LHC constraints and potential on resonant monotop production

Giacomo Cacciapaglia^{1,a}, Eric Conte^{2,b}, Aldo Deandrea^{1,c} , Benjamin Fuks^{3,4,d}, Hua-Sheng Shao^{3,e}

¹ Univ. Lyon, Université Claude Bernard Lyon 1, CNRS/IN2P3, UMR5822 IPNL, 69622 Villeurbanne, France

² Institut Pluridisciplinaire Hubert Curien/Département Recherches Subatomiques, Université de Strasbourg/CNRS-IN2P3, 23 Rue du Loess, 67037 Strasbourg, France

³ Laboratoire de Physique Théorique et Hautes Énergies (LPTHE), UMR 7589, Sorbonne Université et CNRS, 4 place Jussieu, 75252 Paris Cedex 05, France

⁴ Institut Universitaire de France, 103 boulevard Saint-Michel, 75005 Paris, France

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Abstract We discuss the phenomenology associated with a resonant monotop collider signal, i.e. a signal in which a single top quark is resonantly produced in association with missing energy through an s -channel scalar exchange. We study both the bounds originating from dedicated monotop searches performed by the ATLAS and CMS experiments, and the constraints associated with other processes that could be induced by a new physics context favouring monotop production at colliders. The latter class of constraints includes, in particular, the recasting of analyses from the LHC and the TeVatron. All theoretical calculations are performed at the next-to-leading order accuracy in QCD, and we finally combine all results to establish the present limits on the parameter space and test the relevance of the monotop signal at the LHC Run 2.

1 Introduction

Monotop production at colliders consists in the production of a single top quark in association with a large amount of missing transverse energy. This quite peculiar final state has been investigated at the LHC by both the ATLAS and CMS collaborations, and both at Run 1 and 2. As monotop production is heavily suppressed in the Standard Model, its observation would consist in a clear sign of physics Beyond the Standard Model. In a new physics context, there exist two main different monotop production mechanisms [1–3]. In the first of them, the monotop system is produced from a

coloured scalar or vector resonance that decays into a top quark and an invisible neutral fermion, whereas in the second of them, monotops arise from the production of a single top quark in association with an invisible scalar or vector boson via the flavour-changing couplings of the latter to the top and light quarks. After imposing invariance under the full Standard Model gauge symmetry and invoking simplicity, it can be shown that only scalar resonant monotop production and vector flavour-changing monotop production are consistent [3]. Whilst several existing experimental [4–7] and phenomenological [2, 3, 8–10] studies focus on the flavour-changing option, the possibility of resonant production has been less studied, at least comprehensively [2, 5, 7, 11]. On different grounds, monotop signatures have also been considered in the case of compressed supersymmetry [12–16] and models with vector-like fermions [17] or explaining neutrino masses [18].

In this work, we reconsider the resonant production of monotop systems via an intermediate coloured scalar, which consists in the simplest model featuring monotop production as a key new physics signal and that is allowed after imposing invariance under the Standard Model gauge symmetry group [3]. In practice, we embed the generic effective Lagrangian for resonant monotop production [1] within the full gauge symmetry requirements of the Standard Model, which severely constrains the couplings and quantum numbers of the mediator. Hence the coloured scalar mediator has an electric charge of $2/3$ and consists in a colour triplet σ that couples to a pair of different-flavour right-handed down-type quarks. Single mediator production, therefore, occurs via these di-quark couplings, while the decay of the mediator into a (right-handed) top quark plus an invisible neutral fermion χ occurs through an independent coupling parameter. In order for the model to stay monotop-motivated, it is crucial that the fermion χ remains undetected when pro-

^a e-mail: g.cacciapaglia@ipnl.in2p3.fr

^b e-mail: eric.conte@iphc.cnrs.fr

^c e-mail: deandrea@ipnl.in2p3.fr

^d e-mail: fuks@lpthe.jussieu.fr

^e e-mail: huasheng.shao@lpthe.jussieu.fr

duced, and thus decays outside the detector. Remarkably, this model resembles a supersymmetry-inspired simplified model in which the Standard Model field content is supplemented by neutralino and a right-handed top squark featuring R -parity violating couplings to the down-type quark sector.

Besides the direct investigation of monotop probes, this model is also constrained by many other searches for new physics that thus already limit the available parameter space. Hence we will take into account searches from LHC Run 1 and 2 involving jets and top quarks in the final state, as well as dijet searches at the LHC and the TeVatron. Moreover, constraints on the decay length of the invisible (unstable) particle produced in association with the top quark and contributions to the top width also importantly restricts the parameter space. Another important point concerning the resonant production of a coloured spin-0 boson is that, being a QCD process, next-to-leading order (NLO) effects are expected to be important. As the tools allowing for such a calculation at the Monte Carlo level became available recently [19], we employ the full NLO machinery to study the existing bounds, as well as to establish the LHC potential at Run 2 to test the simplest phenomenologically viable monotop model.

The paper is organised as follows. In Sect. 2 we briefly recall the details of the effective Lagrangian describing resonant monotop production and of the theoretical framework allowing for numerical Monte Carlo simulations at the NLO level. In Sect. 3, we analyse the existing bounds coming from different sources. Hence, we reinterpret the results of stop pair searches, include limits from resonance searches using dijet probes and the constraints originating from direct monotop searches at the LHC Run 1. We moreover consider the constraints stemming from the modification of the top quark width and the fact that the neutral fermion χ has to be long-lived or decay invisibly. In Sect. 4, we collect all current constraints on the parameter space and discuss the LHC Run 2 potential, which will be relevant for the ongoing experimental analyses, and hence present our conclusions.

2 Theoretical framework

We consider a general class of scenarios describing the resonant production of monotop systems at hadronic colliders, in which the key model signature consists in the production of a top quark in association with missing transverse energy carried by an invisible Majorana fermion χ of mass m_χ . Imposing electroweak gauge invariance, the resulting Lagrangian is quite simple [1, 3]. The only phenomenologically viable and gauge-invariant production mechanism involves two initial-state right-handed down-type quarks that annihilate into a scalar field σ of mass m_σ lying in the triplet representation of the QCD gauge group, and carrying an electric charge $e_\sigma = 2/3$ in unit of the positron charge. The σ particle then

decays into a monotop state consisting of a right-handed top quark and a χ -particle. This mechanism can be described by a simplified Lagrangian, to be supplemented to the Standard Model one,

$$\Delta\mathcal{L} = D_\mu\sigma^\dagger D^\mu\sigma - m_\sigma^2\sigma^\dagger\sigma + \frac{i}{2}\bar{\chi}\not{D}\chi - \frac{1}{2}m_\chi\bar{\chi}\chi + \left[\lambda\sigma\bar{d}^c P_R d + y\sigma\bar{t} P_L \chi + \text{h.c.}\right], \quad (1)$$

where all indices are left understood for clarity and $P_{L,R}$ denotes the left-handed and right-handed chirality projectors. This Lagrangian includes gauge interactions for all new fields and the Yukawa couplings of the σ field to a pair of down-type antiquarks (parameterised by a 3×3 matrix λ , antisymmetric in flavour space) and to the $t\chi$ monotop state (with a strength y).

Aiming at precision predictions for monotop production at NLO in QCD, the Lagrangian in Eq. 1 must be consistently renormalised to absorb all ultraviolet divergences appearing in the virtual one-loop diagrams. Adopting the on-shell renormalisation scheme, the wave-function and mass renormalisation constants of the five massless quark fields vanish ($\delta Z_q = \delta m_q = 0$), whereas those of the massive top quark (δZ_t and δm_t) and σ field (δZ_σ and δm_σ^2) are given, at the first order in the strong coupling α_s , by

$$\delta Z_t = -\frac{\alpha_s}{3\pi} \left[\frac{3}{\bar{\epsilon}} + 4 - 3 \log \frac{m_t^2}{\mu_R^2} \right] \quad \text{and} \quad \delta m_t = -\frac{\alpha_s m_t}{3\pi} \left[\frac{3}{\bar{\epsilon}} + 4 - 3 \log \frac{m_t^2}{\mu_R^2} \right], \quad (2)$$

$$\delta Z_\sigma = 0 \quad \text{and} \quad \delta m_\sigma^2 = -\frac{\alpha_s m_\sigma^2}{3\pi} \left[\frac{3}{\bar{\epsilon}} + 7 - 3 \log \frac{m_\sigma^2}{\mu_R^2} \right]. \quad (3)$$

In our notation, we denote the regularisation/renormalisation scale by μ_R and the ultraviolet-divergent parts of the renormalisation constants are given in terms of $1/\bar{\epsilon} = 1/\epsilon - \gamma_E + \log 4\pi$, γ_E being the Euler-Mascheroni constant and the number of spacetime dimensions being $D = 4 - 2\epsilon$. The Majorana fermionic field χ is a gauge singlet, so that it does not need to be renormalised at NLO in α_s . In contrast, the wave-function renormalisation constant of the gluon field (δZ_g) reads

$$\delta Z_g = -\frac{\alpha_s}{6\pi} \left[\frac{1}{\bar{\epsilon}} - \log \frac{m_t^2}{\mu_R^2} \right] - \frac{\alpha_s}{24\pi} \left[\frac{1}{\bar{\epsilon}} - \log \frac{m_\sigma^2}{\mu_R^2} \right]. \quad (4)$$

We moreover enforce that the running of the strong coupling constant solely originates from gluons and $N_f = 5$ flavours of light quarks and renormalise it by subtracting, at zero momentum transfer, all contributions from top and σ loops. Any effect induced by these massive fields is hence decoupled and absorbed in the renormalisation of α_s ,

$$\frac{\delta\alpha_s}{\alpha_s} = \frac{\alpha_s}{2\pi\bar{\epsilon}} \left[\frac{N_f}{3} - \frac{11}{2} \right] + \frac{\alpha_s}{6\pi} \left[\frac{1}{\bar{\epsilon}} - \log \frac{m_t^2}{\mu_R^2} \right] + \frac{\alpha_s}{24\pi} \left[\frac{1}{\bar{\epsilon}} - \log \frac{m_\sigma^2}{\mu_R^2} \right]. \quad (5)$$

Finally, we choose to renormalise the λ and y parameters in the $\overline{\text{MS}}$ scheme,

$$\frac{(\delta\lambda)_{ij}}{\lambda_{ij}} = -\frac{\alpha_s}{\pi\bar{\epsilon}} \quad \text{and} \quad \frac{\delta y}{y} = -\frac{\alpha_s}{2\pi\bar{\epsilon}}. \quad (6)$$

The renormalisation group running for the new physics couplings λ_{ij} and y will be performed with the anomalous dimensions as $\beta_{\lambda_{ij}} = -\frac{\alpha_s}{\pi}$ and $\beta_y = -\frac{\alpha_s}{2\pi}$.

In order to handle $2 \rightarrow 2$ resonant processes at the NLO accuracy in QCD, we work in the complex-mass scheme [20–22] where the renormalisation procedure is handled with the complex masses and complex derived parameters. Complications may consequently arise with the choice of proper Riemann sheets when the derivation of the various renormalisation constants is at stake [23]. However, in our simplified model, there is no $\mathcal{O}(\alpha_s)$ contribution to the particle decay widths at tree level, so that such complications are avoided. To achieve the NLO QCD accuracy in the whole phase space, we nevertheless need to evaluate the unstable particle widths by including NLO QCD corrections. For simplicity, we fix in these calculations the renormalisation scale μ_R to the respective particle masses, and include the renormalisation group running of the α_s , λ_{ij} and y couplings. In the context of the width calculations, the potential impact of using a different scale is a pure next-to-NLO effect, and has therefore been ignored.

3 Resonant monotops in the LHC era

The model that has been described in Sect. 2 has six free parameters, *i.e.* two masses and four couplings,

$$\{m_\chi, m_\sigma, y, \lambda_{12} = -\lambda_{21}, \lambda_{23} = -\lambda_{32}, \lambda_{31} = -\lambda_{13}\}. \quad (7)$$

To exhaustively explore its phenomenology, in particular at the LHC, we first simplify the parameter space by assuming that only the λ_{12} parameter dominates,

$$\lambda \equiv \lambda_{12} \gg \lambda_{23}, \lambda_{31} \approx 0. \quad (8)$$

Such a choice allows us not only to avoid the flavour constraints that arise in particular from kaon mixing [11], but also to maximise the monotop production cross section at the LHC by virtue of parton density effects. The monotop signal then solely depends on the scalar mass m_σ and the relative magnitude of the λ and y parameters that control the two branching ratios

$$\begin{aligned} \text{BR}(\sigma \rightarrow t\chi) &\equiv \text{BR}_{t\chi} \quad \text{and} \\ \text{BR}(\sigma \rightarrow d\bar{s}) &\equiv \text{BR}_{jj} = 1 - \text{BR}_{t\chi}. \end{aligned} \quad (9)$$

Our exploration strategy, therefore, consists in slicing the parameter space for fixed values of the $\text{BR}_{t\chi}$ branching ratio and of the coupling λ , and in studying how the constraints evolve for increasing importance of the dijet decay channel. We therefore describe our model in terms of the four parameters

$$\{m_\chi, m_\sigma, \lambda, \text{BR}_{t\chi}\}, \quad (10)$$

where y has been traded with the $\text{BR}_{t\chi}$ branching ratio. We have kept the λ parameter free (and not the y one) as it directly controls the resonant production rate of a σ particle. As the aim of this work is to focus on monotop models, we will exclude from our investigations any parameter space region in which the σ particle cannot decay into a monotop system, *i.e.* regions for which $m_\sigma < m_\chi + m_t$.

In this section, we will consider two classes of constraints, namely those that are respectively independent and dependent on λ . The former ones allow us to directly exclude regions in the (m_χ, m_σ) parameter space at fixed $\text{BR}_{t\chi}$, whilst the latter ones allow us to derive upper limits on the coupling λ for each point of the same mass plane. The first category of constraints includes typical LHC searches for the production of a pair of coloured resonances (see Sect. 3.1), as well as searches capable of targeting the doubly-resonant production of a pair of dijet systems (see Sect. 3.2). These searches are indeed sensitive to the production and decay of a pair of σ particles. We only consider the QCD contribution to pair production, which only depends on the mass m_σ . Additional contribution may arise from the process $d\bar{d} \rightarrow \sigma\sigma^\dagger$ via a t -channel s -quark exchange. However, for the values of λ allowed by the constraints, this process always yields a negligible contribution to the cross section. On different lines, the second category of constraints includes bounds that could originate from dijet (see Sect. 3.3) and monotop (see Sect. 3.4) probes, as these two final states can be induced by the resonant production and decay of a single σ particle. In addition, for parameter space regions in which $m_\chi < m_t$, the top quark can undergo a three-body decay in a χjj final state via an off-shell σ particle. While no direct search is currently dedicated to such a decay, measurements of the top width yield indirect constraints on the model (see Sect. 3.5). The fermion χ is in principle allowed to decay into a $t^{(*)}jj$ system, with the final-state top quark being potentially off-shell, and one must ensure that χ is stable on LHC detector scales (see also in Sect. 3.5). Finally, in the same section, we comment on the usage of the narrow-width approximation for the σ particle.

In our phenomenological investigations, we rely on a numerical evaluation in four dimensions of all loop integrals, so that the numerical results must be complemented by

rational terms related to the ϵ -dimensional pieces of the integrals [24–27]. They consist of the so-called R_1 and R_2 terms, the former being universal and connected to the denominators of loop integrands and the latter being model-dependent and associated with the numerators of loop integrands. In practice, we perform loop-integral computations with the MADLOOP package [28] that automatically estimates the R_1 contributions and makes use of a finite set of special Feynman rules derived from the bare Lagrangian [29] to estimate the R_2 ones. We have computed those R_2 Feynman rules by implementing the Lagrangian of Eq. 1 into the FEYN-RULES package [30], that we have jointly used with the NLOCT [31] and FEYNARTS [32] programs to export the model under the form of a UFO module [33]. Such a module contains here, in addition to tree-level information, all ultra-violet counterterms and R_2 Feynman rules needed for numerical loop-calculations in QCD in the context of our monotop model. In practice, this has allowed us to make use of the MADGRAPH5_aMC@NLO [34] (MG5_aMC) platform for the generation of the LHC signals at the NLO accuracy in QCD.

Moreover, for the resonant processes in which the complex-mass scheme must be employed, we have verified that the NLO widths computed by MG5_aMC agree with independent in-house calculations. We now present the current bounds on the parameter space of the model.

3.1 LHC constraints on σ pair-production from supersymmetry searches

The σ particle, being charged under the QCD gauge group, can be copiously pair-produced at the LHC. Considering a decay mode in which both σ particles decay into a monotop system,

$$pp \rightarrow \sigma\sigma^\dagger \rightarrow t\chi \bar{t}\chi, \quad (11)$$

one obtains a signature that could be probed by typical stop searches in the top-antitop plus missing transverse energy ($t\bar{t} + \cancel{E}_T$) channel as well as by generic supersymmetry searches in the jets plus missing energy mode (assuming that both final-state top quarks decay hadronically). This latter class of searches is also sensitive to signals arising from a mixed decay mode in which one σ particle decays into a monotop system and the other one into a pair of jets:

$$pp \rightarrow \sigma\sigma^\dagger \rightarrow t\chi \, jj + \text{h.c.} \quad (12)$$

This final state can be the more significant one to get a handle on the model for $\text{BR}_{t\chi} \sim 50\%$, and it will also be probed by stop searches in the single-lepton channel such as in Ref. [35]. The remaining decay option in which both σ particles decay into a pair of jets will be addressed in Sect. 3.2. In the two considered cases of Eq. 11 and Eq. 12, the exclusions in the (m_χ, m_σ) planes only depend on the two masses,

Fig. 1 The red region under the curve corresponds to the excluded region at 95% confidence level by the reinterpretation of the results of the CMS-SUS-17-001 search for stops, assuming $\text{BR}_{t\chi} = 100\%$. Our findings are represented in the (m_σ, m_χ) mass plane

with the total rate being rescaled, respectively, by $\text{BR}_{t\chi}^2$ and $2\text{BR}_{t\chi}(1 - \text{BR}_{t\chi})$ factors.

As all LHC stop searches give rise to similar bounds, we reinterpret the results of a single recent search: we thus consider the CMS-SUS-17-001 analysis, which focuses on the dileptonic mode of the top-antitop system [36]. This search is based on an integrated luminosity of 35.9 fb^{-1} of LHC Run 2 proton-proton collision data at a centre-of-mass energy $\sqrt{s} = 13 \text{ TeV}$. It targets a signature made of two opposite-sign leptons (electrons or muons, with a veto on the presence of a reconstructed Z -boson), light and b -tagged jets and a significant amount of missing transverse momentum. The production of the hypothetical stop-pair signal, which in our case is identified with the production of two σ scalars decaying as in the process of Eq. 11, is then efficiently separated from the dominant top-antitop background by a cut on the transverse mass m_{T2} [37,38] reconstructed from the two leptons and the missing transverse momentum. In Fig. 1, we show the LHC bound for the process in Eq. 11, assuming $\text{BR}_{t\chi} = 100\%$ and in the plane of the two masses, (m_σ, m_χ) . We made use of the MADANALYSIS 5 platform [39,40] and its public analysis database [41,42], which contains the corresponding validated reimplementations [43] and the necessary DELPHES 3 configuration cards for handling the simulation of the response of the LHC detectors [44].

We additionally tested the limits arising from two representative ATLAS searches for dark matter and supersymmetry in the multijet plus missing energy channel [45,46]. Both these searches target a small 3.2 fb^{-1} luminosity of proton-proton collisions at $\sqrt{s} = 13 \text{ TeV}$ and consider a signature featuring one very hard jet plus subleading hadronic activity. Whilst the ATLAS-EXOT-2015-03 analysis [45] only allows for a restricted subleading jet activity, the ATLAS-SUSY-2015-06 analysis [46] includes signal regions that require a much more important jet activity. Both these analyses rely on

a large set of signal regions differing by the number of jets, their kinematical configuration and the amount of missing transverse energy. Whilst both these ATLAS analyses only consider a small integrated luminosity of 3.2 fb^{-1} of proton-proton collisions, they are already limited by the systematics. The resulting bounds are consequently not expected to improve with an increased LHC luminosity [47] and the subsequent predictions can be robustly used as the best current constraints originating from multijet plus missing transverse energy production at the LHC. By using the validated MAD-ANALYSIS 5 public implementations [48,49], we have found that they only marginally improve the exclusions at the price of a larger systematic uncertainty due to the fact that the final state in our model differs from the supersymmetric benchmarks. Thus, we will conservatively only use the exclusion from the CMS stop search in the following.

3.2 Searches for a pair of dijet resonances at the LHC

As mentioned in Sect. 3.1, the production of a pair of σ particles can yield a dijet-pair final state when they both decay into quarks,

$$pp \rightarrow \sigma\sigma^\dagger \rightarrow jjjj. \quad (13)$$

Although the relevance of this channel is reduced when monotop production is large (*i.e.*, when $\text{BR}_{t\chi}$ is large), it can lead to important constraints for large $\text{BR}_{jj} = (1 - \text{BR}_{t\chi})$ values as the resulting rate is proportional to BR_{jj}^2 .

A new physics signature featuring a pair of dijet systems originating from a pair of resonances has been searched for at the LHC both during Run 1 and Run 2, and by both the ATLAS and CMS collaborations. One of the investigated benchmark models consists of a simplified model where the SM is supplemented by a light stop decaying into two jets via an R -parity violating interaction with a branching fraction of 100%. The corresponding experimental results can thus be directly reinterpreted as a bound on the $\text{BR}_{t\chi}$ branching ratio, as the signal total rate consists of the stop-pair production cross section rescaled by a $(1 - \text{BR}_{t\chi})^2$ factor.

We include in our study the CMS-EXO-12-052 final Run 1 search [50] dedicated to events exhibiting the presence of at least four jets that are then paired using an algorithm based on their angular distribution. The discrimination from the leading multi-jet background is achieved by relying on a set of kinematical variables including a reduced mass asymmetry between the two pairs of jets. The search has implemented two signal regions. The first region is dedicated to resonance masses larger than 300 GeV and benefits from the full Run 1 dataset with an integrated luminosity of 19.4 fb^{-1} , whereas the second region is restricted to a smaller dataset of 12.4 fb^{-1} and solely considers resonance masses lying in the $[200, 300] \text{ GeV}$ mass window. It has been made possible to access such low masses thanks to a specific trigger that

Fig. 2 Lower bound on $\text{BR}_{t\chi}$ as a function of the mass of m_σ extracted from the search for a pair of dijet resonances in LHC proton-proton data by CMS (blue) at $\sqrt{s} = 8 \text{ TeV}$ (dashed) and 13 TeV (solid), and by ATLAS (solid red) for $\sqrt{s} = 13 \text{ TeV}$. The area under the curves is the excluded region. For the Run 1 CMS search, the discontinuity at $m_\sigma = 300 \text{ GeV}$ stems from the transition between two different search strategies respectively targeting low-mass and high-mass resonances [50]

has been designed especially for this purpose, with a cost in luminosity.

The same signal can also be constrained with Run 2 data. We reinterpret the results of a CMS search focusing on low-mass resonances in the $[80, 300] \text{ GeV}$ regime [51] and on a low-luminosity of data (2.7 fb^{-1}), and the results of ATLAS analysis of the 2015 and 2016 datasets (36.7 fb^{-1}) [52]. Both Run 2 analyses rely on a boosted configuration and select events featuring two back-to-back fat jets and a significant hadronic activity. After a standard cleaning of the signal by means of various grooming and pruning methods, the substructure of the fat jet is employed to improve the quality of the signal and to discriminate it from the background, together with considerations on the mass asymmetry between the two fat jet masses and on other kinematical variables.

Our results are shown in Fig. 2, where we consider the above-mentioned ATLAS and CMS searches and theoretical estimates of the stop pair-production cross section evaluated at the NLO accuracy in QCD [19,53]. The strongest bound arises from the Run 2 ATLAS search for most values of the σ mass, with the exception of two mass points for which the Run 1 CMS search does a better job and the very low mass region that benefits from the dedicated CMS Run 2 analysis. The dip at $m_\sigma = 250 \text{ GeV}$ featured in the CMS Run 1 limit is connected to a small excess of events in a single bin that has however not been confirmed by ATLAS. For fixed $\text{BR}_{t\chi}$, therefore, a lower bound on the mass m_σ can be extracted.

3.3 Dijet searches at the TeVatron and the LHC

Dijet searches constrain the single production of the σ resonance that then decays into jets. Due to the large QCD back-

ground and because of trigger requirements, dijet searches at the LHC typically target high invariant masses, so that dedicated efforts are needed to test the low mass regions that are more relevant for the monotop model under consideration. At Run 2, the most recent ATLAS search based on an integrated luminosity of 37 fb^{-1} has a minimum reach of 1.1 TeV [54], while a low-mass dedicated search [55] (29.3 fb^{-1}) overcomes the trigger limitation by recording only jet trigger information and reaches invariant masses as low as 450 GeV. Similarly, for the CMS analyses of 36 fb^{-1} of 13 TeV data [56], the high mass region starts at 1.6 TeV, whereas a trigger-based low-mass search allows for reaching down to masses of 600 GeV. A special position is reserved for a very low mass CMS search [57] with 36 fb^{-1} , where the trigger limitation is overcome by looking for boosted dijet systems. The boost allows to reduce the background and, therefore, to reach the 50–300 GeV mass range. Similarly, the ATLAS collaboration has performed a search with 15.5 fb^{-1} specifically dedicated to low-mass dijet resonances produced in association with a high- p_T photon [58], which allows for reaching invariant masses in the 200–1200 GeV range at the price of a lower signal cross section. At the TeVatron, the most up-to-date search for new particles giving rise to a dijet final state is from the CDF collaboration [59], where the signal would arise in our case from the

$$p\bar{p} \rightarrow \sigma/\sigma^\dagger \rightarrow jj + X \quad (14)$$

process, for $p\bar{p}$ collisions at $\sqrt{s} = 1.96 \text{ TeV}$. The 95% confidence level limits exclude resonance mass ranging from 260 GeV to 1400 GeV.

The cross section associated with the production of a dijet system originating from the decay of a σ resonance is proportional to $\lambda^2(1 - \text{BR}_{t\chi})$, so that for a fixed $\text{BR}_{t\chi}$ value, an upper limit on λ can be established from the current bounds. In the numerical recast of the relevant searches, we fixed $m_t = 173.3 \text{ GeV}$ and worked in the five light quark flavour scheme. In order to estimate the relevant signal fiducial cross sections and subsequently extract the bounds on the model parameters, we used the MG5_aMC [34] and PYTHIA 8.2 [60] programs to generate NLO-accurate event samples in which the fixed-order results, obtained by convoluting the NLO hard-scattering matrix elements with the NLO set of NNPDF 3.0 parton densities [61], are matched with parton showers following the MC@NLO prescription [62]. We then utilised the MADANALYSIS 5 platform [39] and its interface to FASTJET [63] to impose the same jet reconstruction (based on the anti- k_T algorithm [64]) and kinematical cuts as in the experimental analyses under consideration. In all our computations, the factorisation and renormalisation scales have been set to the σ mass m_σ .

As an example, we focus on the CDF dijet search of Ref. [59]. In Table 1 we provide the predicted NLO fiducial cross sections for a few selected m_σ values. Our results

Table 1 Fiducial cross sections for dijet production at the TeVatron, in $p\bar{p}$ collisions at $\sqrt{s} = 1.96 \text{ TeV}$, after imposing the same jet reconstruction method and signal selection as in the CDF analysis of Ref. [59]. We compare our (normalised) predictions with the CDF limits

m_σ [GeV]	$\sigma_{\sigma \rightarrow jj}^{\text{NLO,CDF}} / (\lambda^2(1 - \text{BR}_{t\chi}))$ [pb]	CDF limit [pb] [59]
260	252	110
300	132	45
400	26.1	7.2
500	5.83	3.9
620	0.960	0.8
700	0.259	0.6

Fig. 3 Upper bound on the λ parameter denoting the coupling strength of the σ resonance to down-type quarks as derived from a variety of dijet searches at colliders. The results are represented in the $(m_\sigma, \text{BR}_{t\chi})$ plane, and the regions under the curves are excluded for λ equal to at least the indicated value. In the upper left panel, we reinterpreted the low mass CMS search in the boosted regime [57], whilst on the upper right panel, we considered the CDF analysis [59]. In the lower left panel, we focused on the trigger-based ATLAS analysis [55], and combined it with the similar CMS search [56] in the lower right panel, recalling that the CMS analysis is insensitive to any m_σ value smaller than 600 GeV

can be compared with the CDF bounds, so that we can extract an upper bound on the λ parameter. Repeating the exercise for all the above-mentioned searches, we present in Fig. 3 the upper limits on λ in the $(m_\sigma, \text{BR}_{t\chi})$ plane that we divide in four kinematical configurations in which different searches dominate.

The low mass regime in which m_σ lies in the 50–300 GeV mass window has been probed by the CMS boosted search of Ref. [57]. The tight requirements of the trigger are over-

come by requiring the presence of at least one broad jet with $p_T > 500$ GeV. Then, various jet substructure techniques are employed to discriminate a signal in which the broad jet is issued from a resonance decaying into a boosted dijet system from the QCD background. The benchmark model used in the search is a Z' model. Even though our model contains a scalar resonance and not a vector one, we do not expect significant kinematical differences in the properties of the dijet system. We have therefore simply reinterpreted the search in terms of our model by a direct comparison of the predicted production cross section with the excluded one. The results are presented in the upper left panel of Fig. 3. The 300–450 GeV mass window is only covered by the CDF analysis, which implies much weaker limits on λ , as shown in the upper right panel of the figure. In the whole parameter space region, the best upper limit is $\lambda < 0.46$, *i.e.* one to two orders of magnitude weaker than any limit that could be derived from the LHC results. For σ masses larger than 450 GeV, the trigger-based low mass search from ATLAS [55] kicks in with limits much stronger than the ones derived from the CDF results, as shown in the lower left panel of the figure. Finally, for masses greater than 600 GeV, the ATLAS limits can be combined with those obtained from the trigger-based search of CMS [56], which is presented in the lower right panel of the figure. In summary, we observe that current low mass dijet searches at the LHC give bounds on the coupling λ of order 10^{-2} , except for the 300–450 GeV mass window where only much weaker CDF limits are applicable.

3.4 Monotop bounds after run 1

Monotop searches have been designed to get hints for new physics in a final state comprised of a single top quark and missing transverse energy. As sketched in Sect. 2, such a final state can originate from the (potentially resonant) production of a σ state followed by its decay into a $t\chi$ system. The first experimental search for monotops has been carried out by the CDF collaboration at the TeVatron [4] and solely focused on the flavour-changing monotop production mode. LHC Run 1 data has been analysed during the last few years, both by the CMS [6] and ATLAS [5] collaborations. Whilst the CMS analysis again focused only on flavour-changing monotop production, the ATLAS results have been interpreted both in the flavour-changing and resonant scenarios. The way in which they are presented however makes their reinterpretation in different theoretical frameworks highly non-trivial. The ATLAS analysis indeed assumes that all model coupling parameters are equal to a common value, and a bound on this value is presented in terms of the resonance mass. It is consequently impossible to make use of the results for model configurations not satisfying this requirement. The first monotop analysis of the LHC Run 2 results has also been recently released by the CMS collaboration [7], but it

targets boosted monotop systems so that it is not relevant for the mass scales probed in this work. For these reasons, we concentrate on the CMS Run 1 monotop analysis that we consider as representative for the most constraining direct LHC search on the resonant monotop model.¹

To this aim, we have implemented the CMS-B2G-12-022 search for monotop production in proton-proton collisions at a centre-of-mass energy of 8 TeV [6] in the MADANALYSIS 5 framework [39–41]. We have validated our reimplementation by making predictions for two monotop scenarios for which the CMS collaboration has provided signal events and associated official results. The first of these scenarios concerns flavour-changing monotop production (with an invisibly-decaying vector boson of 500 GeV and all model couplings fixed to 0.1), whilst the second one addresses resonant monotop production (with a scalar resonance of 1 TeV and an invisible fermion of 50 GeV, all coupling parameters being again set to 0.1). These event samples have been generated with MG5_aMC [34], using the NNPDF 2.3 set of parton densities [65] and a UFO model [33] following the monotop model of Ref. [3]. Parton showering and hadronisation have been simulated with the PYTHIA 8 package [60], and we have included the impact of the detector response by means of the DELPHES 3 programme [44] that internally relies on FASTJET [63] for object reconstruction and on an appropriate detector parameterisation describing the performance of the CMS detector during the LHC Run 1. The validity of our recasting code has been inferred from a comparison between the MADANALYSIS 5 and CMS official results that have been produced from the hard-scattering events that we have provided to CMS. Event files and generator configuration files can be obtained from the public analysis database of MADANALYSIS 5 [66], whilst the recasting C++ code has been additionally submitted to INSPIRE [67].

The CMS-B2G-12-022 analysis relies on a selection that vetoes the presence of isolated electrons (muons) with a transverse momentum p_T larger than 10 GeV (20 GeV) and a pseudorapidity $|\eta| < 2.4$ (2.5), where lepton isolation is imposed by constraining the sum of the transverse momenta of all objects lying in a cone of radius $R = 0.4$ centred on the lepton to be smaller than 20% of the lepton p_T . The analysis next requires the presence of at most three jets with transverse momentum $p_T(j_1) > 60$ GeV, $p_T(j_2) > 60$ GeV and $p_T(j_3) > 40$ GeV respectively, and it additionally forbids the presence of a fourth jet with a p_T greater than 35 GeV. The invariant mass of the three leading jets M_{3j} is then imposed to be compatible with the top mass $M_{3j} < 250$ GeV and the missing energy has to satisfy $\cancel{E}_T > 350$ GeV. Two signal regions S1 and S0 are finally defined, the difference lying in

¹ Due to the lack of publicly available material, we were unable to validate a recasting of the corresponding 8 TeV ATLAS analysis.

Table 2 Comparison of results obtained with our MADANALYSIS 5 reimplementation (MA5) to those provided by the CMS collaboration (CMS-B2G-12-022) in the case of a new physics scenario featuring flavor-changing monotop production. The selection and total

efficiencies are defined in Eq. 15. The official CMS numbers have been derived from the same hard-scattering events entering our simulation chain. Those events have been provided to the CMS collaboration who accepted to produce an official cutflow

	Selection step	CMS	ϵ_i^{CMS}	$\epsilon_{i,\text{tot}}^{\text{CMS}}$	MA5	ϵ_i^{MA5}	$\epsilon_{i,\text{tot}}^{\text{MA5}}$	δ_i^{rel}
0	Nominal	3000			3000			
1	Lepton veto	3000	1.000	1.000	2983	0.994	0.994	0.56%
2	$p_T(j_1) > 60$ GeV	2805	0.935	0.935	2799	0.938	0.933	0.35%
3	$p_T(j_2) > 60$ GeV	1719	0.613	0.573	1900	0.679	0.633	10.8%
4	$p_T(j_3) > 40$ GeV	1116	0.649	0.372	1358	0.715	0.453	10.1%
5	Veto on the fourth jet	598	0.536	0.200	618	0.455	0.206	15.1%
6	$M_{3j} < 250$ GeV	294	0.492	0.098	269	0.435	0.090	11.5%
7	$\cancel{E}_T > 250$ GeV	98	0.333	0.032	109	0.405	0.036	21.6%
8	$\cancel{E}_T > 350$ GeV	27	0.276	0.009	36	0.330	0.012	19.9%
S0	0b-jet	6	0.222	0.002	12	0.333	0.004	50.0%
S1	1b-jet	19	0.704	0.006	23	0.639	0.008	9.2%

Table 3 Same as in Table 2, but for a new physics scenario featuring resonant monotop production

	Selection step	CMS	ϵ_i^{CMS}	$\epsilon_{i,\text{tot}}^{\text{CMS}}$	MA5	ϵ_i^{MA5}	$\epsilon_{i,\text{tot}}^{\text{MA5}}$	δ_i^{rel}
0	Nominal	4000			4000			
1	Lepton veto	4000	1.000	1.000	3989	0.997	0.997	0.28%
2	$p_T(j_1) > 60$ GeV	3932	0.983	1.000	3947	0.989	0.986	0.66%
3	$p_T(j_2) > 60$ GeV	2872	0.730	0.983	3044	0.771	0.761	5.59%
4	$p_T(j_3) > 40$ GeV	1620	0.564	0.718	1944	0.639	0.486	13.2%
5	Veto on the fourth jet	996	0.614	0.405	1006	0.517	0.252	15.8%
6	$M_{3j} < 250$ GeV	536	0.538	0.249	479	0.476	0.120	11.5%
7	$\cancel{E}_T > 250$ GeV	463	0.863	0.134	415	0.866	0.104	0.30%
8	$\cancel{E}_T > 350$ GeV	315	0.680	0.116	284	0.684	0.071	0.59%
S0	0b-jet	104	0.330	0.023	90	0.317	0.023	4.02%
S1	1b-jet	189	0.600	0.040	159	0.560	0.040	6.69%

the presence of either exactly one or exactly zero b -tagged jet.

In Tables 2 and 3, we confront the cutflow charts that have been obtained with MADANALYSIS 5 to the official results of CMS for the two benchmark scenarios under consideration. For each step of the selection, we have calculated the relative (ϵ_i) and cumulative (ϵ_{tot}) efficiencies, as well as the difference δ_i^{rel} between the CMS official and MADANALYSIS 5 relative efficiencies, normalised to the CMS result,

$$\epsilon_i = \frac{n_i}{n_{i-1}}, \quad \epsilon_{i,\text{tot}} = \frac{n_i}{n_0} \quad \text{and} \quad \delta_i^{\text{rel}} = \left| 1 - \frac{\epsilon_i^{\text{MA5}}}{\epsilon_i^{\text{CMS}}} \right|, \quad (15)$$

where n_i and n_{i-1} are the numbers of events after and before the considered selection step, respectively. We have found a very good agreement for the resonant monotop benchmark (see Table 3) reaching a level of 10%–15% for all selection steps. The situation is similar for the flavour-changing sce-

nario case (see Table 2), although a 50% difference between the CMS and the MADANALYSIS 5 results is observed for the final S0 selection. This bin is however populated by a very small number of events so that statistical effects are likely to be important.

We finally also compare the differential distribution in the invariant mass of the three jets M_{3j} after applying all selections but this one for the benchmark scenario presented in the CMS analysis note. The latter consists of a flavour-changing monotop production scenario where an invisible vector state of mass equal to 700 GeV is produced in association with a top quark. The results are shown on Fig. 4 and exhibit once again a good agreement. We therefore consider our reimplementation to be validated.

In order to extract the constraints on the model parameter space, we generated the $pp \rightarrow \sigma^{(*)} \rightarrow t\chi$ monotop signal for various mass configurations and extracted, with MADANALYSIS 5, the number of events N_s populating the two signal regions S0 and S1. For each signal region, we derived

Fig. 4 Distribution of the invariant mass of the three leading jets once all monotop selection steps have been considered, but the one on M_{3j}

an exclusion confidence level by generating 10^5 Monte Carlo toy experiments in which the actual number of background events N_b is computed from a Gaussian distribution of mean $\hat{N}_b$ and width $\Delta\hat{N}_b$, the Standard Model background expectation $\hat{N}_b \pm \Delta\hat{N}_b$ being taken from Ref. [6]. This allowed us to calculate the p -values associated with the background-only (p_B) and signal-plus-background (p_{S+B}) hypotheses. These estimations assume that the number of background events N_b and signal-plus-background events $N_b + N_s$ are Poisson-distributed, and that N_{data} events have been observed (the data numbers being reported in Ref. [6]). We next freely varied the signal production cross section to the smallest value for which

$$1 - \frac{p_{S+B}}{p_B} > 0.95, \quad (16)$$

which corresponds to the cross section σ_{95} excluded at the 95% confidence level. The final exclusion has been obtained by comparing the most stringent constraints originating from the two signal regions with our predicted NLO signal cross section for any given point of the parameter space. Our results are represented, in the (m_σ, m_χ) plane, in Fig. 5 for two specific λ values corresponding to the typical order of magnitude probed by the other LHC analyses potentially constraining the model. Heavy mediator masses are usually excluded, provided that the spectrum is not too compressed. Parameter space configurations in which the mediator is light cannot however be reached with current data. Whilst a stronger exclusion could be derived by combining the two regions, we conservatively ignore this feature as the combination would require correlation information not shared by CMS.

3.5 Width constraints

As we have seen, the resonantly produced scalar σ decays to a monotop system and to a pair of jets more or less impor-

Fig. 5 Excluded regions at the 95% confidence level after reinterpreting the results of the CMS-B2G-12-022 analysis in the considered resonant monotop model. The results are presented in the (m_σ, m_χ) plane, assuming $\text{BR}_{t\chi} = 100\%$. The red area corresponds to the excluded parameter space region for $\lambda = 0.02$, while the green area (plus the red one) corresponds to the excluded region for $\lambda = 0.05$. There is no bound for $\lambda = 0.01$ or smaller. Finally, the blue region (the triangle area on the bottom right) is not considered as the model configuration does not feature any σ decay into a monotop system $\sigma \rightarrow \chi t$

tantly via the relative magnitude of the couplings λ and y respectively. In addition, the y coupling also determines the size of the monotop production cross section. Maximising both the production rate and the branching ratio into a monotop system may thus lead to some tensions in the values of the couplings. The σ branching fraction into a $t\chi$ system, in a simplified limit, only depends on the value of the λ and y couplings. Whilst the full result for the leading order (LO) partial width, $\Gamma(\sigma \rightarrow t\chi)$, is given by

$$\Gamma(\sigma \rightarrow t\chi) = \frac{y^2}{16\pi m_\sigma^3} \left(m_\sigma^2 - (m_\chi + m_t)^2 \right) \times \sqrt{m_\sigma^4 + (m_\chi^2 - m_t^2)^2 - 2m_\sigma^2(m_\chi^2 + m_t^2)}, \quad (17)$$

it simplifies to the simple approximate expression

$$\Gamma(\sigma \rightarrow t\chi) \simeq \frac{m_\sigma}{16\pi} y^2. \quad (18)$$

when the χ and top masses can be neglected. Moreover, in the limit of massless light quarks, the $\Gamma(\sigma \rightarrow d\bar{s})$ partial width reads

$$\Gamma(\sigma \rightarrow d\bar{s}) = \frac{m_\sigma}{2\pi} \lambda^2. \quad (19)$$

Combining these two simplified expressions allows to write the $\sigma \rightarrow t\chi$ branching fraction as

$$\text{BR}(\sigma \rightarrow t\chi) \simeq \frac{y^2}{y^2 + 8\lambda^2}, \quad (20)$$

where it solely depends on the couplings. These formulas allow to see that, in a rough approximation, fixing λ and requiring a given $\text{BR}(\sigma \rightarrow t\chi)$ branching fraction automatically determines the y coupling. For example if we require $\lambda = 0.1$ and $\text{BR}(\sigma \rightarrow t\chi) = 90\%$ (in order to limit the σ decay to dijets, which can give rise to strong bounds on the model, and keep the number of monotop events that are expected at the LHC high), a value of $y \simeq 0.85$ is required. This back-of-the-envelope calculation is not used in the numerical evaluations performed in this work, where we used exact NLO calculations. Nevertheless it allows to qualitatively assess the coupling values that are required and their inter-relations.

The fact that increasing λ while keeping the monotop rate fixed forces us to increase y may lead to parameter space where the total width of σ becomes large. This is however an issue for the reinterpretation of the experimental searches, as we relied on the narrow width approximation (NWA) for the simulation of the resonant signal. This therefore imposes an upper bound on the value of the couplings which may contrast with the requirement of a large monotop production rate. To quantitatively assess where the NWA breaks down, it is convenient to fix the branching ratio in the monotop channel, and study the upper bound on the other coupling λ (responsible for the production rate). The total width of the σ resonance, Γ_{tot} , can thus be written as

$$\Gamma_{\text{tot}} = \frac{\Gamma(\sigma \rightarrow \bar{d}s)}{(1 - \text{BR}_{t\chi})}. \quad (21)$$

The ratio $\Gamma_{\text{tot}}/m_\sigma$ depends on the coupling and the $\text{BR}_{t\chi}$ parameter, so that we can extract an upper bound on λ by imposing the NWA validity. The limit is mass-independent at LO but a slight mass dependence (in m_σ) is present at NLO. For example varying the σ mass between 320 and 1000 GeV induces a variation in the allowed maximal value for λ smaller than 3%, while the difference between the LO and NLO limit generally always lies in the 1–5% range.

The results are shown in Fig. 6, where we show contour lines for $\Gamma_{\text{tot}}/m_\sigma = 10\%$, 20% and 50%. For widths larger than 50% of the mass, the σ can hardly be thought of as a resonance. Also, we see how the larger BR in monotop the smaller λ needs to be, thus suppressing the production rates. For any fixed $\text{BR}_{t\chi}$, we can in this way extract an upper bound for λ beyond which the width of the σ resonance is too large and needs to be fully taken into account in the searches.

Another necessary requirement in order to have the monotop signal as the main feature of the considered new physics model is to make sure that the χ state produced in association

Fig. 6 Upper bound on λ as a function of $\text{BR}_{t\chi}$ to obtain $\Gamma_{\text{tot}}/m_\sigma$ ratios smaller than 10% (blue), 20% (orange) and 50% (green)

with the top quark is long-lived enough to escape the detector. Alternatively, one may always assume that χ dominantly decays into an invisible sector. However, as a fermion, it can only decay into an invisible fermion plus an invisible boson, or into three invisible fermions, so that the dark sector has to be rather involved and one needs to make sure that the couplings to the dark sector are not too large. Here we focus on the minimal setup in which the only allowed decays of χ derive from the two couplings included in the Lagrangian of Eq. 1. We distinguish three kinematic regimes:

1. for $m_\chi > m_{\text{top}}$, the three-body decay $\chi \rightarrow tds$ is open, thus potentially giving very strong bounds on the couplings;
2. for $m_W < m_\chi < m_{\text{top}}$, the decay is four-body, $\chi \rightarrow W^+ bds$ and takes place via a virtual top quark;
3. for $m_b < m_\chi < m_W$, only a five-body decay is allowed via both an off-shell top quark and W -boson.

The issue of the width of the χ fermion has been studied in Ref. [11], where it has been showed that for masses of a few GeV below the W -boson mass m_W , decay lengths in the metre range can be obtained. For all three kinematic regimes defined above, the decay proceeds through a virtual σ exchange, and the width is proportional to the product of the two couplings $(\lambda y)^2$. By requiring that the decay length of χ is larger than the typical scale of an LHC detector, *i.e.* 10 metres, we can obtain an upper bound on λy , as shown in Fig. 7.

Moreover, for $m_\chi < m_{\text{top}}$, the same σ -exchange induces a three-body decay for the top quark, $t \rightarrow \chi \bar{d}s$. The corresponding partial width, which is also proportional to $(\lambda y)^2$, can be constrained by the direct measurement of the top width by the CDF collaboration at the TeVatron [68],

$$\Gamma_t < 6.38 \text{ GeV at the 95\% confidence level.} \quad (22)$$

Fig. 7 Upper bound on λy , represented in the (m_σ, m_χ) plane derived from the width of the χ fermion and the top quark. The black contours correspond to a minimum decay length for the invisible χ of 10 metres, while the red ones correspond to a new physics contribution to the top width of about 5 GeV

Assuming that the Standard Model width is unaffected ($\Gamma_t^{\text{SM}} = 1.41$ GeV), this leads to a bound on the $\Gamma(t \rightarrow \chi \bar{d} \bar{s})$ partial width,

$$\Gamma(t \rightarrow \chi \bar{d} \bar{s}) \lesssim 5 \text{ GeV}. \quad (23)$$

Whilst more precise determinations of the top width exist, as for instance in Ref. [69], these indirect measurements all assume the absence of non-standard decay channels and cannot thus be used here. The above bound is represented by red contours in Fig. 7. Notably, it is complementary to the χ width bounds, as it is more sensitive to the low m_χ region. Nevertheless it does not impose very strong limits on the couplings.

4 Discussion and conclusions

At 13 TeV one can have an idea of the potential of the LHC to observe a monotop signal by considering the production cross section in parameter space regions still allowed by data. A detailed study is not possible at present as this would require the generation of the corresponding Standard Model background, and even this would be only a rough analysis, considered that the 13 TeV environment and background can only be accurately determined using real data. Nevertheless, the signal cross section ranges from a few picobarns for $m_\sigma = 300$ GeV to the femtobarn level for resonance masses lying around $m_\sigma = 2$ TeV, as illustrated in Fig. 8 for

Fig. 8 NLO monotop production cross sections at the LHC operating at a centre-of-mass energy $\sqrt{s} = 13$ TeV for different m_σ values, assuming $m_\chi = 50$ GeV, $\lambda = 0.01$, $\text{BR}_{t\chi} = 90\%$ and $y \approx 0.0977$. Varying the mass of m_χ from a few GeV to 100 GeV barely affects the value of $\text{BR}_{t\chi}$, which is a multiplicative factor entering the monotop production cross section

a specific set of new physics couplings and masses. Those large numbers could in principle motivate the experimental collaborations to attempt a monotop search aiming to discover (or bound in a less optimistic case) the corresponding signal at the LHC Run 2.

We have provided in this work elements allowing to assess this question stronger. We have collected bounds of different origins that range from specific searches at colliders that we have recasted to reinterpret their results in the context of the model under consideration, to limitations implied by the presence of the monotop signal as a key new physics model feature. In the rest of this section, all these bounds will be put together in order to determine which parts of the parameter space are still open and if a monotop signal visible at Run 2 could be expected. On top of the “vanilla” monotop model in which the $\text{BR}_{t\chi}$ parameter is close to 1, we have also considered deviations in which monotop production should be substantial, with the value of $\text{BR}_{t\chi}$ being lowered to 90%, 75%, 50% and 10%. However, in this more general case, a global and detailed analysis including dijet resonance searches, new physics searches in the top quark with missing energy and jets channel as well as in the multijet mode need to be considered more deeply.

In Figs. 9 and 10, we present the bounds on the parameter space in the (m_χ, m_σ) plane for four choices of the σ decay rates into a monotop system, $\text{BR}_{t\chi} = 99\%$, 75%, 50% and 10%. The results include constraints originating from all the searches and features discussed in Sect. 3, namely stop searches, dijet searches, monotop searches and the decay

Fig. 9 Current constraints on the parameter space of the considered model featuring resonant monoton production at colliders. The results are presented in (m_χ, m_σ) planes for various λ and BR_{t_χ} 99% and 75%. Whilst the white region are allowed, the coloured areas are excluded by at least one of the bounds described in Sect. 3. We refer to the text for more details

length of the χ fermion that must be larger than 10 metres. Moreover, we restrict the chosen values of the λ parameter so that the NWA for σ is valid. Imposing $\Gamma_{tot}/m_\sigma < 20\%$, we obtain from Fig. 6,

$$\lambda < 0.35, 0.56, 0.79 \text{ and } 1.06 \quad \text{for} \\ BR_{t_\chi} = 99\%, 75\%, 50\% \text{ and } 10\% \text{ respectively.} \quad (24)$$

Fig. 10 Current constraints on the parameter space of the considered model featuring resonant monoton production at colliders. The results are presented in (m_χ, m_σ) planes for various λ and BR_{t_χ} 50% and 10%. Whilst the white region are allowed, the coloured areas are excluded by at least one of the bounds described in Sect. 3. We refer to the text for more details

The figures include both λ -independent and λ -dependent bounds, so that we indicate the chosen λ value in the upper caption of each subfigure. We moreover recall that, for each of the considered constraint, theoretical predictions have been achieved at the NLO accuracy in QCD thanks to recent developments at the level of the Monte Carlo simulations. This could serve as a starting ground for more

detailed analyses to be performed by the LHC experiments at Run 2.

The blue regions (triangle areas on the bottom right of each subfigure) correspond to parameter space configurations in which a monoton signal cannot be produced resonantly, the $\sigma \rightarrow \chi t$ decay being kinematically closed. We therefore omit it from our analysis. We also represent by rectangular blue areas (bottom left of the two lower subfigures) the regions that are excluded by resonance search in the dijet-pair channel and that we have studied in Fig. 2. As already found in Sect. 3.2, the most powerful searches concern CMS and ATLAS analyses of 13 TeV LHC data, and they only have some constraining power for low σ masses and a large branching ratio associated with the $\sigma \rightarrow jj$ decay (*i.e.* a not too large $\text{BR}_{t\chi}$ value). The light violet regions are excluded by stop searches, as detailed in Sect. 3.1. Typically, the limits presented in Fig. 1 are rescaled down proportionally to the decreasing value of $\text{BR}_{t\chi}$ (that lowers the signal production rate). All figures then feature black horizontal lines that delimitate the green bands of the mass planes that are excluded by dijet searches (see Sect. 3.3). For m_σ in the 300–450 GeV mass window, dijet bounds are weak as this corresponds to a mass configuration only probed by TeVatron searches. In contrast, LHC searches are sensitive for other σ masses and are especially stronger for masses lying in the 200–300 GeV and 450–1000 GeV ranges. Whilst the white areas are in principle reachable by dijet searches, the chosen λ values are too small to yield any constraint. In the upper left panel of the figure (for which $\text{BR}_{t\chi} = 99\%$), the monoton bounds presented in Fig. 5 are overlaid, so that the upper part of the mass plane is excluded (dark red region). Those searches however quickly lose sensitivity with smaller $\text{BR}_{t\chi}$ values in association with a λ value also five times smaller. Finally, this panel also exhibits a trapezoid pink area that corresponds to a region in which the decay length of the χ fermion is smaller than 10 metres, so that there is actually no monoton signal in there. There is no bound for any of the three other (λ , $\text{BR}_{t\chi}$) configurations, as the smaller λ value ensures a larger decay length.

To summarise, current limits severely restrict the model parameter space for light new physics states. Only small specific subregions are still allowed by data, and it is clear that future results will allow one to draw conclusive statements. In contrast, the heavier cases are still viable.

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