

Article

Primordial Black Hole Formation Beyond the Standard Cosmic QCD Transition

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Abstract

We review the role of primordial black holes (PBHs) for illuminating the dark ages of the cosmological evolution and as dark matter (DM) candidates. We elucidate the role of phase transitions for primordial black hole formation in the early Universe and focus our attention on the cosmological QCD phase transition within a recent microscopic model. We explore the impact of physics beyond the Standard Model (SM) on the cosmic equation of state and the probability distribution for the formation of PBHs, which serve as candidates for DM and contribute to present-day binary black hole merger events.

Keywords: Big Bang; QCD transition; black holes; cosmology

1. Introduction

The search for sub-solar mass mergers is of particular interest for cosmology and astrophysics, as such compact objects hint at new physics [1–3]. On 12 November 2025, the LVK collaboration reported a candidate sub-solar mass gravitational wave merger event [4,5]. This single candidate, if confirmed, could be direct evidence that primordial black holes (PBHs) populate the sub-solar mass window—a region inaccessible to standard stellar evolution and therefore a unique fingerprint of early Universe physics [6]. However, another interpretation of this event in terms of a subsolar mass neutron star merger has been discussed as well [7]. PBHs form in the radiation-dominated era when sufficiently large overdensities cross the particle horizon [8]. Unlike other relics, they carry direct information on the Universe before Big Bang Nucleosynthesis (BBN): their mass spectrum is a fossil of the cosmic equation of state (EoS) at their formation [9–12]. The PBH distribution from the thermal history is constrained by many observations [4,6,13–16]. Any modification of the thermal history therefore leaves a direct imprint on the PBH mass distribution and can mitigate the constraints.

The Big Bang scenario describes the early Universe as an extremely dense plasma of fundamental particles, cooling and expanding on a radiation-dominated background. As it cools, a succession of phase transitions (PT) occurs. It is of primary importance to map the nuclear matter phase diagram to understand how the Universe evolved through this period and what kind of PT remnants can be found in observations [17–22]. While the SM predicts smooth crossovers for both the QCD and electroweak (EW) transitions, the motivation for at least one strong first-order PT remains compelling: a strong first-order



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PT can leave imprints in the Stochastic Gravitational Wave Background (SGWB) through bubble nucleation [19,23], and its order has direct consequences for baryogenesis [24,25]. The Sakharov conditions [26] require CP violation, departure from thermal equilibrium, and baryon number violation—conditions that a first-order PT naturally satisfies. The smooth EW crossover suggested by the observed Higgs mass $m_H = 125$ GeV [27] seems to disfavor EW baryogenesis in the SM [28,29], yet the research on beyond-standard-model (BSM) mechanisms remains very active [30–33].

Asymmetries in the lepton sector could also be powering first-order transitions. Schwarz & Stuke [34] showed that an accurate description of the early Universe requires accounting for lepton flavor asymmetries, which generate cosmic trajectories in the QCD phase diagram. A large primordial lepton asymmetry of the Universe (LAU) can suppress sphaleron rates, reconciling a large LAU with the observed small baryon asymmetry (BAU) [35,36], and can trigger a first-order QCD transition with a distinct SGWB signature [37–39]. Crucially, large lepton asymmetries before BBN are observationally allowed, as neutrino flavor oscillations starting at $T \sim 10$ MeV relax them to values consistent with CMB and BBN constraints [40,41]. Bödeker et al. demonstrated that these asymmetries reshape the PBH spectrum in a way potentially consistent with LVK observations [42].

In this work we extend this framework using a microscopic QCD model [43], incorporating BAU and LAU through state of the art Taylor-expanded susceptibilities up to 4th order, following [44]. We compute cosmic trajectories and the EoS across a range of baryon and lepton asymmetry scenarios and translate the results into PBH mass distributions, discussing their implications for gravitational wave and microlensing observations.

2. The Thermal History After the Big Bang

In this section, we develop the cosmic EoS across different transitions. We start by presenting the standard model picture, before discussing effects of Beyond-SM physics.

2.1. The Standard Model

The cosmological equation of state is described through SM interactions, assuming that the laws of physics are rigorously the same as those observed on Earth at present time. The Big Bang scenario describes the early Universe as an extremely dense plasma of fundamental particles in thermal equilibrium on an expanding background. The rate of particle interactions being much larger than the Hubble expansion rate [45], the primordial plasma can be described solely through thermodynamic equilibrium. In its early stage, the “cosmic soup” is radiation-dominated (This study focuses on the radiation-dominated era; later times or the early matter-dominated era are outside the scope of the current study.) and cools down with time according to [46]:

$$t = \sqrt{\frac{90\hbar^3 c^5}{32\pi^3 G g_\epsilon(T)}} (k_B T^{-2}) \approx \frac{2.4 \text{ s}}{\sqrt{g_\epsilon(T)}} T_{\text{MeV}}^{-2} \quad (1)$$

where $g_\epsilon(T) = g_\epsilon = \epsilon/(T^4 \pi^2/30)$ is the effective number of relativistic degrees of freedom for the energy density ϵ and $T_{\text{MeV}} = T/\text{MeV}$. As it cools down, transitions to different phases of matter occur; let us make a brief summary. After the Planck time $t_{\text{Planck}} = 10^{-43}$ s corresponding to a temperature $T_{\text{Planck}} \sim 10^{19}$ GeV, when all four fundamental interactions are unified, the Universe is minuscule, and quantum fluctuations are non-negligible. After t_{Planck} , gravity decouples and the Grand Unification Era begins. It lasts until the strong interaction decouples at $t_{\text{strong}} = 10^{-36}$ s with $T_{\text{strong}} = 10^{16}$ GeV, which triggers the inflation era. According to Guth’s inflationary scenario [47], if the Universe undergoes a first-order phase transition with sufficient supercooling during this early epoch, it can enter

a period of exponential expansion driven by the energy density of a false vacuum state. This inflationary phase addresses two fundamental problems of the standard Big Bang model: the horizon problem and the flatness problem. The exponential expansion during inflation stretches quantum fluctuations, filling the Universe with patches of inhomogeneity in energy density and curvature.

The decay of the inflation field marks the start of the quark-gluon plasma (QGP) phase. The deconfined quarks and ambient gluons provide the bulk of the contribution to the thermodynamics until the temperature drops down to $T_c \approx 156.5$ MeV [48], when a smooth crossover to the hadron resonance gas (HRG) occurs as quarks get confined into hadrons that (except protons) subsequently decay. Such energy regimes can be explored experimentally through large-scale projects such as the LHC [18]. In heavy ion collisions, the fireball formed can be similar to the very early Universe in terms of temperature regime, but it is significantly different when looking at the baryon's net number density and chemical potential. Also, the lifetime of the experimental fireball is much smaller than the Big Bang time scale. Therefore, one cannot straightforwardly apply LHC results to the Big Bang model. It is of primary importance to properly map the QCD phase diagram for cosmology and particle physics experiments; see [20,21] for a recent overview. The preferred way to explore the nuclear matter phase diagram in the Big Bang regime is through lattice QCD simulations. In this study, we rather use the generalized Beth–Uhlenbeck (GBU) approach to the EOS of strongly interacting matter, which we will refer to as the microscopical model.

According to the GBU approach, hadrons appear as correlations in the strong interaction among quarks, which are described by phase shifts with a characteristic jump by π for a bound state and a smooth tail describing the continuum of scattering states. These phase shifts in different hadronic channels are subject to in-medium modifications such as the chiral symmetry restoration (χ SR) and bound state dissociation (Mott effect). When one reverses the cosmic evolution so that the temperature increases and passes T_c , the chiral condensate melts, and the quark masses drop toward their current-quark mass values in the χ SR crossover transition. As a consequence, the continuum edges for two-, three-, and multi-quark states are dramatically lowered and “eat up” the hadronic bound states, which, upon further temperature increases, disappear as fading resonances in the continuum of few-quark scattering states. At the same time, when the quarks lose most of their mass upon χ SR, they appear as dominant degrees of freedom together with the gluons, which are set free by the simultaneous breaking of the center symmetry of the SU(3) color charge of QCD. This aspect of the QCD transition is captured in the behavior of the traced Polyakov loop and its potential. With the emergence of deconfined quarks and gluons in their QGP state, perturbative quark-gluon scattering processes also play an important role and have been included in the microscopical model, which provides an excellent description of QCD thermodynamics in accordance with *ab initio* lattice QCD simulations. For a detailed description and further results of the approach, see [43]. An advantage of the microscopic approach over that of the lattice QCD simulation is its flexibility when applications to cosmology are considered. We can straightforwardly extend the temperature range of the calculation, add photons, leptons, further quark flavors and hadron species, as well as extend the application to nonzero chemical potentials of baryon number and different flavors.

Once the pressure P as thermodynamic potential of the grand canonical ensemble is known, thermodynamic relations can be used to obtain the entropy density s , the number densities n_i and the energy density ε :

$$s = \frac{dP}{dT}, \quad (2)$$

$$n_i = \frac{dP}{d\mu_i}, \quad (3)$$

$$\varepsilon = Ts - P + \sum_i \mu_i n_i. \quad (4)$$

where the temperature T and the set of chemical potentials $\{\mu_i\}$ are the free parameters characterizing the thermodynamic state of the system. The thermal equilibrium allows us to write any summable thermodynamic quantity marked with X as

$$X_{\text{cosmic}}^{(N_f)} = X_{\text{QGP}}^{(N_f)} + X_\gamma + X_\nu + X_{e,\mu,\tau} + X_{\text{bosons}} \quad (5)$$

where the subscripts and their corresponding particle species are self-explanatory. The QCD sector encapsulates deconfined quark and gluon interactions. N_f denotes the number of quark flavors considered; our base data from the microscopic model contains u, d, s quarks. Conveniently, at the QCD transition, heavier quarks have almost fully vanished. Recent studies show that heavy charmed hadrons/mesons are formed at the QCD transition [49,50]. However, since the charm quark abundance is already negligible at this temperature scale, the added contribution to the total thermodynamics is negligible. For consistency, we still added the data found in [50] to the mix, without any significant effect, as expected. Every particle mass was set according to [27]. In higher-temperature regimes, the inclusion of heavy quarks is required for a realistic picture. It is tempting to straightforwardly add a gas of free quarks to the mix on the grounds that the QCD sector at these temperatures is already very close to an ideal gas of free quarks; however, such an approach cannot account for the interaction between the heavy quarks and the gluons. At tree level, we can evaluate the (2+1+1) configuration with (2+1) as follows (from Supplementary Information of [17]):

$$\frac{P^{(2+1+1)}(T)}{P^{(2+1)}(T)} = \frac{\sigma_{\text{SB}}(3) + P_{\text{charm}}(T)}{\sigma_{\text{SB}}(3)}, \quad (6)$$

where $\sigma_{\text{SB}}(N_f)$ is the Stefan–Boltzmann limit for N_f quark flavors considered, and $P_{\text{charm}}(T)$ is the pressure of a gas of free quarks. Note that the Stefan–Boltzmann limit should not only take into account quarks but also gluons, i.e., a non-interacting gas of massive quarks and eight massless gluons. A similar treatment to that shown in Equation (6) can be applied to extend the description of the QCD sector by adding bottom and top quarks (see Appendix A).

Between the start and decay of the QGP phase lies another transition: at $T_{\text{EW}} \approx 100\text{--}200$ GeV, the electromagnetic and weak interactions split, referred to as the electroweak phase transition (EWPT). The Higgs vacuum expectation value v is also set at this time; the observed mass of the Higgs boson $m_H = 125$ GeV is thought to be high enough to produce a smooth crossover rather than a proper phase transition per se [28,29]. In other words, v evolves continuously from the symmetric phase $v = 0$ to the broken symmetry phase with $v = 246$ GeV. In the symmetric phase, the cosmic soup is effectively a mix of massless particles. Such a smooth crossover prevents the formation of cosmological relics like domain walls or cosmic strings, which can leave distinct imprints in the SGWB [23]; see [19] for a review.

Phase transitions, whether first- or second-order or a crossover, leave distinct imprints in the EoS. A deviation from the pure radiation value $w = P/\varepsilon = 1/3$ occurs there, see Figure 1. Here, we corrected the thermodynamics up to the bottom quarks. Given the mass of the top quark and the fact that quarks heavier than the strange do not significantly contribute to the QCD transition, ignoring their contribution in the microscopical temperature range $T \in [1:1300]$ MeV is realistic. We see that the QCD transition is the most prominent feature of the EoS in Figure 1. At the confinement of quarks into hadrons and mesons, the

entropy and energy density drop, as well as the pressure, though less steeply, creating a dip in the EoS parameter w . The same parameter rises to the radiation value $1/3$ after the transition as the leptons and photons become the dominant source of pressure. Then, another dip appears at $T \approx 50$ MeV, corresponding to the annihilation of the newly formed pions.

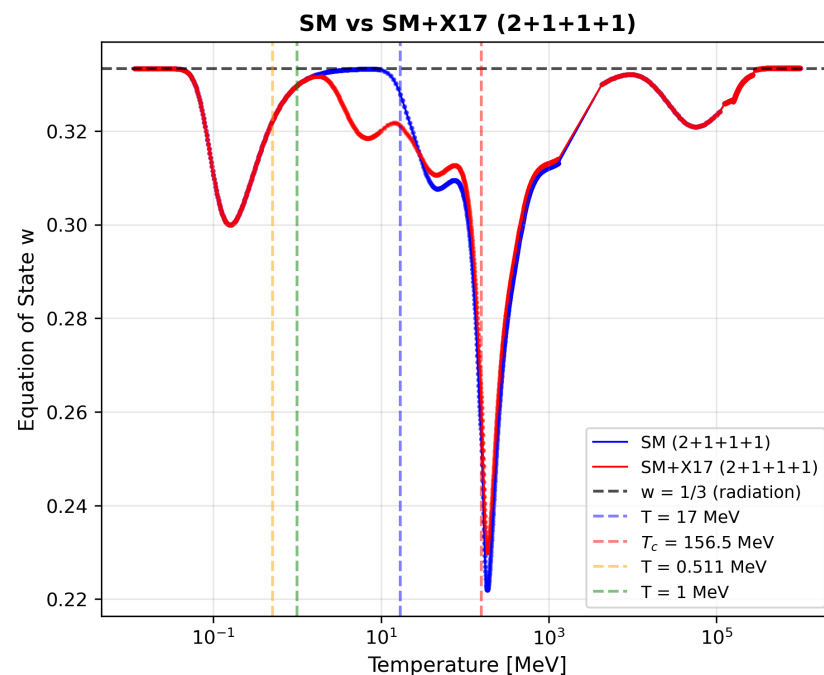


Figure 1. Cosmic EoS tree-level corrected up to the bottom quark, Equation (6), for the SM (blue) and SM+X17 case (red). For $T > 1300$ MeV, at the EWPT scale, we used data from Laine & Meyer [51]. For $T < 10$ MeV, we used the NUDEC BSM code [52,53] to model ν decoupling. Dashed vertical lines denote key temperatures.

Once the QGP has vanished, the other species contributions start to be significant; see Figure 2. Neutrinos and photons are added as a relativistic ideal gas of massless fermions and bosons. For charged leptons, we use a non-degenerate ideal gas of massive particles approximating the Fermi–Dirac integral with Bessel functions; see Appendix of [12] for details. The bump in lepton contributions just after the QCD transition in Figure 2 is due to the muons, which have not yet annihilated. Their mass of $m_\mu = 105$ MeV is close to the critical QCD temperature T_c , so that when the QGP has almost fully decayed, the charged lepton gas is the main pressure contribution for a short time. Then, the muons decay and the charged lepton pressure drops, carried from this point on only by electrons.

At $T \approx 10$ MeV, the neutrinos start oscillating, the individual lepton flavor asymmetries mix with each other according to the mixing angle [40], and are fully decoupled at $T \approx 1$ MeV. Eventually, the electron contribution starts to vanish as electrons annihilate with the positrons at $T \approx 0.511$ MeV, heating up the cosmic components still coupled with photons, mainly hadrons. This succession of processes also leaves a signature in the EoS parameter, modeled here using the publicly available code NUDEC BSM [52,53].

2.2. Beyond the Standard Model

In the previous section, we described the Universe in its most standard way. The picture is consistent with (most) particle physics observations but does not address some fundamental cosmological conundrums. Unknown or overlooked BSM physics could be the culprit. Let us start from the most grounded observations by mentioning “anomalies” in the SM. Experimental research on BSM physics is a particularly active field; so far, many “anomalies” (most often in the form of resonances) have been found in nuclear and

particle physics experiments. The extension in particle physics can occur in the sectors of fermions (spin 1/2), scalar bosons (spin 0), and vector bosons (spin 1). For a recent overview of anomalies in particle physics, see Crivellin & Mellado [54]. In a previous study, we looked at the impact of the putative X17 boson [55,56] on the cosmic EoS and the PBH distribution around the QCD transition, see [12]. In Figure 1, we updated the previous study by including e^+e^- annihilation and the standard crossover at the EW scale. This picture remains close to the standard case, as we do not include new physics per se, but simply add another non-interacting boson to the thermodynamics. It acts as a proof of principle of how particle physics “anomalies” can leave a cosmological imprint.

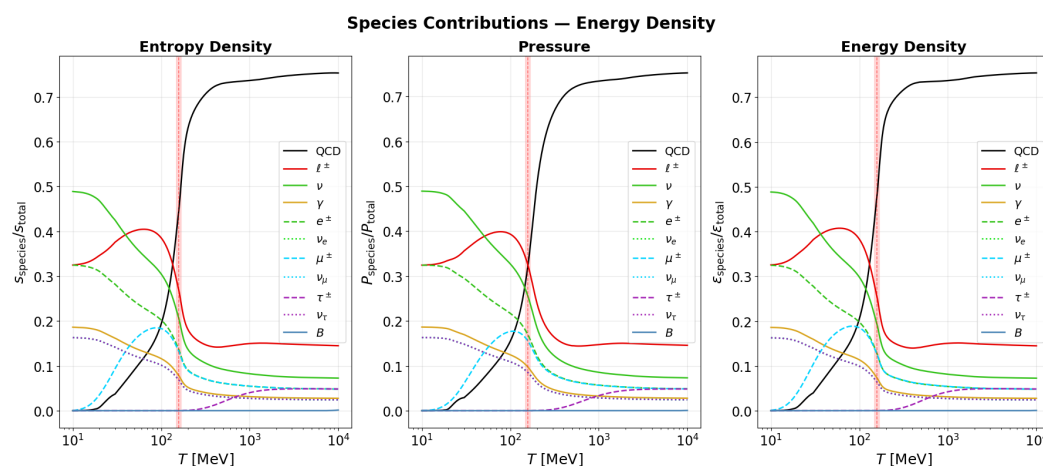


Figure 2. Partial thermodynamic contributions from different particle species and flavors in the standard case “Std b , std ℓ ”. While the model introduces asymmetries, their small values ($b = 8.6 \times 10^{-11}$, $\ell_e = \ell_\mu = \ell_\tau = 5.3 \times 10^{-11}$) make it effectively identical to the blue curve shown in Figure 1. The red shaded region highlights the QCD transition pseudo-critical temperature $T_c = 156.5$ MeV. The different species contributions can be read in the legend; “ B ” denotes bosons. All of the neutrinos contributions, with dotted lines, are the same and the lines overlap.

Turning now to cosmic mysteries, we already mentioned the need for a baryogenesis mechanism able to match the observed baryon asymmetry $b = n_B/s = 8.7 \times 10^{-11}$ inferred from Planck [57]. In the SM, b is supposedly linked to the lepton asymmetry ℓ through the Sphaleron process, allowing violation of baryon B and lepton L numbers while conserving $B - L$. We are not concerned here with the underlying mechanism leading to baryo- or leptogenesis; we refer the interested reader to the reviews and references therein [24,25]. Nevertheless, we want to examine the thermodynamic impact of an asymmetric Universe in baryon and lepton number, i.e., accounting for chemical potentials $\mu_B, \mu_Q, \mu_{L_e}, \mu_{L_\mu}, \mu_{L_\tau}$ —respectively, the baryonic, electric charge, electronic, muonic, and tauonic sectors. Efficient sphaleron processes would yield $\ell = -51/28 \times b$ [58]. Such a scenario would remain very close to the picture described in Section 2.1 and the blue curve of Figure 1, as the asymmetries b and ℓ are very small and the induced chemical potentials do not significantly impact the thermodynamics. It was pointed out in Ref. [34] that an accurate description of the early Universe requires accounting for those asymmetries, as their introduction along with conservation of the associated charges (see Equations (7)–(9) below) generates *cosmic trajectories* in the QCD phase diagram; see Section 4,

$$\ell_\alpha s = n_\alpha + n_{\nu_\alpha} = n_{L_\alpha}, \quad (7)$$

$$bs = \sum_i B_i n_i = n_B, \quad (8)$$

$$qs = \sum_i Q_i n_i = n_Q \quad (9)$$

where n_B is the net baryon number density, n_Q is the net charge number density, n_{L_α} is the net lepton number density, ℓ_α is the lepton flavor asymmetry for the flavor $\alpha = e, \mu, \tau$, and s is the total entropy density from Equation (5). We stress that all the number densities are net: $n_\alpha = n_{\alpha-} - n_{\alpha+}$, and similarly for the neutrinos.

Many ideas and BSM processes have been proposed to generate a baryon asymmetry, often able to generate values above observations. For instance, the Affleck–Dine (AD) mechanism assumes the possibility of inhomogeneous baryogenesis, i.e., there could be patches of the Universe with a high baryon number [59,60]. The AD mechanism also predicts the possibility of a globally high baryon number, which could trigger a 1st-order QCD transition. Understanding how the PT occurs in these regimes is necessary, as unexpected processes could evade the constraints on the baryon asymmetry. Boeckel & Schaffner–Bielich suggested a “little inflation” era (or supercooling) at the QCD crossing for $\mu_B/T > 1$ followed by reheating [30,31]. We want to study how these models could evolve thermodynamically, and to do so, we use QCD thermodynamics with $\mu_B/T = 1, 2$ from [43].

Another pathway for an early Universe with high baryon chemical potential, which does not require a large baryon asymmetry, is through lepton flavor asymmetries, which recent analyses find to be up to $\mathcal{O}(10^{-2})$ [61–65]. The AD mechanism could also be the culprit for such a value of the LAU [66]. The idea is straightforward: large values of lepton asymmetries before BBN are hard to constrain, as neutrino oscillations starting at $T = 10$ MeV can change the lepton flavor asymmetries to values fulfilling the constraints [40]. Some values of lepton flavor asymmetries are preferred by CMB and BBN [41,67]. It should be noted that in the references above only neutrino chemical potentials are taken into account in the calculations, i.e. neutrinos only carries the asymmetries. The present calculations, based on Equation (7), distribute the asymmetry between both charged leptons and neutrinos, resulting in non-vanishing chemical potentials for the former; their contributions are subdominant at $T < 10$ MeV, which makes the CMB and BBN constraints still relevant, even though not exact. Only a direct measurement of the infamously hard-to-measure Cosmic Neutrino Background could shed light on the MeV era. None of the models presented in this study are favored by neutrino oscillations; the aim of the present paper is to examine how the LAU affects the thermal history and to link it to PBH formation in the manner of Bödeker et al. [42]. An exhaustive study of the preferred LAU configuration is planned.

3. Methods

After presenting the thermal history after the Big Bang in various contexts, we now turn to how to solve the conservation equations. In order to solve Equations (7)–(9), the dependence of the entropy density on the baryonic and electric charge chemical potentials must be evaluated. To do so, following the procedure introduced in Refs. [34,44,68,69], we rely on a Taylor expansion over the tree-level correction up to the charm quark. We used state-of-the-art QCD pressure derivatives, called susceptibilities, up to 4th order from [50,70]. The lepton sector was solved using the Fermi–Dirac integrals with the JEL approximation [71]. Even though our approach is very similar to the latest study [44] and we used the same susceptibilities for the Taylor expansion, our base entropy density $s^{QCD}(T, \mu = 0)$ does not come from lattice QCD calculations but rather from the microscopic model [43]. This comes with the advantage of continuous trajectories, rather than 3 distinct regimes. We also extended the calculations to higher temperatures using polynomial extrapolations from Bresciani et al. [72] and interpolation from $T = 1300$ MeV to $T = 2000$ MeV. The interpolation range is chosen for a smooth connection between the two data sets; see Appendix A.

The (2+1) susceptibilities from Ref. [70] range from $T \in [30:800]$ MeV; we used the Thermal-Fist package [73] for computations below 30 MeV. In order to properly include the charm quark in the calculation, we also need the (2+1+1) susceptibilities. Using the susceptibilities from [50], one can derive the relevant susceptibilities $\chi_2^B, \chi_2^Q, \chi_{11}^{BQ}, \chi_4^B, \chi_4^Q, \chi_{13}^{BQ}, \chi_{31}^{BQ}, \chi_{22}^{BQ}$, including the charm quarks. We also extrapolate the susceptibilities from Refs. [50,70] to high temperatures. It can be seen in the figures of the mentioned references that the susceptibilities tend to $\sim 95\text{--}90\%$ of the SB limit at asymptotic temperature. We added a point at 90% of their SB limit at $T = 10$ GeV and performed a polynomial fit to ultimately extend the Taylor expansion to higher regimes. We set the extrapolation limits at 90% of their SB values, since lattice QCD results [72,74] suggest that QCD matter thermodynamics tends to $\sim 90\text{--}85\%$ of the SB limit at high temperatures. It has been demonstrated in Ref. [43] that the $\mathcal{O}(\alpha_s)$ correction to the ideal gas pressure due to quark–quark interaction via one-gluon exchange leads precisely to this 10 – 15% reduction in the SB limit. The tree-level evaluation of the charm and bottom quarks conserves this behavior at the limit; see Appendix A. We show χ_2^B as an example in Figure 3. We observe a similar behavior for the (2+1) and (2+1+1) cases, the former being a reliable pure lattice QCD calculation; the similarities between the red and blue curves show that the inclusion of the charm quark and the polynomial extrapolation are fairly reliable. The comparison with the model of an ideal gas of massive quarks and massless gluons shows how poorly this approximation describes the QCD transition. We stress that the absence of susceptibility data including heavy quarks at high temperature is one caveat of our approach. The present study should motivate the further exploration of this temperature regime.

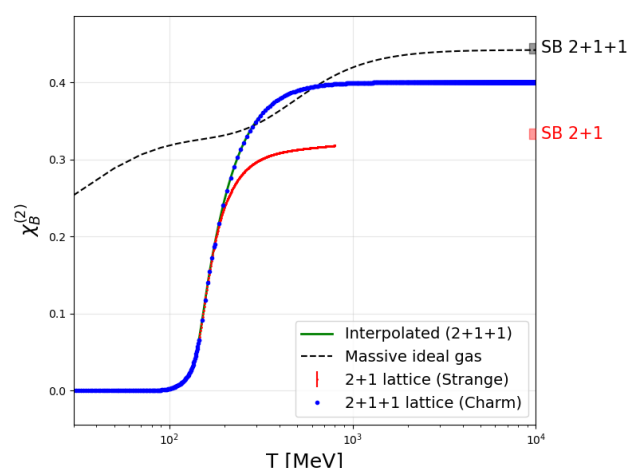


Figure 3. χ_2^B comparison between the ideal QCD sector comprising massive quarks up to the charm sector and gluons, shown as the black dashed line; the data from Ref. [70] in red; and the derived inclusion of the charm quark using Ref. [50] as blue dots. The SB limits are shown as colored squares to the right of the plot.

4. Results and Discussion

4.1. Equation of State with LAU and BAU

After having presented the equations to be solved in Section 2 and how we extended the analysis in Section 3, we now look in detail at the resulting EoS and trajectories of cosmic evolution.

The LAU triggers an increase in the absolute value of the lepton chemical potentials. Then, electric charge neutrality, Equation (9), imposes a correlation between the QCD net electric charge n_Q^{QCD} and the lepton sector net electric charge $n_Q^\ell = n_e + n_\mu + n_\tau$. It follows that the chemical potentials associated with the quarks also increase in absolute value. We can see how the LAU changes the cosmic trajectories in the nuclear matter phase diagram

in Figure 4. The dashed gray and black lines correspond to models presented in Bödeker et al. [42]. The dashed red line, also presented in the same study, corresponds to the case with lepton asymmetry associated with the sphaleron process $\ell_\alpha = -51/28 \times b$, which we call the standard case “Std b , std ℓ ”. The dash-dotted line corresponds to an extreme scenario where $b = -\ell = 0.1$ and the Taylor expansion is no longer reliable. For this dash-dotted model, we only present the massive non-interacting quark and gluon gas approximation. For easier comparison with the other models, the charm quark is the most massive quark of the model. The quarks and leptons have associated chemical potentials:

$$\mu_{up-type} = \frac{1}{3}\mu_B + \frac{2}{3}\mu_Q, \quad (10)$$

$$\mu_{down-type} = \frac{1}{3}\mu_B - \frac{1}{3}\mu_Q, \quad (11)$$

$$\mu_{\nu_\alpha} = \mu_{L_\alpha} = \mu_{\alpha^\pm} + \mu_Q \quad (12)$$

where μ_Q is the electric charge chemical potential, μ_{ν_α} (also written μ_{L_α}) corresponds to the neutrino flavor α chemical potential and $\mu_{\alpha^\pm}$ represents the charged leptons. The gluons are treated without chemical potential as an ultra-relativistic gas of massless bosons. All the dashed lines in the figures throughout this study are determined for $b = 8.6 \times 10^{-11}$ from [57], $q = 0$ consistent with constraints on the electric charge of the Universe [75], and ℓ_α arbitrarily fixed, corresponding to realistic trajectories found by solving the conservation equations Equations (7)–(9).

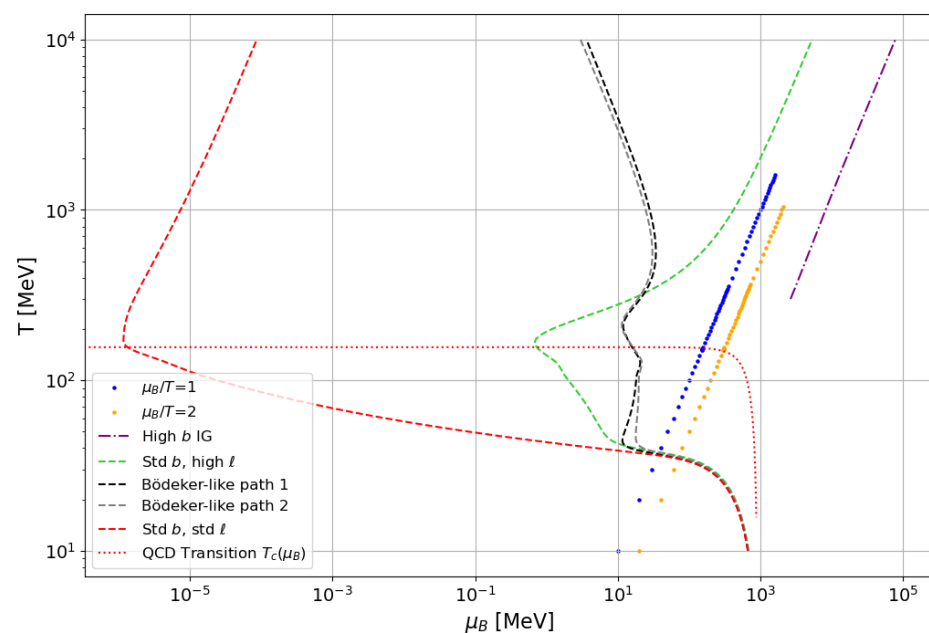


Figure 4. Cosmic trajectories for different scenarios: the dashed lines correspond to realistic scenarios. Colored dots show the trajectories with $\mu_B/T = \text{const}$ imposed. The purple dash-dotted line marked as “High b IG” corresponds to an ideal gas (IG) calculation with $b = 0.1, \ell_e = -0.1, \ell_\mu = -\ell_\tau = 0.1$; the lime green dashed line marked as “Std b , high ℓ ” corresponds to $b = 8.6 \times 10^{-11}, \ell_e = -\ell_\mu = \ell_\tau = -0.1$; the black dashed line marked as “Bödeker-like 1 path”: $b = 8.6 \times 10^{-11}, \ell_e = 0, \ell_\mu = -\ell_\tau = -0.04$; the gray dashed line marked as “Bödeker-like 2 path”: $b = 8.6 \times 10^{-11}, \ell_e = -0.08, \ell_\mu = \ell_\tau = 0.04$; the red dashed line marked as “Std b , std ℓ ”: $b = 8.6 \times 10^{-11}, \ell_e = \ell_\mu = \ell_\tau = 5.3 \times 10^{-11}$. See Appendix B for the cosmic trajectories in the T vs. $\mu_Q, \mu_{L_e}, \mu_{L_\mu}, \mu_{L_\tau}$ planes. The parametrization of $T_c(\mu_B)$ (red dotted) is taken from [76].

Since the Taylor expansion breaks down when $\mu_B/T \approx 0.1$, the extreme dashed-dotted model is computed only at high temperatures using the ideal gas approximation.

Furthermore, a pion condensate can form when $\mu_Q > m_\pi$ or $|\ell_e + \ell_\mu| \gtrsim 0.1$, which could leave an imprint in the SGWB [77–80]. The entry into and departure from the condensate could be a 1st or 2nd order PT depending on the value of $|\ell_e + \ell_\mu|$. We show the pion mass in the T vs. μ_Q plane in the Appendix B Figure A2. None of the realistic trajectories presented form a pion condensate. The cases with $\mu_B/T = \text{const}$ also do not form a pion condensate because $\mu_Q = 0$ in the microscopic model.

Gao & Oldengott [37] used a functional QCD method to explore cosmic trajectories with high lepton flavor asymmetry. They find the signature of a 1st-order QCD transition when looking at the quark chemical potential [22,37,39]. The condition for a 1st-order PT relies on the determination of the critical end point (CEP), where the PT stops being a crossover. They find that if $|\mu_q| > 200$ MeV at $T_{\text{CEP}} = 118$ MeV the PT is 1st-order. None of the realistic cosmic trajectories fulfill this condition.

While it is clear that paths with $\mu_B/T = \text{const}$ are not realistic around the QCD transition, they remain interesting when considering AD baryogenesis or high lepton flavor asymmetries. Such trajectories could approximate the thermodynamic evolution of the models proposed by Boeckel & Schaffner-Bielich [30,31] or the High Baryon Bubbles proposed in [60]. These regions could retain a high baryon number and somehow evolve independently if they are large enough, i.e., on superhorizon scales. We want to keep the number of assumptions to a minimum, so that we do not include lepton asymmetry in the calculations of these models; we take the data from Figures 8 and 12 of [43] and insert them in Equation (5). The paths with $\mu_B/T = \text{const}$ use $(2 + 1)$ quark flavors, while the dashed lines use a tree-level correction to account for the charm quark as well. The reason lies in Equation (6). The accurate determination of P_{charm} in the presence of chemical potentials would have to take into account the electric charge chemical potential; see Equations (10)–(12). Since μ_Q depends on the lepton asymmetry and we chose to use a simplistic model, we shall remain conservative and not add any tree-level corrections. The charm quark is not thought to be very significant at the QCD transition [12,50]; the leptons, however, are, and with μ_B/T fixed, we do not compute the lepton chemical potentials—we discuss the effect of this choice in the following paragraphs. A “zig-zag” feature can be seen at $T \sim 150$ MeV for the case $\mu_B/T = 1$. The parameters of the microscopical model were tuned to unify the QGP and HRG phases at $\mu_B = 0$, where the lattice QCD is the most reliable. With non-vanishing μ_B/T , a discontinuity appears because the microscopic model treats the hadrons and quarks distinctly with a sharp cutoff at T_c . Increasing μ_B/T makes the discontinuity more significant, as seen in [43].

It is worth noting that at $T \gtrsim 500$ MeV, the slopes of the paths with $\mu_B/T = \text{const}$ are similar to those of the “High b IG” and “Std B , high ℓ ” models. This is because, with massless particles, which is a reasonable approximation away from the QCD transition, the charge conservation Equations (8) and (9) are analogous to $\mu_B/T = \text{const}$. In the Appendix of [44], they find $\mu_B/T \approx 65b - 12.5\ell$ valid for $\ell \gg b$.

The EoS of the different models is plotted in Figure 5. The dashed gray, black and red lines reproduce the Bödeker results; the dip in the EoS at the QCD transition is mitigated by the baryon chemical potential but also by the lepton chemical potentials. It is clear from the colored continuous lines that the baryon chemical potential does provide a mitigation of the softening, but the difference with the “lepton powered models” shows that the lepton sector might be even more important when it comes to creating a shallower dip at the QCD transition. We explain this through the muon lepton: in the two Bödeker-like and “Std b , high ℓ ”, ℓ_μ is large; hence, μ_{L_μ} follows, which makes the contribution of the muon sector more significant; see Figure 6. At $T = 156$ MeV, the muons have not yet decayed and can therefore still compensate for the dip created by the QCD transition.

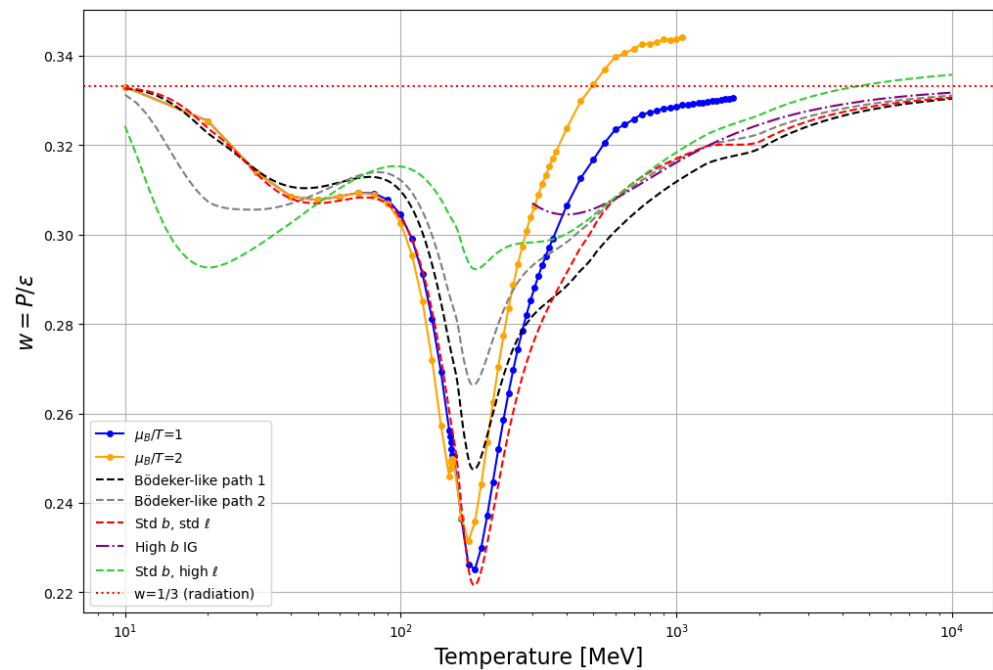


Figure 5. EoS corresponding to the different cosmic trajectories. The same color code from Figure 4 is applied. Between $T = 1300$ MeV and $T = 2000$ MeV, the thermodynamic quantities have been interpolated. Since no physics processes are modeled in this intermediate region, the equation of state parameter remains constant for the case $\mu_B = 0$. The same interpolation method is applied to both the energy density ϵ and pressure P ; consequently, w remains constant across all trajectories. For non-vanishing μ_B , the Taylor expansion makes the stalling of w less significant.

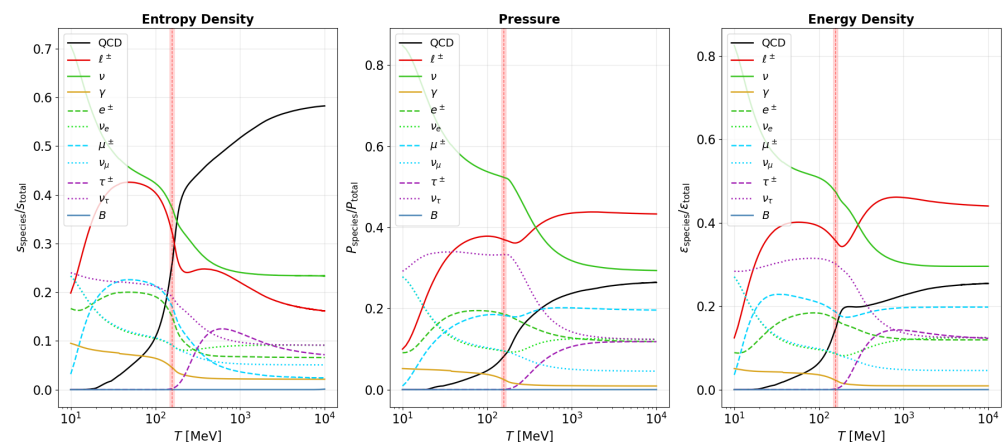


Figure 6. Partial thermodynamic contributions from different particle species and flavors in the case “Std b, high ℓ ” with high lepton flavor asymmetries. $b = 8.6 \times 10^{-11}$, $\ell_e = \ell_\mu = \ell_\tau = 5.3 \times 10^{-11}$.

In Figure 6, the effect of high ℓ_α on the ν_α contribution is clear. Looking closely at the purple lines, we see how the decay of the τ charged leptons is associated with the rise of the ν_τ as the “ τ leptonic” charge must be conserved. Comparing with Figure 2, we see that the QCD sector does not necessarily dominate the pressure or energy density at high T in the presence of LAU. Furthermore, the behavior of the different species contributions is not identical between different thermodynamic quantities. In “Std b, std ℓ ”, charged leptons dominate shortly after the QCD transition; with the “Std b high ℓ ” model, neutrinos are the dominant contributors to the entropy density after the QCD transition. The charged leptons might come close to the neutrino contribution, but the decay of the muon, associated with a rise in ν_μ contribution, leaves the charged lepton contributions subdominant. For the

same model, there is a sharp drop in entropy and energy density at the QCD transition, corresponding to the zig-zag in the associated cosmic trajectory Figure 4.

Increasing μ_B/T also comes with the effect of shifting the minimum of the softening to lower temperatures; this is not observed in the presence of lepton asymmetry because μ_B/T_c is small. The effect is not very significant, but it can be explained by the well-known shift of T_c to a lower value as μ_B increases [48,76], shown in Figure 4 as the red dotted line. $\mu_B/T = 1$ is close to the standard case, and the minimum is barely distinguishable. The dip to the minimum is steeper compared to the Bödeker-like and standard models, which can be partially explained by the difference in the number of quark flavors; see Figure 2 from [12]. The inclusion of heavier quarks would make the rise to $w = 1/3$ slower, because massive quarks such as the bottom are not yet ultra-relativistic.

One more notable feature of the continuous colored lines appears at high temperature: w rises above the radiation value, corresponding to a Universe “stiffer” than what can be expected with pure radiation. Such an effect is not observed in the Bödeker-like models, because the baryon chemical potentials are not high enough for departures from pure radiation at high temperatures. The purple dash-dotted line, where the QCD sector is an ideal gas of massive quarks and gluons, also sits around $w = 1/3$ at high temperature; even though it has a large chemical potential, the Fermi–Dirac distribution imposes $w = 1/3$ in any case. The lime green dashed model also shows $w > 1/3$ at high temperatures, which can be explained by the fact that this model also exhibits high chemical potential at high temperatures, triggering the departure from the ideal gas.

Since w is simply defined as the ratio of pressure over energy density, $w = P/\epsilon = 1/3$ is not a limit indicating incorrect physics. In Figure 7, we plot the pressures of the different models and their associated IG approximations, which act as upper limits. It is clear that all the models sit below the IG limits. The behavior of the energy density is similar. The purple “High b IG” shows the highest pressure, as it is also the model with the highest chemical potential. It might be surprising not to find the fixed μ_B/T models on top of the others, but we are plotting the total pressure from Equation (5), and these models do not account for lepton chemical potential; hence they might not even exhibit higher pressure values than the “Std b , std ℓ ” model. This shows how important it is to take into account the conservation equations Equations (7)–(9).

Having established the cosmic trajectories and their thermodynamic consequences, we now turn to the PBH mass distributions induced by the different LAU.

4.2. The Primordial Black Hole Mass Spectrum

In this section, we briefly reintroduce the formalism to translate the EoS into a PBH distribution. For a detailed discussion, see [10–12] and the references cited in the following paragraph. Primordial black holes are a unique probe of the early Universe; since these objects form in the radiation era, they can provide a picture of the Universe before the CMB. Notably, this was proposed as a probe for lepton flavor asymmetry in [42]. We argue here for a similar idea, computing the PBH mass spectrum from the EoS presented in Section 2 and comparing the induced distributions and their implications for future observations. Fluctuations of energy density in the very early Universe are necessary to explain the current energy distribution of the Universe. Hawking and Carr introduced the idea that some of these early fluctuations could have been large enough to collapse to PBHs when the particle horizon crosses the fluctuation radius at t_{cross} [8]. Although early energy density fluctuations may have many possible origins, the main scenario is cosmic inflation, which typically assumes a Gaussian fluctuation spectrum. No matter what the source of the fluctuations is, they should be larger than the Jeans length at maximum expansion if they are to undergo gravitational collapse. Recently, Carr explored different formation

scenarios [13,14], and various studies provide numerical simulations of the collapse of overdensities [11,81–84]. We refer interested readers to the 2021 review by Escrivà [85] and the references therein. These studies find that a PBH does not form exactly at the horizon mass, which has the notable effect of widening peaks in the distributions; the QCD peak, for instance, sees its width increase. For simplicity, we stay in the purely analytical formulation, as we are interested in drawing a general picture of the PBH spectrum. Collapse can occur at different epochs during the radiation-dominated era, depending on the size of the fluctuation. From Equation 2.2 of Ref. [86], the particle horizon grows as follows:

$$T \approx \sqrt{\gamma} 700 g_{\varepsilon}^{-1/4} \sqrt{M_{\odot}/M_H} \text{ MeV}. \quad (13)$$

where g_{ε} is the number of degrees of freedom of the energy density $g_{\varepsilon} = \varepsilon \frac{30}{\pi^2 T^4}$, M_H is the mass in the cosmic horizon R_H , $\gamma \in [0.2, 1]$ accounts for the critical behavior of a collapse to a PBH, i.e., reflecting the fact that PBHs are not born with exactly M_H . (Note that Equation 2.2 from Ref. [86] seems to partially incorporate the critical behavior of a PBH collapse. The Choptuik’s critical collapse law [87,88] appears more complex. Nevertheless, the equation from Ref. [86] is standard in PBH studies, as similar ones are presented in Bödeker et al. [42] or Carr et al. [15].) For simplicity and because we are mainly interested in describing tendencies, we set $\gamma = 1$. From Equation (13) $M_H \propto \gamma$, so that a change in the value of γ would slightly shift the peaks to lower masses; see [86], where they use $\gamma = 0.8$. The introduction of the γ factor is simplistic, as it is often set as a constant independent of temperature. In fact, one would expect $\gamma(T)$ to change as the matter changes phase and to depend on the actual energy content; for instance, when comparing Figure 2 and Figure 6, we see how the presence of high LAU changes the energy density and pressure content drastically. One should expect the physics inside an overdensity to be different for a Universe whose energy density is dominated by leptons compared to a QCD-dominated one.

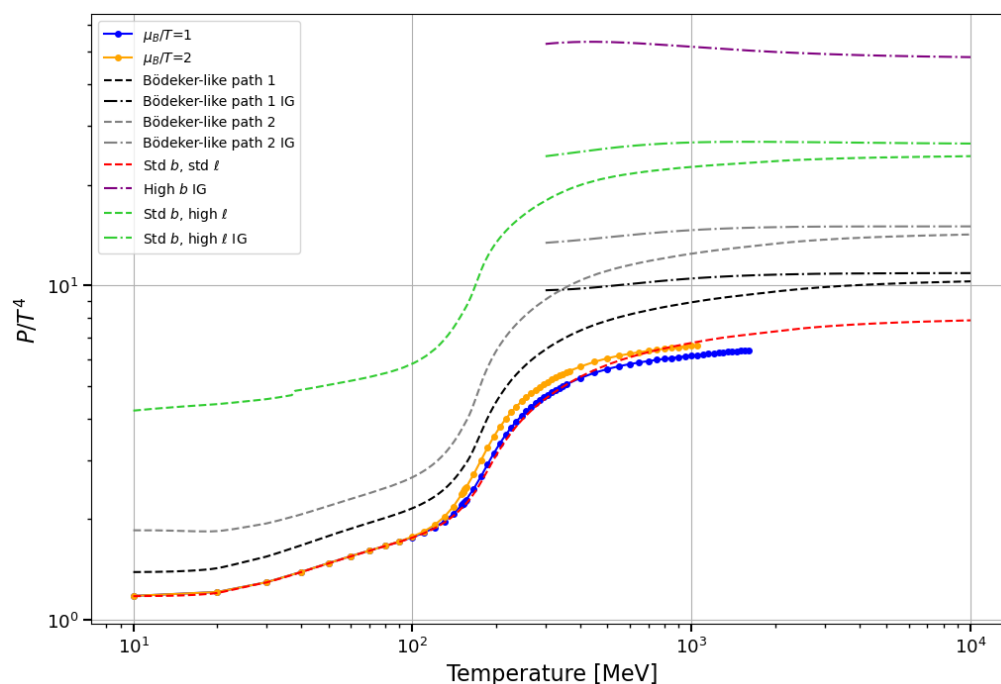


Figure 7. Normalized pressures P/T^4 , the same color code from Figure 4 is applied. The dashed-dotted lines correspond to the ideal gas cases. The slight discontinuity in “Std b, high ℓ ” around $T = 30\text{--}40$ MeV correspond to the concatenation of the microscopical model with Thermal-Fist [73].

In the past five years, new developments have emerged; different statistical methods now exist to evaluate the PBH mass spectrum through compaction function [81,89] or peak theory [90,91]. The evaluation presented here does not compete with these approaches and represents a less accurate representation of the PBH mass spectrum from the thermal history. Nevertheless, regardless of the method used, the overall impact of phase transitions on the PBH spectra is the same.

The size of candidate PBH fluctuations increases with time; a full PBH mass spectrum is possible from Planck's mass (10^{-5} g) if formed at Planck's time (10^{-43} s) to the supermassive range ($10^5 M_\odot$) for those formed as late as 1 s after the Big Bang [14]. It should be noted that Equation (13) depends on g_ϵ , which is a thermodynamic function, so the horizon mass differs slightly between the different models. For a detailed comparison between the SM and SM+X17 scenarios, see the Appendix of [12]. We follow Carr's prescription to fully describe an overdense region by its energy density contrast:

$$\delta = \frac{\epsilon - \epsilon_b}{\epsilon_b}, \quad (14)$$

where ϵ is the energy density in the region, and ϵ_b is the background energy density. Then, the fate of an overdensity is determined by the relation between δ and the threshold δ_c [9]. If $\delta > \delta_c$ at t_{cross} , a PBH is formed; if not, the overdensity is eventually dispersed away by the pressure. Assuming Gaussian fluctuations, the fraction of the Universe collapsing is

$$\beta(M) \approx \text{erfc} \left[\frac{\delta_c(w(T(M)))}{\sqrt{2}\delta_{\text{rms}}(M)} \right], \quad (15)$$

where M is the PBH mass, erfc is the complementary error function and δ_{rms} is the root mean square amplitude of the Gaussian fluctuation; $\delta_c(w(T(M)))$ is taken from [83]. Following [15,42],

$$\delta_{\text{rms}} = A \times (M/M_\odot)^{(1-n_s)/4}, \quad (16)$$

where $n_s = 0.97$ is the spectral index taken at its CMB value [57]. On PBH scales, there is actually more liberty on the shape of δ_{rms} and the value of n_s ; for instance, n_s could be running [86]. The amplitude A is a normalization parameter that expresses the strength of the fluctuations.

The present fraction of DM in a PBH of mass M is then

$$\frac{df_{\text{PBH}}(M)}{d \ln M} \approx 2.4 \beta(M) \sqrt{M_{\text{eq}}/M}, \quad (17)$$

where M_{eq} is the horizon mass at matter–radiation equality. The numerical factor originates from $2.4 = 2(1 + \Omega_b/\Omega_{\text{CDM}})$, with $\Omega_b = 0.0456$ and $\Omega_{\text{CDM}} = 0.245$ being the baryon and CDM density parameters from [57].

$$f_{\text{PBH}} \equiv \int_{M_{\text{min}}}^{M_{\text{max}}} \frac{df_{\text{PBH}}}{d \ln M} d \ln M. \quad (18)$$

The dips in the EoS are of particular interest for PBH formation, as they can be understood as a softening of the Universe: they increase the chance that a given overdense region collapses into a PBH. In Figure 8, we show the PBH mass spectrum for the SM and SM+X17, including now the super-massive black hole (SMBH) production at e^+e^- annihilation, as well as the sub-solar peak at the EWPT. The logarithmic scaling can be deceiving as the SM+X17 seems to produce many more PBHs outside of the QCD peak. The slight mitigation of the QCD softening from X17 does not drastically modify the mass spectrum, but it does show an impact on the entire PBH spectrum. Although we do not

show the PBH spectrum for μ_B/T fixed—since the absence of lepton chemical potential makes the behavior of the QCD transition unreliable—we can predict that the QCD peak would be mitigated, with the rest of the distribution increasing its contribution accordingly. The case would not be as straightforward, since $\mu_B/T = 2$ indicates a stiff EoS prior to the QCD transition; production in this era would be harder, and PBH production would be mitigated. The increased production post-QCD transition seen for SM+X17, “Bödeker-like” and “Std b , high ℓ ”, could have significant observational effects; we discuss this mass range extensively in [12] Sections 3 and 4.

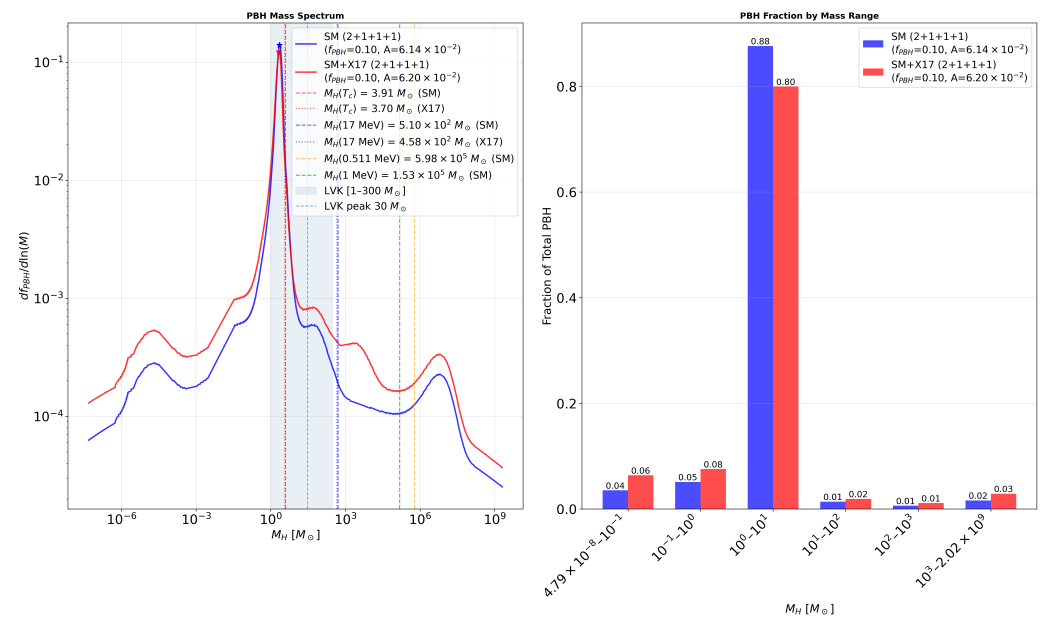


Figure 8. Comparison of SM and SM+X17 PBH mass spectra without chemical potentials. The legend gives f_{PBH} , the fraction of DM in PBHs, and A , the normalization amplitude of δ_{rms} . The asterisks denote the maximum of the distribution; the color code of the vertical lines is the same as in the Figure 1. Note that the vertical lines corresponding horizon mass $M_H(T_c)$ and $M_H(17 \text{ MeV})$ for the SM (dashed) and SM+X17 (dotted) nearly perfectly overlaps.

The standard case, Bödeker-like models and “Std b , high ℓ ” are shown in Figure 9. In Figure 3 of Bödeker et al. [42], similar spectra are plotted, but the normalization of f_{PBH} is not consistent between the models. Here, we plot them again with $f_{\text{PBH}} = 0.1$. Again, the mitigation of the QCD peak triggers increased production in other eras, similarly to the SM+X17 case. “Std b , high ℓ ” might be the most notable one, as it shifts the largest peak to a smaller mass. Looking at the right panel Figure 9, the region $1 - 10 M_\odot$ does not represent the bulk of the PBH population for “Std b , high ℓ ”. The formation of a dome pre-QCD transition at $M_H \sim 2 \times 10^{-1} M_\odot$ is caused by the decay of the τ lepton, leaving a deep imprint in the thermal history; see Figure 6. For models carrying a large tau lepton chemical potential, $\tau^+ \tau^-$ annihilation helps to soften the EoS pre-QCD transition, but the onset of the ν_τ contribution to carry ℓ_τ mitigates the QCD drop and hence the $M_{\text{PBH}} \sim 1 M_\odot$ production. The neutrinos dominate P and ε after the QCD transition, so the behavior of the EoS and PBH production comes from the leptonic sector. At $T \sim 100 \text{ MeV} \sim m_\mu$, the $\mu^+ \mu^-$ annihilate and completely disappeared by $T \sim 20 \text{ MeV}$, creating another peak at $M_H \sim 200 M_\odot$. The number density of ν_μ increases, and the neutrinos become the dominant thermodynamic contribution.

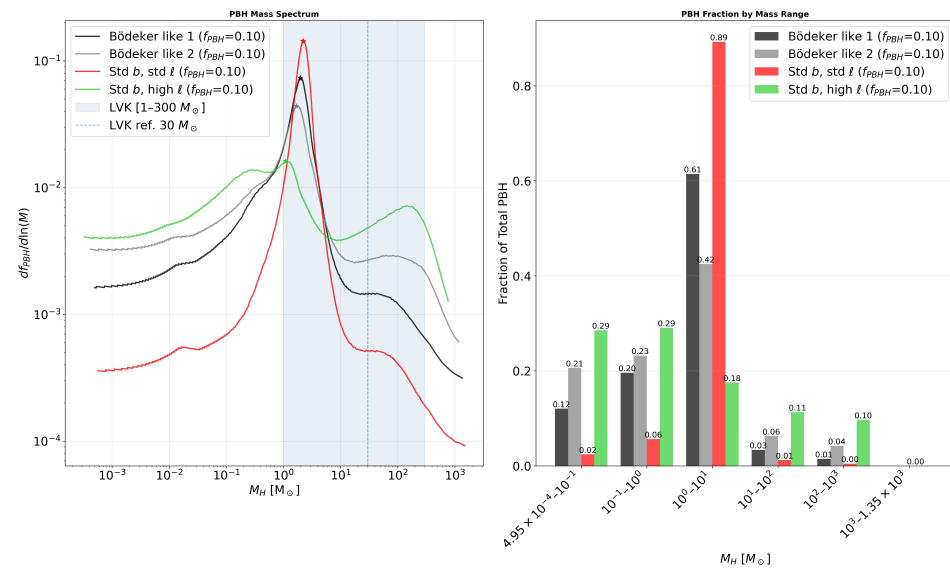


Figure 9. Bödeker-like PBH spectra, “Std b, std ℓ ” and the “Std b, high ℓ ”. Using the same color code as Figures 4, 5 and 7. The shaded region shows the observations of GWTC-4 [92], including the analysis from Ruiz-Rocha et al. [93] up to $M_{\text{BH}} \sim 300 M_{\odot}$. The asterisks denote the maximum of the distribution.

As mentioned in Section 2, a pion condensate can form in the presence of high lepton asymmetry, with impact on the PBH distribution [78]. Furthermore, cosmological relics from first-order transitions could also collapse into PBHs or interact with them [94–96]. PBHs have also been proposed as a probe for a first-order EWPT [97].

4.3. Discussing the Constraints and Positive Evidence

After showing how the PBH mass spectrum is modified in the presence of LAU and BSM anomalies like X17, we discuss how these results compare with current observations. We should start by stating that the aim of the present paper is not to constrain the models presented, nor to make quantitative predictions for future observations. We are mainly interested in describing the EoS behavior around the QCD transition across a diverse variety of models. The PBH mass spectrum is another way to look at the consequences of modifications of the EoS, and would become a direct probe of the Universe pre-BBN should future observations confirm the existence of PBHs. For a thorough review of PBHs and their implications, see Ref. [4] from the LISA theory group, or more generally, the work of B. Carr [6,13].

The stiffening of the EoS pre-QCD transition associated with sub-solar mass PBH formation has implications for microlensing observations. Although it is hard to draw firm conclusions because different collaborations claim different results and constraints, it remains worth mentioning, as it is an active field of research. Moreover, the constraints rely on assumptions that are not fully secured. For instance, the spatial distribution of PBHs could have a significant impact on the constraints; see the recent review by Green [98]. Hawkins in Ref. [99] discussed how the galactic rotation curve can bias microlensing observations. The current discussion, in a nutshell, is as follows: in the late 1990s and early 2000s, the MACHO collaboration reported observations consistent with a significant fraction of the Milky Way halo being composed of sub-solar compact objects [100,101]. EROS [102] and OGLE [103,104] were not able to corroborate the results. Hawkins & Garcia-Bellido [105,106] showed how using the Milky Way rotation curve inferred by GAIA can mitigate the constraints from OGLE. To add to the debate, the Subaru telescope can also be used to search for microlensing events in Andromeda; in their 2017 analysis, they found strong constraints [107] on the mass range $M_{\text{PBH}} \sim 10^{-11} - 10^{-5} M_{\odot}$. A

2026 analysis [108] mitigates this original claim, explicitly stating that it originated from a better treatment of the data. Another way to constrain this mass range is through PBH–star interactions [109,110].

The sub-solar mass range is particularly interesting: on 12 November 2025, the LVK collaboration reported a candidate sub-solar mass merger event [4,5]. It was quickly followed by an analysis finding the merger rate of such an event to be compatible with a slightly modified version of “Bödeker-like 2” (the spectrum they present is tilted compared to the one from this study and Bödeker’s, which typically appears when modifying the value of n_s in Equation (16)) [6]; see also [16]. PBH formation and the peaks associated with thermal history can leave hints in the SGWB [86,111]. In a broader sense, primordial fluctuations are also sources of the SGWB with associated bounds on the fluctuation power spectrum [112,113]; see Ref. [114] for a review.

Stellar mass PBH production has direct implications for GW observations; see Refs. [4,5] for a recent review. As shown by Bödeker et al. in Ref. [42], lepton asymmetry can be a convenient way to make GW observations agree with the predicted PBH mass distribution from thermal history. We showed the current LVK observation mass range in Figures 8 and 9. In Figure 10, we show how the probability density of PBH mergers compares to the actual events from GWTC-4 [92]; to model the PBH merger rate, we assume only late binaries [115]. Comparing with Figures 4–6 from Bödeker et al. [42], it is clear how the lepton asymmetry can shift the peak in the probability density of PBH mergers closer to the observed aggregate in the $q - M_B$ plane, where q is the mass ratio of the binary, and M_B is the heaviest binary component

On the more massive part of the spectrum, Ref. [116] showed that μ distortions (not related to muons nor chemical potentials) in the CMB power spectrum could draw limits on the higher portion of the PBH mass function. In the same intermediate BH mass range 10^2 – $10^4 M_\odot$, Ref. [117] constrains monochromatic PBH abundance through energy deposits in the CMB during the dark ages; see [118]. It is therefore crucial to include the LAU in the PBH mass spectrum estimates and so the SMBH peaks, as it could leave a distinct imprint on the SGWB detected by future experiments [4]. In the absence of a given signal, it would constrain the primordial fluctuations power spectrum, similarly to what is currently inferred from the CMB fluctuation spectrum, for instance. A quantitative study of the SGWB induced by the scenarios presented in this study is outside the scope of the article; we stress that such calculations are of primary importance to decipher the primordial parameters of the universe. On the other hand, it is tempting to understand the newly discovered population of supermassive black holes at high z , the “little red dots”, through PBH seeds [119–121]. Several mechanisms to form the little red dots from PBHs have been considered [122]; the massive seeds $M_{\text{PBH}} > 10^4 M_\odot$ appear to be ruled out by μ -distortion, although the strength of the constraints depends on which statistical framework for PBH formation is adopted [113]. Moreover, supermassive PBHs could impact early structure formation through their seeding effect [123].

We close this section by stressing that the PBH mass spectra produced throughout this study are constrained by various observational channels. Although the results are difficult to compare directly with the extended mass functions, current observations are likely already ruling out the spectra presented in this study. We are particularly concerned with scalar-induced gravitational waves [113,114] and μ -distortion [116]. Nevertheless, our study shows that the relative abundance of different PBH masses can significantly change depending on the LAU, with implications for observations, as shown in Figure 10, which makes PBHs a compelling probe of the primordial Universe. We show that the characteristic peak from the QCD transition does not necessarily make up the bulk of the

PBH population, as seen in “Std b , high ℓ ” in Figure 9. The application of monochromatic constraints to an evenly distributed PBH mass spectrum is, to say the least, questionable.

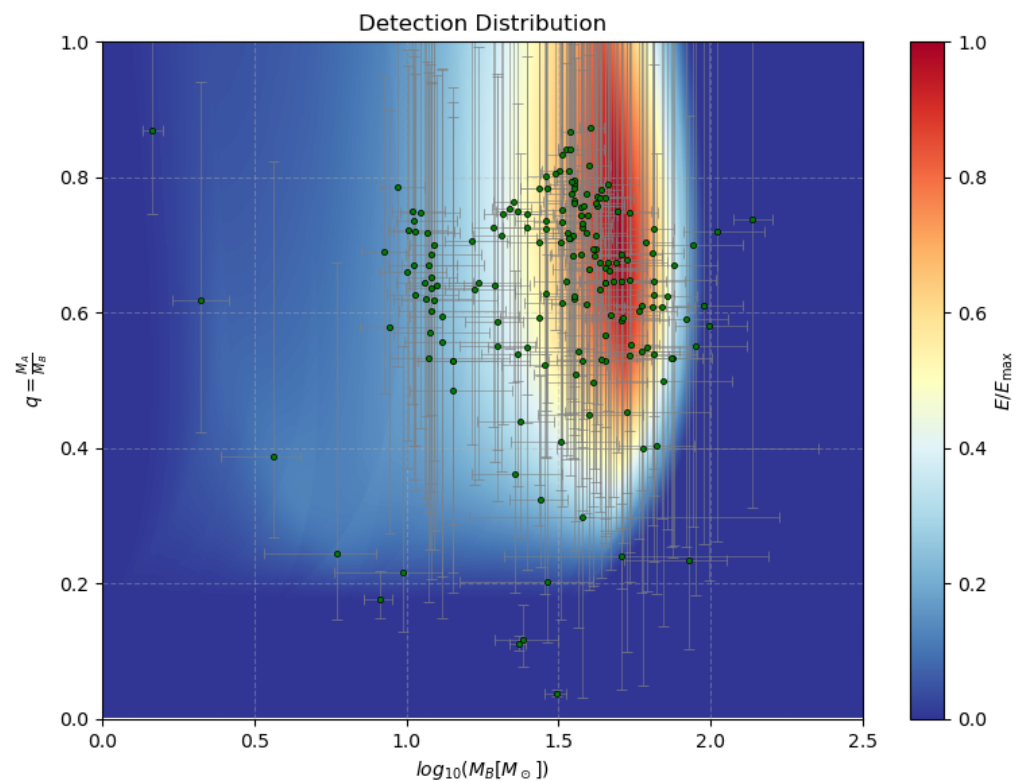


Figure 10. Probability density of PBH mergers with masses $M_A < M_B$. Inspired by Figures 4–6 from Bödeker et al. [42]. Events from GWTC-4 are shown as green dots [92]. We used O4 strain noise from [124].

5. Conclusions

In this study, we explored the cosmic equation of state across the QCD transition for a range of baryon and lepton asymmetry scenarios, extending our previous work [12] to include non-vanishing chemical potentials and to show the complete thermal history in the presence of X17 bosons. Using a microscopic QCD model [43] combined with state-of-the-art lattice QCD susceptibilities up to fourth order [44], we computed cosmic trajectories in the nuclear matter phase diagram and their thermodynamic consequences over a temperature range up to $T = 10$ GeV by including charm quark corrections. We motivate the computation of susceptibilities of heavy quarks in a cosmological context. The extension to $T < 10$ MeV is also of crucial importance, as this epoch corresponds to the formation of SMBHs (see the last peak in Figure 8 or [15]). A future study accounting for LAU during neutrino decoupling is already planned.

Using two simplistic models with high baryonic chemical potential but without LAU, we clarified the role of the baryon and lepton chemical potentials on the thermodynamics, finding the possibility of a stiff universe pre-QCD transition that has not been considered so far. Furthermore, the comparison with realistic models highlights the importance of lepton asymmetry on the EoS. Unlike the previous Taylor expansion and lattice-QCD-based approaches, which rely on distinct hadronic, quark–gluon and non-interacting quark gas regimes, the microscopic model provides continuous trajectories across the transition, offering a smoother and more consistent description of the thermodynamics. We find that the baryon chemical potential μ_B alone does not capture the full picture of the primordial plasma at the QCD transition: lepton flavor asymmetries, through electric

charge conservation, generate comparable or larger chemical potentials in the QCD sector and significantly mitigate the EoS softening. The resulting PBH mass distributions show distinct features—in particular, a pre-QCD stiffening and a modified QCD peak—that differ qualitatively from the standard scenario. The recent LVK sub-solar mass candidate [4,5] and its compatibility with lepton-asymmetry-driven PBH scenarios [6] motivate a deeper quantitative comparison, which we plan to address in a future study including a more systematic exploration of the lepton asymmetry parameter space and a tighter confrontation with gravitational wave observations.

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Conflicts of Interest: The authors claim no conflict of interest.

Appendix A. Base Thermodynamics

In Figure A1, we show different base pressures from the microscopical model and Bresciani in blue, and in orange and green, the tree-level evaluation using Equation (6). We can see how little impact the addition of heavy quarks has on the QCD transition at $T \sim 156$ MeV. The EW scale is not considered here, as it could be the baryogenesis and/or leptogenesis era, making the evolution of the thermodynamic quantities very model dependent [24]. The charge conservation Equations Equations (7)–(9) would not apply with the left-hand side as constant values. Bresciani et al. [72] and the microscopical data rely on different assumptions; a simple concatenation at $T = 1300$ MeV exhibits a discontinuity. At $T > 1300$ MeV, the microscopical model has not been compared with lattice QCD data; $T = 1300$ MeV is therefore the highest reliable temperature of our microscopical model. We tried different interpolation ranges and found out that a concatenation at $T = 2000$ MeV is the smallest temperature for which a smooth connection is possible for the three thermodynamic quantities s , ϵ and P . For consistency, we want to interpolate the three thermodynamic quantities on the exact same temperature range. The orange curve corresponds to the base data used to compute the cosmic trajectories (Sections 3 and 4). The green curve is the one used in Figures 1 and 8.

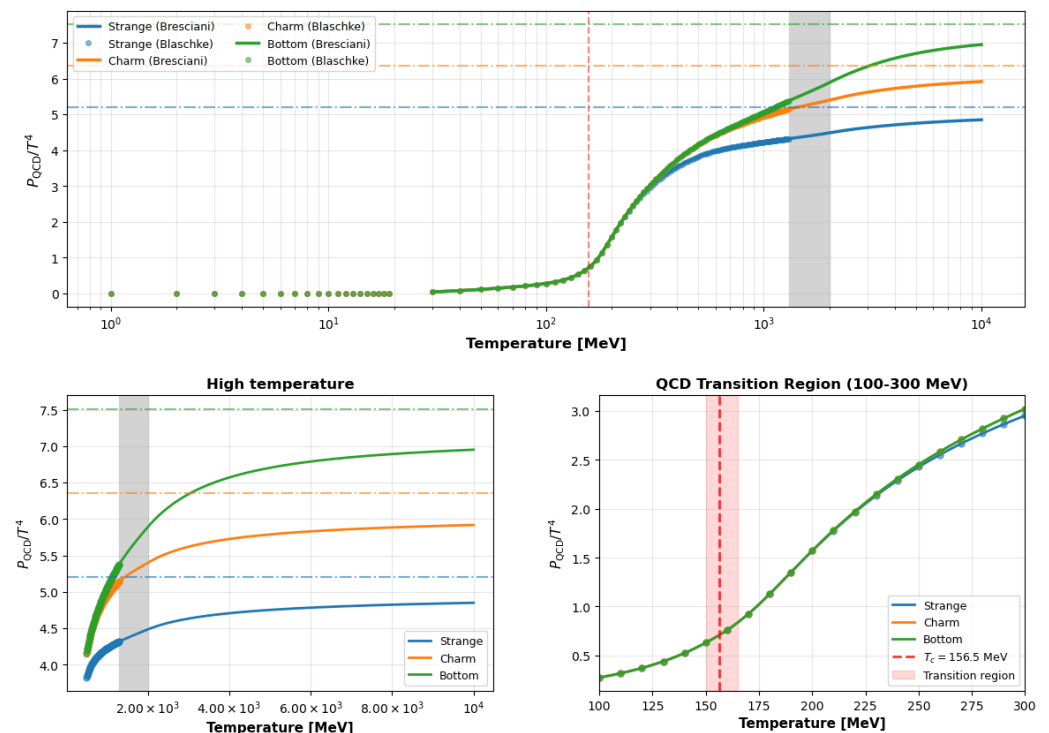


Figure A1. Base QCD pressure for different numbers of quark flavors; the legend shows the heaviest quark of each configuration. Dots represent the base pressure from Ref. [43] down to $T = 1$ MeV. The dash-dotted lines correspond to the SB limits, while the continuous lines show the data used for cosmic trajectory calculations; for $T > 2000$ MeV, we used the polynomial extrapolation from Ref. [72], and the microscopical data extends up to $T = 1300$ MeV. We rely on PCHIP interpolation, marked with the gray shaded region. In the bottom right plot, the three configurations overlaps.

Appendix B. More Cosmic Trajectories

In Section 2, we presented the equations to solve and the induced cosmic trajectories in the T vs. μ_B plane. Since the system of equations is solved in the direction of the five chemical potentials, we show in this Appendix the cosmic trajectories in the other directions. Neutrinos ν_α and their associated charged leptons α are linked according to (12). Comparing the magnitude of μ_Q Figure A3 and μ_{L_e} Figure A2, it becomes obvious that μ_e is small. The same analysis can be done for muon and tau leptons. In “Bödeker-like 2”, μ_e gains a significant value at large temperature, while for “Std b , high ℓ ”, the bounce at the QCD transition lets μ_e gain value before vanishing at low temperature. For completeness, we show the cosmic trajectories in T vs. μ_{L_μ} in Figure A4 and in T vs. μ_{L_τ} Figure A5.

We also observe the difference in behavior between “Bödeker-like 1” and “Bödeker-like 2” in the T vs. μ_{L_e} plane, where the former don’t have asymmetry in the electronic sector, and the latter does, which shows a similar behavior to “Std b , high ℓ ”.

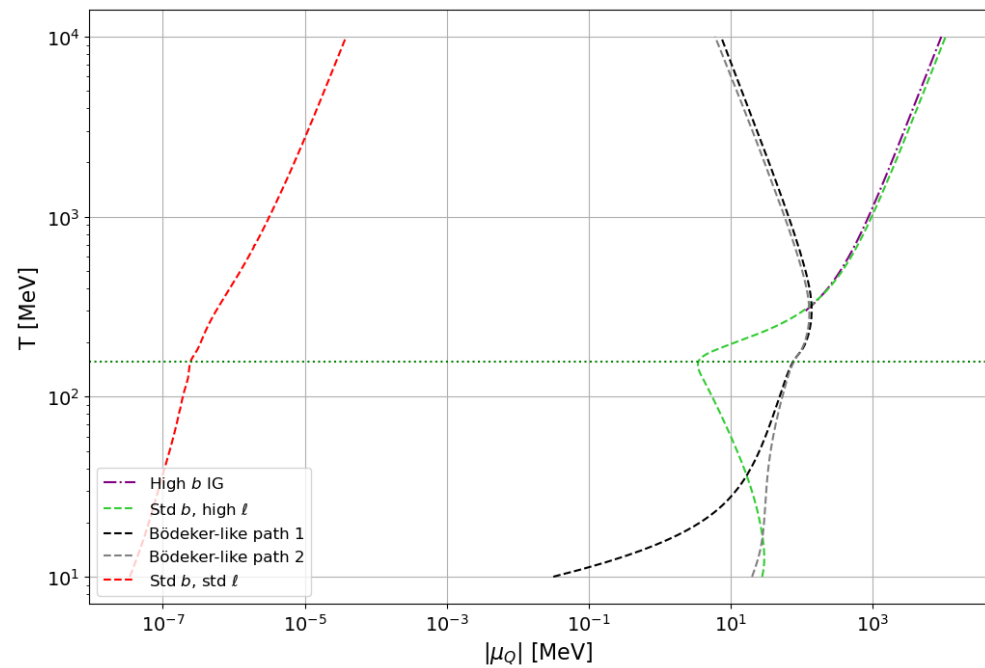


Figure A2. Trajectories in the T vs. μ_Q plane. The trajectories with μ_B/T fixed are not shown, as μ_Q is ignored. If a trajectory were on the right on μ_π , it would indicate the formation of a pion condensate.

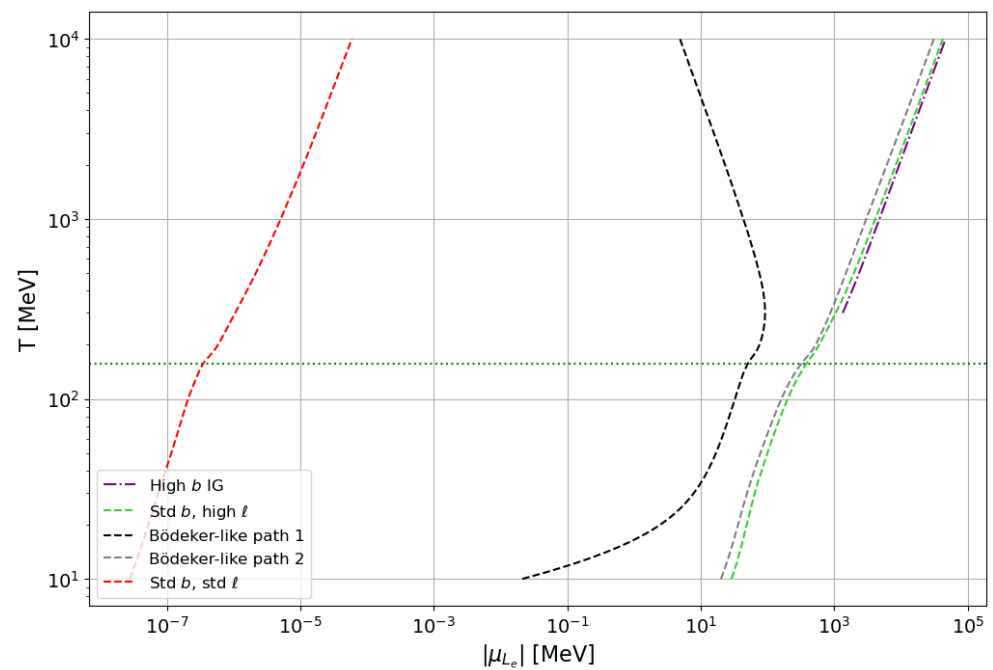


Figure A3. Trajectories in the T vs. μ_{L_e} plane.

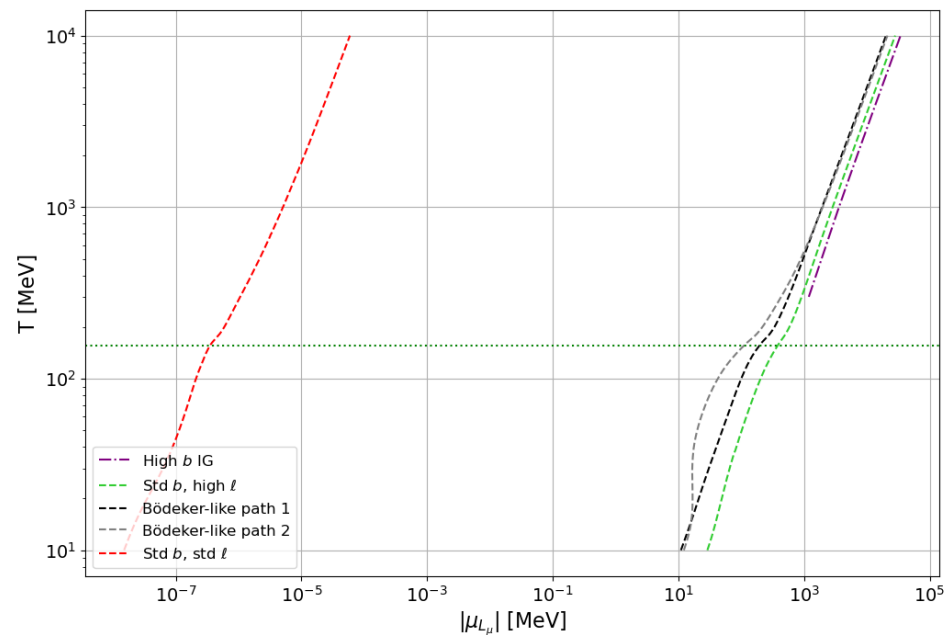


Figure A4. Trajectories in the T vs. $\mu_{L\mu}$ plane. “High b IG”, “Std b , high ℓ ” and “Bödeker-like 2” exhibit positive chemical potential as the muon asymmetry is positive; for “Bödeker-like 1” and “Std b , std ℓ ”, we show $-\mu_{L\mu}$.

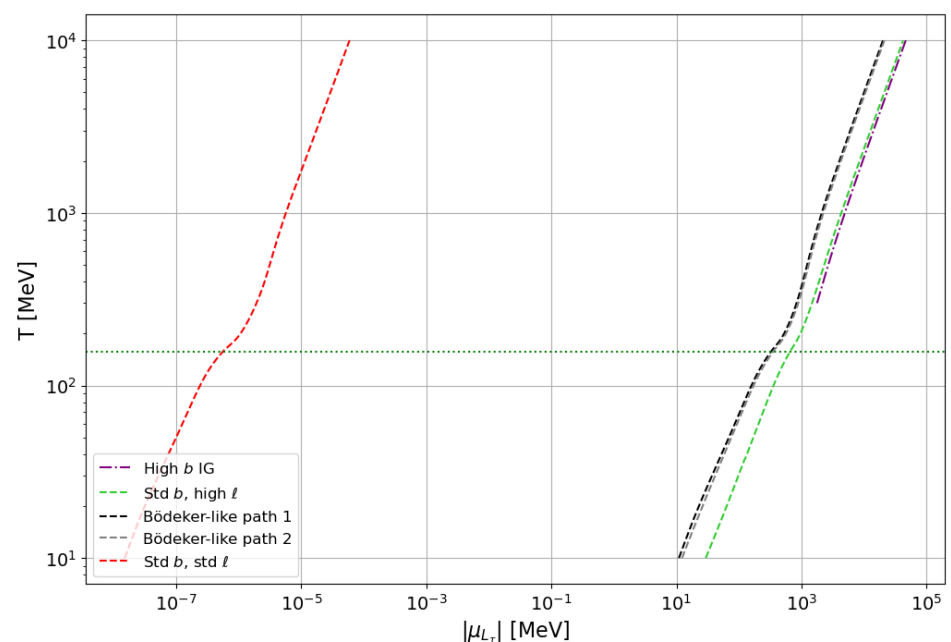


Figure A5. Trajectories in the T vs. $\mu_{L\tau}$ plane. “Bödeker-like 1” and “Bödeker-like 2” exhibit positive chemical potential as the tau asymmetry is positive; for the remaining model, we show $-\mu_{L\tau}$.

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