsilon\phi) \ (l\epsilon\phi).$$

There is however an intrinsic ambiguity in the absolute normalization of a dimensionless spurion Y_ν defined from the dimensionful operator coefficient $\mathcal{Y}_\nu$

whose eigenvalues are experimentally accessible through neutrino mass measurements. The introduction of some scale Λ_ν is required: $Y_\nu \equiv \mathcal{Y}_\nu \Lambda_\nu$.[14] One could maybe argue that it should typically be an electroweak scale like v for the spurion Y_ν to encode the smallness of lepton number violation. A $m_\nu/v \,(\lesssim 10^{-11})$ scaling would then be obtained, similarly to the charged lepton and quark spurions. This is what we will assume in the following. Alternatively, it could be above 10^9 GeV, making Y_ν of order one, at least. As a Majorana mass term is intrinsically a $|\Delta L| = 2$ interaction, this scale could also be assumed of the same order as Λ introduced for baryon- and lepton-number-violating processes. For making apparent the consequences of a different choice, we will retain the Λ_ν/v dependence of all our results.

The new Yukawa spurion $Y_\nu = V_{\nu_L}^T y_\nu V_{\nu_L}$ can then be used in the construction of minimal-flavour-violation expansions of flavoured new-physics couplings. Its symmetric character makes the simplest

$$\epsilon_{l_1 l_2 l_3}\, l^{l_1}\, (Y_\nu^*)^{l_2 l_3} \qquad \text{and} \qquad \epsilon_{e_1 e_2 e_3}\, e^{e_1}\, (Y_e Y_\nu^* Y_e^T)^{e_2 e_3}$$

combinations vanishing. An extra $Y_e^\dagger Y_e$ insertion is therefore required [73, 96]:

$$\epsilon_{l_1 l_2 l_3}\, l^{l_1}\, (Y_e^\dagger Y_e Y_\nu^*)^{l_2 l_3} = \left(\begin{array}{c} \epsilon_{l_1 l_2 l_3}\, \nu_L^{l_1}\, (V_{\mathrm{PMNS}}^\dagger y_e^2 V_{\mathrm{PMNS}} y_\nu)^{l_2 l_3}\, \det(V_{\nu_L}^\dagger) \\ \epsilon_{l_1 l_2 l_3}\, e_L^{l_1}\, (y_e^2 V_{\mathrm{PMNS}} y_\nu V_{\mathrm{PMNS}}^T)^{l_2 l_3}\, \det(V_{e_L}^\dagger) \end{array} \right),$$

$$\epsilon_{e_1 e_2 e_3}\, e^{e_1}\, (Y_e Y_e^\dagger Y_e Y_\nu^* Y_e^T)^{e_2 e_3} = \epsilon_{e_1 e_2 e_3}\, e_R^{e_1}\, (y_e^3 V_{\mathrm{PMNS}} y_\nu V_{\mathrm{PMNS}}^T y_e)^{e_2 e_3}\, \det(V_{e_R}^\dagger),$$

where the rotation from gauge to physical eigenstates writes

$$l = (V_{\nu_L}^\dagger \nu_L,\ V_{e_L}^\dagger e_L)^T, \qquad e = V_{e_R}^\dagger e_R,$$

with the singular values decomposition $Y_e = V_{e_R}^\dagger y_e V_{e_L}$ and the definition $V_{\mathrm{PMNS}} \equiv V_{e_L} V_{\nu_L}^\dagger$. In the vacuum, the neutrinos acquire masses given by the diagonal elements of $v^2 y_\nu/2\Lambda_\nu$.

Numerics As the PMNS matrix is barely hierarchical, the flavour suppressions estimated above are mainly driven by lepton masses and neutrino mass differences. With the PMNS matrix assumed real[15] the minimal flavour sup-

---

[14]In case the above dimension-five operator is the low-energy manifestation of the presence of a heavy right-handed neutrino ν with a Dirac mass term $Y_{\nu_R}\, \bar\nu l$ and a Majorana one $M_{\nu_R}\, \overline{\nu^c}\nu$, we would have $Y_\nu = \Lambda_\nu\, Y_{\nu_R}^T M_{\nu_R}^{-1} Y_{\nu_R}$. The Λ_ν and $|M_{\nu_R}|$ scales could be distinct.

[15]The following estimates are indeed valid as long as the imaginary part of the PMNS matrix can be neglected (or, if the neutrino mass differences are of the same order as the masses themselves). Then,

$$(V_{\mathrm{PMNS}}^\dagger y_e^2 V_{\mathrm{PMNS}} y_\nu)^{l_2 l_3} - (V_{\mathrm{PMNS}}^\dagger y_e^2 V_{\mathrm{PMNS}} y_\nu)^{l_3 l_2}$$
$$= \frac{y_\nu^{l_3} - y_\nu^{l_2}}{2}\, \mathrm{Re}\{V_{\mathrm{PMNS}}^\dagger y_e^2 V_{\mathrm{PMNS}}\}^{l_2}_{\ l_3} + i\frac{y_\nu^{l_3} + y_\nu^{l_2}}{2}\, \mathrm{Im}\{V_{\mathrm{PMNS}}^\dagger y_e^2 V_{\mathrm{PMNS}}\}^{l_2}_{\ l_3}$$

as well as the off diagonal elements of

$$(V_{\mathrm{PMNS}}\, y_\nu V_{\mathrm{PMNS}}^T)^{l_2 l_3} = \max(y_\nu)\, (V_{\mathrm{PMNS}} V_{\mathrm{PMNS}}^T)^{l_2 l_3} - \sum_i (\max(y_\nu) - y_\nu^i)\, V_{\mathrm{PMNS}}{}^{l_2}_{\ i} V_{\mathrm{PMNS}}{}^{l_3}_{\ i}$$
$$= \max(y_\nu) \left[\mathbb{I} + 2i\, \mathrm{Im}\{V_{\mathrm{PMNS}}\} V_{\mathrm{PMNS}}^T \right]^{l_2 l_3} + \mathcal{O}(\Delta y_\nu)$$

are of the order of the neutrino mass differences. The last equality is obtained by using V_{PMNS} unitarity which implies $\mathrm{Re}\{V\}V^T = \mathbb{I} + i\,\mathrm{Im}\{V\}V^T$.

pressions associated with the presence of a single lepton in an interaction are respectively of the order of [73, 96]:

$$
\begin{array}{cc}
\min(m_\nu) = 0 \text{ eV} & \max(m_\nu) = 1 \text{ eV}
\end{array}
$$

$$
\frac{\Lambda_\nu}{v} \times
\begin{array}{c}
\nu_L \\ e_L \\ e_R
\end{array}
\begin{array}{c}
\begin{array}{ccc} e & \mu & \tau \end{array} \\
\left[
\begin{array}{ccc}
10^{-16} & 10^{-17} & 10^{-18} \\
10^{-16} & 10^{-17} & 10^{-19} \\
10^{-21} & 10^{-24} & 10^{-28}
\end{array}
\right]
\end{array}
,
\begin{array}{c}
\begin{array}{ccc} e & \mu & \tau \end{array} \\
\left[
\begin{array}{ccc}
10^{-18} & 10^{-18} & 10^{-20} \\
10^{-18} & 10^{-18} & 10^{-21} \\
10^{-23} & 10^{-26} & 10^{-30}
\end{array}
\right]
\end{array}
. \quad (18)
$$

where we considered two extreme choices of overall neutrino mass scale respectively characterized by $\min(m_v) = 0$ eV and $\max(m_v) = 1$ eV. Those results very mildly depend on whether the neutrino mass hierarchy is *normal* or *inverted*. Approximate expressions are:

$$
\frac{\Lambda_\nu}{v} y_\tau^2 \frac{\Delta m_{\text{atm}}^2}{v \max(m_\nu)} \times
\begin{array}{c}
\nu_L \\ e_L \\ e_R
\end{array}
\begin{array}{c}
\begin{array}{ccc} e & \mu & \tau \end{array} \\
\left[
\begin{array}{ccc}
1 & 1 & \Delta m_{\text{sol}}^2/\Delta m_{\text{atm}}^2 \\
1 & 1 & y_\mu^2/y_\tau^2 \\
y_\tau y_\mu & y_\tau y_e & y_\mu^3 y_e/y_\tau^2
\end{array}
\right]
\end{array}
.
$$

with $\max(m_\nu) \simeq \sqrt{\Delta m_{\text{atm}}^2} \simeq 0.05$ eV [15] in the $\min(m_\nu) = 0$ eV case.

Low-energy constraints

With the flavour suppression arising from minimal flavour violation taken into account, let us now examine what are the constraints on the effective-field-theory scale Λ deriving from the non observation of baryon- and lepton-number violations at low energies. As this specific prescription provides mutually consistent estimates for all the flavour variants of baryon- and lepton-number-violating operators, we need not rely on W emissions and loops to generate one from the other, as in Section II.4.

Without introducing a spurion associated with neutrino Majorana masses, the simplest selection rules allowed by the minimal-flavour-violation prescription are $(\Delta B; \Delta L) = (\pm 1; \pm 3)$ and $(\pm 2; 0)$. As mentioned in Chapter III., operators deriving from the minimal content of six fermions could give rise to neutron–anti-neutron oscillations or (di-)nucleon decays (see Fig. 13). With the Y_ν spurion introduced, $(\Delta B; \Delta L) = (\pm 1; \pm 1)$ processes involving just four fermions become permitted.

Neutron–anti-neutron oscillations   At tree-level, operators inducing $n - \bar{n}$ oscillations necessarily involve only up and down quarks. They would therefore be more suppressed than the one leading to di-nucleon decays for which strange quarks are allowed in the final state. The operators involving this second-generation quarks indeed depart less from the dominant three-generation structures.

The least suppressed combination of six first-generation quarks arises from
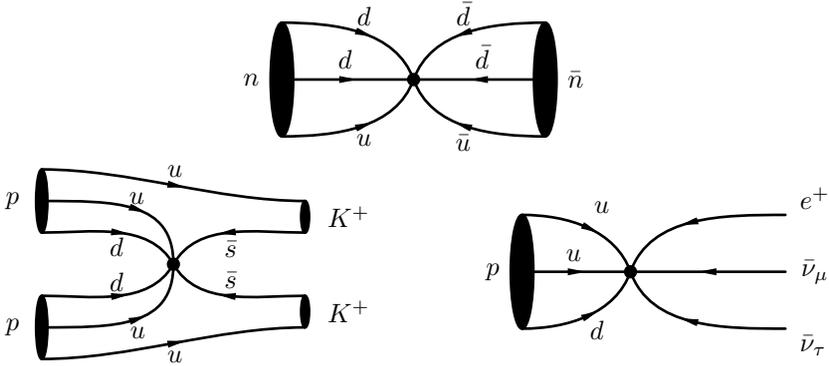
Fig. 13: Examples of diagrams leading to neutron–anti-neutron oscillations and (di-)nucleon decays in the presence of baryon- and lepton-number-violating interactions involving six fermions.

the $qqd\,qqd$ operators through flavour structures similar to

$$\epsilon_{\mathsf{q_1 q_2 q_3}}\, q^{\mathsf{q_1}}(X_u q)^{\mathsf{q_2}}(Y_d^\dagger d)^{\mathsf{q_3}} \ni \det(V_{u_L}^\dagger)\, u_L(y_t^2 V_{td} d_L)(V_{cd} y_d d_R)$$
$$\simeq 10^{-7} u_L d_L d_R,$$

so that minimal flavour violation suppresses the dimension-nine operators leading to neutron–anti-neutron oscillation by a factor of 10^{-14} at least.

The rate of oscillation can then roughly be estimated as

$$\Gamma_{n\bar{n}} \sim (10^{-14})^2 m_n^5 |\psi(0)|^4/\Lambda^{10}$$

where $|\psi(0)|^4 \simeq 10^{-5}$ GeV^6 takes into account wave function effects [98]. A comparison with the limit on the free-neutron oscillation period of the order of 10^8 s (see Section II.1, on p.31) provides a lower bound on Λ of the order of 10^{-3} TeV.

Di-nucleon decay   At tree level, di-nucleon decays to two kaons would be triggered by operators involving two strange, two down and two up quarks. The allowed presence of a second-generation quark renders the flavour suppression relatively milder than for neutron–anti-neutron oscillation operators. A $qqq\,ddq$ combination would, for instance, only suffer a 10^{-6} suppression:

$$\epsilon_{\mathsf{q_1 q_2 q_3}}\, q^{\mathsf{q_1}} q^{\mathsf{q_2}} q^{\mathsf{q_3}} \ni \det(V_{d_L}^\dagger)\, d_L s_L (V_{ub}^* u_L) \simeq 10^{-2} d_L s_L u_L,$$
$$\epsilon_{\mathsf{d_1 d_2 d_3}}\, d^{\mathsf{d_1}} d^{\mathsf{d_2}} (Y_d q)^{\mathsf{d_3}} \ni \det(V_{d_R}^\dagger)\, d_R s_R (y_b V_{ub}^* u_L) \simeq 10^{-4} d_R s_R u_R.$$

With a tentative estimate of the rate for di-nucleon decay given by

$$\Gamma_{NN} \sim (10^{-6})^2\, m_N^{11}/\Lambda^{10}$$

and a bound of the order of 10^{32} year on $pp \to K^+ K^+$ (see Section II.1), an approximate lower limit on Λ of the order of 10^2 TeV is obtained. This estimate, though extremely rough, is probably as trustable as more refined ones [99] since relevant hadronic matrix elements are basically unknown.

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Nucleon decay Without neutrino Majorana masses, minimal flavour violation only permits nucleons to decay through the *uuu ell* and *qqu lll* dimension-nine operators. The following combinations are amongst the least flavour-suppressed ones, relevant for such processes:

$$\epsilon_{\mathsf{u_1u_2u_3}}\, (Y_u X_d Y_u^\dagger u)^{\mathsf{u_1}} (Y_u X_u X_d Y_u^\dagger u)^{\mathsf{u_2}} u^{\mathsf{u_3}}$$

$$\ni \det(V_{u_R}^\dagger)(y_c V_{cb} y_b^2 V_{ub}^* y_u u_R)(y_t^3 V_{tb} y_b^2 V_{ub}^* y_u u_R)u_R$$

$$\simeq 10^{-20} u_R u_R u_R,$$

$$\epsilon_{\mathsf{l_1l_2l_3}}\, (Y_e^\dagger e)^{\mathsf{l_1}} l^{\mathsf{l_2}} l^{\mathsf{l_3}} \ni \det(V_{e_L}^\dagger)(y_\mu \mu_R) e_L \nu_{\tau L} \simeq 10^{-3} \mu_R e_L \nu_{\tau L},$$

$$\epsilon_{\mathsf{q_1q_2q_3}}\, (X_d Y_u^\dagger u)^{\mathsf{q_1}} q^{\mathsf{q_2}} q^{\mathsf{q_3}} \ni \det(V_{d_L}^\dagger)(y_b^2 V_{ub}^* y_u u_R)(V_{cd} u_L) d_L \simeq 10^{-11} u_R u_L d_L,$$

$$\epsilon_{\mathsf{l_1l_2l_3}}\, l^{\mathsf{l_1}} l^{\mathsf{l_2}} l^{\mathsf{l_3}} \ni \det(V_{e_L}^\dagger) e_L \nu_{\mu L} \nu_{\tau L} \simeq 10^0 e_L \nu_{\mu L} \nu_{\tau L}.$$

They lead to $u_R u_R u_R\ \ \mu_R e_L \nu_\tau$ and $u_R u_L d_L\ \ e_L \nu_\mu \nu_\tau$ operators respectively suppressed by a flavour factor of the order of 10^{-23} and 10^{-11}. For a nucleon lifetime larger than about 10^{31} years, lower bounds of the order of 10^{-1} and 10^1 TeV are obtained on the baryon-number-violation scale Λ. Those naive limit estimates just assume the nucleon decay rate scales as m_N^{11}/Λ^{10} and do not even include the significant phase space suppression of four- and three-body decays (of the order of 10^{-7} and 10^{-4} respectively, see footnote 12 on p.45).

When a spurion associated with neutrino Majorana masses is introduced, the most stringent constraint on Λ arises from the limit on $p \to K^+ \bar\nu$ (see Section II.1) through the $O^{(3-4)}$ operators of *qqql* form (see Eq. (14)). Given that

$$\epsilon_{\mathsf{q_1q_2q_3}}\, q^{\mathsf{q_1}} q^{\mathsf{q_2}} q^{\mathsf{q_3}} \quad \ni \epsilon_{\mathsf{q_1q_2q_3}}\, (V_{\mathrm{CKM}}^\dagger u_L)^{\mathsf{q_1}} d_L^{\mathsf{q_2}} d_L^{\mathsf{q_3}} \quad \ni 10^{-2}\ u_L d_L s_L,$$

the minimal-flavour-violation prescription indicates the $u_L d_L s_L \nu_L$ coupling should have a magnitude of the order of $10^{-18}\Lambda_\nu/v$, at most. The partial width for the induced proton decay in that channel would then be of the order of $\Gamma_p \sim (10^{-18}\Lambda_\nu/v)^2 m_p^5/\Lambda^4$. An approximate lower bound of the order of $10^4\sqrt{\Lambda_\nu/v}$ TeV on Λ would then derive. A more refined treatment could include hadronic matrix elements, phase-space factors, etc. and could modify this rough estimate by a couple of orders of magnitude.

Operators with different field contents are in general much more suppressed due to the appearance of first generations Yukawa couplings once spurions are frozen to their background values. The next-to-least suppressed quarks triads include:

$$\epsilon_{\mathsf{u_1u_2u_3}}\, u^{\mathsf{u_1}} (Y_u Y_d^\dagger d)^{\mathsf{u_2}} (Y_u q)^{\mathsf{u_3}} \ni \det(V_{u_R}^\dagger y_u V_{\mathrm{CKM}})\, (V_{ub}^* y_u^{-1} u_R)(y_s s_R) d_L$$

$$\simeq 10^{-8} u_R s_R d_L,$$

$$\epsilon_{\mathsf{q_1q_2q_3}}\, (Y_u^\dagger u)^{\mathsf{q_1}} q^{\mathsf{q_2}} q^{\mathsf{q_3}} \ni \det(V_{d_L}^\dagger)(V_{ub}^* y_u u_R) s_L d_L$$

$$\simeq 10^{-7} u_R s_L d_L,$$

so that the total flavour suppression of $u_R s_R d_L \nu_L$ and $s_L d_L u_R e_R$ interactions are respectively of the order of 10^{-24} and $10^{-28}\Lambda_\nu/v$. The approximate lower limits obtained on Λ in those cases are thus 10^1 and $10^{-1}\sqrt{\Lambda_\nu/v}$ TeV.

So, even though the flavour suppression alone does not seem to be sufficient to lower the bound on the baryon- and lepton-number-violating scale down to the TeV for the *qqql* operators, this is in general realised with other particle contents. The claim made in Section III.3 that baryon-number-violating transitions could be resonant at the LHC is thereby justified. This discussion also substitutes for the one of indirect constraints on four-fermion operators, in Section II.4, when the standard-model flavour symmetry breaking pattern is assumed. In that framework, it should be kept in mind that the very presence of four-fermion operators is conditioned to the existence of a spurion associated to Majorana neutrino masses.

IV.2 R-parity violation

After having described the minimal-flavour-violation framework and its application to baryon- and lepton-number-violating effective operators, let us move on to the specific example of a R-parity violating supersymmetric theory.

Supersymmetry[16]

In brief, supersymmetry is an extension of the Lorentz symmetry (the only non-trivial one, actually [101, 102]) relating bosons and fermions. Were it exact, each standard-model fermion would be associated with a boson (and *vice versa*) of identical mass and internal symmetry properties. The fact that this proliferation of new particles remained so far unobserved requires supersymmetry to be broken and superparticles to be significantly heavier than their standard-model partners.

Soft-breaking To effectively parametrize the supersymmetry breaking necessarily occurring in a *hidden sector*, explicit violation terms are introduced. To avoid spoiling the stability of the Higgs mass parameter protected by an exact supersymmetry, only interactions requiring dimensionful couplings are allowed. Bilinear and trilinear terms for instance provide the new scalar fields with masses much higher than the ones of their standard-model fermionic partners.

Two Higgs fields In the scalar sector, the restrictions imposed by the new symmetry require the up-type quarks, on the one hand, and down-type quarks as well as charged leptons, on the other hand, to receive their masses from two distinct Higgs doublets, ϕ_u and ϕ_d, with the quantum numbers of the standard-model ϕ and ϕ^*, respectively. Their fermionic partners, the *Higginos* $\tilde{\phi}_d$ and $\tilde{\phi}_u$, with opposite quantum numbers, do not introduce any new anomaly in the theory. The ratio of the two corresponding vacuum expectation values is defined as $\tan\beta \equiv v_u/v_d$ and $v = \sqrt{v_u^2 + v_d^2}$.

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[16]See Refs. [100] for general reviews.

R-parity violation[17]   The dramatic doubling (at least) of the available degrees of freedom is sufficient to allow baryon- and lepton-number-violation to occur through operators of dimension as low as three and four. The following interactions involving *squarks* $\tilde{q}$, $\tilde{u}$, $\tilde{d}$, *sleptons* $\tilde{l}$, $\tilde{e}$ and Higginos are indeed permitted by all symmetries

$$l\tilde{\phi}_u,$$
$$\tilde{e}^c ll, \ e^c \tilde{l} l, \ e^c l \tilde{l},$$
$$\tilde{d}^c lq, \ d^c \tilde{l}q, \ d^c l\tilde{q},$$
$$\tilde{u}dd, \ u\tilde{d}d, \ udd\tilde{d}.$$

With squarks and sleptons carrying the same baryon and lepton numbers as their standard-model counterparts, the first three kinds of interactions violate the lepton number by one unit and the last one induces a change of baryon number $|\Delta B| = 1$.

Those interactions would trigger unacceptable matter instability in case the associated couplings were all of order one and the supersymmetric particles got masses below about 10^{16} GeV, the grand-unification scale. The conservation of a new quantum number $R = (-1)^{3B+L+2S}$ is therefore commonly introduced to forbid them. It is positive for all standard-model fields and negative for their supersymmetric counterparts. Imposed *R-parity*, on top of implying baryon and lepton number conservation in the (renormalizable) minimal supersymmetric standard model, forces all supersymmetric particles to be produced in pairs and renders the lightest of them absolutely stable. Cosmology then requires it to be neutral and colourless so that it provides an invisible—*dark*—contribution to the matter density of the universe. R-parity conservation has also radical implications on the expected collider phenomenology of supersymmetric theories. The lightest supersymmetric particle, stable and produced at the end of cascade decays, notably carries out of detectors a significant amount of missing energy.

Minimal-flavour-violation hierarchies

It was however demonstrated in Ref. [96] that R-parity conservation is an unnecessary assumption when the flavoured couplings listed above are aligned with standard-model flavour structures using the minimal-flavour-violation prescription. Natural hierarchies then arise between the first-generation couplings involved in unobserved low-energy processes and the ones involving all three generations that can be much larger. The $\Delta L = \pm 1$ couplings, completely forbidden otherwise, suffer particularly severe suppressions when Majorana mass terms are introduced to account for the tiny observed departure from the assumption of massless neutrinos.

As, in supersymmetrised versions of the standard model, the up- and down-type quarks obtain their masses through interactions with two different Higgs doublets, the Yukawa we will use as spurions get physical values in the vacuum

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[17]See Ref. [103] for a comprehensive review.

that differ from the standard-model ones. Their singular values are:

$$y_u^{\text{susy}} = y_u v_u = y_u \sqrt{1 + \cot^2 \beta}, \qquad y_{d,e}^{\text{susy}} = y_{d,e} \sqrt{1 + \tan^2 \beta},$$
$$y_\nu^{\text{susy}} = y_\nu (1 + \cot^2 \beta).$$

R-parity conserving sector Aligning the flavoured soft-breaking terms with standard-model sources of flavour violation avoids the strong conflicts with flavour observables that would arise in case they were taken generic. Most notably, in the squark sector, the trilinear and bilinear couplings:

$$A_u \ \tilde{u}^c \tilde{q} \, \phi_u, \qquad A_d \ \tilde{d}^c \tilde{q} \, \phi_d,$$
$$M_q^2 \ \tilde{q}^c \tilde{q}, \qquad M_u^2 \ \tilde{u}^c \tilde{u}, \qquad M_d^2 \ \tilde{d}^c \tilde{d},$$

can be expressed in terms of spurions as

$$A_u = a_0 \, Y_u O, \qquad A_d = a_0 \, Y_d O,$$
$$M_q^2 = m_0^2 \, O, \qquad M_u^2 = m_0^2 \, (\mathbb{I} \oplus Y_u O Y_u^\dagger), \qquad M_d^2 = m_0^2 \, (\mathbb{I} \oplus Y_d O Y_d^\dagger),$$

with a_0 and m_0, two characteristic mass scales, and O defined in Section IV.1 as a linear combination of spurion products with order one coefficients at most:

$$O \equiv \mathbb{I} \oplus X_d \oplus X_u \oplus X_u^2 \oplus X_d^2 \oplus \{X_u, X_d\} \oplus i[X_u, X_d] + \ldots$$

A typical supersymmetric spectrum deriving from the minimal-flavour-violation assumption is therefore constituted of almost chiral eigenstates grouped in sets with comparable masses $\{\tilde{u}_L, \tilde{c}_L, \tilde{d}_L, \tilde{s}_L, \tilde{b}_L\}$, $\{\tilde{u}_R, \tilde{c}_R\}$ and $\{\tilde{d}_R, \tilde{s}_R, \tilde{b}_R\}$. The order one coefficients in the minimal-flavour-violation expansions allow the relative separation between the mass scales of those three sets to be sizeable. Due to the large left-right mixing introduced by $[A_u]_{33}$, which scales as the top-quark Yukawa coupling, the two top squark states, $\tilde{t}_1$ and $\tilde{t}_2$, stand apart from those sets. This would also be the case for the bottom squarks, $\tilde{b}_1$ and $\tilde{b}_2$, with $\tan \beta$ approaching or exceeding $1/y_b \simeq 40$.

R-parity violating sector The minimal-flavour-violation prescription generally implies that the lepton-number-violating couplings are incredibly suppressed. They all involve a single (s)lepton doublet l ($\tilde{l}$) whose flavour transformation property cannot be compensated by the one of an anti-(s)lepton. A Majorana neutrino spurion is required to make them formally invariant under the flavour group. In a supersymmetric framework, the flavour suppressions of Eq. (18) only get corrected by a $(1 + \tan^2 \beta)(1 + \cot^2 \beta)$ factor. They are therefore of the order of:

$$\frac{\Lambda_\nu}{v} \times \begin{array}{c} \\ \nu_L \\ e_L \end{array} \overset{\displaystyle \tan\beta = 5}{\begin{array}{ccc} e & \mu & \tau \\ \left[\begin{array}{ccc} 10^{-15} & 10^{-16} & 10^{-17} \\ 10^{-15} & 10^{-16} & 10^{-18} \end{array}\right. \end{array}}, \quad \overset{\displaystyle \tan\beta = 50}{\begin{array}{ccc} e & \mu & \tau \\ \left.\begin{array}{ccc} 10^{-13} & 10^{-14} & 10^{-15} \\ 10^{-13} & 10^{-14} & 10^{-16} \end{array}\right] \end{array}}, \quad (19)$$

where the averages of the suppressions obtained for the normal and inverted mass hierarchies, as well as for $\min(m_\nu) = 0$ eV and $\max(m_\nu) = 1$ eV have been taken. The $\tilde{e}^c ll - e^c \tilde{l} l - e^c l \tilde{l}$ interaction also admits a flavour-invariant structure of the form:

$$\epsilon_{l_1 l_2 l_3} (Y_\nu^* Y_e^T e^c)^{l_1} l^{l_2} l^{l_3} = 2 \det(V_{\nu_L}^\dagger) \, \epsilon_{l_1 l_2 l_3} (y_\nu V_{\mathrm{PMNS}}^T y_e e_R^c)^{l_1} \nu_L^{l_2} (V_{\mathrm{PMNS}}^\dagger e_L)^{l_3},$$

with its least suppressed components involving a right tau (s)lepton and having coefficients of the order of $\Lambda_\nu \max(m_\nu) y_\tau \tan\beta / v^2 \lesssim 10^{-12} \Lambda_\nu / v$. It does however not lead to direct proton decay.

On the other hand, the baryon-number-violating couplings potentially have large entries. The hierarchies between the different flavour variants of the $\tilde{u} dd - u \tilde{d} d - u d \tilde{d}$ coupling can be established using the usual minimal-flavour-violation prescription. When the conservation of none of the three $U(1)_{q,u,d}$ global symmetries is imposed, the three completely antisymmetric tensors of the corresponding $SU(N_g)$s can be used. Since the two down (s)quarks are necessarily of different flavours,[18] the minimal Yukawa insertions are actually sufficient to produce all flavour variants:

$$\epsilon_{q_1 q_2 q_3} (Y_u^\dagger u)^{q_1} (Y_d^\dagger d)^{q_2} (Y_d^\dagger d)^{q_3}$$

$$= \epsilon_{q_1 q_2 q_3} (y_d^{-1} V_{\mathrm{CKM}}^\dagger y_u u_R)^{q_1} d_R^{q_2} d_R^{q_3} \quad \det(V_{d_L}^\dagger y_d) \quad \frac{v^3}{v_d^2 v_u}$$

$$\simeq \epsilon_{q_1 q_2 q_3} \left((1 + \tan^2\beta)\sqrt{1 + \cot^2\beta} \begin{bmatrix} 10^{-10} & 10^{-8} & 10^{-7} \\ 10^{-12} & 10^{-8} & 10^{-8} \\ 10^{-15} & 10^{-11} & 10^{-8} \end{bmatrix} u_R \right)^{q_1} d_R^{q_2} d_R^{q_3},$$

$$\epsilon_{u_1 u_2 u_3} u^{u_1} (Y_u Y_d^\dagger d)^{u_2} (Y_u Y_d^\dagger d)^{u_3}$$

$$= \epsilon_{u_1 u_2 u_3} (y_d^{-1} V_{\mathrm{CKM}}^\dagger y_u^{-1} u_R)^{u_1} d_R^{u_2} d_R^{u_3} \quad \det(V_{u_R}^\dagger y_u V_{\mathrm{CKM}} y_d) \quad \frac{v^4}{v_u^2 v_d^2}$$

$$\simeq \epsilon_{u_1 u_2 u_3} \left((1 + \tan^2\beta)(1 + \cot^2\beta) \begin{bmatrix} 10^{-7} & 10^{-10} & 10^{-14} \\ 10^{-9} & 10^{-11} & 10^{-15} \\ 10^{-12} & 10^{-14} & 10^{-15} \end{bmatrix} u_R \right)^{u_1} d_R^{u_2} d_R^{u_3},$$

$$\epsilon_{d_1 d_2 d_3} (Y_d Y_u^\dagger u)^{d_1} d^{d_2} d^{d_3}$$

$$= \epsilon_{d_1 d_2 d_3} (y_d V_{\mathrm{CKM}}^\dagger y_u u_R)^{d_1} d_R^{d_2} d_R^{d_3} \quad \det(V_{d_R}^\dagger) \quad \frac{v^2}{v_u v_d}$$

$$\simeq \epsilon_{q_1 q_2 q_3} \left(\sqrt{1 + \tan^2\beta}\sqrt{1 + \cot^2\beta} \begin{bmatrix} 10^{-9} & 10^{-7} & 10^{-7} \\ 10^{-9} & 10^{-5} & 10^{-5} \\ 10^{-9} & 10^{-5} & 10^{-2} \end{bmatrix} u_R \right)^{q_1} d_R^{q_2} d_R^{q_3},$$

---

[18] The $\tilde{u} dd$ interactions is antisymmetric under the permutation of the two down-type quarks due to the antisymmetric colour and Lorentz structures, in addition to Fermi statistics. The colour structure alone renders $u \tilde{d} d + u d \tilde{d}$ antisymmetric under the exchange of the two down-type (s)quarks.

where the moduli of each entry have been taken (neglecting the unknown phases of the V_{d_L, d_R, u_R} matrices). Each of the standard-model mass eigenstate stands for either itself or its supersymmetric partner. Combining the three possible structures with order one coefficients, the maximal values obtained for the baryon-number-violating couplings are displayed in Table 13.

Full minimal flavour violation

$$\tan \beta = 5 \qquad\qquad\qquad \tan \beta = 50$$

	u_R	c_R	t_R			u_R	c_R	t_R
$s_R b_R$	10^{-6}	10^{-6}	10^{-6}	,	$s_R b_R$	10^{-4}	10^{-4}	10^{-4}
$d_R b_R$	10^{-8}	10^{-5}	10^{-4}		$d_R b_R$	10^{-6}	10^{-4}	10^{-3}
$d_R s_R$	10^{-8}	10^{-4}	10^{-1}		$d_R s_R$	10^{-7}	10^{-3}	10^{0}

Holomorphic restriction

$s_R b_R$	10^{-8}	10^{-6}	10^{-6}	,	10^{-6}	10^{-4}	10^{-4}
$d_R b_R$	10^{-10}	10^{-7}	10^{-6}		10^{-8}	10^{-5}	10^{-4}
$d_R s_R$	10^{-14}	10^{-10}	10^{-6}		10^{-12}	10^{-8}	10^{-4}

Table 13: Maximal values of the baryon-number- and R-parity-violating couplings $\tilde{u}dd - u\tilde{d}d - udd\tilde{d}$, obtained using the minimal-flavour-violation prescription, for two different values of $\tan \beta$. All three possible $SU(N_g)_{q,u,d}$ completely antisymmetric tensors are combined in the *full* minimal flavour violation while only the $SU(N_g)_q$ one is retained in its *holomorphic* restriction.

Holomorphic restriction In case the Yukawa spurions are interpreted as fields that are dynamical at a high scale and condense in the vacuum to their measured values, supersymmetry imposes an additional restriction on the type of spurion insertions allowed. No interaction involving simultaneously a spurion and its conjugate is permitted (from the *holomorphicity* of the superpotential). Examination of the above flavour structures indicates that only the ones built upon $\epsilon_{q_1 q_2 q_3}$ satisfy this condition. In the *holomorphic restriction* [104] of minimal flavour violation, it is therefore the only one allowed. The resulting couplings for $\tan \beta = 5$ and 50 are displayed in Table 13.

For lepton-number-violating interactions, the introduction of a neutrino Majorana mass term only allows for the flavour structures based on

$$\epsilon^{l_1 l_2 l_3} (Y_e^T e^c)_{l_1} (Y_\nu l)_{l_2} (Y_\nu l)_{l_3}.$$

Those are incommensurably suppressed provided Λ_ν is somewhat smaller than 10^9 GeV.

Low-energy constraints As we have seen, the transposition of standard-model flavour symmetries on the R-parity-violating couplings produces pro-

nounced hierarchies and suppresses most significantly the interactions involving light quarks that are the most relevant in low-energy baryon- and lepton-number-violating processes.

With two baryon-number-violating vertices, $\Delta B = \pm 2$ processes can be triggered. At energies much lower than the supersymmetric particles masses, they would appropriately be modelled in an effective field theory. The $u_R d_R s_R$ $u_R d_R s_R$ operator could for instance lead to di-nucleon decay into kaons (see diagram in Fig. 14) and would have a flavour suppression of the order of $(10^{-9} \tan \beta)^2$. The rough estimation procedure of Section IV.1 (on p.67) then leads to a bound on the supersymmetric particles masses of the order of 10^0 $\tan^4 \beta$ TeV. As the two down (s)quarks involved in a $\Delta B = \pm 1$ and R-parity-violating interaction are necessarily of different flavours, the neutron–anti-neutron oscillation limit is expected to be less constraining.

With a Y_ν Majorana spurion allowing for variations of the lepton number by one unit, the tree-level exchange of down-type squarks $\tilde{d}$ for instance produces the four-fermion operator $udql$ at energies well below $m_{\tilde{d}}$ (see Fig. 14). It is now clear that the flavour suppression of the resulting $u_R d_R s_L \nu_L$ interaction in the physical basis is at least of the order of $10^{-25} \tan^3\beta \, \Lambda_\nu/v$. The bound on $p \to K^+ \nu$ therefore translates to an approximate lower limit of $\sqrt{10^{-25} \, \tan^3\beta \, \Lambda_\nu/v} \, (m_N^5 \tau_N)^{1/4} \simeq 10^0 \sqrt{\tan^3\beta \, \Lambda_\nu/v}$ TeV on $m_{\tilde{s}}$.
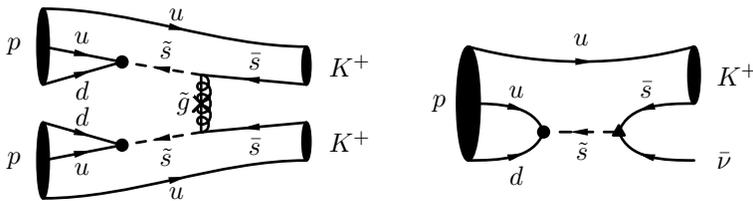
Fig. 14: Examples of diagrams leading to (di-)nucleon decay in a R-parity-violating theory.

So, with minimal flavour violation imposed, the indirect low-energy bounds obtained within a supersymmetric theory are looser than the ones obtained in Section IV.1 in an effective-field-theory framework. The least suppressed $qqql$ operator is for instance not directly generated at low-energy by R-parity-violating interactions. A characteristic scale for superparticle masses of the TeV order is therefore allowed.

IV.3 LHC phenomenology

At colliders, lepton-number- and R-parity-violating interactions of the order of the ones obtained by using the minimal-flavour-violation prescription are too small to give any observable signal. Some baryon-number-violating couplings, especially the $t_R d_R s_R$ one (or $t_R s_R b_R$ in the holomorphic restriction) can, on the other hand, have sizeable magnitudes. The low-energy constraints examined in the previous section proved not to forbid superparticles with masses at the TeV. Let us therefore study the potential phenomenology of this setup at the LHC.

As reminded above, in R-parity-conserving settings, supersymmetric partners are only produced in pairs and the lightest one is absolutely stable. Cosmology therefore imposes it to be neutral and colourless. All decay chains of supersymmetric particles produced at colliders would eventually end with such a particle that would carry away of detectors some amount of missing energy.

The presence of appreciable R-parity-violating couplings changes dramatically this phenomenology. Coloured superparticles are still mostly produced in pairs if the strength of a R-parity-violating coupling is comparable or smaller than the strong coupling constant. Though, the lightest supersymmetric particle is now unstable and needs not be electromagnetically neutral or colourless.

With sufficiently small R-parity-violating couplings, a long lifetime for the lightest supersymmetric particle can be expected. Displaced vertices may be identifiable. Nearly-stable hadrons or electrically charged particles could then be seen flighting through detectors. As will be argued below, the magnitude of R-parity-violating couplings prescribed by minimal flavour violation is generally too large for this situation to be realised.

Prompt R-parity-violating decays would, on the other hand, trade the missing-transverse-energy signatures of standard supersymmetry for additional visible activity. In a minimal-flavour-violation setup, the baryon-number-violating couplings would produce extra hadrons. Because couplings involving a top quark or squark are then the largest ones (see Table 13), a significant amount of tops can be expected.

With R-parity violation, all processes involving supersymmetric states produce (and are initiated by) standard-model particles. Two baryon-number-violating interactions are therefore required. The prompt processes inducing no net change of baryon number have fairly unspecific collider signatures like jets or $t\bar{t}$ plus jets. Unless a resonance can be identified, separating the standard-model and supersymmetric contributions to those final states may be difficult. On the other hand, the processes in which the baryon number changes by two units may be much more peculiar. In the following, we will show that same-sign tops is a widely produced final state. As was highlighted in Section III.3, this signature of baryon-number-violation could actually be constituted of more anti-top than top pairs. After their semi-leptonic decays, a striking predominance of negatively-charged leptons would then develop.

Same-sign tops

Let us argue in this section that same-sign tops are a widely present collider signature of R-parity-violating supersymmetric scenarios, when minimal flavour violation is imposed.

At hadron colliders, the lightest coloured particles would be the most copiously produced. Their R-parity-violating decays would also be kinematically favoured and often produce top quarks since the largest R-parity-violating coupling is the $\tilde{t}_R d_R s_R - t_R \tilde{d}_R s_R - t_R d_R \tilde{s}_R$ one, denoted λ''_{tds} (or λ''_{tbs} in the holomorphic restriction of minimal flavour violation).[19]

---

[19]In case the top quark is heavier than the lightest supersymmetric particles, sub-leading R-parity-violating couplings may be required in the decay of the latter. This possibility will

Let us therefore focus on the production and decay of squarks or gluinos and neglect the interactions of neutralinos, charginos and sleptons (were they much heavier than the other particles, they would actually be irrelevant). With the typical squark mass hierarchies appearing in a minimal-flavour-violation context, several cases can be distinguished, depending on the nature of the lightest coloured particles. A summary of the more detailed discussion that follows is provided in Table 14 and in the adjacent paragraphs on p.80.

Light gluino   When the gluino is lighter than squarks, its strong pair production would be followed by three-body R-parity-violating decays through virtual $\tilde{d}_R$, $\tilde{s}_R$, $\tilde{t}_1$ or $\tilde{t}_2$. The $\tilde{g} \to tds$ and $\bar{t}\bar{d}\bar{s}$ modes would have branching fractions close to $1/2$. (In the holomorphic restriction, b quarks would almost always replace the d's in gluino decays.) Using a constant matrix element approximation obtained by dimensional analysis and the phase-space volume of three massless bodies,[20] the corresponding partial width scales as:

$$\Gamma_{\tilde{g} \to tds, \bar{t}\bar{d}\bar{s}} \sim \frac{\alpha_s |\lambda''_{tds}|^2}{4(4\pi)^2} \frac{M_{\tilde{g}}^5}{M_{\tilde{d}_R, \tilde{s}_R, \tilde{t}_{1,2}}^4}.$$

For $M_{\tilde{g}} = 800$ GeV, $M_{\tilde{d}_R, \tilde{s}_R, \tilde{t}_1} = 1$ TeV and $\lambda''_{tds} = 0.1$, a partial width of about 10^{-3} GeV is obtained.

The squarks, produced in more modest proportions, would then either undergo a direct R-parity-violating decay, e.g.: $\tilde{d}_R \to \bar{s}\,\bar{t}$, $\tilde{t}_1 \to \bar{d}\,\bar{s}$, etc. ($\tilde{d}_R \to \bar{s}\,\bar{t}$, $\bar{b}\,\bar{t}$, $\tilde{t}_1 \to \bar{b}\bar{s}$ in the holomorphic restriction of minimal flavour violation) or produce gluinos, e.g.: $\tilde{u}_L, \tilde{u}_R \to u\tilde{g}$. For first-generation squarks, corresponding partial widths,

$$\Gamma_{\tilde{d}_R \to \bar{s}\bar{t}} = \frac{|\lambda''_{tds}|^2}{8\pi} M_{\tilde{d}_R} \left(1 - \frac{m_t^2}{M_{\tilde{d}_R}^2}\right)^2, \qquad \Gamma_{\tilde{Q} \to Q\tilde{g}} = \frac{2\alpha_s}{3} M_{\tilde{Q}} \left(1 - \frac{M_{\tilde{g}}^2}{M_{\tilde{Q}}^2}\right)^2,$$

scale respectively as 10^{-1} and 10^1 GeV for $M_{\tilde{Q}} = 1$ TeV while for stops, they write

$$\Gamma_{\tilde{t}_1 \to \bar{d}\bar{s}} = \frac{|\lambda''_{tds}|^2}{8\pi} M_{\tilde{t}_1},$$

$$\Gamma_{\tilde{t}_1 \to t\tilde{g}} = \frac{2\alpha_s}{3} M_{\tilde{t}_1} \sqrt{\lambda\left(1, \frac{M_{\tilde{g}}}{M_{\tilde{t}_1}}, \frac{m_t}{M_{\tilde{t}_1}}\right)} \left(1 - \frac{M_{\tilde{g}}^2}{M_{\tilde{t}_1}^2} + 2\sin 2\theta_t \frac{m_t M_{\tilde{g}}}{M_{\tilde{t}_1}^2} - \frac{m_t^2}{M_{\tilde{t}_1}^2}\right),$$

however not be considered here.

[20]With a centre of mass energy M and one single particle of non-vanishing mass xM, the two-, three- and four-body phase space volumes are:

$$\frac{1}{8\pi} \left[1 - x^2\right], \qquad \frac{M^2}{4(4\pi)^3} \left[1 - x^4 + 2x^2 \ln x^2\right],$$

$$\frac{M^4}{24(4\pi)^5} \left[1 + 9x^2 - 9x^4 - x^6 + 6x^2(1 + x^2) \ln x^2\right].$$

For $x = 173$ GeV/800 GeV, the corrections due to a non-vanishing x are of about -4.7%, -29% and -50%.

where $\lambda(a, b, c) \equiv [a^2 - (b+c)^2][a^2 - (b-c)^2]$ and (with mixing between generations neglected) θ_t determines the proportion of left and right gauge eigenstates in $\tilde{t}_1 = \tilde{q}^{a=3,i=1} \cos \theta_t + \tilde{u}^{a=3} \sin \theta_t$. For $M_{\tilde{g}} = 800$ GeV, $M_{\tilde{t}_1} = 1$ TeV and $\sin 2\theta_t = -1$, both partial widths scale as 10^0 GeV.

In this situation, no displaced vertices ($\Gamma \lesssim 10^{-12}$ GeV or $hc/1\,\mathrm{mm}$) or long-lived particles at detector scales ($\Gamma \lesssim 10^{-16}$ GeV or $hc/10\,\mathrm{m}$) would be obtained. The tagging of b-jets arising from top decays and demanding $\Gamma \gtrsim 10^{-14}$ GeV [105] poses moreover no problem. The prompt superparticle decays give rise to many tops and same-sign pairs are dominantly produced through $gg \to \tilde{g}\tilde{g}$. As gluino decays involving a top or an anti-top are equally likely, no significant $t\,t - \bar{t}\,\bar{t}$ asymmetry is expected.

Light stop A light stop would quickly decay through $\tilde{t}_1 \to \bar{s}\,\bar{d}$ (or $\bar{s}\,\bar{b}$ in the holomorphic restriction of minimal flavour violation) with a rate,

$$\Gamma_{\tilde{t}_1 \to \bar{s}\bar{d}} = \frac{|\lambda''_{tds}|^2}{8\pi} M_{\tilde{t}_1},$$

of the order of 10^0 GeV for $M_{\tilde{t}_1} = 800$ GeV and $\lambda''_{tds} = 0.1$.

The right down and strange squarks would, again, have fast R-parity-violating decays ($\Gamma_{\tilde{d}_R \to \bar{s}\bar{t}} \sim 10^{-1}$ GeV) while the other squarks could undergo a three-body $\tilde{Q} \to Q\,\bar{t}\,\tilde{t}_1$ decay if kinematically allowed ($M_{\tilde{t}_1} < M_{\tilde{Q}} - m_t$). An estimate for the rate of such processes is provided by:

$$\Gamma_{\tilde{Q} \to Q\,\bar{t}\,\tilde{t}_1} \sim \frac{\alpha_s^2}{16\pi} \frac{M_{\tilde{Q}}^3}{M_{\tilde{g}}^2} \, \mathrm{ps}_3,$$

where the three-body phase-space volume with two massive decay products gives rise to a ps_3 factor. For $M_{\tilde{Q}} = 1$ TeV, $M_{\tilde{t}_1} = 800$ GeV (and $M_{\tilde{g}} = 1$ TeV), it amounts[21] to about 2.7×10^{-4} so that the partial width is of the order of 10^{-5} GeV. A four-body decay is otherwise required (see below).

Gluinos would dominantly decay to top-stop pairs with a partial width,

$$\Gamma_{\tilde{g} \to t\,\tilde{t}_1^*} = \frac{\alpha_s}{8} \sqrt{\lambda\left(1, \frac{M_{\tilde{t}_1}}{M_{\tilde{g}}}, \frac{m_t}{M_{\tilde{g}}}\right)} \left(1 - \frac{M_{\tilde{t}_1}^2}{M_{\tilde{g}}^2} + 2\frac{m_t}{M_{\tilde{g}}} \sin 2\theta_t + \frac{m_t^2}{M_{\tilde{g}}^2}\right),$$

which amounts to about 10^{-2} GeV for $M_{\tilde{g}} = 1$ TeV, $M_{\tilde{t}_1} = 800$ GeV and $\sin 2\theta_t = -1$.

In addition to many dijets arising from stops, the gluino, up and down squark decay chains would result in some prompt same-sign top pairs provided the production of those states is kinematically accessible.

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[21] The three-body phase-space volume for a centre-of-mass energy M and two particles of non-vanishing masses xM and yM is given by:

$$\frac{M^2}{4(4\pi)^3} \left[\quad \sqrt{\lambda}\,(1 + x^2 + y^2)/2 \quad + 4x^2(1 - y^2) \ln \frac{2x}{1 + x^2 - y^2 + \sqrt{\lambda}} \quad + x \leftrightarrow y \quad \right]$$

where $\lambda \equiv [1 - (x + y)^2]\,[1 - (x - y)^2]$.

Light left squarks If amongst the lightest supersymmetric particles, the left squarks would have somewhat involved decay patterns. First, a four-body decay could proceed through a virtual gluino and another virtual squark, *e.g.*:
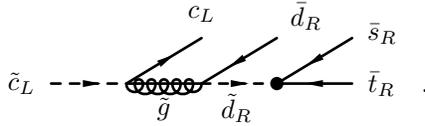
It was stressed in Ref. [5] that a gluino mass insertion, required in the $\tilde{c}_L \to c\, \bar{t}d\bar{s}$ decay, enhances its rate by a factor of approximately $M_{\tilde{g}}^2/M_{\tilde{c}_L}^2$ with respect to the $\tilde{c}_L \to c\, tds$ one ($\tilde{c}_R \to c\, tds$ is subject to a similar enhancement with respect to $\tilde{c}_R \to c\, \bar{t}\bar{d}\bar{s}$). An order of magnitude estimate for the corresponding partial width is given by:

$$\Gamma_{\tilde{c}_L \to c\, \bar{t}d\bar{s}} \sim \frac{\alpha_s^2 |\lambda''_{tds}|^2}{24(4\pi)^3} \frac{M_{\tilde{c}_L}^7}{M_{\tilde{g}}^2 M_{\tilde{d}_R,\tilde{s}_R,\tilde{t}_{1,2}}^4}.$$

It amounts to about 10^{-7} GeV for $\lambda''_{tds} = 0.1$, $M_{\tilde{c}_L} = 800$ GeV and $M_{\tilde{g}} = 1$ TeV$= M_{\tilde{d}_R,\tilde{s}_R,\tilde{t}_{1,2}}$. For these choices of masses, a baryon-number-violating couplings smaller than about 10^{-4} and 10^{-6} would be needed to give displaced vertices ($\Gamma_{\tilde{c}_L} \lesssim 10^{-12}$ GeV) and long-lived squarks ($\Gamma_{\tilde{c}_L} \lesssim 10^{-16}$ GeV). Alternatively, for $\lambda''_{tds} = 0.1$, the gluino, stop and right squark masses would have needed to be raised to about 7 and 32 TeV.

The tiny right component of light left squarks could also allow them to decay through a R-parity violating coupling: $\tilde{u}_L \rightsquigarrow \tilde{u}_R \to \bar{s}\bar{b}$, $\tilde{d}_L \rightsquigarrow \tilde{d}_R \to \bar{t}\bar{s}$. With minimal flavour violation imposed, the rates of these transitions would scale as:

$$\Gamma_{\tilde{u}_L \to \bar{s}\bar{b}} \sim \frac{|y_u \lambda''_{usb}|^2}{8\pi} M_{\tilde{u}_L}, \qquad \Gamma_{\tilde{d}_L \to \bar{t}\bar{s}} \sim \frac{|y_d \tan\beta\, \lambda''_{tds}|^2}{8\pi} M_{\tilde{d}_L},$$
$$\simeq 10^{-20} \text{ GeV}, \qquad\qquad\qquad \simeq 10^{-10} \text{ GeV},$$

for $M_{\tilde{d}_L,\tilde{u}_L} = 800$ GeV, $\lambda''_{usb} = 10^{-6}$ and $\tan\beta$ of order one.

The mixings amongst generations could also help up-type squarks to find their way to decay: $\tilde{u}_L \rightsquigarrow \tilde{t}_R \to \bar{d}\bar{s}$. The associated partial width,

$$\Gamma_{\tilde{u}_L \to \bar{d}\bar{s}} \sim \frac{|y_t\ V_{ub} y_b^2 V_{tb}^* \tan^2\beta\ \lambda''_{tds}|^2}{8\pi} M_{\tilde{u}_L},$$

would scale as 10^{-12} GeV with $M_{\tilde{u}_L} = 800$ GeV and $\tan\beta$ of order one.

The decays of light left up squarks therefore appear fairly suppressed but only the holomorphic restriction of minimal flavour violation with small $\tan\beta$ (and $\lambda''_{tsb} \simeq 10^{-6}$) could actually give rise to some long-lived squarks. For $\tan\beta$ of the order of 50, the dominant $\lambda''_{tsb} \simeq 10^{-4}$ coupling would only produce displaced vertices. On the contrary, with the full minimal-flavour-violation hierarchies of Table 13 the four-body decays of left squarks would all be prompt, provided the gluino, stop or right down squarks are lighter than about 7 TeV $(M_{\tilde{c}_L,\tilde{u}_L}/800 \text{ GeV})^{7/6}$.

The four-body decay of light left down squarks $\tilde{d}_L \to d\,\bar{t}\bar{d}\bar{s}$ would dominate the two-body one $\tilde{d}_L \to \bar{t}\bar{s}$ as long as the gluino, right down squarks and stops are lighter than about $3.8 M_{\tilde{d}_L}$ ($1.0 M_{\tilde{d}_L}$ for $\tan\beta = 50$) which amounts to about 3 TeV for $M_{\tilde{d}_L} = 800$ GeV.

Produced gluinos would decay to quark-squark pairs like $\tilde{g} \to \bar{u}\,\tilde{u}_L$ with a rate,

$$\Gamma_{\tilde{g} \to \bar{u}\,\tilde{u}_L} = \frac{\alpha_s}{8} M_{\tilde{g}} \left(1 - \frac{M_{\tilde{u}_L}^2}{M_{\tilde{g}}^2} \right)^2,$$

of the order of 10^0 GeV for $M_{\tilde{u}_L} = 800$ GeV and $M_{\tilde{g}} = 1$ TeV. As for right squarks, they could undergo three- or two-body decays like $\tilde{u}_R \to u\,\bar{d}\bar{d}_L$ and $\tilde{d}_R \to \bar{s}\,\bar{t}$ of rates,

$$\Gamma_{\tilde{u}_R \to u\,\bar{d}\bar{d}_L} \sim \frac{\alpha_s^2}{16\pi} \frac{M_{\tilde{u}_R}^3}{M_{\tilde{g}}^2} \, \mathrm{ps}_3, \qquad \Gamma_{\tilde{d}_R \to \bar{s}\,\bar{t}} = \frac{|\lambda''_{tds}|^2}{8\pi} M_{\tilde{d}_R} \left(1 - \frac{m_t^2}{M_{\tilde{d}_R}^2} \right)^2,$$

scaling as 10^{-3} and 10^0 GeV for $M_{\tilde{d}_L} = 800$ GeV, $M_{\tilde{u}_R, \tilde{d}_R} = M_{\tilde{g}} = 1$ TeV and a phase space volume factor ps_3 of about 0.019.[20] The $\tilde{u}_R \to u\,d\bar{d}_L^*$ decay rate would rather scale with $M_{\tilde{u}_R}^5/M_{\tilde{g}}^4$. A weak decay of the stop, if sufficiently split from the left bottom squark, with a rate,

$$\Gamma_{\tilde{t}_1 \to \tilde{b}_L W^+} = \frac{\alpha_W \cos^2\theta_t}{8} \frac{M_{\tilde{t}_1}^3}{m_W^2} \lambda\left(1, \frac{M_{\tilde{b}_L}}{M_{\tilde{u}_L}}, \frac{m_W}{M_{\tilde{u}_L}} \right)^{3/2},$$

of the order of 10^1 GeV would compete with the R-parity violating one $\Gamma_{\tilde{t}_1 \to \bar{s}\bar{d}} \sim 10^0$ GeV, for $M_{\tilde{t}_1} = 1$ TeV, $M_{\tilde{b}_L} = 800$ GeV and $\cos\theta_t = 1$.

A significant production of same-sign tops then proceeds through $pp \to \tilde{u}_L\tilde{u}_L,\ \tilde{u}_L\tilde{d}_L,\ \tilde{d}_L\tilde{d}_L$ after dominant four-body decays preferably involving anti-tops (the two-body $\tilde{d}_L \to \bar{t}\bar{s}$ decay appears subleading). Gluino and right first-generation squark production ($\tilde{g}\tilde{u}_L,\ \tilde{d}_R\tilde{u}_L,\ \tilde{d}_R\tilde{d}_L,\ \tilde{g}\tilde{g},\ \tilde{d}_R\tilde{d}_R,\ \tilde{u}_R\tilde{u}_L$, etc.) would also provide sizeable contributions provided they are not much heavier than left squarks.

Light right up squarks Light $\tilde{u}_R$ or $\tilde{c}_R$ would dominantly undergo a four-body decay like $\tilde{u}_R \to u\,tds$, provided the $\tilde{d}_R$, $\tilde{s}_R$, $\tilde{t}_{1,2}$ squarks and gluino are not much heavier. The decays to a pair of light jets $\tilde{u}_R \to \bar{s}\,\bar{b}$, $\tilde{c}_R \to \bar{s}\,\bar{d}, \bar{b}\,\bar{d}$ would otherwise become competitive. With the following expressions for partial widths:

$$\Gamma_{\tilde{u}_R \to u\,tds} \sim \frac{\alpha_s^2 |\lambda''_{tds}|^2}{24(4\pi)^3} \frac{M_{\tilde{u}_R}^7}{M_{\tilde{g}}^2 M_{\tilde{d}_R, \tilde{s}_R, \tilde{t}_{1,2}}^4}, \qquad \Gamma_{\tilde{u}_R \to \bar{s}\bar{d}} = \frac{|\lambda''_{usb}|^2}{8\pi} M_{\tilde{u}_R},$$

as well as $M_{\tilde{u}_R} = 800$ GeV and $\lambda''_{usb} = 10^{-6}$ (10^{-4} for $\tan\beta = 50$), the four-body decay is seen to dominate for $M_{\tilde{d}_R, \tilde{s}_R, \tilde{t}_{1,2}} = M_{\tilde{g}} \lesssim 5.4 M_{\tilde{u}_R} \simeq 4$ TeV ($2.5 M_{\tilde{u}_R} \simeq 2$ TeV). At that point, both partial widths are of the order of 10^{-10} GeV (10^{-6} GeV). For $M_{\tilde{g}} = M_{\tilde{d}_R, \tilde{s}_R, \tilde{t}_{1,2}} = 1$ TeV, the four-body decay rate rises to about 10^{-7} GeV.

Like for light left squarks, the extra suppression of their four-body decays arising in the holomorphic restriction would cause the $\tilde{u}_R$ to have detector-scale lifetimes at small $\tan\beta$. Displaced vertices would also occur at $\tan\beta$ of the order of 50.

Produced gluinos would undergo prompt two-body decays like $\tilde{g} \to \bar{u}\,\tilde{u}_R$. A three-body decay of the left squarks (*e.g.*, $\tilde{d}_L \to d\,\bar{u}\tilde{u}_R$) would be required if the two-body one (*e.g.*, $\tilde{d}_L \to d\tilde{g}$) is not not kinematically allowed.

Provided $\tilde{d}_R$, $\tilde{s}_R$, $\tilde{t}_1$ and $\tilde{g}$ are not much heavier than $\tilde{u}_R$, a significant amount of same-sign tops would therefore also be produced in this situation.

Light down right squarks Light $\tilde{d}_R$ would almost exclusively decay as $\tilde{d}_R \to \bar{s}\,\bar{t}$, with rate of the order of 10^{-1} GeV, for $M_{\tilde{d}_R} = 800$ GeV. Their same-sign production $dd \to \tilde{d}_R\tilde{d}_R$ through the t-channel exchange of a gluino would then lead to many more anti-top pairs than tops.

The three-body decays of the other squarks (like $\tilde{Q} \to Q\,\bar{d}\tilde{d}_R$) would have sizeable branching fractions provided the two-body ones ($\tilde{Q} \to Q\tilde{g}$) are not open. On the other hand, the produced gluinos would tend to decay to $\bar{d}\tilde{d}_R$, $\bar{s}\tilde{s}_R$ and $\bar{b}\tilde{b}_R$ pairs (or charge conjugates), with subsequent R-parity-violating decays of the squarks.

Almost all superparticle decays would therefore lead to tops in this situation and many same-sign pairs would be produced.

Summary For each mass hierarchy, the dominant decay channels of the gluino, stop and first-generation squarks are summarized in Table 14. For definiteness, the largest R-parity violating coupling is notably assumed to be $\lambda''_{tds} = 0.1$. *Light* superparticles are taken to have masses of 800 GeV and the masses of other coloured states are set to 1 TeV (colourless superparticle interactions are neglected).

	Light gluino $\tilde{g}$	Light stop $\tilde{t}_1$	Light left squarks $\tilde{Q}_L$	Light right up squarks $\tilde{U}_R$	Light right down squarks $\tilde{D}_R$
$\tilde{g} \to$	$tds,\ \bar{t}\bar{d}\bar{s}$	$t\tilde{t}_1^*,\ \bar{t}\tilde{t}_1$	$Q\tilde{Q}_L^*,\ \bar{Q}\tilde{Q}_L$	$U\tilde{U}_R^*,\ \bar{U}\tilde{U}_R$	$D\tilde{D}_R^*,\ \bar{D}\tilde{D}_R$
$\tilde{t}_1 \to$	$t\tilde{g},\ \bar{s}\bar{d}$	$\bar{s}\bar{d}$	$\tilde{b}_L W^+,\ \bar{s}\bar{d}$	$\bar{s}\bar{d}$	$\bar{s}\bar{d}$
$\tilde{u}_L/\tilde{d}_L \to$		$Q\,t\tilde{t}_1^*$	$u/d\ \bar{t}\bar{d}\bar{s}\ (-/\bar{t}\bar{s})$	$u/d\ \bar{U}\tilde{U}_R$	$u/d\ \bar{D}\tilde{D}_R$
$\tilde{u}_R \to$	$Q\tilde{g}$	$Q\,\bar{t}\tilde{t}_1$	$u\,\bar{Q}\tilde{Q}_L$	$u\,tds\ (\bar{s}\bar{b})$	$u\,D\tilde{D}_R^*$
$\tilde{d}_R \to$		$\bar{t}\bar{s}$	$\bar{t}\bar{s}$	$\bar{t}\bar{s}$	$\bar{t}\bar{s}$

Table 14: Dominant decay modes of the gluinos, stop and first generation squarks for the different mass hierarchies discussed in the text, when minimal flavour violation is imposed. With rates $\Gamma(\tilde{d}_L \to \bar{t}\bar{s}) \sim 10^{-10}$ GeV $\sim \Gamma(\tilde{u}_R \to \bar{s}\bar{b})$, the two-body decays of light $\tilde{d}_L$ and $\tilde{u}_R$ could compete with the four-body ones, if $\tilde{d}_R$, $\tilde{s}_R$, $\tilde{t}_1$ and $\tilde{g}$ are about four and five times heavier (one and three times, for $\tan\beta = 50$).

The productions of gluinos and first-generation squarks are seen to initiate decay chains ending with top quarks. The $pp \to \tilde{g}\tilde{g}$, $\tilde{g}\tilde{u}_{R,L}$, $\tilde{u}_{R,L}\tilde{d}_{R,L}$, etc. processes would therefore produce same-sign top pairs. The same-sign squark pair production processes proceed through the t-channel exchange of a Majorana gluino.

As the dominant two-, three- and four-body decays have rates of the order of $10^{[+1,-2]}$, $10^{[-3,-5]}$ and 10^{-7} GeV for our choices of parameters, very few displaced vertices (arising for $\Gamma \lesssim 10^{-12}$ GeV) are produced. With the holomorphic restriction of minimal flavour violation imposed, R-parity-violating two- and four-body decays get suppressed by an extra factor of 10^{-6} for $\tan\beta = 50$ (10^{-10} with $\tan\beta = 5$). Four-body decays can then lead to displaced vertices (and, for $\tan\beta = 5$, superparticles of width $\Gamma \lesssim 10^{-16}$ GeV that are long-lived at detector scales).

Same-sign leptons

Same-sign tops production was just seen to be a fairly generic feature of R-parity-violating supersymmetric scenarios satisfying the minimal-flavour-violation hypothesis. As stressed before, this is a clear signature especially in its same-sign lepton channel. Let us in this section provide quantitative estimates and relate model parameter with LHC data.

Simplified hierarchy   The gluino and right down squark productions (see Fig. 15) are the simplest processes we will focus on:

$$dd \to \tilde{d}_R \, \tilde{d}_R \to \bar{s}\,\bar{t} \ \ \bar{s}\,\bar{t},$$
$$dg \to \tilde{d}_R \, \tilde{g} \to \bar{s}\,\bar{t} \ \ \bar{d}\,\bar{s}\,\bar{t},$$
$$gg \to \tilde{g} \, \tilde{g} \to \bar{d}\,\bar{s}\,\bar{t} \ \ \bar{d}\,\bar{s}\,\bar{t}.$$

Conjugated final states are understood. In the holomorphic restriction, the gluino decays involve bottoms instead of down quarks and d_R's also decay to $\bar{b}\bar{t}$ pairs on top of $\bar{s}\bar{t}$ ones.
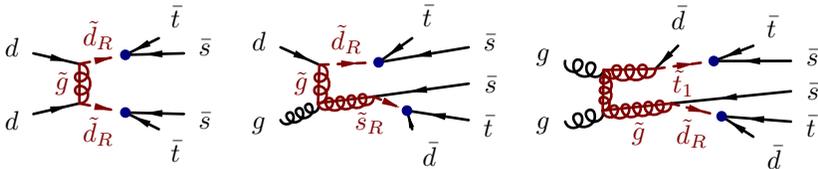
Fig. 15: Especially relevant contributions to same-sign tops plus jets production in a R-parity violation scenario satisfying the minimal-flavour-violation prescription.

Other processes may obviously also contribute to same-sign top pair production. The signal estimates obtained through right down squark and gluino resonant intermediate states only will therefore be rather conservative. To get an idea of the potential increase of rates caused by other processes, a scenario where the contribution of $\tilde{d}_L \to \bar{s}\bar{t}$ exactly equals to the one of $\tilde{d}_R$ will also be considered. A realistic situation is expected to yield results somewhere between the ones obtained in these conservative and (probably) optimistic settings.

The rates of the simplest processes listed above are rather independent of the details of the mass spectrum and mixings. Only the right down squark and gluinos masses are of primary relevance. A very simplified setup characterised by those two parameters, $M_{\tilde{Q}}$ and $M_{\tilde{g}}$, therefore captures the most relevant physics. For simplicity, the other squark masses are also fixed to $M_{\tilde{Q}}$. Although the stop contributes to the gluino $\tilde{g} \to tds, \bar{t}\bar{d}\bar{s}$ decays, its mass would only affect the associated partial width, not much the corresponding branching fraction. It could therefore have been chosen at will, instead of being fixed to $M_{\tilde{Q}}$, without altering much the following conclusions.

For definiteness, the λ''_{tds} coupling is fixed to 0.1 in the general case while, for the holomorphic restriction, we set $\lambda''_{tbs} = 10^{-3}$ and $\lambda''_{tds,tdb} = 10^{-4}$. As will be clear from the results obtained in those two cases, the exact magnitude of the R-parity-violating couplings is actually very mildly relevant. Like the stop mass, they mainly impact particles widths and lifetimes but barely modify branching fractions.

Monte Carlo simulation The superparticle productions and decays are simulated at leading order and parton level using the FEYNRULES-MADGRAPH5 software chain [67, 68]. All superparticle widths are also computed iteratively with MADGRAPH5. Event samples are produced for the 8 and 14 TeV LHC, with $M_{\tilde{Q}}$ and $M_{\tilde{g}}$ varying between 200 and 1100 GeV. In this preliminary study, radiative QCD corrections have not been taken into account. Signal rates are therefore somewhat underestimated. Gluons-initiated only, our gluinos pair production cross sections at 8 TeV are for instance about 3.5, 4.1, 4.9 and 5.8 times smaller than the ones provided by the LHC SUSY Cross Section Working Group [106] for $M_{\tilde{g}} = 550$, 650, 750 and 850 GeV.

Superparticles and same-sign top production rates are displayed in Fig. 16. In the upper-left half of the $M_{\tilde{Q}} - M_{\tilde{g}}$ plane, the production of same-sign tops is dominated by processes involving resonant down squarks. The $\tilde{d}_R \to \bar{t}\bar{s}$ branching fraction is close to 100% in that region. In the bottom-right half, resonant gluino production provides the leading contribution and the branching fractions for $\tilde{g} \to tds, \bar{t}\bar{d}\bar{s}$ both approach 50%. The intermediate region receives significant contributions from $\tilde{d}_R\tilde{g}$ production.

Experimental data Both CMS [74–77] and ATLAS [107–109] studied the same-sign dilepton signature and set generic constraints on new physics contributions with 7 and 8 TeV LHC data. Signal regions characterized by moderate missing energy, relatively high hadronic activity or jet multiplicity and one or two b tags are expected to be the most sensitive to the same-sign tops plus jets signature.

Backgrounds In these searches, irreducible and instrumental backgrounds have comparable magnitudes. Irreducible backgrounds with isolated same-sign leptons and b jets arise from $t\bar{t}Z$ and $t\bar{t}W$ production processes. Their 8 TeV LHC cross sections at next-to-leading-order in QCD amount respectively to 208 and 232 fb [78, 110]. The di- and tri-boson plus jets productions also

$\sigma(\, p\, p \to \tilde{d}\, \tilde{d},\ \tilde{d}\, \tilde{g},\ \tilde{g}\, \tilde{g}\,)$ [fb]

$\sigma(\, p\, p \to \tilde{d}\, \tilde{d},\ \tilde{d}\, \tilde{g},\ \tilde{g}\, \tilde{g} \to t\, t,\ \bar{t}\, \bar{t} + X\,)$ [fb]
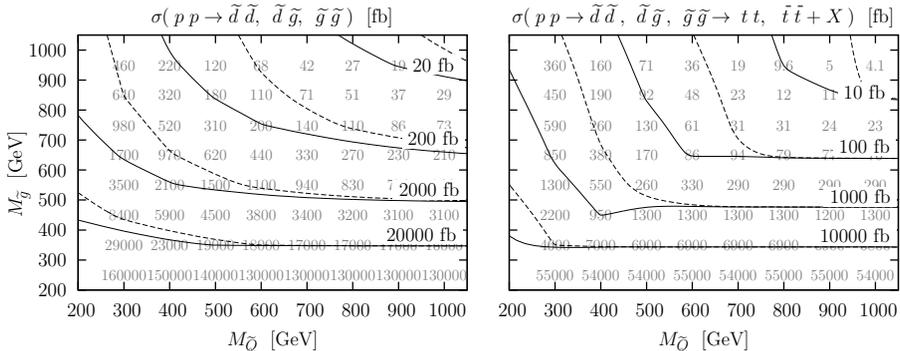
Fig. 16: Leading-order and parton-level superparticles and same-sign tops production rates at the 8 TeV LHC. The solid curve (and grid numbers) includes the contributions from the gluino and right down squark production only while the dashed one assumes $\tilde{d}_L$ contributes as much as $\tilde{d}_R$.

contribute, generally without hard b and sometimes with a third opposite-sign lepton arising from Z boson leptonic decays. Positively charged dileptons dominate over negatively charged ones at the LHC. As mentioned already in Section III.3 (on p.52) and to be further examined in Section IV.5, this feature is generic in the standard model which communicates the proton-proton initial-state charge asymmetry to the final state.

Instrumental backgrounds arise from the mis-reconstruction, mainly in $t\bar{t}$ events, of:

- (heavy) mesons decaying leptonically within jets,
- hadrons as leptons,
- asymmetric conversions of photons,
- electron charges (if a hard bremsstrahlung radiation converts to a $e^+ e^-$ pair in which the electron with a charge opposite to the initial one dominates).

The first three sources are often collectively referred to as *fake leptons* and the last one as *charge flips*. The important contribution of b quark semi-leptonic decays is significantly reduced when b tags are demanded [76].

Selection criteria The selection of events follows CMS one whose collaboration provides information (including efficiencies) and guidelines for constraining any model in an approximate way [75–77]. Tops are decayed semi-leptonically to electrons or muons, and, at the partonic level, we require:

- two same-sign leptons with $p_T > 20$ GeV and $|\eta| < 2.4$,
- jets to have $p_T > 40$ GeV and $|\eta| < 2.4$,
- at least two of them to be b-tagged.

As in CMS searches, eight signal region (SR) are defined (see Table 15). They are characterized by different requirements on the number of jets, b-tags, amount of missing transverse energy $\not{E}_T$ and transverse hadronic activity H_T. The selection of an isolated lepton is taken to have an efficiency of 60%

Lepton efficiency in SR0

Lepton efficiency in SR8

b-tagging efficiency in SR0

b-tagging efficiency in SR8
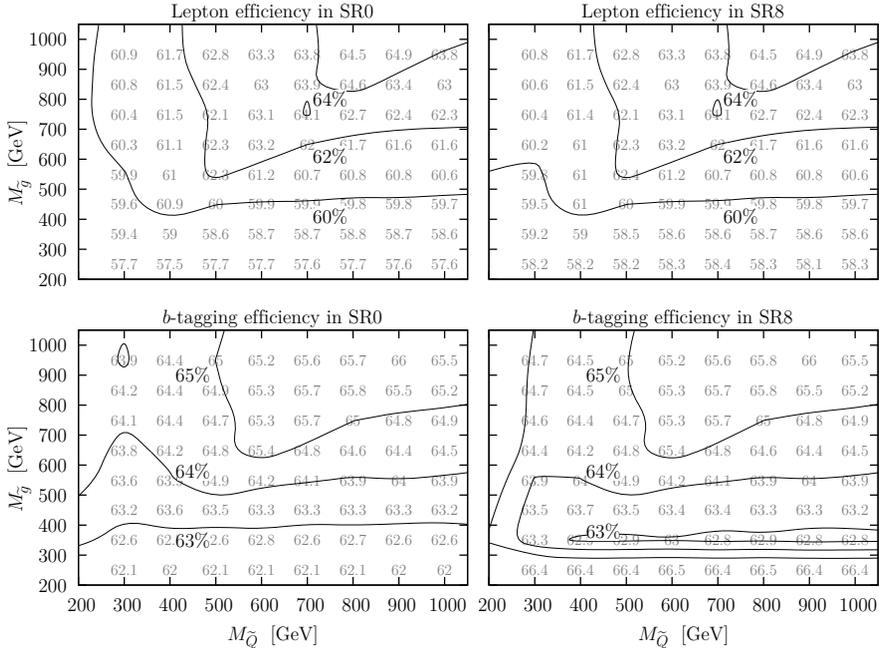
Fig. 17: Isolated lepton and b-jet identification efficiencies (in percent) obtained using the p_T-dependent parametrization provided by the CMS experiment [75–77], at 8 TeV, in signal regions SR0 and SR8 (defined in Table 15). As they are very much constant over the considered range of masses, we approximate them by a flat 60% value.

and the tagging of a parton-level b quark as a b jet is fixed to be 60% efficient too. These values have been chosen in view of the almost constant efficiencies obtained (see Fig. 17) using the p_T-dependent parametrizations provided by CMS. Note that, for b tagging, the value chosen is a few percent lower than those estimated in this way. With backgrounds under control, a higher number of isolated leptons from signal events could be selected by lowering the cut on their p_T or by modifying the isolation requirement [111]. On the other hand, the p_T-dependent parametrization of the isolated lepton selection efficiency might not be reliable in regions where the tops can be boosted or when the hadronic activity of a typical event is important [112]. For our preliminary limit setting we will however keep efficiencies constant.

The goodness of our parton-level approximate selection is assessed by comparison (relaxing the same-sign condition for leptons) to the total acceptance in SR1 quoted by CMS for standard-model $t\bar{t}$ events with semi-leptonic top decays. Our total acceptance of 0.20% (including top branching fractions) is compatible but lower than the $(0.29 \pm 0.04)\%$ quoted by CMS [75, 77]. So, again, the strength of our signal is probably conservatively estimated.

	SR0	SR1	SR4	SR3	SR8	SR5	SR6	SR7
Min. num. of b tags	2	2	2	2	2	2	2	3
Min. num. of extra jets	0	0	2	2	2	2	2	0
Cut on H_T [GeV]	80	80	200	200	320	320	320	200
Cut on $\not{E}_T$ [GeV]	0	30	50	120	0	50	120	50
Limit on BSM events	30.4	29.6	12.0	3.8	10.5	9.6	3.9	4.0

Table 15: Definitions of the signal regions used by CMS [75, 77] for same-sign dilepton searches. For each of them, a 95% CL upper limit on beyond-the-standard-model (BSM) events is derived from 10.5 fb^{-1} of 8 TeV data, assuming a 30% uncertainty on signal efficiency and using the CLs method.

IV.4 Current constraints and prospects

For several choices of squark and gluino masses, the number of events in each signal region is compared with the 95% CL limit set by CMS assuming a conservative 30% uncertainty on the signal selection efficiency and using 10.5 fb^{-1} of 8 TeV data [75, 77]. The corresponding exclusion contours in the $M_{\tilde{Q}} - M_{\tilde{g}}$ plane are displayed in Fig. 18.
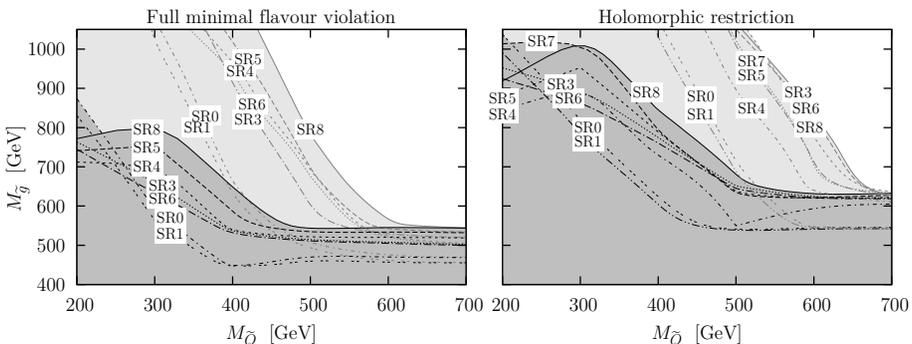
Fig. 18: Exclusion regions at the 95% CL in the $M_{\tilde{Q}} - M_{\tilde{g}}$ plane derived from the CMS same-sign dilepton search [75, 77]. The light-grey region is excluded in scenarios where resonant $\tilde{d}_R$ and $\tilde{d}_L$ contributions are identical while the dark-grey region is excluded when $\tilde{d}_R$ only is contributing. Although $M_{\tilde{Q}}$ is common to all squarks, only the down squarks masses are actually relevant. In particular, these limits are basically independent of the top squark (or neutralino) one.

In the full minimal-flavour-violation case, signal regions with low H_T cuts perform well in the lowest mass ranges where jets are softer. Everywhere else, SR8 characterized by no $\not{E}_T$ cut and a relatively high $H_T > 320$ GeV requirement provides the best sensitivity. As expected, in the presence of R-parity violation, the supersymmetry searches requiring a large amount of missing energy are not the best suited. This can be understood from the shapes of the R-parity-violation signal and $t\bar{t}W + t\bar{t}Z$ background in the $H_T - \not{E}_T$ plane (see Fig. 19). For squark and gluino masses close to the exclusion contour of SR8 (Fig. 19, left), the two missing energy distributions are very similar while, for higher superparticle masses (Fig. 19, right), the average $\not{E}_T$ is only slightly higher in signal events. On the other hand, a relatively good discrimination

between signal and background is provided by the transverse hadronic activity. This is especially true in the $M_{\tilde{g}} < M_{\tilde{Q}}$ region (Fig. 19, bottom) where the H_T distribution starts to rise at values about 200 GeV higher than in the $M_{\tilde{Q}} < M_{\tilde{g}}$ case (Fig. 19, top). A H_T cut higher than 320 GeV could probably increase the signal over background ratio. The jet multiplicity or highest jet p_T may also provide powerful handles [111].
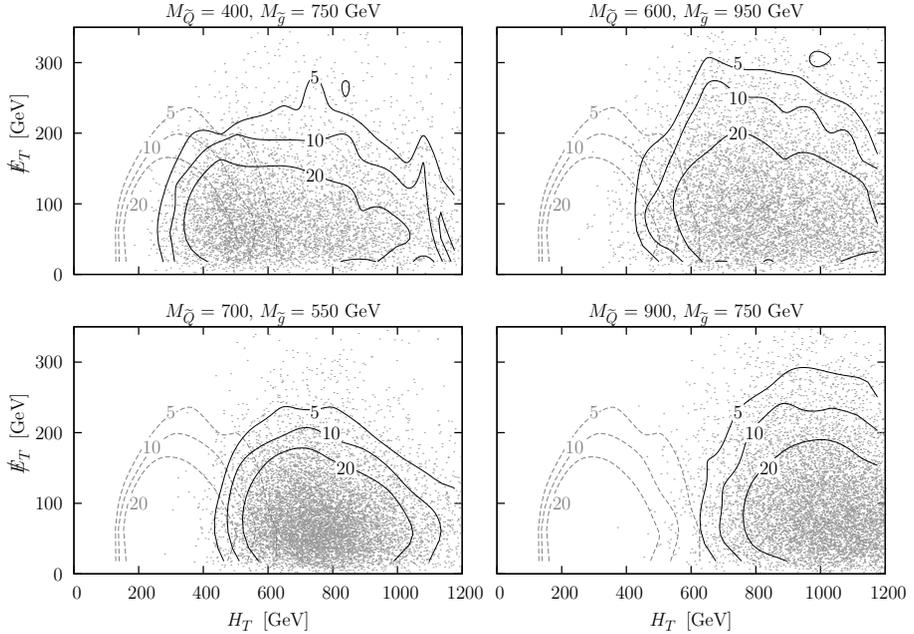
Fig. 19: Shape $1/\sigma \times d^2\sigma/dH_T d\!\!\!/E_T$ $[100 \text{ GeV}]^{-2}$ of the R-parity-violating signal (solid contours, scattered points) and standard-model $t\bar{t}W$ and $t\bar{t}Z$ backgrounds (dashed contours). Full minimal flavour violation and SR0 selection criteria are imposed; only the $\tilde{d}_R$ and $\tilde{g}$ contributions to top pair production are included; leading-order and parton-level results obtained using MadGraph5 [67].

In the whole squark mass range, the SR8 limit excludes gluino masses below roughly 550 GeV. In the low- and mid-range squark mass region however, the bound varies significantly depending on the contributions of first-generations squarks to the same-sign tops signal. In the most unfavourable situation where only $\tilde{d}_R$ contributes, the gluino mass limit saturates around 800 GeV while it rises well above the TeV when the $\tilde{d}_L$ contribution equals the $\tilde{d}_R$ one. This saturation arises from the gluino mass dependence of the $\tilde{d}_R\tilde{d}_R$ production rate.

With the holomorphic restriction imposed, the final-state b multiplicity increases on average. Obtaining at least two b-tagged jets is therefore more likely and the limits improve slightly. SR7 where three b tags are required then also provides competitive bounds. Overall, this pushes the limit on superparticle masses higher, towards regions where the average $\!\!\!/E_T$ of signal events increases. There, SR3 and SR6 characterized by a higher $\!\!\!/E_T > 120$ GeV cut and very small backgrounds perform more and more efficiently. This is especially visible when the contributions of first-generations squarks like $\tilde{d}_L$ are significant and

further enhance the signal rate. For moderate superparticle masses though, SR8 still leads to the best limit.

Our exclusion regions are somewhat more conservative than the $M_{\tilde{g}} \gtrsim 800$ GeV limit obtained in Ref. [112], with the holomorphic restriction of minimal flavour violation. The scenario analysed there decouples all sparticles except the gluino and a lighter top squark. Same-sign top pairs are produced through $p\,p \to \tilde{g}\,\tilde{g}$ with the gluino decaying as $\tilde{g} \to t\,b\,s, \bar{t}\,\bar{b}\,\bar{s}$ via on-shell top squarks. Such a scenario corresponds to our $M_{\tilde{Q}} \to \infty$ region where gluino pair production only contributes to the same-sign top final state. There, a $M_{\tilde{g}} \gtrsim 630$ GeV limit is obtained. It was checked not to change significantly when the stop is light enough to be on-shell in the gluino decay. Even though the kinematics is different, the selection criteria are broad enough to prevent a significant change of sensitivity. The differences in results could arise from the QCD corrections and quarks-initiated gluinos production modes not included here, the finite-width effects not included in Ref. [112] or our simpler simulation procedure with, for instance, constant efficiencies.

Still for a holomorphic scenario in this region of the plane, the last same-sign dilepton searches by CMS [74] and ATLAS [107] set $M_{\tilde{g}} \gtrsim 900$ GeV limits (that are actually independent of the stop mass) with the full 8 TeV data set which contains almost twice as much collisions as the one considered here. The multijet search of Ref. [113] appear to have a comparable sensitivity (see the comparison in Fig. 5d of Ref. [107]).

The perspectives of improvement on the mass bounds can be assessed from the fiducial 8 TeV cross sections in SR8 (currently providing the best sensitivity in most cases) and SR0 (the baseline selection) displayed in Fig. 20. A comparison of the top row of Fig. 20 with the right panel of Fig. 16 shows the overall acceptance, including top branching fractions, in both SR0 and SR8 is of about $(0.3 \pm 0.04)\%$, once the lowest masses $M_{\tilde{Q},\tilde{g}} \leq 300$ GeV are excluded. The signal rate is actually not much affected by the additional hadronic requirements in SR8 with respect to SR0. Improving the limits by a factor of ten could lead to an increase of the absolute bound on the gluino mass of the order of a couple of hundred GeV. The improvement would be the much more significant in the low squark mass region. Similar gains would be obtained at the 14 TeV LHC if a bound comparable to the one obtained so far at 8 TeV is achieved. It is however important to stress that the characteristics of the signal change as sparticles get heavier. With increasing mass bounds, signal regions with higher H_T and $\slashed{E}_T$ requirements would have better sensitivities. Adequate techniques should then also be used to identify the higher proportion of boosted top quarks (see for instance Ref. [114]).

IV.5 Charge asymmetries

As already mentioned, the standard-model $t\bar{t}W$ background features a predominance of positively charged dileptons over negative ones. More quantitatively, MADGRAPH5 [67] leading-order estimates for the lepton charge asymmetry of
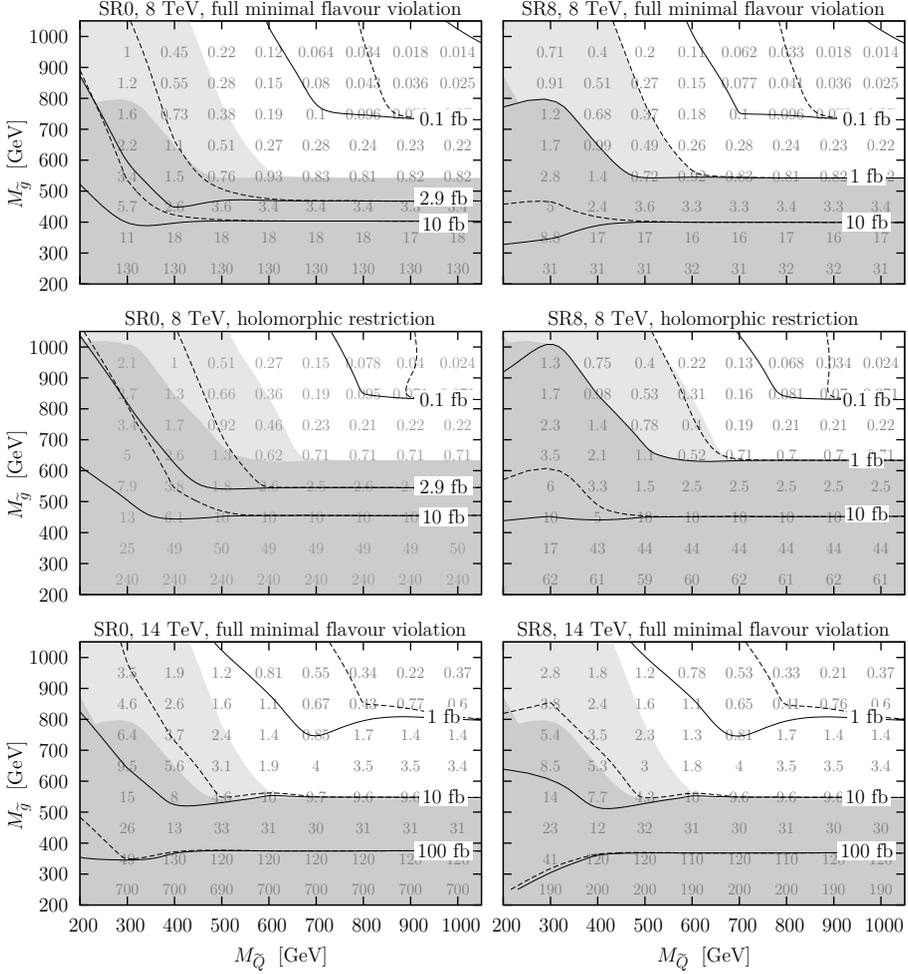
Fig. 20: Fiducial cross sections [fb] in SR0 and SR8 signal regions for the same-sign dilepton signature of R-parity violation. Currently excluded regions are shaded (see Fig. 18). The contributions of $\tilde{d}_R$ and $\tilde{g}$ production only are used to obtain the solid contours (and grid numbers) while the dashed ones moreover assume a $\tilde{d}_L$ contribution of the same magnitude as the $\tilde{d}_R$ one.

are:

$$A_{\ell\ell'}^{t\bar{t}W+t\bar{t}Z} = \begin{array}{ccccccc} \text{SR0} & \text{SR1} & & \text{SR4} & \text{SR3} & & \text{SR8} & \text{SR5} & \text{SR6} \\ +26\% & +28\% & & +29\% & +36\% & & +26\% & +28\% & +35\% \end{array}.$$

A good agreement is obtained in SR1 with the value of $+29\%$ deriving from the central value of CMS background estimates [75]. This is the only signal region in which CMS distinguished positively and negatively charged pairs. Instrumental backgrounds are symmetric and therefore dilute the global asymmetry expected in the absence of signal. A central value of $+9.3\%$ is obtained including all backgrounds while -26% is observed in data. An unreasonably large $\pm53\%$ uncertainty on the background expectation is obtained when neglecting presumably important correlations (not publicly available) and performing a quadratic combination of the uncertainties on positive and negative lepton pair yields. A proper measurement of the asymmetry by the experimental collaboration therefore appears needed.
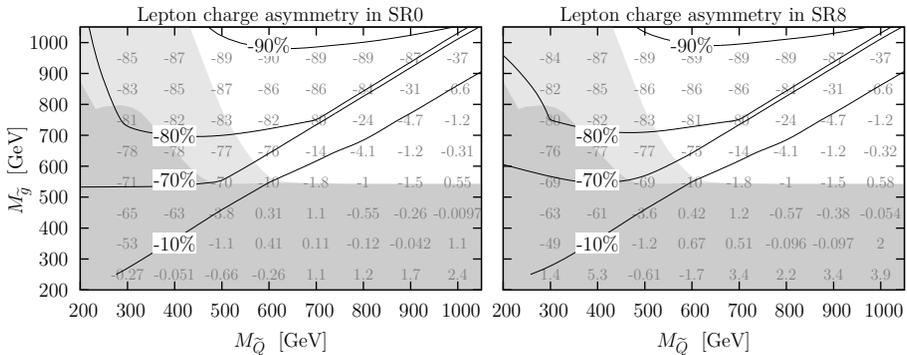
Fig. 21: Lepton charge asymmetry of the same-sign dilepton signal of R-parity violation (electrons or muons only). Full minimal flavour violation is imposed and only $\tilde{d}_R$ and $\tilde{g}$ contributions to the top pair production are included. Excluded regions are shaded (see Fig. 18).

On the contrary, the R-parity-violating processes contribute negatively to the asymmetry. The ones initiated by down valence quarks are indeed significantly more probable than their conjugates, initiated by anti-down quarks. As, they dominate when squarks are lighter than gluinos, much more anti-top than top-quark pairs are expected in the upper-left part of the $M_{\tilde{Q}} - M_{\tilde{g}}$ plane. This leads to a predominance of negatively charged dileptons and $A_{\ell\ell'}^{\mathrm{RPV}}$ approaches -1 for all $\ell, \ell' = e, \mu, \tau$ (see Fig. 21, where only electrons and muons are considered).

As already emphasized in Section III.3, such a negative asymmetry is a smoking gun for new physics and an important evidence for baryon number violation. It is indeed almost impossible to obtain in other realistic new-physics scenarios. On the experimental side, a precise measurement of this asymmetry, in which systematic uncertainties cancel, could provide important constraints on new-physics models, including R-parity-violating ones. In addition, a limit on the production rate of negatively charged lepton pairs only, for which

standard-model irreducible backgrounds are smaller, could in principle be used to improve the current bounds in the upper-left half of the $M_{\tilde{Q}} - M_{\tilde{g}}$ plane.

IV.6 Conclusions

This chapter studied in more details a specific new-physics scenario where new resonant states give rise to observable baryon number violation at the LHC, without conflicting with low-energy bounds. R-parity violation in a minimal-flavour-violation context was shown to give rise to a clear same-sign top signature whose presence and strength depend fairly mildly on many of the model parameters (*e.g.* R-parity-violating coupling strengths, top squark or neutralino masses). A quantitative description of this signal in terms of the few relevant model parameters (*i.e.* the down squark and gluino masses) was provided and the constraints arising from a same-sign dilepton search were derived. Distinctive features of the signal that could be exploited in future searches were identified. Beside a hadronic activity significantly higher than the background (but not a much more important missing transverse energy), the predominance of negatively charged lepton pairs over positively charged ones in half of the parameter space is a striking characteristic of the signal. A measurement of the associated asymmetry in which uncertainties cancel could almost unambiguously point at baryon number violation.

Conclusions

In this dissertation, some means of describing physics beyond the standard model have been explored with the aim of contributing theoretically to the present quest for new physics that is actively pursued experimentally, notably at the LHC.

An effective-field-theory view at the standard model naturally leads to the construction of higher-dimensional interactions which are not *sensu stricto* considered as parts of our standard paradigm while systematically deriving from the same symmetry requirements and field content. Lepton number violation is the very first consequence of such an extension (left-handed neutrino masses and mixings derive after the spontaneous breaking of the electroweak symmetry). Next, comes a wide variety of new interactions with significant implications at colliders.

Amongst those, baryon number violation is an especially striking kind of phenomenon that is theoretically expected to occur but remained so far intriguingly unobserved. The LHC offers a new domain in which it could be probed and lifts the restriction to first light generations affecting lower-energy experiments. The copiously producible top quark in particular allows for tests at the quark level. Its comparatively high but actually natural mass imparts it with a special status amongst matter fields.

The effective-field-theory consistency in principle prescribes that all operators of identical dimension should be considered simultaneously. At some definite order in gauge coupling constants, only a limited set of operators are usually relevant for given kinds of processes. Dimension-six baryon-number-violating ones moreover form an isolated subclass which we considered separately. A four-fermion operator description of single top production or decay involving one lepton only was provided. This work served as a basis for a baryon-number-violating top decay search performed by the CMS collaboration.

The puzzling absence of low-energy signals for baryon number violation could have its origin in the standard-model flavour structures from which no departure has been seen experimentally. The minimal-flavour-violation prescription allows for their explicit extrapolation to beyond-the-standard-model

interactions. Baryon-number-violating four-fermion operators are then only allowed if neutrino Majorana masses provide a fundamentally new flavour structure to the standard-model. The smallness of the associated Yukawa couplings would lead to significant suppressions. The intrinsic flavour antisymmetry of quark triads would make the coefficients of operators involving only light first-generation fermions even smaller. The existing bounds on nucleon decays would then only require the characteristic scale of most four-fermion baryon-number-violating operators to be in the TeV region. Somewhat higher values may be required in particular for the *qqql* operator involving weak doublets only. Unobserved neutron–anti-neutron oscillations and (di-)nucleon decays impose similar constraints on six-fermion interactions.

Baryon- and lepton-number-violating transitions could therefore be resonant at the LHC. An effective field theory would then no longer provide a suitable description. The global flavour symmetry requirement would nonetheless still favour interactions involving all three generations. A generic characterisation of such processes was established and the most phenomenologically promising signatures singled out. Tops and charged leptons are ideally suited to identify the special nature of the baryon- and lepton-number-violating signals. The asymmetry in production of same-sign positively and negatively charged leptons was stressed to be a deciding observable.

This generic treatment of the consequences of new resonances was unable to describe quantitatively the signal characteristics. Specific models were required for that purpose. R-parity-violating supersymmetric scenarios actually predict the dominant baryon-number-violating interaction should involve a top quark, or its supersymmetric partner, when minimal flavour violation is imposed. The supersymmetry moreover restricts the diversity of interactions allowed so that the most strongly constrained *qqql* operator is not generated directly at low energies. Superparticle masses in the TeV region are therefore in agreement with matter stability constraints. The resonant production of same-sign tops, identified generically beforehand, was shown to be a widely present signature in this family of scenarios. Its rate and the asymmetry between tt and $\overline{tt}$ plus jets production would primarily depend on a couple of parameters only. A comparison with existing LHC data illustrated the reach of LHC searches and some characteristics of the signal that could be used in the future were highlighted.

At the TeV scale and beyond, pp colliders appeared as particularly well suited laboratories for probing baryon-number-violating production processes (*e.g.* $\Delta B = \pm 2$). Together with (future) ep colliders, they are also able to search for simultaneous baryon and lepton-number-violating transitions (*e.g.* $\Delta B = \Delta L = \pm 1$ or $N_c \Delta B = \pm \Delta L = \pm 3$). Hypothetical machines colliding same-sign lepton—possibly of different flavours—are potentially very special experiments for testing lepton number violation alone (*e.g.* $\Delta L = \pm 2$), through neutrinoless same-sign W production [115].

In conclusion, several descriptions have been combined with the important guidance provided by the constraints on allowed flavour structures to identify how new physics could manifest itself at the LHC. This approach has been applied to a peculiar though expected kind of new phenomena: baryon-number-

violating processes. They could be visible at colliders without conflicting with observed matter stability, once a standard-model pattern of flavour-symmetry breaking is imposed. Flavour and phenomenological considerations indicated the top quark is an outstanding system for probing baryon number violation at the LHC.

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The astonishing stability of ordinary matter is a mysterious fact established experimentally far beyond common wisdom. A conservation of "baryon number" would actually make absolutely stable the protons of atomic nuclei. This is, however, a doubtful theoretical hypothesis. Violations are indeed required to explain the apparent matter-antimatter asymmetry in the universe and occur "non-perturbatively" in our standard description of fundamental interactions.

The present dissertation investigates the forms a violation of baryon number could take at energy scales a thousand times larger than the proton mass. By accelerating particles against each other at such "TeV" energies, the LHC is currently exploring the frontiers of our understanding of elementary particle physics.

An effective description of baryon-number-violating interactions is first adopted. It is built upon the same ingredients (fields, Lorentz and gauge symmetries) as our standard model and is therefore particularly general. In addition, transposing the standard-model pattern of global "flavour" symmetry breaking to baryon-number-violating interactions can simultaneously render them elusive at the proton mass scale and observable at the TeV scale. Their possible manifestations at the LHC are studied, generically and in a specific "supersymmetric" framework.

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