Revisiting monotop production at the LHC

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ABSTRACT: Scenarios of new physics where a single top quark can be produced in association with large missing energy (monotop) have been recently studied both from the theoretical point of view and by experimental collaborations. We revisit the originally proposed monotop setup by embedding the effective couplings of the top quark in an SU(2)_L invariant formalism. We show that minimality selects one model for each of the possible production mechanisms: a scalar field coupling to a right-handed top quark and an invisible fermion when the monotop system is resonantly produced, and a vector field mediating the interactions of a dark sector to right-handed quarks for the non-resonant production mode. We study in detail constraints on the second class of scenarios, originating from contributions to standard single top processes when the mediator is lighter than the top quark and from the dark matter relic abundance when the mediator is heavier than the top quark.

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1 Introduction

The first phase of the LHC experiments has given two important messages: a scalar resonance closely resembling the Standard Model Higgs boson has been discovered, and new physics beyond the Standard Model has not been found. The latter result would imply that new states or effects beyond the Standard Model predictions may be much more difficult to spot at the LHC than we previously thought. In fact, very strong bounds have been posed on easy–catch models, like the constrained version of the Minimal Supersymmetric Standard Model [1, 2].

Many theorists have therefore recently turned their attention on a more signature-based strategy, focusing on unusual final states which are difficult to detect or have not been considered yet by experimental collaborations. One such final state that has been gaining popularity among phenomenologists [3–17] and experimentalists [18, 19] is the monotop signature: a single top quark produced in association with a large amount of missing energy. Although the production of this final state is very suppressed in the Standard Model, it is however not easy to obtain this kind of events in realistic and complete models of new physics. Two main production mechanisms can lead to a monotop state [8, 15], arising either from the resonant production of a coloured bosonic state which further decays into a top quark plus an invisible neutral fermion; or via the production of a single top quark in association with an invisible boson that has flavour-changing couplings to top and light quarks. Examples of the first class of models include $R$-parity violating...
supersymmetry, where the produced resonance is a top squark decaying into a top plus a long-lived neutralino [3, 5, 6, 12]. The second class of models has been described in scenarios of dark matter from a hidden sector that couples to the Standard Model via flavour-violating couplings of a bosonic mediator [7, 9, 14, 20].

All such models can be described in terms of a simple effective Lagrangian [8], which contains all the possible couplings giving rise to a monotop signal. A very general analysis of this framework can be found in Ref. [15], the limiting case of higher-dimensional operators has been discussed in Ref. [10], while monotop production via flavour-changing interactions of quarks with an invisible Z-boson has been detailed in Ref. [4]. Although the effective description has the advantage of being complete, it has the drawback of containing too many free parameters to be efficiently scanned by an experimental search. Furthermore, the included couplings do not respect the symmetries of the Standard Model, as they are intended to describe the model dynamics after the breaking of the electroweak symmetry. In this way, this approach ignores other interactions needed to restore gauge invariance which can give rise to new physics signals in different search channels, the latter possibly implying stronger constraints on the parameters of the model than the monotop search itself. In this work, we revisit the effective parametrisation originally proposed in Ref. [8] by paying particular attention to the embedding of the Lagrangian description into $SU(2)_L \times U(1)_Y$ invariant operators. Our analysis allows us to restrict the number of “interesting” scenarios, i.e., the cases where the monotop signal is genuinely the main signal of new physics to be expected at the LHC. Equivalently, this reduces the number of free parameters to a manageable number. Finally, we discuss in detail how the effective model could be completed in order to guarantee that the missing energy particle produced in association with the top quark is indeed either long-lived or decaying into invisible states.

The rest of this work is organized as follows. In Section 2, we describe how to embed the effective monotop description of Ref. [8] in the Standard Model gauge structure, considering separately the resonant and the flavour-changing monotop production modes. We then focus on non-resonant scenarios which turn out to be less “standard” and investigate, in Section 3, the conditions under which the invisible state is effectively invisible, and other experimental observations, which can further constrain the model. Our conclusions are presented in Section 4.

2 Gauge-invariant descriptions of the monotop dynamics

2.1 Resonant monotop production

In the first class of scenarios yielding the production of a monotop system at colliders, the produced top quark recoils against an invisible fermionic state $\chi$. Note that, being singly produced, $\chi$ cannot be stable, thus it is either long-lived or it decays into a pair of stable particles: in either case, it has to be a neutral and colour-singlet state. Both particles in the final state arise from the decay of a heavy scalar $\varphi$ or vector $X$ field, lying in the fundamental representation of $SU(3)_c$, that is resonantly produced from the fusion of two
down-type (anti-)quarks. The effective Lagrangian describing those scenarios is given by [8]

\[ \mathcal{L}_{\text{res}} = \mathcal{L}_{\text{kin}} + \left( \varphi \bar{d}_i^C \begin{bmatrix} (a_{SR}^q)^{ij} + (b_{SR}^q)^{ij} \gamma^5 \end{bmatrix} d^C_j + \varphi \bar{t} \begin{bmatrix} a_{SR}^{1/2} + b_{SR}^{1/2} \gamma^5 \end{bmatrix} \chi + \text{h.c.} \right) \]

\[ + \left( X_\mu \bar{d}_i^C \begin{bmatrix} (a_{VR}^q)^{ij} \gamma^\mu + (b_{VR}^q)^{ij} \gamma^{\mu \gamma^5} \end{bmatrix} d^C_j + X_\mu \bar{t} \begin{bmatrix} a_{VR}^{1/2} \gamma^\mu + b_{VR}^{1/2} \gamma^{\mu \gamma^5} \end{bmatrix} \chi + \text{h.c.} \right), \]

(2.1)

where \( i, j \) are flavour indices and where we omit all colour indices for clarity. The Lagrangian \( \mathcal{L}_{\text{kin}} \) includes kinetic and mass terms for all new fields, while the other terms focus on their interactions with the Standard Model quarks. As the colour contraction in the interactions with two down-type quarks is antisymmetric, necessarily the scalar, pseudoscalar and axial-vector coupling matrices \( (a_{SR}^q), (b_{SR}^q) \) and \( (a_{VR}^q) \) are antisymmetric under the exchange of the flavour indices, while the vector coupling strength matrix \( (a_{VR}^q) \) is a symmetric matrix in flavour space. Consequently, parton density effects enhance the production modes \( dd \rightarrow X \) and \( ds \rightarrow \varphi \) at hadron colliders (with the relevant coupling strengths being non-vanishing), as already pointed out in previous works [8, 12, 15]. Finally, the parameters \( a_{VR}^{1/2} \) and \( b_{VR}^{1/2} \) appearing in Eq. (2.1) represent the strengths of the interactions of the resonant states with the monotop \( t \chi \) system.

All these interactions are completely generic, and in particular no assumption is made on the chirality of the Standard Model quarks that are involved. However, \( SU(2)_L \) gauge invariance will necessarily constrain such couplings and force the invisible state \( \chi \) and the extra coloured fields \( \varphi \) and \( X \) to lie in possibly non-trivial representations of the group, as already briefly shown in Ref. [11] for the scalar case. This implies the existence of additional component fields whose masses cannot be much larger than those of the \( \varphi, X \) and \( \chi \) fields. In the simplified picture above, any mass splitting can indeed only be generated by the vacuum expectation value of the Higgs field, so that larger mass differences will induce sizable corrections to the electroweak precision observables [21, 22] and are thus strongly disfavoured.

We show in the rest of this section that studying the \( SU(2)_L \) embedding of the effective Lagrangian of Eq. (2.1) can allow us to derive precious constraints on viable and realistic scenarios that deserve to be further studied in high-energy physics experiments. We start our analysis with the scalar case. The \( \varphi \) field, as any scalar field, can only couple to two fermions with opposite chiralities. As a consequence, its coupling to down-type quarks can only have the form

\[ \lambda_1^S \varphi_1 \bar{d}_R^C d_R^C + \lambda_2^S \varphi_2 \bar{d}_L^C d_L^C + \text{h.c.} , \]

(2.2)

recalling that the charge conjugate of the right-handed quark \( d_R^C \) is left-handed while the one of the left-handed quark \( d_L^C \) is right-handed. Moreover, we have introduced generic couplings \( \lambda_S \) and numbered indices to distinguish between both terms in the following discussion. The product \( \bar{d}_R^C d_R^C \) transforms as a singlet of \( SU(2)_L \) with an hypercharge quantum number of \(-2/3\). This implies that the charge of the \( \varphi_1 \) field under \( U(1)_Y \) is \( 2/3 \) and that this field is not charged under the weak gauge group. Analogously, the \( \bar{d}_L^C d_L^C \) product of quark fields belongs to a combination of two left-handed doublets of \( SU(2)_L \) lying in the adjoint representation of the group and whose hypercharge is \( 1/3 \). This enforces the \( \varphi_2 \) field to belong to a weak triplet of fields with an hypercharge of \(-1/3\).
The above discussion demonstrates that the $\varphi_1$ and $\varphi_2$ fields are necessarily two different fields, whose representation under the Standard Model gauge group is given by

$$\bar{d}_R^C \bar{d}_L = (\bar{3}, 1, -2/3) \Rightarrow \varphi_1 = (3, 1, 2/3) \equiv \varphi_2^{2/3},$$

$$\bar{d}_L^C \bar{d}_R = (\bar{3}, 3, 1/3) \Rightarrow \varphi_2 = (3, 3, -1/3) \equiv \left(\varphi_t^{2/3}, \varphi_t^{-1/3}, \varphi_t^{-4/3}\right)^t,$$

(2.3)

where in the right-hand side of both relations, the subscripts $s$ and $t$ stand for the singlet and triplet representation of $SU(2)_L$, respectively, and the superscripts indicate the electric charge of all component fields. Equivalently, the possible couplings of the new scalar field $\varphi$ of Eq. (2.1) to right-handed and left-handed down-type quarks have very different gauge structure, and must arise from two different scalar fields $\varphi_1$ and $\varphi_2$ whose interactions are given instead by Eq. (2.2). In general, both new scalar fields can mix with the Standard Model Higgs doublet. However, the resulting mass splitting is constrained to be small by the perturbativity of the couplings [23] and corrections to the $S$ and $T$ parameters [21, 22]. Any new trilinear and quartic scalar interaction is thus neglected.

The last term of the first line of the Lagrangian of Eq. (2.1) describes the couplings of the $\varphi$ field to the top quark and the invisible fermion $\chi$. It must be modified accordingly when considering that the $\varphi$ field has to lie in either the singlet or the triplet representation of the weak isospin group. In the singlet case, a gauge-invariant interaction term can easily be written down as

$$\lambda_3^S \varphi_1 \bar{t} R \chi_R + \text{h.c.}$$

(2.4)

When we consider instead an $SU(2)_L$ triplet of fields $\varphi_2$, no coupling with an electroweak singlet $\chi$ is allowed so that $\chi$ must be embedded into a larger representation of $SU(2)_L$. We have two such possibilities,

$$\lambda_4^S \varphi_2 \bar{t} R \chi_{R,t}, \quad \text{where} \quad \chi_{R,t} = (1, 3, 1) \equiv \left(\chi_t^1, \chi_t^1, \chi_t^0\right)^t,$$

$$\lambda_5^S \varphi_2 \bar{t} L \chi_{L,d}, \quad \text{where} \quad \chi_{L,d} = (1, 2, 1/2) \equiv \left(\chi_d^1, \chi_d^0\right).$$

(2.5)

As in the above syntax, the subscripts $d$ and $t$ indicate the representation under $SU(2)_L$ and the superscripts refer to the electric charges of the various component fields. Under both setups, additional single top processes where the top quark is resonantly produced in association with one of the charged component fields of the $\chi$ multiplet are expected,

$$\bar{u}_L \bar{d}_L \rightarrow \varphi_t^{-1/3} \rightarrow t_{L/R} \chi_{L,d}^1 \quad \text{and} \quad \bar{u}_L \bar{u}_L \rightarrow \varphi_t^{-4/3} \rightarrow t_R \chi_t^2.$$

(2.6)

The mass splitting between the various components of $\chi$ is expected to be small (see above), so that these extra processes will accompany any hint for new physics in the monotop channel and cannot therefore be ignored. New charged long-lived particles are however heavily constrained by current searches [24, 25], which renders these scenarios unlikely to be realized.

Finally, additional constraints arise when the new scalar field $\varphi_2$ couples to a left-handed top quark and a fermionic field $\chi_d$. Gauge invariance induces a coupling to the
left-handed bottom quark too, which leads to a fast decay of the neutral \( \chi \) field via a virtual \( \varphi_2 \) scalar,

\[
\chi_d^0 \rightarrow b_L (\varphi_1^{1/3})^* \rightarrow b_L u_L d_L .
\] (2.7)

As a consequence, the initial monotop signal has to be traded with new physics contributions to more standard single top production in association with jets.

We now turn to cases where monotop systems are produced from the decay of a spin-1 resonance. Vector fields couple to spinors of the same chiralities, so that the surviving couplings of the \( X \)-field to down-type quarks in Eq. (2.1) are of the form

\[
\lambda_v^1 X^\mu \bar{d}_L \gamma^\mu d_R + \lambda_v^2 X^\mu \bar{d}_R \gamma^\mu d_L + \text{h.c.} ,
\] (2.8)

where we denote by \( \lambda_v \) generic coupling strengths. In order for such couplings to be \( SU(2)_L \)-invariant, the \( X \)-boson must belong to a weak doublet with hypercharge \( 1/6 \),

\[
X^\mu = (3, 2, 1/6) \equiv (X_{2/3}^\mu, X_{-1/3}^\mu)^t ,
\] (2.9)

the superscripts referring again to the electric charge of the component fields. In this case, the decay into a monotop system can only be generated by the coupling of the \( X \)-field to an electroweak fermionic singlet \( \chi \) and a left-handed top quark,

\[
\lambda_v^3 X_{2/3}^\mu t^L \gamma^\mu \chi + \text{h.c.}
\] (2.10)

Weak isospin gauge invariance enforces such a Lagrangian term to be accompanied with the interaction of a left-handed bottom quark to the second component of the \( X \)-doublet. This induces the fast decay of the neutral \( \chi \) fermion via an off-shell \( X \)-state,

\[
\chi \rightarrow b_L (X_{1/3}^\mu)^* \rightarrow b_L u_L d_R / b_L d_L u_R ,
\] (2.11)

so that this setup does not predict any monotop signal.

From all the above considerations, a monotop signature arising from the decay of a new coloured resonance can only be generated via scalar mediator \( \varphi_{2/3}^s \), singlet under the weak isospin group and coupling only to right-handed fermions,

\[
\mathcal{L}_{\text{res}} = \mathcal{L}_{\text{kin}} + \lambda_S \varphi_{2/3}^s d_R^C d_R + \lambda_S' \varphi_{2/3}^s t_R \chi + \text{h.c.}
\] (2.12)

In terms of the effective model of Refs. [8, 15] shown in Eq. (2.1), this implies that

\[
a_{VR}^q = b_{VR}^q = a_{V}^{1/2} = b_{V}^{1/2} = 0 , \quad a_{SR}^q = b_{SR}^q \quad \text{and} \quad a_{SR}^{1/2} = b_{SR}^{1/2} .
\] (2.13)

This case corresponds to an \( R \)-parity violating supersymmetric scenario where the only \( R \)-parity violating terms of the superpotential are the so-called \( UDD \) interactions violating the baryon number. The scalar field \( \varphi_s \) is then identified with a right-handed stop and \( \chi \) with a neutralino (bino).

As \( \varphi_s \) is a colour triplet, it will also be pair produced via standard QCD interactions and could be directly searched for in non-monotop processes,

\[
pp \rightarrow \varphi_s \varphi_s^* \rightarrow jj jj , \quad jj t\chi \quad \text{and} \quad t\chi \bar{t}\chi .
\] (2.14)
As a consequence, existing LHC searches at a center-of-mass energy $\sqrt{s} = 8$ TeV could further constrain the model. We however assume the new scalar field to be heavy, so that the new physics contributions to the three channels above are phase-space suppressed. This allows one to evade the bounds possibly arising from the analysis of data recorded during the LHC run I. However, the situation will change with run II at $\sqrt{s} = 13$ TeV. In this case, combining monotop searches to analyses of events whose final states feature paired dijet and top-antitop plus missing energy signatures may further constrain the model parameters.

### 2.2 Non-resonant monotop production

In the second class of scenarios yielding the production of monotop states, the top quark is produced in association with an invisible bosonic field that couples in a flavour-changing way to top and light up-type (up or charm) quarks. Denoting these new states by $\phi$ (in the scalar case) and $V_\mu$ (in the vector case), this series of models can be described by the effective Lagrangian \[8\]

$$L_{\text{non-res}} = L_{\text{kin}} + \left( \phi \bar{u}_i \left[ (a_{FC}^0)^{ij} + (b_{FC}^0)^{ij} \gamma^5 \right] u_j + \text{h.c.} \right) + \left( V_\mu \bar{u}_i \left[ (a_{FC}^1)^{ij} \gamma^\mu + (b_{FC}^1)^{ij} \gamma^\mu \gamma^5 \right] u_j + \text{h.c.} \right),$$

(2.15)

where $L_{\text{kin}}$ contains kinetic and mass terms for all new fields and the coupling strengths $a_{FC}$ and $b_{FC}$ are symmetric matrices under the exchange of flavour indices. The states $\phi$ and $V$ are in general not stable since they couple to quarks. A missing energy signature is therefore enforced by requiring these fields either to be long-lived so that they decay outside of the detector, or to decay predominantly into a pair of additional neutral stable particles. In particular, the latter possibility has been proposed in the framework of flavourful dark matter models \[9\], where the extra boson ($\phi$ or $V$) is a mediator of the interactions of the dark matter candidate with the Standard Model particles.

The main issue of this class of models is to make sure that the new boson leads to a missing energy signature in a detector. In this work, we address it by assuming that the $\phi/V$ field dominantly decays into a pair of dark matter candidate particles.\(^1\) In this case, extra constraints arise from the requirement that the particle the boson decays into is a good candidate for dark matter, or that at least it does not overpopulate the Universe.

As already stated in Section 2.1, the interactions of a scalar field to quarks involve both the right-handed and left-handed components of the fermions. There are thus two possible structures allowed for the couplings. Focusing on contributions leading to a monotop hadroproduction rate enhanced by parton density effects, the associated Lagrangian (the first line in Eq. (2.15)) can be rewritten as

$$\phi \left( y_1 \bar{t}_R u_L + y_2 \bar{u}_R t_L \right),$$

(2.16)

\(^1\)One could also consider that neither the boson nor its decay products are stable, but instead long-lived. Although this is a viable assumption, this implies further complications in the building of the model. We therefore stick with the minimal case.
where the $y$ parameters denote generic coupling strengths. Consequently, the scalar $\phi$ field must transform as a doublet of $SU(2)_L$ with an hypercharge quantum number of $1/2$,

$$\phi = (1, 2, 1/2) \equiv (\phi^+, \phi^0)^t.$$  \hspace{1cm} (2.17)

Rendering the Lagrangian of Eq. (2.16) gauge invariant implies the addition of interactions between the charged component field $\phi^+$ and quarks. The $\phi^+$ field will therefore always promptly decay into two-body final states, $\phi^+ \rightarrow u\bar{b}$ or $t\bar{d}$. Analogously, the neutral component $\phi^0$ could also decay into an associated particle pair comprised of a top and an up quark, $\phi^0 \rightarrow u\bar{t} + \bar{t}u$, as well as into a three-body final state via the exchange of a virtual charged scalar field, $\phi^0 \rightarrow W^-[\phi^+]^* + W^+[\phi^-]^* \rightarrow W^-\bar{b}u + W^+\bar{b}u$ or $W^-\bar{d}t + W^+\bar{d}t$. All these decay channels are however assumed to be negligible when compared to a decay into a pair of dark matter particles. In this case, no minimal coupling to a single stable state is achievable since $\phi$ is a doublet of $SU(2)_L$, and one must design an interaction of the $\phi$ state to two extra fields whose combination forms a doublet of $SU(2)_L$. If we restrict ourselves to $\phi^0$-decays into fermionic particles, the most minimal option is given by the Lagrangian

$$L_{\phi-\text{decay}} = y_\chi \bar{\chi}_d \chi_s + \text{h.c.}, \hspace{1cm} (2.18)$$

where $\chi_s$ is an electroweak singlet and $\chi_d$ a weak doublet with an hypercharge of $1/2$. This term induces decays of both components of $\phi$

$$\phi_0 \rightarrow \chi_s \chi_0^d \quad \text{and} \quad \phi^+ \rightarrow \chi_d^+ \chi_s \rightarrow [W^+]^* \chi_0^d \chi_s,$$  \hspace{1cm} (2.19)

the charged component $\chi_d^+$ being taken heavier than, but close in mass to, the neutral component $\chi_0^d$ so that both neutral fields $\chi_s$ and $\chi_d$ can be seen as viable dark matter candidates.

As a consequence of this non-minimal dark sector of the model, monotop production via flavour-changing interactions of up-type quarks with a new invisible scalar field will always be accompanied by an extra single top production mode

$$pp \rightarrow t\phi^- \rightarrow t\chi_0^d \chi_0^s [W^-]^*.$$  \hspace{1cm} (2.20)

The nature and magnitude of the associated effects are very benchmark dependent. For instance, a small mass splitting between the component fields of $\chi$ leads to very soft $W$-boson decay products, so that the process of Eq. (2.20) would imply new contributions to monotop production. On the other hand, in the case of larger mass splittings, related new physics scenarios feature an LHC signature comprised of a single top quark and an isolated lepton.

Nevertheless, we choose to keep the focus on minimal models, and therefore ignore, in the rest of this work, scenarios where monotop states are produced from flavour-changing interactions of up-type quarks with a scalar particle mediating dark matter couplings to the Standard Model.

\footnote{Due to reasons already stated in Section 2.1, the $\phi^+$ and $\phi^0$ states are assumed to have similar masses.}
When the mediator is a vector boson $V$, one can design very simple models since it can be singlet under the electroweak group. In this setup, the associated couplings (shown in the second line of the Lagrangian of Eq. (2.15)) involve either right-handed or left-handed quarks and take the form

$$
\left( a_R V_\mu \bar{t}_R \gamma^\mu u_R + a_L V_\mu (\bar{t}_L \gamma^\mu u_L + \bar{b}_L \gamma^\mu d_L) + \text{h.c.} \right),
$$

(2.21)

where the $a_{L,R}$ parameters denote the strengths of the interactions of the $V$-field with up and top quarks. As in the rest of this section, we have restricted ourselves to interactions focusing on the monotop hadroproduction modes enhanced by parton densities.

The Lagrangian terms of Eq. (2.21) open various decay channels for the $V$-field. First, the left-handed couplings allow the mediator to always promptly decay into jets, $V \to bd + db$. Next, the importance of the decays into top and up quarks (this time both in the context of left-handed and right-handed couplings) depends on the mass hierarchy between the mediator and the top quark, the tree-level decay $V \to t\bar{u} + u\bar{t}$ being only allowed when $m_V > m_t$. Furthermore, when $m_V < m_t$, a triangle loop-diagram involving a $W$-boson could also contribute to the decay of the $V$-field into a pair of jets, $V \to d_i \bar{d}_j$. Finally, when $m_W < m_V < m_t$, the three-body decay channel $V \to bW^+u + \bar{b}W^-u$ is open, mediated by a virtual top quark. A monotop signal is thus expected only when the $V$-field is invisible and dominantly decays into a pair of dark matter particles. Since $V$ is an electroweak singlet, the associated couplings can be written, in the case of fermionic dark matter, as

$$
\mathcal{L}_{V-\text{decay}} = V_\mu \left( g_R \bar{\chi}_R \gamma^\mu \chi_R + g_L \bar{\chi}_L \gamma^\mu \chi_L \right),
$$

(2.22)

where $\chi$ is a Dirac fermion, singlet under the Standard Model gauge symmetries. The consistency of the model, i.e., the requirement that $V$ always mainly decays into a pair of $\chi$-fields and not into one of the above-mentioned visible decay modes, implies constraints on the Lagrangian parameters. They will be studied in details in the next section, together with other requirements that can be applied to viable non-resonant monotop scenarios.

Summarising all the considerations above, the minimal gauge-invariant Lagrangian yielding monotop production in the flavour-changing mode is given by

$$
\mathcal{L}_{\text{non-res}} = \mathcal{L}_{\text{kin}} + \left( a_R V_\mu \bar{t}_R \gamma^\mu u_R + a_L V_\mu (\bar{t}_L \gamma^\mu u_L + \bar{b}_L \gamma^\mu d_L) + \text{h.c.} \right)
+ V_\mu \left( g_R \bar{\chi}_R \gamma^\mu \chi_R + g_L \bar{\chi}_L \gamma^\mu \chi_L \right),
$$

(2.23)

In the notations of Ref. [8, 15] employed in Eq. (2.15), the above choice corresponds to setting

$$
a_{FC}^0 = b_{FC}^0 = 0 ,
$$

(2.24)

Moreover, the parameter basis of Eq. (2.15) ($a_{FC}^1, b_{FC}^1$) is fully equivalent to the one of Eq. (2.23), ($a_L, a_R$).

3 Monotop phenomenology specific to non-resonant models

Some features of the resonant models mediated by a scalar, like the lifetime of the invisible fermion produced in association with the top, have been studied in details in Ref. [11].
In the following, we therefore focus on various features of non-resonant spin-1 models by studying the effective lifetime of the invisible vector, associated single top signals, and the dark matter relic density.

We separately consider two regions of the parameter space which have very different phenomenology: the case where the mediator is lighter than the top quark (its mass $m_V$ being smaller than the top mass $m_t$) and the case where it is heavier, with $m_V > m_t$.

3.1 Mediators lighter than the top quark

When the spin-1 mediator $V$ is lighter than the top quark, its possible decay modes into a pair of top and lighter quarks are kinematically forbidden. At tree-level, $V$ can therefore only decay into a multibody final state such as $V \to u\bar{b}W^-$ or $\bar{u}bW^+$, where the $W$-boson is virtual when $m_V < m_W$ ($m_W$ denoting the $W$-boson mass). In this mass range, loop-induced decays must however be considered too. For instance, a triangle loop-diagram with a $W$-boson exchange generates couplings to down-type quarks, which consequently opens a dijet decay channel. As the decay channels in this region are either kinematically or loop-suppressed, one may wonder whether $V$ may be long-lived without the need for an additional invisible decay channel. Another interesting property of this mass region is that a new decay of the top quark is allowed, $t \to uV$, and extra constraints on monotop scenarios could therefore be extracted from, e.g., top width measurements or the analysis of $t\bar{t}$ events when one of the top quarks decays into a jet plus missing energy.

3.1.1 Loop-induced decays of the mediator

Light mediators, below the top mass threshold, may decay dominantly into two jets via loop-induced interactions. The structure of the loop crucially depends on the chirality of the monotop couplings, and we can study separately the two limiting cases $a_L = 0, a_R \neq 0$ and $a_L \neq 0, a_R = 0$, the two setups being prevented from interfering in the limit of massless light quarks.

We start with the left-handed scenario, $a_L \neq 0$ and $a_R = 0$. It has been shown in Section 2.2 that embedding this class of monotop effective theories within $SU(2)_L$ implies that the mediator $V$ couples to down-type quarks, and therefore always decays to two jets at tree-level. As a consequence, an extra invisible fermion $\chi$ had been added to the theory, allowing one to tune the partial width related to the process $V \to \chi\chi$ to be dominant and preserve in this way the monotop signature. The $Vdd$ vertices play also an important role for the consistency of the theory. Assuming an anomalous coupling approach, one could imagine an effective model where the left-handed couplings of the mediator to up and top quarks are allowed and those to down-type quarks neglected. The latter will however be regenerated via triangle-loop diagrams involving a $W$-boson that are logarithmically divergent in the ultraviolet limit. Following a standard procedure, these divergences must be treated with appropriate counterterms that naturally appear after renormalization of the (complete) Lagrangian of Eq. (2.23). This consequently motivates the use of an $SU(2)_L$-invariant Lagrangian from the beginning. In a similar fashion, those couplings will generate mediator interactions to all combinations of up-type quarks at the loop level, so that the initial hypothesis of a unique coupling between up and top quarks is unphysical and extra
decay channels must be considered too. All these higher-order contributions to the total width are however loop- and/or CKM-suppressed and can thus be neglected, in particular as we recall that the invisible partial decay width is tuned to be dominant. Consequently, hints for new physics are still expected to occur in monotop events, although additional quark-mediator couplings could induce other observable effects that may imply stronger bounds on the parameter space. For instance, it is not unlikely that a monotop signal could be accompanied with a monojet signal in scenarios with a non-vanishing $a_L$ parameter.

We now turn to the study of right-handed scenarios with $a_L = 0$ and $a_R \neq 0$. The Lagrangian of Eq. (2.23) simplifies and there are no more interactions of down-type quarks with the mediator. However, they are as above generated at the loop-level via $W$-boson triangle diagrams. Since weak interactions are left-handed, the chiralities of the quarks involved in these diagrams must be flipped, which implies that the loop-induced couplings are proportional to the product of the up and top masses $m_u m_t$. Contrary to setups where monotops are produced from left-handed interactions of the mediator with quarks, the loop-induced $V d_d d_d$ couplings are this time finite, in line with the fact that no associated counterterm appears after renormalization. The interaction strength reads, in the limit of small light quark masses,

$$g_{V d d d}^{1\text{-loop}}(a_R) = \frac{\alpha a_R}{4\pi s_W} \frac{m_u}{m_t} (V^*_{ud} V_{td} + V^*_{td} V_{ud}) \tilde{c}_0,$$

where $\alpha$ stands for the electromagnetic coupling constant, $s_W$ for the sine of the weak mixing angle and $V_{ij}$ for the elements of the CKM matrix. In addition, the loop factor

$$\tilde{c}_0 = m_t^2 C_0(p_1, -(p_1 + p_2); m_W, m_t, 0)$$

depends on the Passarino-Veltman three-point function $C_0$ where $p_1$ and $p_2$ are the momenta of the external down-type quarks. We can therefore calculate the partial width associated with the decay $V \to \bar{d} d d$ which reads, after summing over all down-type quark flavours,

$$\Gamma(V \to jj) = \frac{\alpha^2 a_R^2}{64\pi^3 s_W^2} \frac{m_V m_u^2}{m_t^2} |\tilde{c}_0|^2.$$

We observe that it exhibits both a loop-suppression and a $(m_u/m_t)^2$ factor, so that it is expected to be numerically small.

In Figure 1, we show the partial width in Eq. (3.3) as a function of the mediator mass for $a_R = 0.04$ (left panel). On the right panel of the figure, the partial width is translated as an upper bound on the value of $a_R$ in order for $V$ to have a mean decay length of at least 50 metres so that it is long-lived enough to decay outside of typical hadron collider detectors. The figure shows that the lifetime of $V$ would be long enough only for values of the coupling satisfying $a_R \lesssim 10^{-2}$. Such small values may however challenge the possible observation of a monotop signal at the LHC by reducing the associated production cross section. It should also be mentioned that above the $W$-boson threshold, a tree-level three-body decay is kinematically open, which further shortens the decay length of $V$. Finally,
Figure 1. Partial decay width associated with the loop-induced decay of the $V$-field into down-type quarks as a function of the $V$-boson mass, as given by Eq. (3.3) (left panel). We consider scenarios in which $a_R = 0.04$ and $m_V$ is kept smaller than the top mass. The results are translated, in the right panel, as a bound on $a_R$ that ensures that the $V$-boson has a decay length of at least 50 m.

Figure 2. Partial width associated with the $t \rightarrow Vu$ decay mode of the top quark as a function of the $V$-boson mass and the $a_R$ coupling. We show curves for a partial width of 1, 0.5, 0.1 and 0.01 GeV. The solid black curve corresponds to the upper bound on $a_R$ from a partial width less than 0.5 GeV.

in cases where the model features a coupling of the $V$-field to top and charm quarks, the partial width of Eq. (3.3) would exhibit an enhancement proportional to $(m_c/m_u)^2$.

In summary even for monotop scenarios in which the mediator cannot decay into a top quark, its lifetime is generally too short and one needs to complete the model by adding a decay channel into an invisible state. Although the class of minimal scenarios described in this section features a light extra vector boson, the setup is compatible with current Tevatron and LHC bounds on monotop production as the latter are always derived under the assumption of very large coupling values of $O(0.1)$ [18, 19]. They could however be constrained by other observations, as will be shown in the next subsections.

3.1.2 Single top constraints on monotop scenarios

Motivated by minimality principles, we have discussed, in the previous section, appealing monotop scenarios in which the mediator $V$ is lighter than the top quark. In this case, the former couples to up and top quarks via right-handed couplings and one needs to add an
invisible decay channel to potential dark matter particles $\chi$ to guarantee a monotop signature, unless the coupling strength $a_R$ is very small. On different grounds, these scenarios feature a new decay channel for the top quark, $t \to uV$. This observation can be used to further restrict the viable regions of the parameter space by imposing that new physics contributions to the top width do not challenge the measured value of $\Gamma_t = 2.0 \pm 0.5$ GeV [26]. Assuming a good agreement between the Standard Model expectation and the top width measurement, the partial width $\Gamma(t \to V u)$ can thus be enforced to be of at most 0.5 GeV. On Figure 2, we present the dependence of this partial width on the coupling $a_R$ and the mediator mass $m_V$. We observe that for couplings smaller than 0.01, new physics effects in the top width are predicted to be very small, except when the mediator is almost massless. This consequently disfavors such setups in which the mediator is very light, even in cases with coupling strengths of $\mathcal{O}(0.001)$.

Kinematically allowed $t \to V u$ decays also imply that monotop events can be issued from the production of a top-antitop pair when one of the top quarks decays into a $V$-boson and a light quark,

$$ pp \to t\bar{t} \to t\bar{u}V \quad \text{or} \quad pp \to t\bar{t} \to \bar{t}uV . $$

This process induces additional contributions to the production of a monotop system ($tV$ or $\bar{t}V$) in association with an additional jet, a signature already accounted for in the LHC monotop analysis of Ref. [19]. How much this new channel will contribute to the monotop signal depends on the cuts employed in the experimental analysis. However, due to the large $t\bar{t}$ cross section, these effects cannot be neglected.

Complementary constraints on this channel could be deduced from Standard Model single top analyses whose signal regions could capture monotop events as above. For instance, both CMS [27] and ATLAS [28] have analyses dedicated to the measurement of the single top cross section in the $t$-channel which contain a region that could be populated by monotop events as above.\(^3\) In the CMS analysis, events are selected by requiring one single isolated electron or muon and exactly two jets, one of them being $b$-tagged. The background is reduced by requiring an important amount of missing energy and by imposing that the transverse mass computed after combining the lepton transverse momentum with the missing transverse momentum is large. A final selection is performed by means of an advanced multivariate technique. We have nevertheless to ignore this last step of the selection as the amount of information provided in the experimental publication is not sufficient for satisfactorily recasting it (see Ref. [29] for more information on this aspect).

We simulate our new physics signal by using the monotop model [8] implemented in the FeynRules package [30, 31], tuning the model parameters to the setup of Eq. (2.23), so that we can export the model to a UFO library [32] that is then linked to MADGRAPH5\_aMC@NLO [33]. The generated parton-level events have subsequently been processed by PYTHIA [34] for parton showering and hadronization and by DELPHES [35] for detector simulation, making use of the recent ‘MA5Tune’ [36] of the CMS detector.

\(^3\)Other single top analyses could be considered. However, they in general use multivariate techniques that cannot be employed in the reinterpretation framework pursued below.
description of DELPHES. The CMS analysis of Ref. [27] has finally been implemented in the MADANALYSIS5 framework [37, 38], which has allowed us to derive exclusion bounds at the 95% confidence level in the \((m_V, a_R)\) plane, as shown on Figure 3. The figure also shows the constraint from the top width, and from the dedicated CMS monotop search [19]. The monotop search is currently more sensitive. However, the bound from the single top is a rough estimate, and the bound may be much stronger once the full analysis, including the multivariate selection, is taken into account. Nevertheless, our result shows that the constraints from single top searches can play an important role in constraining monotop scenarios.

### 3.1.3 Dark matter constraints

We have argued that, even for mediator masses below the top threshold, an invisible decay channel is typically needed in order for the monotop signature to be present. The simplest way out is to couple \(V\) to a fermionic stable dark matter candidate \(\chi\). However, in a minimal scenario where \(V\) is the only mediator for the interactions of the dark matter candidate, one needs to ask whether the relic abundance of \(\chi\) is enough to fulfill the bounds from observations. Below the top threshold, the main annihilation process \(\chi\chi \rightarrow V \rightarrow t\bar{u}\) and \(\bar{t}u\) is kinematically forbidden, so that the annihilation of dark matter particles can only proceed to a three-body or four-body final state (via a virtual top quark), or via loop-diagrams \(\chi\chi \rightarrow V \rightarrow d_i\bar{d}_j\). As discussed in Section 3.1.1, the loop contributions are suppressed by the mass of the light up-type quark that the mediator couples to, so that the \(\chi\chi\) annihilation rate may be too slow for the stable particle \(\chi\) not to overpopulate the Universe.

We are therefore left with either a scenario of small \(a_R\) coupling, where the mediator \(V\) is a long-lived particle, or with a non-minimal model with an invisible decay channel where \(\chi\) is either the next-to-minimal odd particle or long-lived itself. In any case, the constraint from the dark matter abundance plays a crucial role for the viability of the model and should be carefully verified.
3.2 Mediators heavier than the top quark

Scenarios exhibiting mediator masses above $m_t$ are very different from the light case discussed in Section 3.1: the mediator $V$ can always decay into a top quark. Including in the model a $V$-decay channel into an invisible state to be considered as a dark matter candidate is thus always necessary. Moreover, the top quark cannot decay into the mediator, which allows one to avoid constraints from standard single-top signatures. Focusing on the minimal case, we study below the interesting interplays between the requirement that the invisible channel dominates and bounds originating from the relic density of the dark matter candidate.

3.2.1 Tree-level decays of the mediator

When the $V$-boson is heavier than the top quark, it can decay into either a pair of down-type quarks, an associated pair comprised of a top quark and a lighter quark or a pair of dark matter particles, as already discussed in Section 2.2. Since the first two decay modes are driven by the same interaction vertices allowing one for monotop production, we need to make sure that the invisible decay channel always dominates. The relevant partial widths are given by

$$
\Gamma(V \rightarrow bd + \bar{b}d) = \frac{m_V}{4\pi} |a_L|^2 ,
$$

$$
\Gamma(V \rightarrow tu + \bar{t}u) = \frac{m_V}{4\pi} \left( |a_R|^2 + |a_L|^2 \right) \left( 1 - \frac{m_t^2}{2m_V^2} - \frac{m_t^4}{2m_V^4} \right) ,
$$

$$
\Gamma(V \rightarrow \chi\chi) = \frac{m_V}{24\pi} \sqrt{1 - 4 \frac{m_\chi^2}{m_V^2}} \left[ \left( |g_{L\chi}|^2 + |g_{R\chi}|^2 \right) \left( 1 - \frac{m_\chi^2}{m_V^2} \right) + 6 \frac{m_\chi^2}{m_V^2} \Re \{g_{L\chi}g_{R\chi}^*}\right] ,
$$

where we neglect all quark masses but the top mass. In addition, we denote by $m_\chi$ the mass of the dark matter candidate.

Focusing on the simplest subclass of scenarios where the couplings of the $V$-boson to left-handed quarks are all vanishing ($a_L = 0$), we study typical constraints that can be imposed on ratios of the $g_{L\chi}$, $g_{R\chi}$ and $a_R$ parameters when they are all assumed to be real quantities. Since ratios of branching ratios are equivalent to ratios of partial widths, we use this latter quantity and show, in Figure 4, the maximum value of the $a_R$ coupling strength in units of the $\chi V$ coupling that ensures the $V$-field to decay invisibly in at least 99% of the cases. In the left panel of the figure, we consider scenarios where $g_{R\chi}$ vanishes (the same result holds for vanishing $g_{L\chi}$), while in the right panel of the figure, we assume vector-like couplings, $g_{L\chi} = g_{R\chi} = g_{V\chi}$. In general, the coupling to the top $a_R$ (that is responsible for the monotop signal) has to be quite small compared to the coupling to the dark matter candidate in order for the mediator $V$ to be invisible, unless the mass of the mediator $V$ is close to the top mass. On the contrary, if the mass of $V$ is close to the $\chi\chi$ threshold, the invisible decays are suppressed. This study shows that it is not straightforward to have $V$ to decay invisibly, and this constraint may play an important role in the interpretation of the signal, especially when associated with the study of the properties of $\chi$ as a dark matter candidate. We study more in detail this question in the next subsection.

$^4$The results have been checked using the decay module of FeynRules [39].
3.2.2 Dark matter constraints

We have seen that, in order to avoid visible decays of the mediator $V$, it has to be coupled to a stable particle $\chi$ and the decay $V \to \chi \chi$ must always dominate. If $\chi$ is stable, and if the model is minimal in the sense that $V$ is the only mediator of interactions between the dark sector and the Standard Model, then the only annihilation process that will determine the thermal relic abundance of $\chi$ is $\chi \chi \to V \to t\bar{u}$ and $\bar{t}u$. Such process is proportional to the same coupling that gives rise to the monotop signature at the LHC, and also to the coupling of $V$ to dark matter. By studying the relic abundance of $\chi$ one can therefore derive interesting constraints on the couplings, especially when imposing that the relic abundance is smaller than the measured density of dark matter. Those restrictions can in principle always be evaded by assuming that there are additional mediators, or that $\chi$ is not a stable particle but rather a long-lived one that decays on cosmological time scales. In the rest of the section, we nevertheless focus on the minimal case of $\chi$ being the only dark matter candidate.

As the relic abundance decreases with increasing annihilation cross sections, one can calculate an upper bound on the product of $a_R$ with the couplings of $V$ to the dark matter. This has been computed by implementing the model described by the Lagrangian of Eq. (2.23) in CALCHEP [40] and using approximate formulas for the relic abundance. We consider, for concreteness, a vectorial model with $g_{L\chi} = g_{R\chi} = g_{V\chi}$. The results of the calculation are shown in Figure 5, where we present the lower bound on $a_R \times g_{V\chi}$ as a function of the mediator mass $m_V$ and the dark matter mass $m_\chi$. We restrict ourselves to values of the $\chi$ mass above the top threshold, $2m_\chi > m_t$, so that a two-body process is kinematically allowed. Below the top threshold, the dark matter candidate can only
Figure 5. Lower bound on $g_{V\chi} \times a_R$ from the dark matter relic abundance as a function of $m_V$ and $m_\chi$.

Annihilate into three-body final states or via loop-induced processes, so that the annihilation cross section is too small and the $\chi$ particle overpopulates the Universe. The figure shows that the product of couplings is bound to be larger than about 0.1, with the lower bound increasing towards the top threshold as the phase space closes down, and becomes smaller towards the $V$ threshold $2m_\chi = m_V$ where the resonant $V$ exchange enhances the annihilation. We recall that the $V$-boson mass must be at least twice as large as the dark matter candidate mass to allow invisible decays for $V$. The corresponding regions of the parameter space are tagged as kinematically inaccessible.

This result, very interesting per se, can be combined with other constraints to better determine the viable regions of the parameter space of the model. The requirement that the invisible $V$-decay dominates has allowed us, in Section 3.2.1, to calculate a lower bound on the ratio $g_{V\chi}/a_R$ which depends on the mediator and dark matter masses (see Figure 4). Multiplying it with the limits derived from the relic abundance predictions, we extract a lower bound on $g_{V\chi}$ independently of the value of $a_R$: the results are shown in Figure 6. The lower bound on $g_{V\chi}$ is found to grow with smaller values of the $\chi$ mass. Moreover, near the top threshold, it reaches values well above unity, tending hence to the non-perturbative regime.

Under the assumption that $\chi$ is the only dark matter candidate of the theory, we can further restrict our analysis to parameter space regions where the values of the couplings are such that the bound from the dark matter abundance is saturated. We first reinterpret, as a function of the masses, the limits calculated in the CMS monotop search [19] by accounting for an invisible branching ratio of the mediator that may not be 100%. Next, we correlate these to the dark matter results: for increasing values of $a_R$, the coupling $g_{V\chi}$ has to be smaller to satisfy the dark matter constraints. This indicates that an enhancement of the $tV$ production rate (by increasing $a_R$) is accompanied by a reduction of the invisible branching.
Figure 6. Lower bound on $g_{V\chi}$ obtained combining the dark matter relic abundance constraints with the requirement that the mediator $V$ decays invisibly in 99% of the cases.

The result is shown on the right panel of Figure 7. Above the blue curve, the argument of the square root is negative and the inequality of Eq. (3.7) has no solution, therefore there is no bound that can be applied on $a_R$. Below the blue line, near the top threshold, the Dark Matter constraint requires larger couplings and therefore larger monotop rates are allowed, thus a bound on $a_R$ can be calculated. Naturally, larger portions of the parameter space are expected to be covered with the upcoming run II of the LHC.

The region where the monotop signal is suppressed can have interesting additional features. The boson $V$ may dominantly decay into top and lighter quarks, yielding at the same time a signature comprised of same-sign top quark pairs ($tV \to tt\bar{u}$) and extra...
Figure 7. Reinterpretation of the CMS monotop limits of Ref. [19] in terms of $a_R$ (left) for a 100% invisible mediator. The results are then used to determine the viable regions of the parameter space when enforcing dark matter and LHC constraints as shown by Eq. (3.7) (right). The region above the blue line is found not to be bounded by current searches. Below the blue line, limits on $a_R$ deduced from the left panel of the figure are in order.

contributions to top-antitop production ($tV \rightarrow t\bar{t}u$) that may be difficult to observe due to the overwhelming $t\bar{t}$ Standard Model background. These extra channels deserve a particular attention, in particular in upcoming data from LHC collisions at $\sqrt{s} = 13$ TeV.

4 Conclusions

Monotop final states comprised of a single top quark produced in association with missing energy can be a striking sign of new physics at the LHC. The main production mechanisms can be divided into two classes: resonant production, where a heavy coloured boson is first produced in the $s$-channel and further decays via its couplings to a single top quark and an invisible neutral fermion, and non-resonant production where the top quark is produced in association with an invisible boson that couples to top and up (or charm) quarks. A complete and model independent parametrisation of the two channels has been provided in Ref. [8]. In the present work, we have revisited this description by embedding the effective interactions in an $SU(2)_L \times U(1)_Y$ invariant formalism. In doing so, we have shown that, depending on the chirality of the tops, a complete model contains necessarily extra states and couplings that may spoil the monotop signal, or add more new physics signatures that should be studied in association with the monotop one.

We have identified two minimal setups. In the first case, a scalar field is resonantly produced by the fusion of a pair of down-type quarks and couples to a right-handed top quark and a new invisible fermion, like a right-handed stop in $R$-parity violating supersymmetry. In the second case, a vector state couples to right-handed top and up quarks and decays dominantly into new invisible fields, like in models of dark matter where the dark
sector couples to the Standard Model via a flavour-sensitive mediator. We have further investigated the phenomenology of the second class of models that can be split into two subclasses, depending on the mass of the mediator.

For mediators lighter than the top quark, their visible decay modes are either loop-suppressed or CKM-suppressed, or both. Nevertheless, one always needs to add (and tune the couplings of) an invisible field to prevent the mediator from decaying inside a typical hadron collider detector as this would otherwise spoil the monotop signature originally motivating the model. An important feature of these scenarios is that they allow for the top quark to decay into the mediator and an extra jet. This feature can enhance the monotop production rate, as the monotop system can be produced in association with an extra jet from $t\bar{t}$ events when one of the top quarks decays in the exotic channel. Such events could also be searched for in standard typical single-top searches, as they are expected to populate signal regions of associated analyses. We have indeed observed that a CMS analysis of single top events could imply significant constraints on the mediator couplings, competitive and sometimes stronger than those obtained from monotop searches.

Scenarios with a mediator mass above the top threshold have a very different phenomenology as the mediator decays significantly into top quarks and jets. One needs a large coupling to the invisible sector in order to preserve the monotop signature. Describing the dark sector with a new fermion $\chi$, we have found that the latter could be a viable dark matter candidate if heavier than half the top quark mass, with a correct relic abundance driven by its annihilation via an $s$-channel mediator into a top and an up quark. We have used relic abundance constraints to derive lower bounds on the product of the couplings of the mediator to quarks and to the dark matter candidate. We have then further restricted the monotop parameter space by combining cosmological and collider results and enforcing the mediator to decay mostly invisibly. We have found that the issue of the perturbativity of the model could be raised for dark matter masses close to the top mass and that the parameter space turns out to be largely constrained when the $\chi$ fermion is demanded to reproduce the observed relic density. However, a large portion of the parameter space is still left unconstrained by current data and future experimental results are in order, in particular analyzing a same-sign top quark pair final state arising from the visible decays of the mediator.

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