PROCEEDINGS OF SCIENCE

Searching for dark radiation at the LHC

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There is an interesting connection between early universe cosmology and searches for long-lived particles (LLPs) at the LHC. Light particles can be produced via freeze-in and act as dark radiation, contributing to the effective number of relativistic species N_{eff}. The parameter space of interest for future CMB missions points to LLP decay lengths in the mm to cm range. These decay lengths lie at the boundary between prompt and displaced signatures at the LHC and can be comprehensively explored only by combining searches for both. We consider a model where the LLP decays into a charged lepton and a (nearly) massless invisible particle. By reinterpreting searches for promptly decaying sleptons and for displaced leptons at both ATLAS and CMS we can then directly compare LHC exclusions with cosmological observables. Our results show how in this model the target value of CMB-S4 is already excluded by current LHC searches.

ArXiv ePrint: 2204.01759

41st International Conference on High Energy physics - ICHEP2022
6-13 July, 2022
Bologna, Italy

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PoS(ICHEP2022)105

1. Introduction

In recent years particle colliders have emerged as one of the primary tools used to look for new physics beyond the Standard Model (SM). The LHC in particular, given its high center of mass energy and luminosity, has the potential to probe new physics at the electroweak scale. However, there has been no detection of particles beyond the SM so far and this has led to an increased interest in unusual LHC signatures. Displaced signatures coming from long-lived particles (LLPs) have consequently become a very popular object of study at the LHC. LLPs are particles whose proper decay length is comparable with the LHC geometric length and hence are in the millimetre to metre range.

LLPs are not only interesting from an experimental point of view, but they are justified theoretically by looking at their connection with early universe cosmology. We know that to understand whether an interaction is efficient in the early universe, we need to compare its rate with the Hubble expansion rate. It is very intriguing to notice that the Hubble rate around electroweak temperatures ($T_{\rm EW} \sim 100\,$GeV) would correspond to decay lengths at approximately the cm scale. This essentially means that decays whose proper lifetime is macroscopic are effective in the early universe.

It would then be natural to assume these interactions were the ones responsible for the production of dark matter (DM) in the early universe. Unfortunately, that cannot be the case as these decays would lead to a large overabundance of dark matter. In our work [1] we investigate the possibility that these interactions are indeed present, but they produce an almost massless relic, rather than a massive one as in the case of DM. What is being produced in this model is not dark matter, but rather dark radiation. The production of new relativistic particles in the early universe has consequences on big bang nucleosynthesis (BBN) and CMB formation, so we need to take into consideration cosmological observations. The future mission CMB-S4 will be able to constrain the cosmological parameter of interest $\Delta N_{\rm eff}$ with a target sensitivity of $\sigma(\Delta N_{\rm eff}) = 0.03$. In our work we illustrate how the parameter space accessible by future CMB missions can be probed by a combination of prompt and displaced signatures at LHC.

2. Dark radiation model and freeze-in production

The minimal setup for this scenario requires the presence of a decaying bath particle B whose mass is at the electroweak scale and an (almost) massless daughter particle χ. For the sake of clarity, we will focus on a specific model identified by the spin nature and quantum charges of the new particles, but most of the discussion will generalize to other models. In our work the bath particle $B = (B_e, B_\mu, B_\tau)^{\rm T}$ is a scalar with a flavour structure and hypercharge $Y_B = -1$: it is analogous to the right-handed sleptons of SUSY. The massless relic χ is instead a Majorana fermion, singlet under the SM gauge group. Assuming a $\mathbb{Z}_2$ symmetry under which the new particles are odd, the renormalizable interaction with the SM is given by

$$\mathcal{L}_{\rm int} = B^{\rm T} \cdot y_\ell \cdot (\bar{\ell}_{\rm R}\, \chi) \; + {\rm h.c.} \tag{1}$$

In this contribution we will rely on some simplifying assumptions. First of all, we consider a flavour universal scenario so that the coupling matrix is $y_\ell = {\rm diag}(y, y, y)$. Then, we will consider the infrared freeze-in production of dark radiation through decay, which requires that $y \ll 1$ in the parameter space of interest. If the coupling y was too large, we would have the simple scenario of χ thermalized with the SM in the early universe. Finally, we assume that the new light particle χ stays completely relativistic through both BBN and CMB formation, so $m_\chi \ll 1\text{eV}$. The specific value of the mass will have no consequences on the phenomenology of our model, provided it is small enough.

The cosmological observable of interest for this model is the shift in the effective number of relativistic species ΔN_{eff}. This is defined as the ratio between the radiation energy density of our new particle ρ_χ and the energy density of a massless neutrino. It is then convenient to introduce the comoving energy density Z_χ in analogy with the well known comoving number density used for DM phenomenology:

$$Z_\chi(x) \equiv \frac{\rho_\chi(x)}{s^{4/3}(x)}, \tag{2}$$

where $s(x)$ is the entropy at an arbitrary time $x \equiv m_B/T$. Just like the comoving number density for DM calculations, this quantity will stay constant after the particle χ has decoupled. Deriving ΔN_{eff} then becomes just a matter of tracking the comoving energy density Z_χ as they are related by

$$\Delta N_{\text{eff}}(x) = \frac{Z_\chi(x)\, s_0^{4/3}}{\frac{7}{8}\left(\frac{4}{11}\right)^{4/3} \rho_{\gamma,0}}\ , \tag{3}$$

where $\rho_{\gamma,0}$ and s_0 are the photon energy density and entropy today, respectively.

We can proceed in analogy with DM calculations and calculate $Z_\chi(x)$ by solving a Boltzmann equation. The relevant (integrated) Boltzmann equation is given by

$$x\tilde{H}(x)s^{4/3}(x)\frac{\mathrm{d}Z_\chi(x)}{\mathrm{d}x} = \frac{m_B\,\Gamma_B}{8\pi^4} I(x, T_\chi, \text{spins}) \tag{4}$$

with $\tilde{H}(x) = H(x)(1 - \frac{1}{3}\frac{\mathrm{d}\log g^*(x)}{\mathrm{d}\log x})$ and $\Gamma_B = \frac{y^2 m_B}{16\pi}$. Here $I(x, T_\chi, \text{spins})$ indicates an integral which in the most general case depends on the SM temperature, the dark sector temperature T_χ and the spin nature of the bath and daughter particles. The complete expression for this integral is derived by the integration of the collision operator in the Boltzmann equation and reads

$$I = \frac{m_B\,\Gamma_B}{8\pi^4}\int \frac{d^3 p_\chi}{2E_\chi}\int \frac{d^3 p_\ell}{2E_\ell}\frac{2E_\chi}{E_B}\left(1 - \frac{f_\chi(E_\chi)}{f_\chi^{\text{eq}}(E_\chi)}\right) f_B^{\text{eq}}(E_B)\delta(E_B - E_\chi - E_\ell), \tag{5}$$

where E_i, p_i and f_i are the energy, momentum and phase space distribution of particle i. For DM calculations this integral would simplify considerably by relying on two well justified assumptions. The first one is that the particle B decays while completely non-relativistic. This is a common approximation which does not introduce qualitative errors. The second assumption is that the dark matter abundance is small enough that the backreaction effect can be completely neglected. This second approximation is well justified for the parameter space of interest of DM, but not in our case. So we will not employ these approximations and use the fully relativistic treatment already developed in Refs. [2, 3] for DM calculations. This integral can then be tabulated[1] and the

[1]The integral tables are provided in Ref. [1].

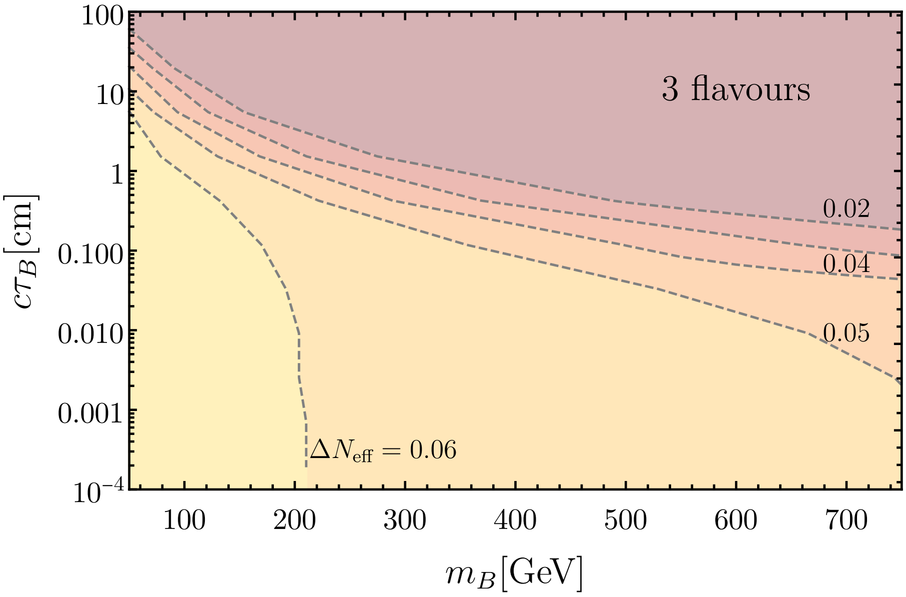

Figure 1: Shift in the effective number of relativistic species as a function of the mass and proper decay length of B. Shorter lifetimes correspond to larger couplings and for short enough lifetimes we reach the equilibrium densities and the upper bound on ΔN_{eff}.

Boltzmann equation can be solved. Solving the Boltzmann equation gives us the cosmologically interesting parameter space shown in Fig. 1.

This figure illustrates several points of interest. First of all, the target sensitivity of CMB-S4 is $\sigma(\Delta N_{\text{eff}}) = 0.03$, which corresponds to a 95% exclusion bound of $\Delta N_{\text{eff}} \leq 0.06$. We have highlighted in the plot several reference values of ΔN_{eff} which can be probed by future CMB missions. Then we can notice how for large couplings (small lifetimes) there is a saturation due to strong backreaction effects, which is also the reason why we cannot rely on the usual assumption that the backreaction is negligible. Finally, the parameter space of interest encompasses decay lengths from 10^{-4} cm up to a metre. This means that we will need to rely on both prompt and displaced LHC signatures to properly constrain the parameter space.

3. LHC constraints

A crucial part of our work is deriving the constraints on the model coming from collider experiments. We will find these by recasting the analysis of prompt and displaced searches at the LHC. It will also be important to consider the LEP constraints, which are the most relevant in the case of short lifetimes and low masses. These are derived from a combination of different LEP experiments and are summarized in Ref. [4].

3.1 Prompt signatures

The optimal prompt searches for our model are already implemented at the LHC, since our model can be probed by SUSY slepton searches [5, 6]. To recast the analysis we simulate our events and apply the event cuts. The most relevant cuts for what concerns us are the ones on the transverse

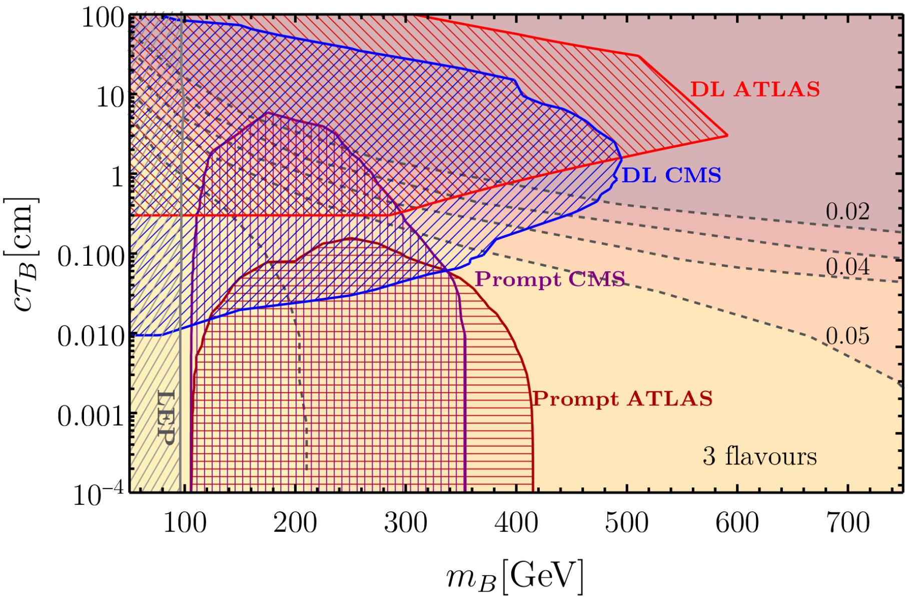

Figure 2: Collider constraints on the parameter space of interest for the 3-flavour case.

and longitudinal impact parameters. The reason for this is that the cut efficiency strongly depends on the proper lifetime τ_B. We expect the search to impose weaker constraints on particles with a longer lifetime τ_B. It is worth mentioning that ATLAS and CMS impose different cuts on the impact parameters, leading to a different sensitivity to the particle lifetime.

3.2 LLP signatures

In the case of LLP signatures we can rely on searches for displaced leptons [7, 8]. Deriving the constraints for our model is straightforward since the ATLAS and CMS collaborations provide limits on the production cross section for right-handed sleptons. These limits are given as a function of the mass and lifetime of the parent particle and directly apply to our model. More subtleties in the recasting would arise if we had different couplings to different flavours. Similarly to the case of prompt searches, also here the cuts imposed on the impact parameters by ATLAS and CMS are different. As a consequence ATLAS will put stronger constraints on longer lifetimes with respect to CMS.

4. Conclusions

We can now combine the cosmologically interesting parameter space with the collider constraints and obtain Fig. 2. Here the complementarity between prompt and displaced signatures is made clear. This is particularly true when we consider that the 95% target value $\Delta N_{\rm eff} = 0.06$ can only be excluded by applying a combination of both. For what concerns smaller values of $\Delta N_{\rm eff}$ within reach of future CMB experiments, they are more sensitive to LLP searches rather than prompt ones. Given that these searches have low backgrounds, we expect the constraints to improve considerably with higher luminosities.

Overall, our work shows how the limits on the model will improve when more data become available and the complementarity between prompt and displaced signatures. This complementarity highlights the importance of a correct reinterpretation of prompt searches in the case of macroscopic decay lengths.

Acknowledgments

We thank Torsten Bringmann and Jan Heisig for helpful discussions. FK and AM acknowledge funding from the Deutsche Forschungsgemeinschaft (DFG) through the Collaborative Research Center TRR 257 "Particle Physics Phenomenology after the Higgs Discovery" under Grant 396021762 – TRR 257 and the Emmy Noether Grant No. KA 4662/1-1. ML is funded by the Alexander von Humboldt Foundation. This manuscript has been authored by Fermi Research Alliance, LLC under Contract No. DE-AC02-07CH11359 with the U.S. Department of Energy, Office of Science, Office of High Energy Physics.

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