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Baryogenesis and gravitational waves from runaway bubble collisions

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Abstract. We propose a novel mechanism for production of baryonic asymmetry in the early Universe. The mechanism takes advantage of the strong first order phase transition that produces runaway bubbles in the hidden sector that propagate almost without friction with ultra-relativistic velocities. Collisions of such bubbles can non-thermally produce heavy particles that further decay out-of-equilibrium into the SM and produce the observed baryonic asymmetry. This process can proceed at the very low temperatures, providing a new mechanism of post-sphaleron baryogenesis. In this paper we present a fully calculable model which produces the baryonic asymmetry along these lines as well as evades all the existing cosmological constraints. We emphasize that the Gravitational Waves signal from the first order phase transition is completely generic and can potentially be detected by the future eLISA interferometer. We also discuss other potential signals, which are more model dependent, and point out the unresolved theoretical questions related to our proposal.

Keywords: baryon asymmetry, cosmological phase transitions, particle physics - cosmology connection, gravitational waves / sources

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Contents

1 Introduction

Theories of baryogenesis (for a review see [1, 2]) aim at explaining the tiny difference between the number density of baryons and antibaryons, about 10^{-10} (in units of the entropy density) we observe in our universe. Until now, many mechanisms for the generation of the baryon asymmetry have been proposed. Grand Unified Theories unify the strong and the electroweak interactions and predict both baryon and lepton number violation at the tree level. They may be considered perfect candidates for a theory of baryogenesis. The out-of-equilibrium decay of superheavy particles can explain the observed baryon asymmetry, even though there remain problems strictly related to the dynamics of reheating after inflation and to the fact that the baryon asymmetry generated by the decays may be erased by the nonperturbative sphaleron configurations present already in the Standard Model (SM).

This latter problem is in fact made a virtue in the leptogenesis mechanism where the cosmic baryon asymmetry originates from an initial lepton asymmetry generated in the decays of heavy sterile neutrinos in the early universe and then converted into baryon asymmetry by the sphalerons themselves (for a review see [3]).

If one wishes to resort to much lower energy scales, one should consider the electroweak baryogenesis (for reviews see [4, 5]), where the baryon number violation takes place during first-order phase transitions happening at temperatures of the order of the electroweak scale and the baryon number is preserved against the SM sphalerons if the first-order transition is sufficiently strong. Of course in the SM as it stands with $m_h \approx 125\,\text{GeV}$ there is no electroweak phase transition (EWPT) and instead one finds a smooth crossover [6–9]. Therefore in order to have a viable electroweak (EW) baryogenesis, one should necessarily augment the SM with new degrees of freedom that strongly couple to the Higgs. Moreover, new CPV sources at the EW scale are needed. These features make electroweak baryogenesis scenarios attractive because they could be tested or ruled out at the LHC and future colliders. However, current LHC results and, in particular, the Higgs precision data already corner the parameter space of the EW baryogenesis [10, 11] and even more gains are expected at HL

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LHC and in future colliders FCC-ee and FCC-pp (see [12] for a recent review of the FCC-pp potential reach).

Can one generate a baryon asymmetry at temperatures lower than about 100 GeV, when the baryon number violation provided by the SM sphalerons is switched off? The three required ingredients for baryogenesis are of baryon number violation, C and CP violation and out-of-equilibrium dynamics [13]. It is not easy to generate the baryon asymmetry in a universe that reheats after inflation to a low temperature because the first and third ingredients are hard to come by: it is difficult to introduce baryon number violation at low temperatures without contradicting laboratory bounds on baryon number violation, and the universe is expanding so slowly at low temperatures that it is very close to equilibrium. However, several examples of such mechanisms exists, noticeably in the context of baryon-number violating SUSY at the EW scale. In ref. [14] such a scenario demands a notoriously low reheating scale of order of dozens of MeV. This unappealing feature can be potentially avoided in a scenario of [15] if very heavy gravitinos are assumed.[1]

Another interesting example of low-energy baryogenesis, that we will later return to, was introduced in the context of "darkogenesis" [17]. In the latter scenario the asymmetry is first produced in the hidden sector, which harbors new particles with the baryon number, and later on shared with the visible sector. The asymmetry in the hidden sector is produced due to the strong 1st order phase transition (PT), which provides low-temperature departure from equilibrium.

In this paper we propose a testable mechanism that overcomes these difficulties and naturally produces the baryon asymmetry at the low temperatures, when the sphalerons are not active anymore. The scenario that we propose crucially relies on the existence of a hidden valley that is endowed with a very non-trivial thermal dynamics and is capable of producing heavy particles out of equilibrium, which further decay to the SM via interactions that violate both the baryon number and the CP. The hidden valleys as beyond the SM (BSM) particle sector with potentially low masses, significantly below the EW scale, but with weak couplings to the SM via higher-dimensional operators (suppressed, say, by the TeV scale) have been proposed in refs. [18–20]. The motivation for this kind of physics was dominantly signature-driven, because production of these new particles promised rare but spectacular events at LHC. Nonetheless, it was soon understood that the hidden-valleys scenarios can be highly theoretically motivated, for example, being a necessary ingredient of various scenarios, addressing the naturalness problem of the SM [21–23].

In this paper we show that the hidden valleys can also be closely related to the problem of the baryogenesis. The basic idea of our mechanism is that the strong 1st order PT in the hidden sector can produce non-thermally particles, which are themselves much heavier than the temperature of the PT, closely following the old idea of ref. [24]. This can happen only in very strong 1st order PT, when the friction of surrounding plasma is not big enough to stop the bubble walls of the broken phase from propagating with the speed on light. Collisions of these ultra-relativistic bubble walls can potentially produce particles much heavier than the temperature of the PT itself. The decays of these heavy particles, however, dominantly proceed into the SM rather than the hidden sector due to the accidental symmetries of the model. The observed baryon asymmetry is produced in these decays. Thus an appealing and unique feature of this scenario is that one should neither necessarily reheat the Universe to the temperature of the decaying particle, nor produce the decaying particles during the

[1]This first of these scenarios is now heavily disfavored by the LHC searches for the RPV SUSY and displaced vertices [16].

inflation. Reheating to the temperatures above the temperature of the hidden PT should be sufficient to trigger the mechanism. In this paper we will both discuss the mechanism in detail and introduce a particular model, which is fully calculable and perturbative, satisfies all the existing cosmological constraints and illustrates all the features of the mechanism, being its existence proof.[2]

Of course, because one can probably construct lots of pretty different models along the lines of our proposal, potentially including the models that take advantage of non-perturbative dynamics, it is hard to talk about generic signatures of this scenario, for example any kind of guaranteed signature for the LHC or future colliders (although particular realizations might have interesting predictions both for the LHC and the di-nucleon decay experiments). However, one signature of this scenario is completely generic and virtually inevitable: the gravitational waves. The strong first order PT which must proceed in the hidden sector necessarily produces gravitational waves. We will show in our paper, that these gravitational waves are of right intensity and frequency to be observed or excluded by eLISA experiment, providing a unique opportunity for this future interferometer.

Our paper is structured as follow. In the next section we discuss the basic mechanism that we propose in details, discussing qualitatively the conditions for the runaway bubbles during the PT and why one might expect heavy particle production from these collisions. In section 3 we present an existence proof to the mechanism, described in section 2. We show that this particular model can satisfy all the cosmological constraints, identify particular spots in parameter space where phase transition with the runaway bubbles occur and estimate possible baryonic abundance. We show that the baryonic abundance produced can be easily be in agreement with the observed one. In section 4 we discuss the experimental signature with a special emphasis on the primordial gravitational waves and the prospects of eLISA. Finally, in the last section we conclude and discuss some open questions and possible future directions.

2 The mechanism description

The mechanism that we propose relies on an assumption, that heavy particles, ψ, which decay out of equilibrium via CPV and baryon-number violating (BNV) or lepton-number violating (LNV) operators, are produced *non-thermally* at temperatures much lower than their masses. This will allow us to produce these particles abundantly even if the Universe is reheated to the temperature, which is significantly lower than the mass of the particle ψ, or, alternatively if the ψ is annihilated too efficiently.

If the decays of the particle ψ proceed at high temperatures, above the temperature of the EWPT (which, in the absence of the new physics, is simply a crossover), one can take advantage of, for instance, LNV (or of a production of any other charge asymmetry as long as is not orthogonal to the baryon number), as long as this does not contradict the low-energy experimental test. If the asymmetry is formed in the leptonic sector, it will efficiently moved to the baryonic sector by the SM sphalerons. On the other hand, we will be in particular interested in the very *low temperature baryogenesis*, both because this region

[2]It is not the first time when the runaway bubbles are invoked in the model-building to produce the Universe baryonic asymmetry. Refs. [25, 26] take advantage of the reheating process which that follows the runaway bubbles collisions to produce the baryon asymmetry at the EW scale. Ref. [27] advocates low-scale baryogenesis from the delayed EWPT with the runaway bubbles. In spite of these similarities, as we will see, our proposal is quite different.

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of the parameter space is less explored and because this is where we can expect interesting gravity wave signature in the detection range of eLISA. If these decays proceed after the EWPT, the SM sphalerons are not active anymore and the decay must proceed via the BNV operator. This is why we will mostly concentrate on the BNV decays, since they can generically work at any temperature scale, however we should bare in mind that for the high-temperature decays LNV is also a viable option.

On top of this, the decays of the ψ must proceed via CPV operators in order to produce the desired baryonic asymmetry. We will show that in particular realizations that we have in mind this CPV can be very efficiently hidden from the low-energy probes like electron and neutron EDMs even if the baryogenesis scale is very low. Similar to the common lore of the leptogenesis and the decays of the "WIMP-like particles", the sector of the new fermions should have at least two independent CPV phases.

Many mechanisms are known to produce stable or metastable particles non-thermally. Usually they rely on decays of another heavy particles out of equilibrium. However, here we are interested here in low temperature scenario with potentially very low reheating temperature. A non-thermal particle production mechanism, which can efficiently produce the particles much heavier than the equilibrium temperature, we will take advantage of is the runaway bubble collisions in the strong first order cosmological phase transition. This mechanism has been discovered lots of time ago [24, 28] and was mostly discussed in the context of the SM and the potentially first order EWPT.

The non-thermal low-temperature production of the heavy particles crucially demands a strong first-order cosmological PT. Note, that this requirement is very different from the usual EW baryogenesis requirement of the strong first order PT during the electroweak symmetry breaking (EWSB). The EW baryogenesis requires quasi-stationary bubbles of the broken phase, namely that the bubbles do not propagate "too fast" through the plasma [29–31]. Practically for successful EWBG one demands that v_w is smaller than the speed of sound ~ 0.3, while the baryon asymmetry is almost v_w-independent for the values of v_w between 10^{-3} and the speed of sound (see ref. [32] and review [5] for detailed discussions).[3] In the our case we do not desire to produce an *asymmetry* via the phase transition, but rather *an abundant population of heavy states*. Therefore we will require bubbles which propagate almost with the speed of light through the plasma, namely $\gamma \gg 1$, with the regular definitions

$$\gamma \equiv \left(1 - v_w^2\right)^{-1/2} \tag{2.1}$$

It can be intuitively explained why do we expect the collisions of ultra-relativistic bubble walls to produce particles that might be much heavier than the temperature of the PT. In the case of the steady bubbles, the bubble wall carries a typical momentum $p \sim T$, therefore their collisions produce only particles with masses of order $m \lesssim T$, which further promptly thermalize in the primordial plasma. In the case of the runaway bubbles the incoming momentum is in fact much bigger, $p \sim \gamma T$. This in principle allows non-negligible production of particles of order γT, which is much heavier than the temperature of the surrounding plasma if $\gamma \gg 1$ [24]. Of course in oder for the mechanism to work, one should make sure that the particles that are formed do not thermalized in the surrounding plasma and stay out-of-equilibrium.[4]

[3]This requirement is needed because the excitations in plasma must diffuse efficiently in front of the bubble wall. See also [33] for a related discussion.

[4]Although the collision of the ultra-relativistic bubbles is a highly non-equilibrium process, in the proposed mechanism it is responsible for production of heavy particles which do not carry any asymmetry by themselves.

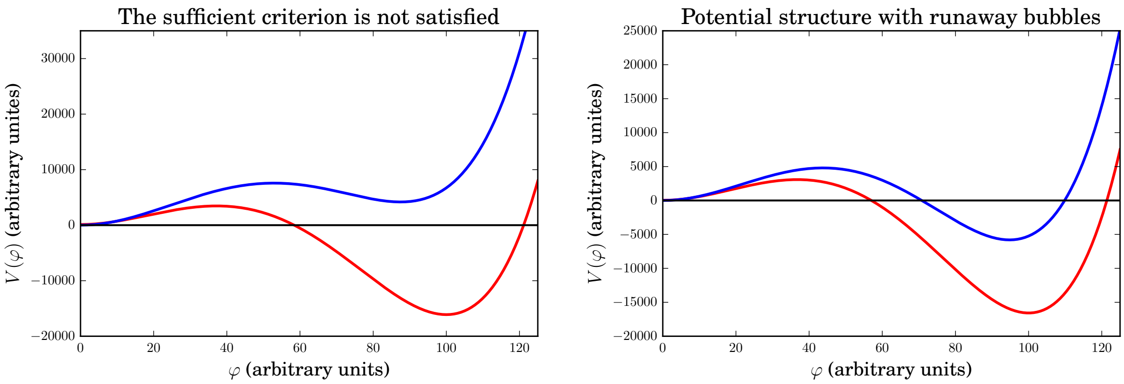

Figure 1. Runaway bubbles in potentials which satisfy Bödeker-Moore criterion. The red line stands for the full one-loop thermal potential, while the blue line stands for the mean free field approximation. If the potentials look like on the right panel, the bubbles during the first order phase transition clearly have runaway behavior. The picture on the left panel does not satisfy the criterion: although the local symmetry breaking minimum exists both in the full potential and in the mean free field approximation, the vacuum energy of the symmetry breaking minimum in the mean field approximation is bigger than the vacuum energy of the symmetry preserving phase.

A condition for sustaining the runaway bubbles in certain cosmological phase transitions was nicely formulated in ref. [34]. This reference shows, that in a gauge symmetry breaking phase transition the pressure on the interface between the phases for $\gamma \gg 1$ can be found by replacing the standard thermal potential by a mean field thermal potential $V_{\rm MF}$. The latter can be approximated as

$$V_{\rm MF} = \frac{T^2}{24} \sum_a m_a^2(\phi) \tag{2.2}$$

when the sum runs over all the particles which are in thermal equilibrium during the phase transition. This in turn sets a simple self-consistency criterion for sustaining the runaway bubbles if they beforehand reach the ultra-relativistic velocity: if at the nucleation temperature one finds that the mean field-approximated potential at the symmetry breaking minimum point (as determined by the full thermal potential!) has lower energy than the symmetry preserving local minimum, the bubbles are expected to be runaway.[5]

We will further use this criterion to establish the parameter space of the model. We illustrate this criterion on figure 1.

From the very formulation of this criterion it is clear that we will further have to focus on the calculable models of the symmetry breaking PT. This does not mean that our mechanism cannot work, for example, in systems with confinement PT. However, since these systems are inherently uncalculable and the criterion that would be analogous to the Bödeker-Moore criterion is not known for these systems, we cannot make a firm point that such a system produces the runaway bubbles. We believe that some of these systems do, but we cannot show this with our tools and leave this as an open question. It is also worth noticing that

[5] Another criterion for the runaway bubbles, based on the calculation of [34], was introduced in ref. [35]. We find, in full agreement with [36], that this criterion is stronger than the Bödeker-Moore criterion, although the models that we consider differ significantly from those, considered in [36]. Given the current theoretical uncertainties in the treatment of the runaway bubbles, we are going to adopt the Bödeker- Moore criterion which is the most conservative.

because our model will rely on the higgs-like phase transition, it will not be natural in the same sense as the SM is unnatural. Since we would like to present here an existence proof, we will not try to solve this problem here. However, it can be relatively easily rendered natural either via supersymmetrization of the hidden sector or via taking advantage of the confinement PT as we have previously discussed.

Notice, that from the point of view of hidden valley thermal behavior our mechanism shares some similarities with the idea of the "darkogenesis" [17]: both demand strong order phase transition in the hidden sector at temperatures that are smaller than the temperature of the EWPT. However, the "darkogenesis" demands stationary bubbles, because the asymmetry is produced in during the dark PT itself, and the dark sector must be augmented with the sufficient sources of the CPV. Our mechanism takes advantage of the different regime of this kind of models, with the runaway bubbles and demands no CPV sources in the dark sector itself.

Let us now concentrate on the last part of our mechanism, namely the heavy particle that we have introduced. In order to be produced abundantly, they should be coupled sufficiently strong to the particles of the hidden sector that undergoes the strong first order PT. On the other hand, they should not decay to the particles of that sector. We will concentrate on the models, in which these heavy particles possess an approximate $\mathbb{Z}_2$ symmetry, which is exact with respect to the particles of the hidden sector and broken by a small coupling to the SM matter to allow out-of-equilibrium decays. Therefore, these heavy particles can be either in singlet or adjoint representations of the SM. For simplification purposes we will assume them to the the SM singlet fermions. The models with the SM adjoints can be potentially much more challenging, because they can imply other particles with the SM charges in the hidden sector.

From the structural point of view the sector of the heavy decaying particles very closely resembles the analogous structure in the baryogenesis for weakly interacting massive particles [37], which can by itself viewed as an incarnation of "WIMPy baryogeneis" [38]. The crucial difference between our scenario and these works, is that by assuming a light hidden valley, we can make the entire mechanism operate at much lower temperature scales that the standard WIMP-inspired scenario. Moreover, because we can produce even at low scales much bigger abundances of the heavy decaying particles that one expects from the thermal scenarios, we will be able to get the observed baryonic asymmetry with much smaller values of the CPV, than one needs in the "WIMPy" scenarios.

3 A model

Let us now describe a model that can serve as an existence proof to the mechanism, described in section 2. Consider an SU(2) gauge symmetry in the hidden sector with "dark higgs" Φ in the fundamental of the dark SU(2) and a pair of fundamental fermions L_i with $i = 1, 2$. We allow a standard symmetry-breaking potential for Φ

$$V = -m^2|\Phi|^2 + \lambda|\Phi|^4 \; . \tag{3.1}$$

We will further use Φ for the full dark-SU(2) doublet and φ for the dark higgs excitation, using following parametrization

$$\Phi = \begin{pmatrix} 0, & \frac{f+\varphi}{\sqrt{2}} \end{pmatrix} \tag{3.2}$$

and f standing for the VEV d the dark higgs.

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Let us also introduce a pair of singlet fermions e_i, again with $i = 1, 2$. The most generic couplings involving the fermions that we can write down are

$$\mathcal{L} \supset y_{ij} \Phi L_i e_j + m_L \epsilon_{ij} \epsilon^{ab} L_a^i L_b^j + (m_e)_{ij} e^i e^j \ . \tag{3.3}$$

We imagine that the couplings $y_{ij} \sim \mathcal{O}(1)$, but non-degenarate, while all the rest of the bare masses, m_L and m_e, are of the same order of magnitude as $y\langle\Phi\rangle$. As previously, i is a flavor index and a is an SU(2) index. After the hidden SU(2) breaking and diagonalization of the fermionic mass matrix one gets 6 Majorana fermion states roughly at the scale $\sim y\langle\Phi\rangle$. Note that these particles generically will inherit $\mathcal{O}(1)$ couplings to the dark higgs, however the lightest fermionic states *cannot decay into the hidden sector* because of the accidental $\mathbb{Z}_2$ symmetry, under which they are odd and the rest of the hidden sector is even. Namely they perfectly satisfy the criteria, that we have laid out in the previous section: heavy particles with $\mathcal{O}(1)$ couplings to the dark sector but unable to decay into it. We will later couple these particles weakly to the SM particles to ensure their decays via BNV and CPV operators to produce the visible sector baryonic asymmetry.

The fermions that we have described above have to be produced in the bubbles collisions during the dark SU(2) breaking phase transition. As we will see later explicitly, this will be possible if we choose following regime

$$m_{\text{fermions}} \gg m_{\text{gauge-bosons}} \gg m_\varphi \ . \tag{3.4}$$

Of course this model of the hidden sector is by no mean novel. It strongly resembles the EW model (without gauging the $U(1)_Y$), as well as handful of the "darkogenesis" models. The trick is that in the regime (3.4), which has not been studied in detail beforehand, the model exhibits a very non-trivial thermal behavior that we will take advantage of.

The structure of the rest of this section is as follow. First, in 3.1 we establish the parameter space for the runaway bubbles and discuss the production of the heavy fermions in the collisions. However, because the hidden sector might populate very low mass ranges, one might worry about undesired relics and late decays in this sector. We address these questions of the cosmological safety in 3.2. In the subsection 3.3 we discuss the couplings of the exotic heavy fermions to the SM and the production of asymmetry in decays. Finally, in 3.4 we will return back to the heavy fermions sector, define the couplings to the SM and calculate the expected baryonic asymmetry that we produce.

3.1 Runaway bubbles in the hidden sector

We concentrate now on the higgs and gauge bosons part of the hidden sector. As we have alluded beforehand, we will be interested in regime where the gauge bosons are much heavier than the dark higgs, triggering gauge bosons driven PT. This in turn, determines the hierarchy of couplings $g^2 \gg \lambda$, where g is a dark gauge coupling and λ is a dark quartic. It is well known from the studies of the SM with a very light higgs, that the phase transition in this case is strong first order, and driven by the gauge bosons. We further calculate the thermal potential for the dark higgs, which in our case becomes

$$V_{\text{th}} = \frac{g_i T^4 (-1)^{F_i}}{2\pi^2} \int_0^\infty dx\, x^2 \log\left(1 - (-1)^{F_i} \exp\left(\sqrt{x^2 + \frac{m_i^2(\varphi)}{T^2}}\right)\right) \tag{3.5}$$

It is well known that although qualitatively this expression is sufficient, it misses lots of numerical effects that have to do with large two-loop corrections, non-perturbative effects and further uncertainties that we will latter list in detail. However, since from the SM lesson we know that it does give qualitatively correct answer, we will proceed with this expression. Because we are merely trying here to give a existence proof, we will be less worried about numerical factors of $\mathcal{O}(1)$ that we will be likely to miss with our procedure. We also perform "daisy" resummation [39, 40] by replacing in the thermal potential

$$m_a^2(\varphi) \to m^2 + \Pi_a(T)\,, \tag{3.6}$$

with $\Pi_a(T)$ being the thermal masses squared of the particles, usually of order of the plasma temperature.

In order to check, if the potential satisfies the Bödeker-Moore criterion, we compare it at temperature T to the mean field potential, which in our case takes a compact form

$$V_{\text{th,MF}} = \frac{3\,g^2T^2h^2}{32}\,. \tag{3.7}$$

Of course, in both cases we add the thermal potential to the tree level potential and the one-loop Coleman-Weinberg potential.

It is worth now explaining why do we originally choose the regime (3.4) and why does the system in this regime often satisfy the sufficient criterion for the runaway bubbles. Ref. [34] pointed out that runaway bubbles in the context of the EWSB is reachable with at least two scalar fields in the higgs sector and can happen in singlet-catalyzed EWSB. The reason is very straightforward: one usually gets the system with two different vacua from the *cubic* higgs term in high-temperature expansion of eq. (3.5). Because the mean field thermal potential only cares about the leading term in T^2, it contains no cubic term in the higgs field, essentially ruling out the possibility of getting two-vacuum structure in the mean field approximation with a cubic term that is not produced already at the zero temperature. In the concrete example of ref. [34] the extra singlet has been used to produce a tree level effective cubic term.

However, our scenario is qualitatively different from the EWSB. Because we insist on the mass hierarchy (3.4), our quartic coupling in the higgs sector is much smaller than the gauge coupling g. In this situation we might discover, that in the $T = 0$ potential the 1-loop Coleman-Weinberg terms are as important, as the tree level terms. This is indeed the case when

$$\lambda \sim \frac{g^2}{16\pi^2} \tag{3.8}$$

and this is precisely a regime that we are going to further focus on. For this choice of parameter space, the $T = 0$ potential has a much more complicated structure, because it also includes the Coleman-Weinberg term

$$V_{\text{CW}} = \frac{N_g}{64\pi^2} m_g^4(\varphi) \left(\log\frac{m_g(\varphi)^2}{v^2} - \frac{5}{6}\right) \tag{3.9}$$

where N_g is the total number of heavy gauge bosons degrees of freedom, which is 9 in the case of our simple SU(2) toy model and $m_g = \frac{g\varphi}{\sqrt{2}}$ is the higgs-dependent mass of the heavy gauge bosons. Of course, in order to make sure, that the values of the $T = 0$ higgs mass and its VEV do not move from the values that we have chosen, we add appropriate one-loop counter-terms to the λ and m^2.

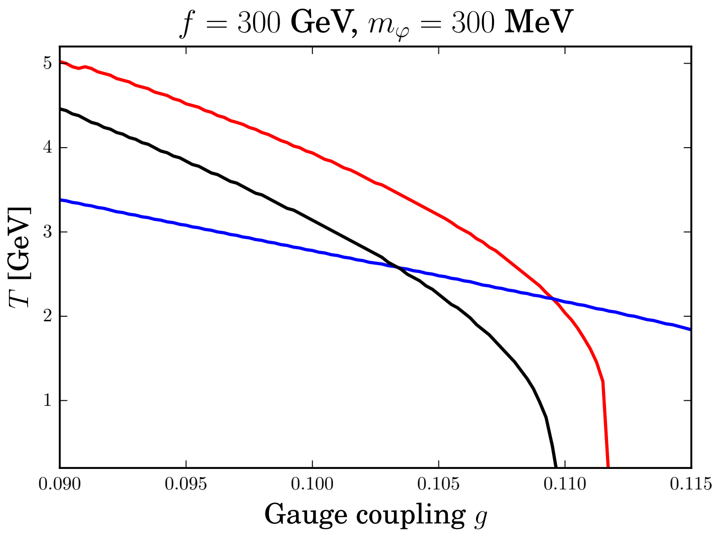

Figure 2. The upper bound on the nucleation temperature (red line) and the highest temperature at which BM criterion for the runaway bubbles is satisfied (blue line). Below the black line the barrier between the real vacuum and the false vacuum disappears. In the most pessimistic scenario the relevant parameter space lies to the right of the intersection between the blue and the red curves, while in the most optimistic scenario it is to the right of the intersection between the black and the blue curves.

This is not the first time that it noticed that significant changes in the structure of $T = 0$ potential might have an important impact on the nature of the symmetry breaking. Similarly to our structure, it was noticed in [41] that the changes in the *zero-temperature potential* can trigger the first order PT in the EWSB. This picture is very different from the standard EW first order PT, driven by the thermal loops of the new particles. In the picture of [41] the barrier between the symmetry-breaking and the symmetry-preserving vacuum may persist even at the zero temperature, while the thermal loops play relatively a minor role in shaping the thermal potential (the effect was also discussed in detail in [42]). In the SM case this is not that easy to achieve because for the $m_h \approx 125\,\text{GeV}$ the quartic of the SM is $\lambda_{\text{SM}} \approx 0.12$. If one would like to have a Coleman-Weinberg (CW) potential that is numerically comparable to the tree level one, he would need essentially non-perturbative couplings of the new particles to the higgs. Ref. [41] suggests overcoming this problem by putting large number of new scalars at the EW scale with $\mathcal{O}(1)$ couplings to the SM higgs. However, it is not clear weather the model, presented by this reference is in fact calculable in 't Hooft sense, namely $N\lambda \ll 1$.

Nevertheless, one can easily envision this scenario without dubiously large couplings or high-multiplicity fields in the hidden sector. Moreover, the situation emerges naturally in the regime (3.8) where the corrections to the 1-loop CW potential at $T = 0$ from the gauge bosons are so significant, that they in fact reproduce the picture of the zero-temperature produced barrier.

To check further for the candidate points in the parameter space we proceed as follow. We identify the points which satisfy the Bödeker-Moore criterion for a certain temperature and calculate the upper bound on the nucleation temperature for these points. If the $O(3)$ solution has the least action, the tunneling probability per unit time per unit volume is given by [43, 44]

$$\frac{\Gamma}{\nu} = A(T)e^{-S_3/T} \tag{3.10}$$

where S_3 is the three-dimensional bounce action and $A(T)$ is a prefactor of order $\sim T^4$. In order to calculate the bounce action, we should solve for ϕ the Euclidian equation of motion

$$\frac{d^2\varphi}{dr^2} + \frac{2}{r}\frac{d\varphi}{dr} = \frac{dV(\varphi;T)}{d\varphi}\,, \tag{3.11}$$

which is a boundary value problem with the boundary conditions $\varphi(\infty) = 0$ and $\dot{\varphi}(0) = 0$. One further calculates the 3D bounce action using this approximation.

One can make a further approximation by assuming that the thickness of the bubble wall at the moment of its formation is much smaller than the radius of the bubble [43, 44]. In this case the 3D actions can be approximated as

$$S_3 \approx -\frac{4\pi}{3}r^3\Delta V + 4\pi r^2 S_1, \quad \text{where} \quad S_1 \equiv \int d\varphi\sqrt{2V(\varphi;T)} \tag{3.12}$$

where r stands for the radius which extremizes the 3D action, ΔV is the energy difference between the false minimum and the global one, and the integration in the surface term S_1 runs from zero to the point at which $V(\varphi_{\rm end}) = V(\varphi = 0)$. In fact, this approximation is decent if ΔV is much smaller than the barrier between the minima. This corresponds to the *weak first order PT*, which is not our case. Therefore, strictly speaking the approximation (3.12) never works in our case. However, it had been shown in ref. [45] that this approximation always overestimates the temperature of the PT, henceforth it always underestimates the strength of the PT. The underestimation is in fact as big as $\mathcal{O}(1)$.

Finally after the we have calculated the 3D action as a function of temperature, the probability for a single bubble to nucleate within the horizon volume is order-one when following condition holds [46]:

$$\xi \equiv \int_{T_{\rm nuc}}^{\infty} \frac{dT}{T}\left(\frac{2\zeta M_{\rm pl}}{T}\right)^4 e^{-\frac{S_3(T)}{T}} \sim \mathcal{O}(1)\,, \tag{3.13}$$

where

$$\zeta^{-1} \equiv 4\pi\sqrt{\frac{\pi g_*(T)}{45}}\,. \tag{3.14}$$

We show the bound on the nucleation temperature vs the minimal temperature at which the Bödeker-Moore criterion for the runaway bubbles is satisfied for a particular choice of m_φ and f on figure 2. Clearly we see that the upper bound on the nucleation temperature is lower than the highest possible temperature at which the bubbles run away for an appropriate choice of the gauge coupling g. Note also that because the PT is very strong first order in the entire range presented on the figure, the bound on the nucleation temperature is expected to overestimate the real nucleation temperature vastly. Therefore, it is not unlikely that the entire parameter space, or much bigger portion of it, where the minimal PT temperature does not exceed the bound on the BM criterion ($g \gtrsim 0.103$), is suitable for the runaway bubble collisions. Note also, that the bound on the nucleation temperature that we find, is in fact significantly bigger than the dark higgs mass, while at least naively we expect it to scale as the higgs mass.

3.2 Cosmological safety of the hidden sector

Although the model that we propose in this paper is merely an existence proof to the larger new mechanism, and probably should not be taken too seriously for the phenomenological purposes, one should still check that it is not in conflict with the basic cosmological observations. Here we show that all the potential problems can be resolved in the relevant regions of the parameter space.

The potential problems are twofold:

1. The dark higgs is the lightest dark particle without any obvious efficient annihilation channels. We should introduce new couplings, to make sure that the higgs decays into the SM faster than within 1 sec to ensure the Big Bang Nucleosynthesis (BBN) safety

2. The SU(2) gauge bosons themselves are dark-stable. A-priori this might be a worry. We will show, that in the full theory, where we include the fermions and the interactions with the SM, they decay to the SM states faster than within 1 sec.

The first problem on this list can be addressed in several different ways. The most straightforward one that we will focus on is simply via introducing a tiny higgs-portal coupling to the SM higgs portal:

$$\mathcal{L} = \kappa |H|^2 |\Phi|^2 \tag{3.15}$$

One should be careful with this kind of coupling because it also induces an undesired mixings between the visible and the dark higgs, as well as it induces invisible higgs decays, which are now constrained by the LHC data. Moreover, while such a coupling induces new mass parameters both to the visible and the hidden higgses. We explicitly check that the induced $T = 0$ mass does not exceed the bare mass, triggering a potential instability of the higgses potentials.

In the mass insertion approximation the coupling (3.15) induces following decay rates of ϕ into the SM fermions

$$\Gamma(\varphi \to f\bar{f}) = \frac{\kappa^2 f^2 m_f^2 m_\varphi}{4\pi m_h^4}\left(1 - \frac{4m_f^2}{m_\varphi^2}\right)^{3/2} \tag{3.16}$$

while the invisible decay width of the SM higgs into the hidden higgses is

$$\Gamma(h \to \varphi\varphi) = \frac{1}{8\pi}\frac{\kappa^2 v^2}{m_h}\sqrt{1 - \frac{4m_\varphi^2}{m_h^2}}\ . \tag{3.17}$$

The invisible higgs width has been looked for both by ATLAS and CMS collaborations [47, 48], reporting exclusion of the SM-like higgs invisible branching ratios of 65% and 75% respectively at 95% confidence level. Of course a proper combination of these results is beyond the scope of this work and we will crudely assume for the further analysis that the invisible higgs BR should not exceed 50%, providing us with an upper bound on κ.

The lower bound on κ comes from the requirement that φ should decay with at 1 sec in order not to spoil the BBN predictions, or in other words $\Gamma(\varphi \to f\bar{f}) < 6.58 \times 10^{-25}$ GeV. Although this constraint looks pretty innocuous at the first glance, it turns out to set a meaningful and important bound on the light scalar mass. Note that below the mass of 1 GeV

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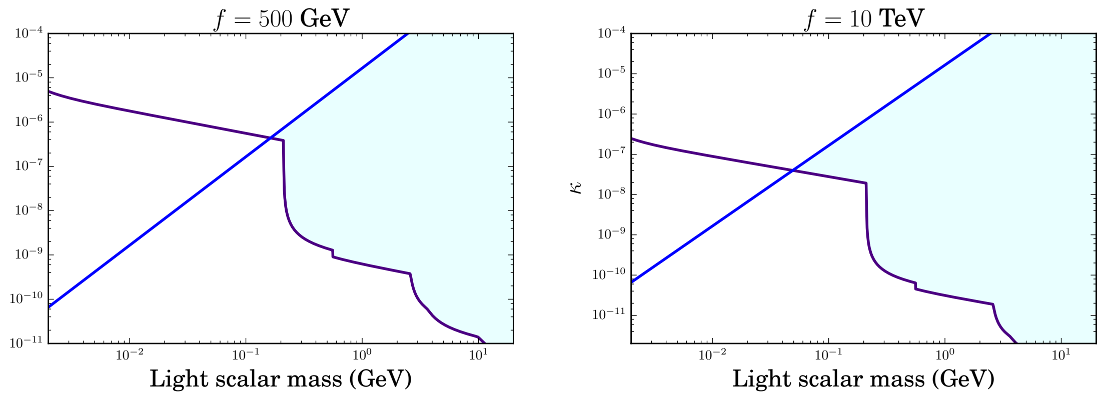

Figure 3. Constraints on the value of κ in the higgs portal coupling for different values of f. Above the blue line the structure of the hidden sector potential may be altered without appropriate fine-tuning due to too large induced dark higgs mass. Below the purple line the decay time of the dark higgs exceeds 1 sec. The shaded region is allowed. The constraints from the invisible higgs decays $\kappa \lesssim 0.01$ are completely subdominant and they are not shown on these plots.

the scalar is allowed to decay only into the light fermions, suffering both from the small κ suppression and from the SM Yukawa suppression.[6] On top of this, one gets an upper bound on κ from the demand of the stability of the dark higgs potential. We summarize all these constraints for representative choices of f on figure 3.

As we see, the constraints coming from the higgs invisible width are completely subdominant and the constraints are dominated by considerations of the structure of the dark higgs potential and the fast enough decay of the dark higgs. The parameter space widely opens up around $m_\varphi \sim 100\,\mathrm{MeV}$. Here we emphasize that the bound that we get here is only relevant for this particular model, otherwise the scale of the *mechanism, rather than the model* can be as low as the BBN scale 1 MeV.

Another potential cosmological worry is the dark-stability of the massive gauge bosons. However, even in the dark sector itself the accidental $\mathbb{Z}_2$ symmetry which protects the stability of the dark gauge bosons, will be broken by the couplings to the dark fermions. Further, when the interactions with the SM fermions will be taken into account, we will see that these dark massive gauge bosons can potentially decay way faster than 1 sec into the SM fermions and therefore pose no real cosmological threat.

3.3 Couplings to the SM

After establishing the parameter space of the runaway bubbles PT and setting the scales which do not contradict the existing cosmological observations, we are now ready to introduce the heavy fermions sector and later to estimate the final baryon asymmetry.

Let us see how can we couple it to the BNV operators in the SM and estimate the baryon asymmetry that they produce. Clearly, we cannot couple L_i particles to the SM, because such a coupling would necessarily be non-gauge invariant under the hidden SU(2). However, the gauge singlets e_i come here to our rescue. The most generic couplings to the

[6] Although the decays into di-photons are kinematically allowed, their branching ratio (BR) never exceeds the the BR into the fermions, as long as the decays into e^+e^- are kinematically allowed. The same is true for the decay into the gluons above the Λ_{QCD} scale.

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BNV SM operators one can introduce are

$$\mathcal{L} \supset \frac{1}{\Lambda^2}\left(\lambda_{\alpha ijk} e_\alpha u_i^c d_j^c d_k^c + \eta_{\alpha ijk} e_\alpha Q_i Q_j (d_k^c)^\dagger\right) \ . \tag{3.18}$$

It does not really matter at this stage what is precisely the operator that we further proceed with. Note, however, that these operators have a slightly different flavor structure, which will be important for our further discussion in section 4. While the first operator is antisymmetric in the color space, and therefore demand also anti-symmetrization over j and k indices in the λ coefficients, the second operator is anti-symmetrized in the color space and in the SU(2) space, allowing non-zero $i = j$ coefficients inside η.

We further analyze the decays of the heavy hidden fermions, which eventually produce the baryon asymmetry. Let us collectively call all the Majorana fermions, that we get from the diagonalization of the fermion matrix ψ. For concreteness we concentrate on the couplings $\lambda_{\alpha ijk}$ in eq. (3.18) and set all the coupling η to zero.[7]

To simplify further the discussion let us UV-complete the four-fermi effective operator (3.18) via introducing an exotic colored scalar Δ with the hypercharge $-2/3$:

$$\mathcal{L} \supset \lambda'_{\alpha i} \Delta e_\alpha u_i^c + \lambda''_{jk} \Delta^* d_j^c d_k^c + M_\Delta^2 |\Delta|^2 \ . \tag{3.19}$$

In general, all the coefficients λ' and λ'' are complex. Note also that the fermions e_α are *real* due to the Majorana mass term (cf. eq. (3.3)) and therefore they cannot be re-phased.

If we forget for a moment about the SM Yukawa couplings, all the imaginary phases of the coupling λ'' can be absorbed into the re-phasing of the RH down-type fermions. If we had been forced to rely only on those couplings, the only CPV effects would be proportional to the Jarlskog invariant, forcing our CP-efficiency to be as small as $\epsilon_{\rm CP} \lesssim 10^{-8}$.

To avoid this marginal possibility, one can rely on the coefficients λ'. If we had had only one generation of e_α, we would have faced the same problem here, because all the imaginary phases could be absorbed into redefinitions of u_i^c. However, if we can two generations of the fermions e_α we have three irremovable phases, very similar to the common lore of the standard leptogenesis scenario. Also, similar to the standard leptogenesis scenario, one gets the CPV violating effect from the interference between the tree level and the one-loop diagram as schematically shown on figure 4. With the assumption that the phases can be as big order-1 and the splittings between various Majorana states are comparable to the masses of the fermions ψ, the maximal value of the CP-efficiency we can get in a non-resonant perturbative scenario is

$$\epsilon_{CP} \lesssim \frac{1}{8\pi}\left(\frac{m_\psi}{m_\Delta}\right)^2 \ . \tag{3.20}$$

As we will see in the next subsection, this allows us lots of freedom with choosing parameters, before we even get close to the collider bounds imposed by the LHC. First, we will see that we do not need very high value of $\epsilon_{\rm CP}$, and the values as small as $10^{-4} \ldots 10^{-7}$ can be perfectly sufficient in order to get an observed relic abundance. This safely allows

[7]Our following discussion has a handful of similarities to TeV-scale model baryogenesis of refs. [49, 50] because of the very similar couplings structure. However, there are important differences. First, it will be much simpler, because we neither involve the neutrino masses in our discussion nor invoke the resonant processes. Second, more important, we introduce *two generation of Majorana fermions* e_α, which allows us, similar to the leptogenesis scenario, to get irremovable phases in the exotic fermions sector, without relying on the tiny Jarlskog invariance.

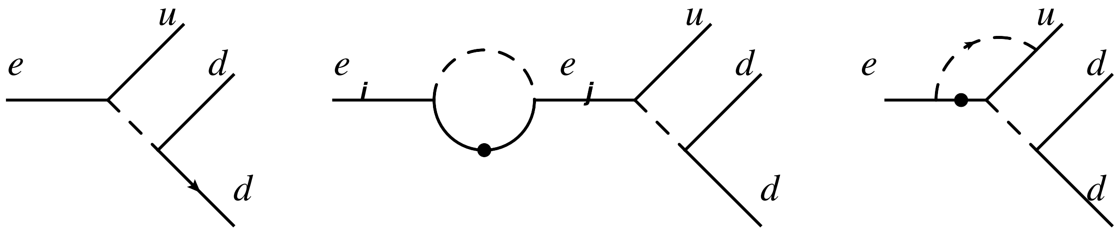

Figure 4. CP violating decays of the heavy fermions. In order to produce the asymmetry, one needs the interference between the (first) tree level diagram and the loop level diagrams, which are sensitive to the irreducible phases of the couplings of the scalar Δ. The blobs stand for the mixings between the Majorana particles, that we necessary have in the most generic model. These mixings are unavoidable because all the phases in the single-generation couplings are removable.

us to keep the colored scalar Δ more than an order of magintude heavier than the fermions ψ. Second in order not to suppress $\epsilon_{\rm CP}$ further, it is sufficient to keep the couplings λ' to be $\sim \mathcal{O}(1)$. The couplings λ'' essentially play no role in contributing to the CP-aymmetry. On the other hand the scale Λ from the eq. (3.18) maps onto $\Lambda^2 \sim M_\Delta^2/(\lambda'\lambda'')$. Hence by keeping λ'' small, we can make the scale Λ almost arbitrarily large without suppressing the CPV effects, easily above the generic bounds of $\sim 100\,\text{TeV}$, which apply from the double-nucleon decay (see section 4 for more details).

At this point we can return to the dark gauge bosons and briefly comment on their lifetime. We will explicitly see that it poses no cosmological worry. When the effects of the heavy fermions and the scalar Δ are properly taken into account, we find that the dark gauge bosons can decay to a pair of the SM fermions already at the one loop level. The decay proceeds via the gauge coupling and the couplings $\lambda'_{\alpha i}$ in eq. (3.19) (see figure 5). The decay rate is parametrically

$$\Gamma(W' \to u\bar{u}) \sim \frac{1}{8\pi}\left(\frac{g\lambda'^2}{16\pi^2}\right)^2 \frac{m_{W'}^5}{m_\Delta^4} \tag{3.21}$$

Note, that although we assume that the gauge coupling is small, say $g \sim 10^{-2}$, the coupling λ' is not necessarily small, and it can easily be $\mathcal{O}(1)$, while the scale Λ can be easily small because λ'' is small. For representative values $m_\Delta \sim 1\,\text{TeV}$, $m_{W'} \sim 10$ GeV and $g \sim 10^{-2}$ we get that the decay width is approximately 10^{-16} GeV, clearly faster than 1 sec, posing no cosmological problem.

3.4 Production of heavy matter in bubble collisions and estimation of the asymmetry

Now we can finally calculate the abundance of the heavy fermions produced in the bubble collisions. Because we expect dominantly post-sphaleron decays and no important sources of the washout, we expect final produced baryon asymmetry to be $\Omega_B h^2 \approx \epsilon_{\rm CP} \times \Omega_\psi h^2$.

The abundance of heavy particles produced in runaway bubble have been explicitly calculated in ref. [51] with assumptions of either purely elastic or purely inelastic collisions.[8] In general, the energy produced in these collisions per unit area in the form of particle a is

[8]In this subsection, when we discuss the elastic vs. inelastic collisions of the DM, we mean the collisions of the bubble walls rather than the particles.

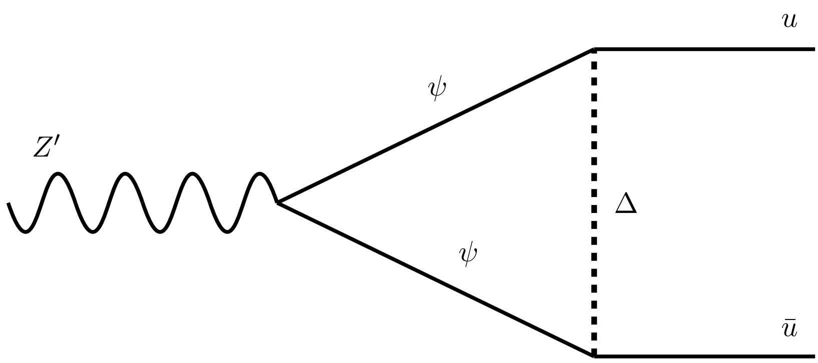

Figure 5. Dark gauge boson decay at one-loop, induced by the couplings of the heavy fermions and the scalar Δ to the SM. The coupling to the fermions is gauge coupling suppressed, but assuming $\lambda' \sim 1$ there are no suppressions in the Yukawa interactions.

given by

$$\frac{\mathcal{E}}{A} = \frac{1}{4\pi^2}\int_{\chi_{\min}}^{\infty} d\chi f(\chi)\sqrt{\chi}\int d\Pi_a \left|\bar{\mathcal{M}}(\varphi\to\alpha)\right|^2 = \frac{1}{2\pi^2}\int_{\chi_{\min}}^{\infty} d\chi f(\chi)\sqrt{\chi}\;\mathrm{Im}\tilde{\Gamma}^{(2)}(\chi)\,, \tag{3.22}$$

where the function $f(\chi)$ carries the information about the details of the bubble collisions and the efficiency of production of various particles in the runaway bubble collisions depending on the nature of the collision (elastic or inelastic respectively). $d\Pi_\alpha$ is a Lorentz-invariant phase space and $\tilde{\Gamma}^{(2)}$ is a two-point 1PI Green function in the momentum space. $|\mathcal{M}(\varphi\to\alpha)|^2$ is spin-averaged squared amplitude of the dark higgs decay into a set of particles α.

In expression (3.22) the integration runs over all the available invariant masses squared of the collisions, parametrized by variable $\chi \equiv \omega^2 - \vec{p}^2$. $\chi_{\min}$ stands for the minimal available invariant mass at which the production of the final state α is kinematically possible and it is equal to the square of the sum of all the particles masses in the final state α, namely $\chi_{\min} = (\sum_{\text{state }\alpha} M_i)^2$.

We will further analyze two different cases, differentiating the elastic bubbles collisions from the inelastic ones. The explicit expression for the function $f(\chi)$ in the elastic case, obtained in ref. [51], is

$$f(\chi)_{\mathrm{EL}} = \frac{16 f^2 \log\left(\frac{2(\frac{\gamma_w}{l_w})^2 - \chi + 2\frac{\gamma_w}{l_w}\sqrt{(\frac{\gamma_w}{l_w})^2-\chi}}{\chi}\right)}{\chi^2}\Theta\left(\left(\frac{\gamma_w}{l_w}\right)^2 - \chi\right) \tag{3.23}$$

where γ_w is defined in eq. (2.1), Θ stands for the step function and l_w is the thickness of the bubble wall. Practically one usually gets $l_W \sim \mathcal{O}(10)/T_{\mathrm{PT}}$. The above written approximation assumes the thin wall approximation. The reader can find the full expression, without a thin wall approximation in ref. [51], however we made an explicit numerical check and found that (3.23) is sufficient for our purposes and does not introduce intolerable numerical imprecisions. The analogous expression for the purely inelastic collisions is

$$f(\chi)_{\mathrm{IN}} = 4 f^2 m_\varphi^4 \frac{\log\left(\frac{2(\frac{\gamma_w}{l_w})^2 + \chi + 2\frac{\gamma_w}{l_w}\sqrt{(\frac{\gamma_w}{l_w})^2+\chi}}{\chi}\right)}{\chi^2\left((\chi - m_\varphi^2)^2 + m_\varphi^6(\frac{\gamma_w}{l_w})^{-2}\right)} \tag{3.24}$$

Parenthetically we note that this expression is also an approximation, although it is also adequate for us in the entire parameter space.

The work of Falkowski and No have analyzed the efficiency of the elastic and inelastic collisions in the case of the EWPT, namely where one naturally has $T_{\rm PT} \sim m_h \sim v$. They concluded that there is no efficient production above the scale m_h (already in striking disagreement with the original conjecture of Watkins and Widrow [24], which talks about efficient production of particles as heavy as $m \sim \gamma T$), while the situation is more subtle with the elastic collisions. The paper by Falkowski and No emphasizes that while scalars heavier than the higgs mass cannot be produced in the EW theory runaway PT, it is possible that heavy exotic gauge bosons and fermions can still be produced abundantly. While we agree with the statement on the scalars, we find that such a production in the SM is impossible also for the gauge bosons and the exotic fermions. The problem is that in order to ensure sufficient production of the exotic particles above $T_{\rm PT}$, ref. [51] assumed the tree level couplings of the exotic fermions to the Higgs $\mathcal{L} \supset \lambda H f \bar{f}$ and the tree level coupling of the higgs to the exotic vector bosons $\mathcal{L} \supset \lambda_V h M_V V_\mu V^\mu$, with H standing for the SM higgs doublet and h is standing for the physics higgs boson. In the former case it is impossible to talk about production of particles much heavier than the m_h because this would imply non-perturbative Yukawa coupling to the higgs.[9] In the latter case an efficient production of the heavy exotic gauge boson will jeopardize the unitarity of the $VV \to VV$ scattering.

After the short overview of the EWPT, let us return to our case that allows for a separation of scales m_h, f and $T_{\rm PT}$, which does not happen in the SM. We find that expressions (3.23) and (3.24) behave almost identically at $\sqrt{\chi} < m_\varphi$, with an inelastic solution having a characteristic resonance at $\sqrt{\chi} \approx m_\varphi$. However the behavior at higher masses is very different: while the $f(\chi)_{\rm IN}$ drops very fast above the scale m_φ, rendering the production of heavy particles at $\sqrt{\chi} > m_\varphi$ virtually impossible, the function $f(\chi)_{\rm EL}$ descends much less steeply. We find that in purely elastic collisions one can still efficiently produce in the runaway bubble collisions particles as heavy as f. Above this scale, both elastic and inelastic collisions turn out to be highly inefficient. Therefore, it is clear, that in order for our mechanism to work, we will need the collisions to be fairly elastic.

Another worry that we should emphasize here is the possibility of elastic collisions and the possibility of production of particles in the elastic collisions in general. It is entirely possible that in the case of perfect elasticity the bubbles simply travel through one another without producing any particles in the collisions.[10] This is the case, for instance, of the Sine-Gordon model, as noticed already in [24], for which the production becomes possible only at γ_w^{-1} order. Numerical simulations though show that the dynamics of bubble collisions is extremely complicated [52]: in models with a φ^4 models bubbles collide and indeed may reflect each other by also emitting soft scalar waves. We leave the decisive conclusion on this point to future studies.

Now we can use all these formula and, assuming the elastic collisions, estimate the total abundance of the baryonic matter in the Universe produced by our mechanism. The relevant two-point 1PI Green function for the *Dirac fermions* (note the factor of 2 because of the two fermion generations) is

$$\text{Im}\, \tilde{\Gamma}^{(2)}(\chi) = \frac{y^2}{4\pi}\chi \left(1 - \frac{4m_f^2}{\chi}\right) \Theta(\chi - 4m_f^2)\ . \tag{3.25}$$

[9] A theoretical possibility that this exotic fermion can gain its mass partially from the non-higgs sources would imply additional sources of the SM SU(2) × U(1), the possibility which is cornered by the EW precision tests.

[10] We are grateful to Thomas Konstandin for an illuminating discussion on this issue.

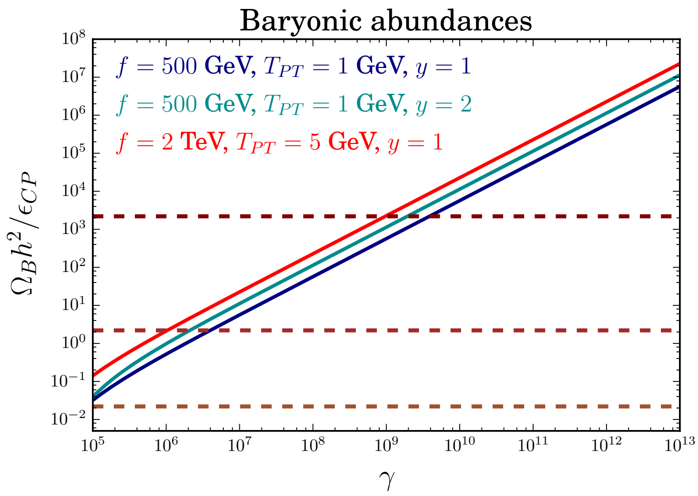

Figure 6. Predicted baryonic abundances (reweighed by the CP-efficiency of the process) as a function of γ. The dashed lines bottom-up show the measured $\Omega_B h^2$ (namely requiring perfect CP-efficiency), $100 \times \Omega_B h^2$ and $10^5 \times \Omega_B h^2$ respectively. In all the cases we have assumed $\beta^{-1} = 10^{-3}\, H^{-1}$.

A small subtlety is that due to the mixings resulting from the couplings (3.3) we are producing Majorana rather than Dirac fermions. However, if we are interested in total production of all the available states, the final result will still be given by expression (3.25).

After picking up together all the relevant terms, one finds the produced number of particles per comoving volume

$$Y_\alpha = \frac{225}{4\pi^2} \frac{1}{\sqrt{g_*} T M_{\rm Pl}} \left(\frac{\beta}{H}\right) \left(\frac{\mathcal{E}}{A}\right), \tag{3.26}$$

with β^{-1} being the duration of PT, typically $\beta \sim (100 \dots 1000) \times H$ [53].

We show the expected abundances for various points in the parameter space on figure 6. The maximal possible γ that we can expect is

$$\gamma \lesssim \left(\frac{\beta}{H}\right)^{-1} \frac{M_{\rm pl}}{f}, \tag{3.27}$$

namely can potentially be as large as $(10^{14} \dots 10^{15})$. In practice this value is smaller than the maximum possible velocity, because this naive estimation neglects the $\log\gamma$ friction terms, which might also become important for these values of γ. However, even for relatively mild values of $\gamma \sim 10^9$ we can get a correct baryonic abundance with pretty low CP-efficiency, in fact as small as $\epsilon_{\rm CP} \sim 10^{-5}$ which can very easily achieved perturbatively without any need invoking resonant baryogenesis.

4 Experimental signatures

Because the entire mechanism is crucially based on the strong first order PT at relatively low temperature with runaway bubbles, the prediction of the primordial gravity waves is completely model independent. Primordial gravitational waves are generic signatures of the strong cosmological first order PT, and their effects were studied in detail in the context of the non-standard EWPT [36, 54–59], in the case of SM QCDPT with the lepton asymmetry [60], and, recently, in the case of low-temperature hidden sector PT [61, 62]. The latest case is in fact similar to our study.

We analyze the predicted Gravity waves in the first subsection of this section. We later turn to comment on more model dependent signature, like $n-\bar{n}$ oscillations, double nucleon decay and even possible collider signatures.

4.1 Gravitational waves

The gravitational wave signal is especially interesting and intriguing in this particular mechanism not only because that it is very generic, but also because of its potential detectability in the forseeable future. Indeed, strong cosmological phase transitions at temperatures in range $(1\dots 50)$ GeV can be precisely in the preferred range of the future eLISA detector. At this point it is little bit difficult to estimate the sensitivity of the eLISA exactly, because the sensitivity studies are now under way. However lots of future designs of the eLISA will be able either provide evidence to our mechanism, or falsify big parts of its parameter space.

During a strong first order phase transition the primordial gravitational wave is produced due to three different processes which all generically coexist: collisions of bubble walls, sounds wave in the plasma after the bubbles collided and magnetohydrodinamic (MHD) turbulence in plasma forming after the bubbles collided. Therefore

$$\Omega_{\rm GW}h^2 = \Omega_\varphi h^2 + \Omega_{\rm sw}h^2 + \Omega_{\rm turb}h^2\,, \tag{4.1}$$

where Ω_φ stands for the first mentioned contribution. In our following analysis we will closely follow the discussion of [63], from where we also take the expressions for all these distributions. We provide here only the final formula for our estimations, and for more details the interested reader is referred to [63] and references therein.

The fundamental parameter, which largely determines the intensity of the gravitational wave emitted by the strong 1st order phase transition is the ratio between the vacuum energy density and the radiation energy density during the PT:

$$\alpha \equiv \frac{\rho_{\rm vac}}{\rho_{\rm rad}} = \frac{30}{\pi^2}\frac{V_{\rm th}(0,T_{\rm nuc}) - V_{\rm th}(f(T_{\rm nuc});T_{\rm nuc})}{g_* T^4} \tag{4.2}$$

With this parameter and the duration of the PT β^{-1} we can express the intensity of the gravity waves as a function of the frequency. Note, that the discussion in ref. [63] is pretty generic and applies to different velocities of the bubble walls. Here we assume runaway bubble walls, $v_w = 1$, in which case the relevant expressions for intensity are:

$$\Omega_\varphi h^2(f) = 0.13\times 10^{-5}\left(\frac{H}{\beta}\right)^2\left(\frac{\kappa_\varphi\alpha}{1+\alpha}\right)^2\left(\frac{100}{g_*}\right)^{1/3} S_{env(f)}\,, \tag{4.3}$$

$$\Omega_{\rm sw} h^2(f) = 2.65\times 10^{-6}\left(\frac{H}{\beta}\right)\left(\frac{\kappa_v\alpha}{1+\alpha}\right)^2\left(\frac{100}{g_*}\right)^{1/3} S_{\rm sw}(f)\,, \tag{4.4}$$

$$\Omega_{\rm turb} h^2(f) = 3.35\times 10^{-4}\left(\frac{H}{\beta}\right)\left(\frac{\kappa_{\rm turb}\alpha}{1+\alpha}\right)^{3/2}\left(\frac{100}{g_*}\right)^{1/3} S_{\rm turb}(f)\,. \tag{4.5}$$

The functions $S(f)$, which determine the dependence of the gravitational wave intensity on the frequency are:

$$S_{\rm env}(f) = \frac{3.8(f/f_{\rm env})^{2.8}}{1+2.8(f/f_{\rm env})^{3.8}}, \tag{4.6}$$

$$S_{\rm sw}(f) = \left(\frac{f}{f_{\rm sw}}\right)^3 \left(\frac{7}{4+3(f/f_{\rm sw})^2}\right)^{7/2}, \tag{4.7}$$

$$S_{\rm turb}(f) = \frac{(f/f_{\rm turb})^3}{(1+(f/f_{\rm turb}))^{11/3}(1+8\pi f/h_*)} \tag{4.8}$$

with

$$h_* = \left(\frac{T}{100\ {\rm GeV}}\right)\left(\frac{g_*}{100}\right)^{1/5} 1.65\times 10^{-5}\ {\rm Hz}\,. \tag{4.9}$$

In these expressions the frequencies $f_{\rm env}$, $f_{\rm sw}$ and $f_{\rm turb}$ are the frequencies at which each of the functions peaked. In general we expect them to be of order $\sim \beta$ (inverse duration of the phase transition) and further redshifted with h_*, given by eq. (4.9). Taking all this into account we get the expected peak frequencies:

$$f_{env} = \left(\frac{\beta}{H}\right)\left(\frac{T}{100\ {\rm GeV}}\right)\left(\frac{g_*}{100}\right)^{1/6} 3.8\times 10^{-6}\ {\rm Hz} \tag{4.10}$$

$$f_{\rm sw} = \left(\frac{\beta}{H}\right)\left(\frac{T}{100\ {\rm GeV}}\right)\left(\frac{g_*}{100}\right)^{1/6} 1.9\times 10^{-5}\ {\rm Hz} \tag{4.11}$$

$$f_{\rm turb} = \left(\frac{\beta}{H}\right)\left(\frac{T}{100\ {\rm GeV}}\right)\left(\frac{g_*}{100}\right)^{1/6} 2.7\times 10^{-5}\ {\rm Hz} \tag{4.12}$$

Finally the parameters $\kappa_{env,v,turb}$ stand for the fraction of the vacuum energy converted into the gradient energy of the dark higgs field, bulk motion of the fluid and into the MHD turbulence respectively. These efficiencies for the bulk motion and turbulence in $v_w \sim 1$ limit are given by

$$\kappa_v = \alpha(0.73+0.083\sqrt{\alpha}+\alpha)^{-1}, \quad \kappa_{\rm turb} = \epsilon\kappa_v\,, \tag{4.13}$$

where ϵ is of order $0.05\ldots 0.1$ based on the numerical simulations. The efficiency for the gradient energy of the higgs is

$$\kappa_{\rm env} = \frac{\alpha-\alpha_\infty}{\alpha} \quad \text{where} \quad \alpha_\infty = \frac{30}{24\pi^2}\frac{\sum_a c_a(m_a^2(f(T))-m_a^2(0)}{g_*(T)T}\,, \tag{4.14}$$

with $c_a = N$ (number of degrees of freedom) for the bosons and $c_a = N/2$ for the fermions. Note also that condition of [35] for the runaway bubbles that we have already mentioned in section 2 is exactly $\alpha > \alpha_\infty$, namely that non-negligible amount of energy density is damped into the gradient energy of the higgs field.

Because full analysis of the parameter space of the toy model is well beyond the scope of our work, we do not study here, which parts of the parameter space can be explicitly probed by eLISA (given, also, the uncertainties in future eLISA sensitivity). But to illustrate the point we plot on figure 7 the signal expected spectrum of one of the points in the parameter space, that we have already analyzed explicitly in subsection 3.1.

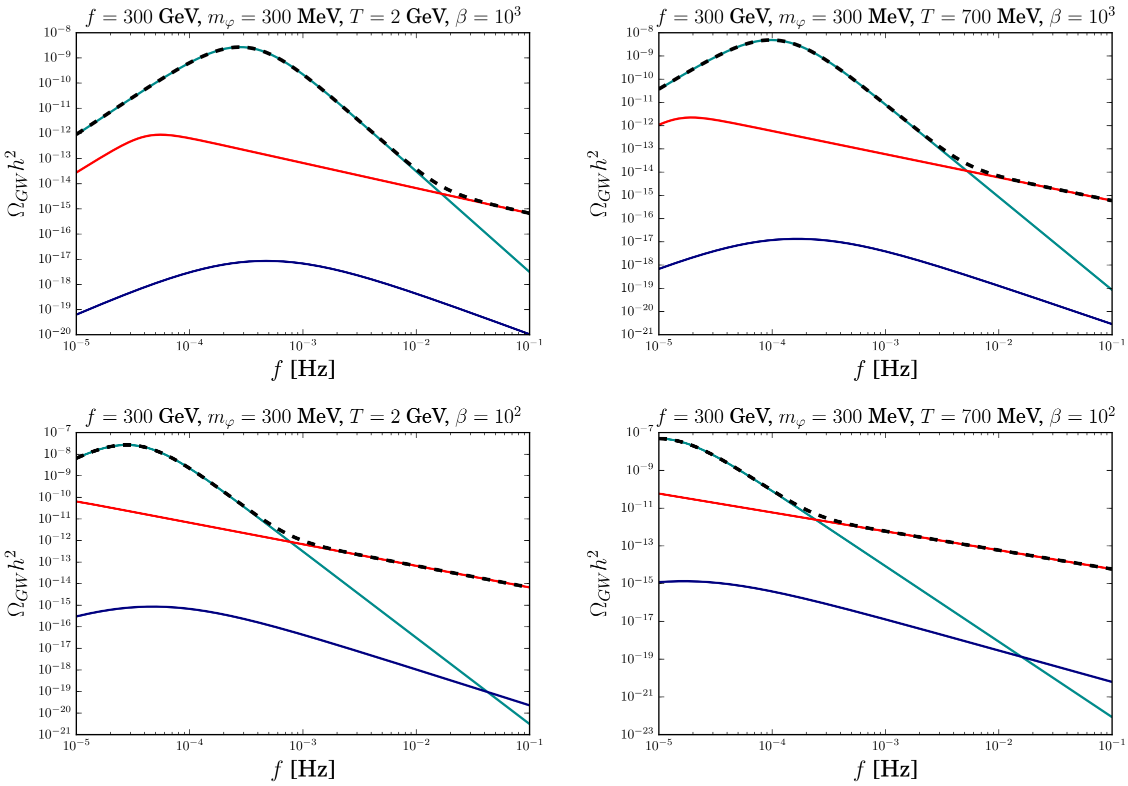

Figure 7. Gravitational wave spectrum of one particular point point in parameter space of a model, described in section 3. On the left panels we assume $T_{\rm nuc} = 2\,\text{GeV}$, just below the bound on the nucleation temperature that we derive in subsection 3.1. On the right panel we assume $T_{\rm nuc} = 0.7\,\text{GeV}$, significantly below the upper bound. On the upper panels we show the spectrum for relatively high β, and lower its value on the lower panel. The dark blue line stands for the MHD turbulence contribution, the red line stands fro the direct contribution from bubble collisions $\Omega_\varphi h^2$ and the cyan line stands for the sound waves contribution. The total intensity is indicated by the dashed black line. Everywhere $\epsilon = 0.05$ is assumed.

Note that all these points that we illustrate on figure 7 can be well within the reach of eLISA, at least if we rely on the old LISA configuration.[11] Not surprisingly, the intensity of the signal is peaked at very low frequency and the highest intensity point might be inaccessible to eLISA. This is especially true for low nucleation temperature and low β, such that the intensity distributions are peaked at frequencies as low as 10 μHz, where eLISA almost looses its sensitivity. However, due to very big energy release in the PT in this region, the signal from the tail is strong enough to yield $\Omega_{\rm GW} h^2 \sim (10^{-13} \ldots 10^{-12})$ for frequencies of order 10 mHz, where the eLISA sensitivity is maximized. The α gets values which are as high as 500 for $T \approx 0.7\,\text{GeV}$. The most optimistic eLISA sensitivity projected sensitivity will be able to probe in this region $\Omega_{\rm GW} h^2$ as low as 10^{-14}, while its sensitivity drops to values as low as $\Omega_{\rm GW} h^2 \sim 10^{-10}$ for frequencies of order 10 μHz. The contribution to the GW intensity is heavily dominated by the sounds waves in the low frequency regions and by the direct bubble collisions contribution at higher frequencies, while the MHD turbulence is always negligibly small.

[11]This corresponds to 6 links, 5 million km arm length and duration of 5 years. For more details on possible configurations of eLISA and the expected noise the reader is referred to [63] and especially to [64].

4.2 Brief comments on other possible signatures

It is tempting to conclude, based on the interactions (3.18) that the model that we have introduced, and largely the entire mechanism, inevitably predict $n - \bar{n}$ oscillations and the respective bounds apply. Indeed, by taking the first term of the equation (3.18) on its face value and integrating out the heavy sterile fermion e_α one gets

$$\mathcal{L} \supset \frac{\lambda_{\alpha ijk}\lambda^{\alpha}_{lmn} u^c_i d^c_j d^c_k u^c_l d^c_m d^c_n}{m_e \Lambda^4} \tag{4.15}$$

which looks like a perfect $n - \bar{n}$ oscillations operator. Nonetheless, the color contraction in this expression forces antisymmetrization on the flavor indices in the down-sector. Namely, at least two of the down-type quarks must be either s or b-quarks. We get a qualitatively different behavior if we consider the second operator in the (3.18), namely the $QQ(d^c)^\dagger$. Here we must antisymmetrize both on the color and the SU(2) indices of the Q fields, and therefore there is no problem with getting a non-zero coefficient for the same generations fields.

Although the $n - \bar{n}$ operator does not form at the leading order, if only the coupling λ is invoked, with democratic flavor structure of the coupling and sizable $\lambda_{\alpha 112}$ one does get a double nucleon decay $pp \to KK$. The lifetime of this process is constrained to be bigger than 1.7×10^{32} years [65], which was inferred from non-observation of ^{16}O nucleus decay into the $^{14}CK^+K^+$. Rephrasing the calculation of [66] in our variables we get

$$\tau \approx \frac{\pi}{8}\frac{\Lambda^8 m_e^2 m_N^2}{\rho \Lambda_{\rm QCD}^{10}}\,, \tag{4.16}$$

where $\rho \approx 1.91 \times 10^{-3}$ GeV is a nuclear matter density and Λ is the suppression scale from eq. (3.18) assuming that $\lambda = 1$. In fact, one gets precisely the same expression for the lifetime if $\eta = 1$ is assumed. This translates to the bound in the range of PeV scale, and in fact it is the most stringent bound that one can obtain even in the case of the unsuppressed $n - \bar{n}$ oscillations rate. However, as we discussed in section 3, it is not difficult to satisfy this bound by taking the coupling of the exotic colored scalar to the d^c to be of order $\mathcal{O}(10^{-3})$ or smaller for $m_\Delta \sim$ TeV. In lots of senses prediction of the dinucleon decay is an interesting feature of a model, that we presented, but the exact scale unfortunately depends on the model details and cannot be predicted firmly, also because the exact flavor structure of the new operators are not known, and the couplings to the light generations can in principle be highly suppressed.

Let us now again consider the $u^c d^c d^c$ coupling. The contribution to the $n-\bar{n}$ oscillations is of course not identically zero, and, if the coupling to the strange quark is non-zero, one gets it from the strange quark component of the neutron, which is of order 1%. In order to estimate this constraint we rely on the estimations of [49]. Although the model of baryogenesis that that reference proposes is very different, the fermion sector shares some clear similarities with ours and the estimation of $n - \bar{n}$ oscillation rate goes along the same lines. With the bound on the neutron oscillation time $\tau_n > 1.3 \times 10^8$ years [67] and assuming that $\lambda_{\alpha 112} \sim \mathcal{O}(1)$ we get

$$m_e \Lambda^4 \gtrsim 10^{19}\ \mathrm{GeV}^5\ . \tag{4.17}$$

For $m_e \sim 1\,$TeV, this translated to a very minor bound $\Lambda \gtrsim 10\,$TeV, very subdominant to the double nucleon decay.

In case of the $QQ(d^c)^\dagger$ coupling, assuming that one have $\eta_{111} \sim \mathcal{O}(1)$, one gets a bound from the $n-\bar{n}$ oscillations in the vicinity of 100 TeV. Note though that in this case it does not directly compare to the dinucleon decay, because different flavor operators are responsible for these processes.

Finally, let us make a brief comment about the collider phenomenology. Although the mechanism makes absolutely no firm prediction in this area, one can still expect the the bosonic operators in the hidden valley couple to the SM higgs operator $|H|^2$, because it is the lowest dimension SM gauge invariant scalar operator. This can lead to the rare exotic higgs decays, and intriguing possibility, which has been studied both in theoretical and experimental context [68]. In the model that we presented here such an operator is Eq (3.15). Although in our case the invisible rate is way too small to be observed at the LHC, it might still be visible in the future leptonic and hadronic colliders. And, not unlikely, in other models which take advantage of our mechanism, the rate can be even bigger.

Another potentially interesting aspect of the collider phenomenology has to do with the production and decay of the colored scalars Δ. Of course the existence of these particles is strongly model-dependent, however, if exist, they can be abundantly produced by the LHC. Even though these particles are colored, the LHC potential to discover them is modest, mainly because they decay into pair of anonymous jets, similar to the RPV stops in SUSY with the baryon number violation. The constraints on these particles are exactly the same as the constraints on the BNV stops. The ATLAS search [69] bounds the mass of this particle to be above $\gtrsim 320\,\text{GeV}$ in the case that one of the resulting jets is b-tagged, and puts no constraints in case the decay proceeds into the light jets. The companion CMS search [70] is also looking for the resonances in a decay channel, where one of the final state jets is heavy, and they constrain such particles in the mass window $200\ \text{GeV} < m_\Delta < 380\ \text{GeV}$. In the most general case, if Δ decays into a pair of light jets only CDF bound applies [71] $m_\Delta > 100\,\text{GeV}$. In general, it is a notoriously hard search, mostly due to the trigger issues and even after the 13 TeV run the bound is not expected to exceed 600 GeV [72].

5 Conclusions and outlook

In this paper we have introduced a novel mechanism to produce the Universe baryonic asymmetry via out-of-equilibrium decays of heavy particles which are produced non-thermally. The mechanism can work at the temperatures that are much lower than the masses of the decaying particles and therefore there is conceptually no need for the Universe to be reheated after inflation up to the masses of the decaying particles. We have also presented a simple model that serves an existence proof that the mechanism can be embedded into a realistic framework, which satisfies all the known cosmological, collider and other constraints.

Although we have concentrated on the low energy regime and post-sphaleron baryogenesis, even this part of parameter space does not necessarily provide us with robust LHC signatures. In fact, one might hope for such signatures (like long-living color triplet at the TeV scale), however they all appear to be highly model-dependent. On the other hand, the prediction of primordial gravitational waves background is model independent, because it relies on the crucial component of the mechanism: low-energy strong first order cosmological PT with the runaway bubble walls. This opens intriguing perspectives for the future eLISA interferometer, which can potentially provide first supportive evidence of this mechanism.

There are still several open questions on the theory side which should still be answered to address the feasibility of our mechanism. We have mentioned them already in the text, but it would probably be useful to summarize them here. Our mechanism relies on the possibility of efficiently producing heavy states during bubble collisions. This requires large value of the Lorentz factor γ. We have made use of the fact that such large values are potentially achievable during a PT in which bubbles runaway. The exact criterion under which runaway bubbles are present is still under debate we have decided to stay on the safe side by adopting the most conservative criterion. We have not computed ourselves the value of γ achieved during the PT, but rather indicated an upper bound for which our mechanism can work. Of course, a crucial step would be to quantitatively check if such large values of the Lorentz boost can be achieved. Also we note, that we have not analyzed this theory in the limit of resonant baryon asymmetry production. This might be interesting to do in order to understand whether lower values of γ can be relevant for our mechanism. Finally, a more thorough, most probably numerical, study will be required to estimate the role of elasticity during bubbles collisions and the efficiency of particle production.

Our proposal also opens a handful of new directions in the studies of the Universe baryon asymmetry and leaves several interesting questions open? Are there any *natural* scenarios that can accommodate the mechanism that we propose? For example, can it be confinement PT, which does not require a fine-tuned higgs mechanism? Or whether this mechanism can emerge in a hidden sector theory which is naturally higgsed, either by virtue of the hidden valley being naturally supersymmetric or because the higgs is emergent (composite higgs scenario)? Moreover, in this paper we have just brought one model, kind of existence proof that the mechanism can work, but it is still unclear how generic the mechanism is. It would be interesting to see more models along these lines, understand their parameter space, and more important see what are the experimental implications of these models.

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