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Gravitational portal to Reheating, Leptogenesis and Dark Matter

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Considering graviton as the only messenger that exists between the inflaton, the dark sector and the Standard Model (SM), we show, it is possible to simultaneously generate the observed relic density of dark matter (DM), the baryon asymmetry via leptogenesis, as well as successful reheating after inflation. Assuming the inflaton ϕ oscillates in a monomial potential $V(\phi) \propto \phi^k$ at the end of inflation, reheating via minimal gravitational interaction turns out to be excluded due to excessive production of gravitational wave (GW) energy density from inflation. We thus (a) extend the minimal model by considering a non-minimal gravitational contribution to radiation and (b)) propose an explanation for the PeV excess observed by IceCube, when the DM has an explicit Yukawa coupling to the SM leptong and Higgs. We also propose a novel scenario, where the gravitational production is a two-step process, through the production of a pair of scalars, that eventually decay into fermionic final states. Our framework gives rise to detectable primordial GW signal, that falls within the sensitivity of several proposed GW detectors.

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1. Introduction

The weakly interacting massive particles (WIMPs) are one of the most popular candidates [1–4] for dark matter (DM), a non-baryonic entity that makes up around 24% of the matter-energy budget of the Universe [5]. In standard WIMP paradigm, DM particles thermalize with the baryon-photon plasma in the early Universe with a weak-scale interaction strength, and then freeze out when the interaction rate falls below the Hubble expansion rate, producing the observed relic density $\Omega_{\rm DM}h^2 \simeq 0.12$ [5]. This scenario, albeit appealing, currently being severely constrained from the non-observation of potential signature at the experimental frontier (see, e.g., Refs. [6, 7] for a review). As opposed to WIMPs, feebly interacting massive particles (FIMPs), which couple to the SM sector very feebly evading the stringent experimental bounds, can as well be viable candidates for DM [8–10]. In the early Universe, FIMPs can be generated from either the decay or annihilation of states in the visible sector. When the SM temperature falls below the typical mass scale of the interaction, the generation process becomes Boltzmann suppressed, giving rise to a constant comoving DM number density, dubbed as freeze-in [9].

In contrast to the “WIMP-miracle" which produces the observed relic abundance with weak-scale couplings and masses, a “FIMP-miracle" happens when one considers renormalizable couplings of strength $\sim \mathcal{O}(10^{-11})$, in order to ensure out-of-equilibrium production of the FIMPs. Such a feeble interaction strength occurs quite naturally when the only effective coupling between the DM and the visible sector is gravity. Indeed, the minimal irreducible interaction that should exist between DM and the Standard Model (SM) is mediated by graviton exchange [11–29]. Such an interaction can lead to the observed amount of DM through the scattering of the particles in the thermal bath or directly through the gravitational transfer of the energy stored in the inflaton condensate [23–27]. It has also been argued that the thermal bath itself may be generated from gravitational interactions [26, 30, 31]. However, reheating the Universe from graviton exchange processes alone requires a steep inflaton potential during reheating, resulting in a low reheating temperature and a massive enhancement of tensor modes after inflation. Hence, the minimal scenario of gravitational reheating is excluded by an excessive generation of dark radiation in the form of gravitational waves (GWs) during BBN, as already noted in [32]. This limitation of minimal gravitational reheating is one motivation to introduce, as a natural generalization, non-minimal couplings of fields with gravity.

As it is well-known, the visible or baryonic matter content of the Universe is asymmetric. While the SM is not capable of explaining the observed matter-antimatter asymmetry, the baryon asymmetry of the Universe (BAU) can be produced via leptogenesis [33], where, instead of creating a baryon asymmetry directly, a lepton asymmetry is generated first from the CP-violating decay of heavy right handed neutrinos, and that asymmetry gets converted into baryon asymmetry via electroweak sphaleron transitions [34]. In thermal leptogenesis [35–38], the decaying particles, typically right-handed neutrinos (RHNs), are produced thermally from the SM bath. However, the lower bound on the RHN mass in such scenarios (known as the Davidson-Ibarra bound), leads to a lower bound on the reheating temperature $T_{\rm RH} \gtrsim 10^{10}$ GeV [39] so that the RHNs can be produced from the thermal bath. This bound might be in conflict with an upper bound on the reheat temperature that applies in supersymmetric models with a “gravitino problem" [40]. One way to avoid this is to consider the non-thermal production of RHNs [41–45]. This interaction is necessarily model dependent as it depends on the Yukawa interaction between the inflaton and the RHNs.

In the present work, we have presented a simultaneous solution for the DM abundance, baryon asymmetry, and origin of the thermal bath after inflation, from purely gravitational interactions. Our scenario can be considered as the most minimal possible, since we do not introduce any new interactions for any process beyond the SM, except for gravity. The only new fields beyond the SM that are required, are the RHNs. We show, the present framework can give rise to a detectable primordial GW background with an inflationary origin, that in turn excludes the minimal gravitational reheating scenario. However, a large part of the parameter space still remain within the reach of several futuristic GW detection facilities.

2. The Set-up

On expanding the metric around the Minkowski space-time: $g_{\mu\nu} \simeq \eta_{\mu\nu} + \frac{2h_{\mu\nu}}{M_P}$, one can obtain possible gravitational interactions between the graviton field and the stress-energy tensor as [46, 47]

$$\sqrt{-g}\,\mathcal{L}_{\rm int} = -\frac{1}{M_P}\,h_{\mu\nu}\left(T^{\mu\nu}_{\rm SM} + T^{\mu\nu}_{\phi} + T^{\mu\nu}_{X}\right), \tag{1}$$

where ϕ is the inflaton and X denotes field beyond the SM. The form of the stress-energy tensor $T^{\mu\nu}_i$ depends on the spin of the field. In the present scenario we are interested in Majorana spin-1/2 fermions, for which

$$T^{\mu\nu}_{1/2} = \frac{i}{8}\left[\bar\chi\gamma^\mu\overset{\leftrightarrow}{\partial^\nu}\chi + \bar\chi\gamma^\nu\overset{\leftrightarrow}{\partial^\mu}\chi\right] - g^{\mu\nu}\left[\frac{i}{4}\bar\chi\gamma^\alpha\overset{\leftrightarrow}{\partial_\alpha}\chi - \frac{m_\chi}{2}\,\overline{\chi^c}\chi\right], \tag{2}$$

while for a scalar S,

$$T^{\mu\nu}_0 = \partial^\mu S\partial^\nu S - g^{\mu\nu}\left[\frac{1}{2}\partial^\alpha S\,\partial_\alpha S - V(S)\right]. \tag{3}$$

For the inflaton potential $V(\phi)$, we consider the class of α-attractor T-models [48]

$$V(\phi) = \lambda\,M_P^4\left[\tanh\left(\frac{\phi}{\sqrt{6\,\alpha}\,M_P}\right)\right]^k \simeq \lambda\,M_P^4 \times \begin{cases} 1 & \text{for } \phi \gg M_P, \\ \left(\frac{\phi}{\sqrt{6\,\alpha}\,M_P}\right)^k & \text{for } \phi \ll M_P\,. \end{cases} \tag{4}$$

The overall scale of the potential, parameterized by the coupling λ, can be expressed in terms of the amplitude of the scalar perturbation power spectrum $A_S \simeq (2.1\pm0.1)\times10^{-9}$ [49] as $\lambda \simeq \frac{18\,\pi^2\,\alpha\,A_S}{6^{k/2}\,N_\star^2}$, where $N_\star$ is the number of e-folds measured from the end of inflation to the time when the pivot scale $k_\star = 0.05~\text{Mpc}^{-1}$ exits the horizon. Here onward we will also fix $\alpha = 1/6$. In addition to the inflationary sector and the SM, in order to explain DM abundance and the baryon asymmetry via leptogenesis, we consider the following renormalizable interaction Lagrangian

$$\mathcal{L} \supset -\frac{1}{2}\,M_{N_i}\,\overline{N_i^c}N_i - (y_N)_{ij}\,\overline{N}_i\,\widetilde{H}^\dagger\,L_j + \text{h.c.}, \tag{5}$$

where H and L are the SM Higgs and lepton doublet respectively. We consider two cases, first, $(y_N)_{1i} = 0$ and thus N_1 is a stable DM candidate. Later, we will relax this condition and consider a metastable DM candidate, allowing it to decay into neutrinos that could be observed at IceCube. The

DM candidate N_1 can be produced during reheating from inflaton scattering $\phi\phi \to N_1 N_1$ as well as from the thermal bath (mediated by a massless graviton in both cases). On the other hand, for the generation of the baryon asymmetry, we will cater to non-thermal leptogenesis, where the RHNs $N_{2,3}$ are too weakly coupled to reach thermal equilibrium. Hence they are predominantly produced only during reheating from gravitational inflaton scattering. We thus consider the following production via graviton exchange

- $\phi\,\phi \to N_1\,N_1$, SM SM $\to N_1\,N_1$ for production of the DM candidate N_1.
- $\phi\,\phi \to N_{2,3}\,N_{2,3}$ for production of $N_{2,3}$ that will lead to non-thermal leptogenesis.
- $\phi\phi \to$ SM SM for gravitational reheating.

3. Reheating via gravity portal

The evolution of inflaton and radiation energy densities (ρ_ϕ, ρ_R respectively) can be tracked with the help of the following set of coupled Boltzmann equations (BEQ)

$$\begin{aligned}
&\frac{d\rho_\phi}{dt} + 3H\,(1+w_\phi)\,\rho_\phi = -(1+w_\phi)\,\Gamma_\phi\,\rho_\phi\,,\\
&\frac{d\rho_R}{dt} + 4H\,\rho_R = +(1+w_\phi)\,\Gamma_\phi\,\rho_\phi\,,
\end{aligned} \tag{6}$$

where $w_\phi \equiv \frac{p_\phi}{\rho_\phi} = \frac{k-2}{k+2}$ [50] is the general equation of state parameter and $\rho_R = \frac{\pi^2 g_*}{30}\,T^4 \equiv c_*\,T^4$, and g_* is the effective number of degrees of freedom for the thermal plasma at the temperature T. Since during reheating, the total energy density is dominated by the inflaton, it is possible to analytically solve Eq. (6)

$$\rho_\phi(a) = \rho_{\rm end}\left(\frac{a_{\rm end}}{a}\right)^{\frac{6k}{k+2}}\,. \tag{7}$$

For reheating, we consider the Higgs channel and expand the inflaton potential energy in terms of the Fourier modes [26, 28, 31, 50–52]

$$V(\phi) = V(\phi_0)\sum_{n=-\infty}^{\infty}\mathcal{P}_n^k e^{-in\omega t} = \rho_\phi\sum_{n=-\infty}^{\infty}\mathcal{P}_n^k e^{-in\omega t}\,, \tag{8}$$

obtaining the production rate of radiation as [26, 27, 31]

$$(1+w_\phi)\,\Gamma_\phi\,\rho_\phi = R_H^{\phi^k} \simeq \frac{N_h\rho_\phi^2}{16\pi M_P^4}\sum_{n=1}^{\infty}2n\omega|\mathcal{P}_{2n}^k|^2 = \alpha_k M_P^5\left(\frac{\rho_\phi}{M_P^4}\right)^{\frac{5k-2}{2k}}\,, \tag{9}$$

with $N_h = 4$ being the number of internal degrees of freedom for one complex Higgs doublet. The frequency of oscillations of ϕ is given by [50]

$$\omega = m_\phi\sqrt{\frac{\pi k}{2(k-1)}}\frac{\Gamma(\frac{1}{2}+\frac{1}{k})}{\Gamma(\frac{1}{k})}\,, \tag{10}$$

with $m_\phi^2 = \frac{\partial^2 V(\phi)}{\partial \phi^2}$ being the inflaton mass squared. The definition of α_k follows the analysis in [31]. The evolution of the radiation energy density, on the other hand, reads[50]

$$\rho_R(a) = \rho_{\rm RH} \left(\frac{a_{\rm RH}}{a}\right)^4 \left[\frac{1-(a_{\rm end}/a)^{\frac{8k-14}{k+2}}}{1-(a_{\rm end}/a_{\rm RH})^{\frac{8k-14}{k+2}}}\right]. \tag{11}$$

The evolution in Eq. (11) is valid when $a_{\rm end} \ll a \ll a_{\rm RH}$ where $a_{\rm end}$ marks the end of inflation (or the onset of reheating), while $a_{\rm RH}$ indicates the end of reheating defined as $\rho_\phi(a_{\rm RH}) = \rho_R(a_{\rm RH}) = \rho_{\rm RH}$. The reheating temperature can be determined by solving the Friedmann equation (6) for the radiation energy density. This yields [31]

$$\rho_R(a) \simeq \alpha_k \frac{k+2}{8k-14} \sqrt{3} M_P^4 \left(\frac{\rho_{\rm end}}{M_P^4}\right)^{\frac{2k-1}{k}} \left(\frac{a_{\rm end}}{a}\right)^4, \tag{12}$$

and evaluating this at $a_{\rm RH}$ we have

$$T_{\rm RH}^4 = \frac{30}{\pi^2 g_{\rm RH}} M_P^4 \left(\frac{\rho_{\rm end}}{M_P^4}\right)^{\frac{4k-7}{k-4}} \left(\frac{\alpha_k \sqrt{3}\,(k+2)}{8k-14}\right)^{\frac{3k}{k-4}}. \tag{13}$$

From Eq. (13) we find $T_{\rm RH} \simeq 60$ MeV for $k = 10$ and $\rho_{\rm end} \simeq (4.8 \times 10^{15}\ {\rm GeV})^4$.

The non-minimal coupling of Higgs bosons to gravity provides an additional channel to reheat the Universe through gravitational processes, with the following rate [31]

$$(1+\omega_\phi)\Gamma_\phi = R_H^{\phi,\xi} \simeq \frac{\xi_h^2 N_h}{4\pi M_P^4} \sum_{n=1}^{\infty} 2n\omega \left|2 \times \mathcal{P}_{2n}^k \rho_\phi + \frac{(n\omega)^2}{2} \phi_0^2 |Q_n|^2\right|^2 = \alpha_k^\xi M_P^5 \left(\frac{\rho_\phi}{M_P^4}\right)^{\frac{5k-2}{2k}}, \tag{14}$$

where Q_n has been defined in Eq. (23). On solving Eq. (6) for ρ_R, the reheating temperature in the presence of the non-minimal coupling turns out to be

$$\left(T_{\rm RH}^\xi\right)^4 = \frac{30}{\pi^2 g_{\rm RH}} M_P^4 \left(\frac{\rho_{\rm end}}{M_P^4}\right)^{\frac{4k-7}{k-4}} \left(\frac{\alpha_k^\xi \sqrt{3}\,(k+2)}{8k-14}\right)^{\frac{3k}{k-4}}. \tag{15}$$

The maximum temperature in this case is determined from

$$\rho_{\rm max}^\xi \simeq \sqrt{3}\,\alpha_k^\xi\, M_P^4 \left(\frac{\rho_{\rm end}}{M_P^4}\right)^{\frac{2k-1}{k}} \frac{k+2}{12k-16} \left(\frac{2k+4}{6k-3}\right)^{\frac{2k+4}{4k-7}} \equiv c_*\,(T_{\rm max}^\xi)^4. \tag{16}$$

4. Dark matter via gravity portal

The DM production rate via 2-to-2 scattering of the bath particles, mediated by graviton exchange reads [17, 24, 26, 53]

$$R_{N_i}^T = \frac{1}{2} \times \beta_{1/2} \frac{T^8}{M_P^4}, \tag{17}$$

where we have $\beta_{1/2} = 11351\pi^3/10368000 \simeq 3.4\times 10^{-2}$ [26]. The evolution of RHN number density n_{N_i} (with $i = 1, 2, 3$) is governed by the BEQ

$$\frac{dn_{N_i}}{dt} + 3\,H\,n_{N_i} = R^T_{N_i}\,, \tag{18}$$

where $H = \dot{a}/a$ is the Hubble parameter. Introducing the comoving number density $Y_{N_i} = n_{N_i}\,a^3$, we can re-cast the Boltzmann equation as

$$\frac{dY^T_{N_i}}{da} = \frac{a^2}{H}\,R^T_{N_i}\,, \tag{19}$$

where $i = 1$ for DM production. The DM number density at the end of reheating turns out to be

$$n^T_{N_i}(a_{\rm RH}) \simeq \frac{\beta_{1/2}\,(k+2)\,\rho_{\rm RH}^{\frac{3}{2}}}{12\,\sqrt{3}M_P^3\,c_*^2}\,\frac{2(7-4k)^2}{(k+5)(k-1)(5k-2)}\,r^{\frac{10+2k}{k+2}}\,. \tag{20}$$

The contribution of gravitational scattering of the particles in the primordial plasma to the DM relic abundance can then be determined using [54]

$$\Omega^T_{N_1}\,h^2 = 1.6\times 10^8\,\frac{g_0}{g_{\rm RH}}\,\frac{n^T_{N_1}(a_{\rm RH})}{T^3_{\rm RH}}\,\frac{M_{N_1}}{\rm GeV}\,, \tag{21}$$

which gives

$$\Omega^T_{N_1}\,h^2 \simeq 1.6\times 10^8 \times \frac{g_0\,\beta_{1/2}}{g_{\rm RH}} \times \frac{M_{N_1}}{\rm GeV}\,\frac{c_*^{-\frac{5}{6}-\frac{5}{3k}}\,(7-4k)^2\,(k+2)}{6\,\sqrt{3}\,(k+5)(k-1)(5k-2)}\left(\frac{T_{\rm RH}}{M_P}\right)^{\frac{5k-20}{3k}}\left(\frac{\rho_{\rm end}}{M_P^4}\right)^{\frac{k+5}{3k}}\,, \tag{22}$$

where $g_0 = g_*(T_0) = 43/11$ and $g_{\rm RH} = g_*(T_{\rm RH}) = 427/4$ are the number of relativistic degrees of freedom at present and at the end of reheating respectively.

The DM candidate N_1 can also be produced directly from inflaton scattering. The oscillating inflaton field with a time-dependent amplitude can be parametrized as

$$\phi(t) = \phi_0(t)\cdot \mathcal{Q}(t) = \phi_0(t)\sum_{n=-\infty}^{\infty}\mathcal{Q}_n e^{-in\omega t}\,, \tag{23}$$

where $\phi_0(t)$ is the time-dependent amplitude that includes the effects of redshift and $\mathcal{Q}(t)$ describes the periodicity of the oscillation. Furthermore, we assume a mass hierarchy $M_{N_{1,2,3}} < m_\phi$ such that the graviton mediated process is kinematically allowed. The production rate for N_i from inflaton scattering mediated by gravity is given by [26]

$$R^{\phi^k}_{N_i} = \frac{\rho_\phi^2}{4\pi M_P^4}\,\frac{M^2_{N_i}}{m^2_\phi}\,\Sigma^k_{N_i}\,, \tag{24}$$

where

$$\Sigma^k_{N_i} = \sum_{n=1}^{+\infty}|\mathcal{P}^k_{2n}|^2\,\frac{m^2_\phi}{E^2_{2n}}\left[1-\frac{4\,M^2_{N_i}}{E^2_{2n}}\right]^{3/2}\,, \tag{25}$$

accounts for the sum over the Fourier modes of the inflaton potential, and $m_\phi^2 = \lambda\, k\,(k-1)\,\left(\rho_\phi/(\lambda\, M_P^4)\right)^{\frac{k-2}{k}}$. Here $E_n = n\omega$ is the energy of the n-th inflaton oscillation mode. Then, the number density of RHN reads

$$\frac{dY_{N_i}^{\phi^k}}{da} = \frac{\sqrt{3}M_P}{\sqrt{\rho_{\rm RH}}}a^2\left(\frac{a}{a_{\rm RH}}\right)^{\frac{3k}{k+2}} R_{N_i}^{\phi^k}(a). \tag{26}$$

Integration of Eq. (26), leads to the following expression for the RHN density [26, 31]

$$n_{N_i}^{\phi^k}(a_{\rm RH}) \simeq \frac{M_{N_1}^2\,\sqrt{3}\,(k+2)\,\rho_{\rm RH}^{\frac{1}{2}+\frac{2}{k}}}{24\,\pi\,k(k-1)\lambda^{\frac{2}{k}}\,M_P^{1+\frac{8}{k}}}\left(\frac{\rho_{\rm end}}{\rho_{\rm RH}}\right)^{\frac{1}{k}}\,\Sigma_{N_1}^k\,, \tag{27}$$

evaluated at the end of reheating. In order to obtain the DM relic abundance, one can again follow Eq. (21), but now replacing $n_{N_1}^T(a_{\rm RH})$ with $n_{N_1}^\phi(a_{\rm RH})$, and obtain [26]

$$\begin{aligned}\frac{\Omega_{N_1}^{\phi^k}h^2}{0.12} &= \frac{\Sigma_{N_1}^k}{2.4^{\frac{8}{k}}}\frac{k+2}{k(k-1)}\left(\frac{10^{-11}}{\lambda}\right)^{\frac{2}{k}}\left(\frac{10^{40}\text{GeV}^4}{\rho_{\rm RH}}\right)^{\frac{1}{4}-\frac{1}{k}}\\ &\times\left(\frac{\rho_{\rm end}}{10^{64}\text{GeV}^4}\right)^{\frac{1}{k}}\left(\frac{M_{N_1}}{1.1\times10^{7+\frac{6}{k}}\,\text{GeV}}\right)^3.\end{aligned} \tag{28}$$

The total DM relic abundance is therefore a sum of the gravitational contribution from thermal bath ($\Omega_{N_1}^T\,h^2$) and from inflaton scattering ($\Omega_{N_1}^{\phi^k}\,h^2$).

5. Leptogenesis via gravity portal

Since N_1 is the stable DM candidate, in the present scenario the lighter of $N_{2,3}$ can undergo out-of-equilibrium decay to SM final states. The resulting CP asymmetry from the decay of N_2 is given by [40, 55–58]

$$\epsilon_{\Delta L} = \frac{\sum_\alpha[\Gamma(N_2\to l_\alpha+H)-\Gamma(N_2\to\bar{l}_\alpha+H^*)]}{\sum_\alpha[\Gamma(N_2\to l_\alpha+H)+\Gamma(N_2\to\bar{l}_\alpha+H^*)]}\,. \tag{29}$$

The CP asymmetry can be expressed as [31, 58, 59]

$$\epsilon_{\Delta L} \simeq \frac{3\delta_{\rm eff}}{16\,\pi}\,\frac{M_{N_2}\,m_{\nu,\rm max}}{v^2}\,, \tag{30}$$

where $\langle H\rangle \equiv v \approx 174$ GeV is the SM Higgs doublet vacuum expectation value, $\delta_{\rm eff}$ is the effective CP violating phase in the neutrino mass matrix with $0\le\delta_{\rm eff}\le1$, and, we take $m_{\nu,\rm max} = 0.05$ eV as the heaviest light neutrino mass. Here we are interested in non-thermal leptogenesis [41, 43–45, 60–62]. The produced lepton asymmetry is eventually converted to baryon asymmetry via electroweak sphaleron processes leading to

$$Y_B = \frac{n_B}{s} = \frac{28}{79}\epsilon_{\Delta L}\,\frac{n_{N_2}^\phi(T_{\rm RH})}{s}\,, \tag{31}$$

where $n^{\phi}_{N_2}(T_{\rm RH})$ is the number density from Eqs. (20) and (27) at the end of reheating and $s = 2\pi^2 g_{\rm RH} T_{\rm RH}^3/45$ is the entropy density. The final asymmetry then becomes

$$Y_B \simeq 3.5 \times 10^{-4}\, \delta_{\rm eff} \left(\frac{m_{\nu,\rm max}}{0.05\,\rm eV}\right) \left(\frac{M_{N_2}}{10^{13}\,\rm GeV}\right) \frac{n^{\phi}_{N_2}}{s}\bigg|_{T_{\rm RH}}, \tag{32}$$

while the observed value, as reported by Planck [63], is $Y_B^{\rm obs} \simeq 8.7 \times 10^{-11}$. The lepton asymmetry is not washed out because the lepton-number violating process involving the Yukawa scattering and the electroweak sphaleron processes are never in equilibrium at the same time.

6. Spectrum of primordial gravitational wave

The ratio of the gravitational wave (GW) energy density to that of the radiation bath is given by [64]

$$\frac{\rho_{\rm GW}}{\rho_R} = \frac{1}{32\pi G \rho_R} \frac{k_{\rm GW}^2}{2} \mathcal{P}_T(k_{\rm GW}) \quad \text{with} \quad \mathcal{P}_T(k_{\rm GW}) \equiv \frac{2H_I^2(k_{\rm GW})}{\pi^2 M_P^2}, \tag{33}$$

where $k_{\rm GW}$ is the momentum mode of the GW, $\mathcal{P}_T$ is the primordial tensor power spectrum, $H_I(k_{\rm GW})$ is the Hubble scale during inflation when the mode $k_{\rm GW}$ exits the horizon, $T_{\rm hc}$ is the horizon-crossing temperature when the mode re-enters the horizon at $k_{\rm GW} = H(T_{\rm hc})$, and the factor of $1/2$ accounts for the time average of the rapidly oscillating metric perturbations. As one can see, if horizon crossing occurs during radiation domination $k_{\rm GW}^2 = H^2(T_{\rm hc}) = \rho_R/(3M_P^2)$, then the GW spectrum becomes scale invariant. On the other hand, if horizon crossing occurs during the inflaton-dominated era, the GW strength is enhanced by a factor of ρ_ϕ/ρ_R evaluated at $T_{\rm hc}$. As a result, the largest enhancement is for the mode that re-enters the horizon right after inflation at $T_{\rm max}$. For minimal gravitational reheating ($\xi_h = 0$), the enhancement in this mode is $\rho_{\rm end}/\rho_R(T_{\rm max}) \simeq (4-6) \times 10^{13}$ for $k \in [6, 20]$. These values are excluded by the BBN bound of $\Omega_{\rm GW} h^2 \simeq 1.3 \times 10^{-6}$ [65]. Therefore, the case with minimal gravitational reheating is ruled out [32]. The constraint is relaxed when $T_{\rm max}$ is increased, e.g., by non-minimal gravitational interactions via ξ_h [cf. Eq. (16)], because the GW energy density relative to that of radiation is smaller in this case. In addition to setting a constraint, such enhanced gravitational waves offer an exciting signature to search for [66]. The amount of enhancement depends on ρ_ϕ/ρ_R at the time of horizon crossing, implying that the GW spectrum depends on k via ρ_ϕ. By analyzing modes that re-enter the horizon after $T_{\rm max}$ and using $\rho_\phi \propto a^{-6k/(k+2)}$ from Eq. (7), we find the GW spectrum scales with the frequency as $\Omega_{\rm GW} h^2 \propto f^{\frac{k-4}{k-1}}$, which is consistent with Ref. [32]. The GW frequency is obtained by redshifting the initial momentum mode at $T_{\rm hc}$ to today's photon temperature $T_{\gamma,0}$

$$f = \frac{k_{\rm GW}}{2\pi} \frac{T_{\gamma,0}}{T_{\rm hc}} \left(\frac{g_*(\rm eV)}{g_*(T_{\rm hc})}\right)^{\frac{1}{3}}, \tag{34}$$

with the maximum possible frequency being

$$f_{\rm max} = \frac{H(T_{\rm max})}{2\pi} \frac{a_{\rm end}}{a_0}, \tag{35}$$

corresponding to the modes that re-enter the horizon right at the end of inflation.

7. Results and discussion

7.1 Case of stable dark matter

In the left panel of Fig. 1, we show contours of total relic abundance $(\Omega_{N_1}^T + \Omega_{N_1}^{\phi^k})\, h^2$. We clearly see that the desired relic density $(\Omega_{N_1} = 0.12)$ is obtained *twice*: (i) at a lower reheating temperature, where inflaton scattering dominates, and (ii) for a higher reheating temperature, when we are in the thermal production regime. For $M_{N_1} > 3 \times 10^8$ GeV, there are no values of $(T_{\rm RH}, k)$ that result in an acceptable density of DM, and the allowed range in $T_{\rm RH}$ is larger with lighter DM. This is understandable as, the thermal relic requiring a *larger upper bound* on $T_{\rm RH}$ for lighter DM, while the inflaton scattering requires a *smaller lower bound* on $T_{\rm RH}$ for lighter DM. Low values of $T_{\rm RH}$ are excluded by BBN. The gray-shaded region in the lower right corner of this panel is also excluded since minimal gravitational interactions would produce a reheating temperature larger than the values in that region. For DM of masses very close to 1 PeV, there exists a viable parameter space for $k \geq 9$ (along the boundary of the excessive GWs region), requiring $\xi_h \simeq 0.5$. For larger masses, the range in k extends to lower values, and higher reheating temperatures are possible and require larger non-minimal coupling to gravity.

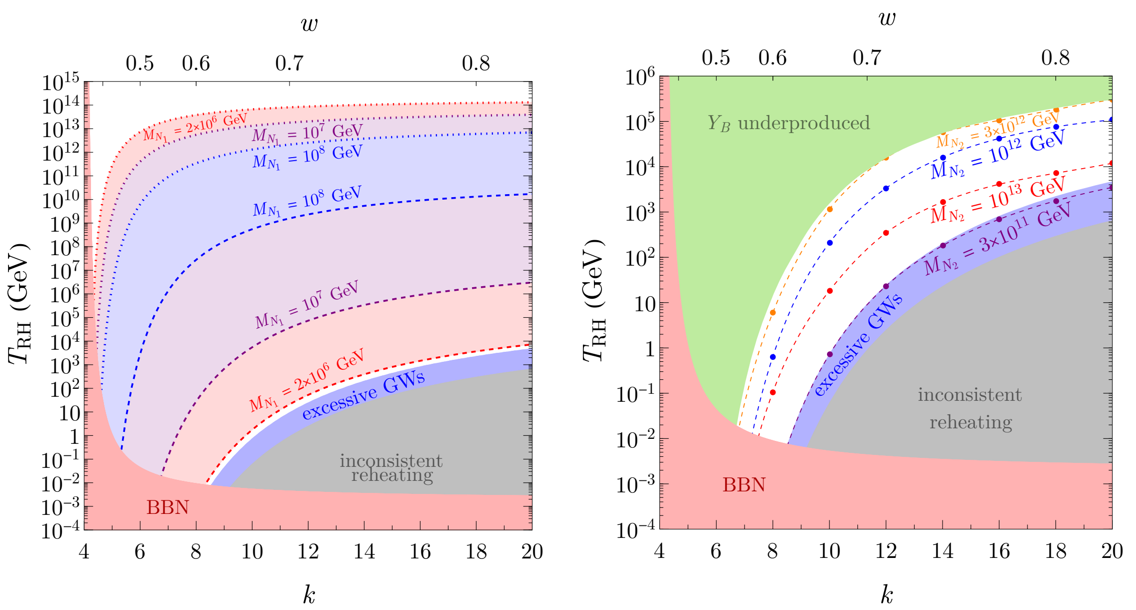

Figure 1: *Left:* Coloured regions correspond to values of $(k, T_{\rm RH})$ with $(\Omega_{N_1}^T + \Omega_{N_1}^{\phi^k})h^2 \leq 0.12$ for the three choices of M_{N_1}. *Right:* Contours of M_{N_2} corresponding to the observed baryon asymmetry [cf. Eq. (32)] in the $(k, T_{\rm RH})$ plane. The red-shaded region correspond to the lower bound on $T_{\rm RH}$ from BBN, and the green region leads to underproduction of Y_B due to the kinematic suppression in inflaton scattering when M_{N_2} approaches m_ϕ.

In the right panel of Fig. 1, we show contours of some benchmark values of the mass of N_2 that reproduce the observed baryon asymmetry $Y_B^{\rm obs}$. We find that the gravitational contribution to the baryon asymmetry is essentially entirely due to inflaton scattering rather than the thermal particles in the SM bath. Since minimal gravitational interactions are excluded by excessive GWs, non-minimal interactions are required to produce a sufficiently large thermal bath so that GW fractional energy is consistent with BBN. Leptogenesis via N_2 is therefore possible above the border of the blue-shaded region in Fig. 1, indicating a mass $M_{N_2} \gtrsim 3 \times 10^{11}$ GeV is required. Larger values of M_{N_2} can produce the correct asymmetry so long as $\xi_h > 0$. Once again, the bottom red region is forbidden by BBN because of an excessive inflaton energy density during BBN. In summary, we observe that, saturating the bound on GWs from BBN, together with the right DM abundance and successful leptogenesis requires $\xi_h \simeq 0.5$, $M_{N_2} \simeq 3 \times 10^{11}$ GeV and $M_{N_1} \simeq 10^6$ GeV.

We project the viable parameter space in the (M_{N_1}, M_{N_2}) plane in Fig. 2 for different values of ξ_h, allowing k to vary within $k \in [6, 20]$. In each coloured line segment, gravitational interactions are responsible for the observed DM relic abundance, the baryon asymmetry and reheating. Different coloured slanted line segments in this figure correspond to different choices of the non-minimal coupling ξ_h, with $\xi_h = 0$ being ruled out from overproduction of GWs. The maximum possible value for ξ_h is around 13.5, above which the mass M_{N_2} necessary to reproduce the observed baryon asymmetry gets too close to m_ϕ and kinematic suppression becomes significant, as can be seen from the right panel of Fig. 1. Note that for each ξ_h, the allowed parameter space satisfying all the constraints, is rather restricted. This is better seen from the right panel figure, where we have zoomed in to the $\xi_h = 1$ case. Interestingly, this shows that the viable parameter space is approximately independent of k, while $k = 6$ and 8 are excluded by BBN.

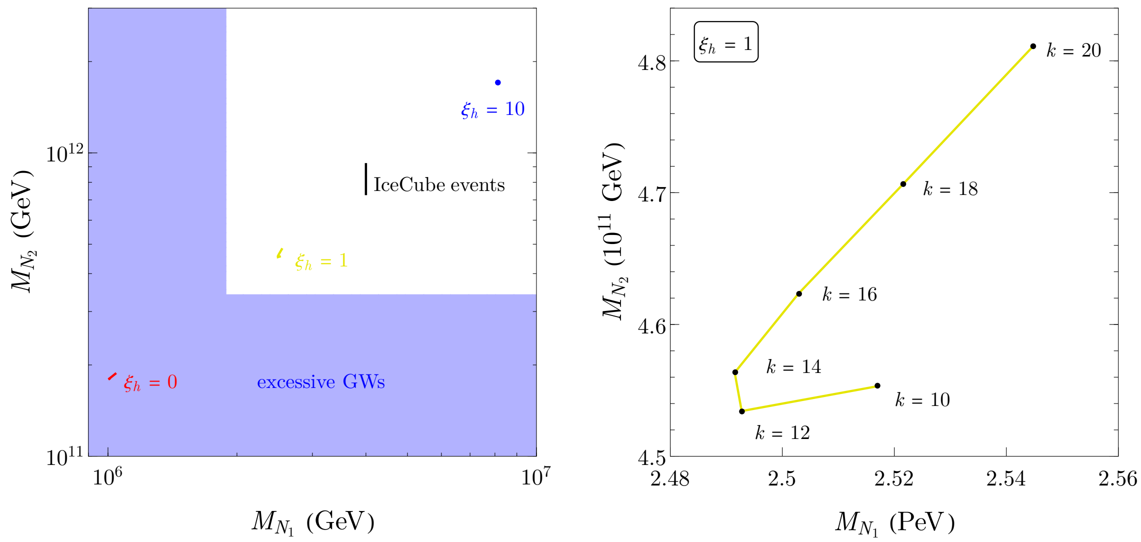

Figure 2: Viable parameter space in the (M_{N_1}, M_{N_2}) plane for which gravitational interactions are responsible for the observed DM relic abundance (in N_1), the baryon asymmetry (produced from N_2 decays), and reheating for $k \in [6, 20]$. In the left panel, different colours correspond to $\xi_h = \{0, 1, 10\}$ diagonally from bottom left (red) to top right (blue). The vertical black segment indicates the range in M_{N_2} for $M_{N_1} = 4$ PeV for the range in k considered, where the connection to the IceCube high-energy neutrino excess will be discussed in the next subsection. In the right panel, we magnify the parameter space for a fixed non-minimal coupling $\xi_h = 1$. The dots correspond to even values of k as indicated.

7.2 The case of decaying dark matter

Until now, we have assumed that the DM candidate, N_1, is absolutely stable. If it is not, and N_1 has non-zero Yukawa components, y_{1i}, N_1 can decay to SM final states. In this case, one necessary (but not sufficient) constraint on the DM mass and Yukawa coupling arises from the requirement of having a lifetime larger than the age of the Universe $\tau_{N_1} \gtrsim \tau_{\rm univ} \simeq 4.35 \times 10^{17}$ s. On the other hand, the IceCube detector has reported the detection of three PeV neutrinos, a roughly 3σ excess above expected background rates [67–69]. The three highest energy events correspond to deposited energies of 1.04 PeV, 1.14 PeV and 2.0 PeV. Although the origin of these very high energy events is still unclear, it has been shown in Refs. [70–84] that such events could be sourced from the decays of superheavy DM particles. Our minimalistic framework contains a natural avenue to reconcile both the DM abundance and IceCube events, through the gravitational production of decaying PeV neutrinos in the early Universe. We show in Fig. 3 contours for $\Omega_{N_1} h^2 = 0.12$ for $M_{N_1} = 4$ PeV in the $(k, T_{\rm RH})$ plane. In the left panel of Fig. 2, we show, by the black vertical line segment, the range in M_{N_2} obtained from varying k while fixing $M_{N_1} = 4$ PeV.

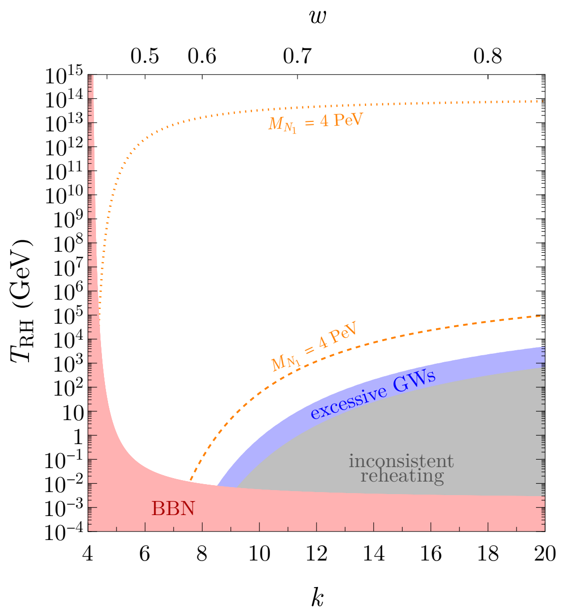

Figure 3: Contours of fixed relic density, $\Omega_{N_1} h^2 = 0.12$ for $M_{N_1} = 4$ PeV. The upper dotted contour corresponds to production from gravitational scattering in the thermal bath (and requires a large value of ξ_h) and the lower dashed contour corresponds to production from inflaton scattering (and requires a relatively low value of ξ_h) Between the two contours $\Omega_{N_1} h^2 < 0.12$ for $M_{N_1} = 4$ PeV.

8. Dark matter & leptogenesis with a Majoron

The relevant Lagrangian of this extension can be written as

$$\mathcal{L}_\Phi = (-y_R^i\, \Phi\, \overline{N_i^c}\, N_i + {\rm h.c.}) + \frac{1}{2}\mu_\Phi^2\, \Phi^2 - \frac{1}{4}\lambda_\Phi\, \Phi^4\,. \tag{36}$$

After symmetry breaking, the real part of Φ acquires a non-zero vacuum expectation value, around which one can expand the field as: $\Phi = \frac{1}{\sqrt{2}}(S + v_S)e^{iJ/v_S}$, and J is the Majoron. This expectation value is the origin of the RHN Majorana masses, $M_{N_i} = y_R^i\, v_S/\sqrt{2}$. Then $m_S = \mu_\Phi < m_\phi$ and the gravitational production rate of the real scalar, S is

$$R_S^{\phi^k} = \frac{2 \times \rho_\phi^2}{16\pi M_P^4} \Sigma_S^k \,, \tag{37}$$

where the factor of two accounts for the fact we produce two scalar particles per scattering, with [23, 26]

$$\Sigma_S^k = \sum_{n=1}^{\infty} |\mathcal{P}_{2n}^k|^2 \left[1 + \frac{2\mu_\Phi^2}{E_{2n}^2}\right]^2 \sqrt{1 - \frac{4\mu_\Phi^2}{E_{2n}^2}} \,. \tag{38}$$

Since each scalar decays into 2 right-handed neutrinos, we obtain for the density of N_i after integration of the Boltzmann equation [26],

$$n_{N_i}^{S\phi^k}(a_{\rm RH}) \simeq {\rm Br}_i \times \frac{\sqrt{3}\rho_{\rm RH}^{3/2}}{4\pi M_P^3} \frac{k+2}{6k-6} \left(\frac{\rho_{\rm end}}{\rho_{\rm RH}}\right)^{1-\frac{1}{k}} \Sigma_S^k, \tag{39}$$

where we assumed $a_{\rm RH} \gg a_{\rm end}$, and here ${\rm Br}_i = \frac{(y_R^i)^2}{\sum (y_R^i)^2}$ so ${\rm Br}_i = \frac{M_{N_i}^2}{M_{N_1}^2 + M_{N_2}^2 + M_{N_3}^2}$ if $N_{1,2,3}$ are all lighter than S. The relic abundance of N_1 is then given by

$$\begin{aligned} \frac{\Omega_{N_1}^{S\phi^k} h^2}{0.12} &\simeq {\rm Br}_1 \times \left(\frac{\rho_{\rm end}}{10^{64}{\rm GeV}^4}\right)^{1-\frac{1}{k}} \left(\frac{10^{40}{\rm GeV}^4}{\rho_{\rm RH}}\right)^{\frac{1}{4}-\frac{1}{k}} \left(\frac{k+2}{6k-6}\right) \\ &\times \; \Sigma_S^k \times \frac{M_{N_1}}{2.5 \times 10^{\frac{24}{k}-8}{\rm GeV}} \,, \end{aligned} \tag{40}$$

whereas the baryon asymmetry follows from Eq. (32). We show in Figs. 4 and 5 respectively, the parameter space allowed by the relic abundance and the baryogenesis constraint in the $(k\,, T_{\rm RH})$ plane. We notice that the DM mass respecting Planck constraint is much lower if the branching fraction to N_1 is large. For smaller branching fraction, the density of N_1 through this channel is suppressed and the effect is milder and proportional to ${\rm Br}_1$, as one can see in Fig. 4 right panel.

Finally, we can combine all of the preceding results, adding the possibility for a gravitational reheating with non-minimal coupling. We illustrate this in Fig. 6, which is the analogue of Fig. 2 but with the scalar S as an intermediate state. Each color segment in Fig. 6 assumes a fixed ξ_h and allows values of $k \in [6, 20]$ that are consistent with the BBN bound on $T_{\rm RH}$. The black dot indicates the $M_{N_{1,2}}$ masses, independent of k, required to explain the IceCube high-energy neutrino excess. The green region is inaccessible because $m_S > M_{N_2} > m_\phi/2$ forbids the production of S via ϕ scattering.

9. Conclusions

We have shown that there exists the possibility that inflationary reheating, dark matter and the baryon asymmetry can be generated solely gravitational interactions. The baryon asymmetry is produced through the decay of a right-handed neutrino M_{N_2} (leading first to a non-zero lepton asymmetry). For minimal gravitational interactions, $\xi_h = 0$, a large amount of dark radiation is created in the form of gravitational waves and is inconsistent with BBN. Thus, we allow for a

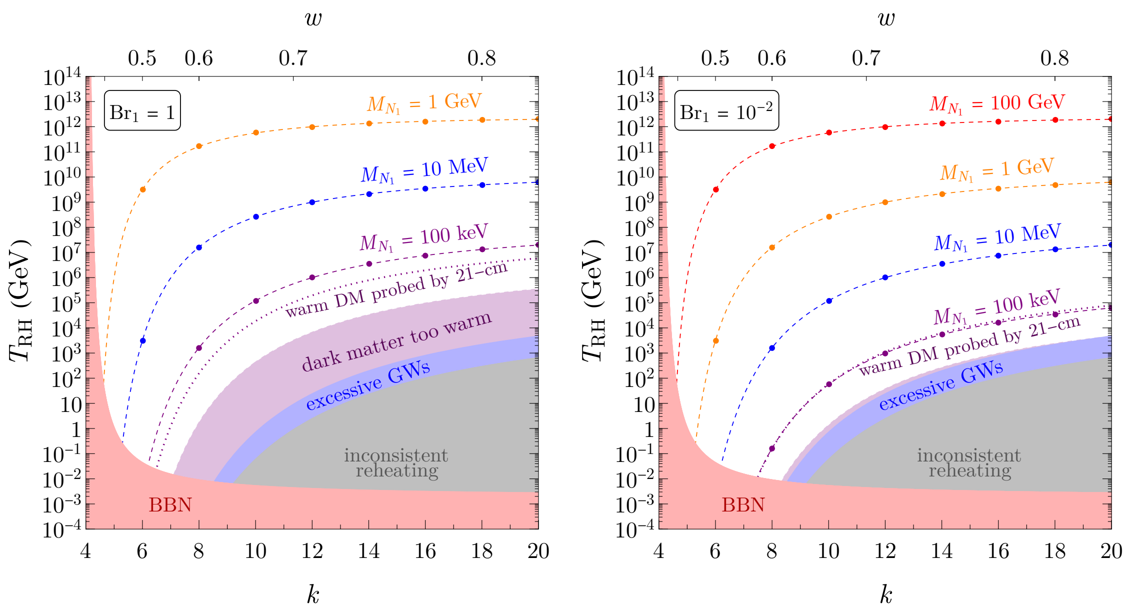

Figure 4: *Contours of observed relic abundance assuming $Br_1 = 1$ (left) and $Br_1 = 10^{-2}$ (right) for different choices of the DM mass, considering only Majoron contribution. The purple-shaded region is disallowed from the warm DM limit (see text).*

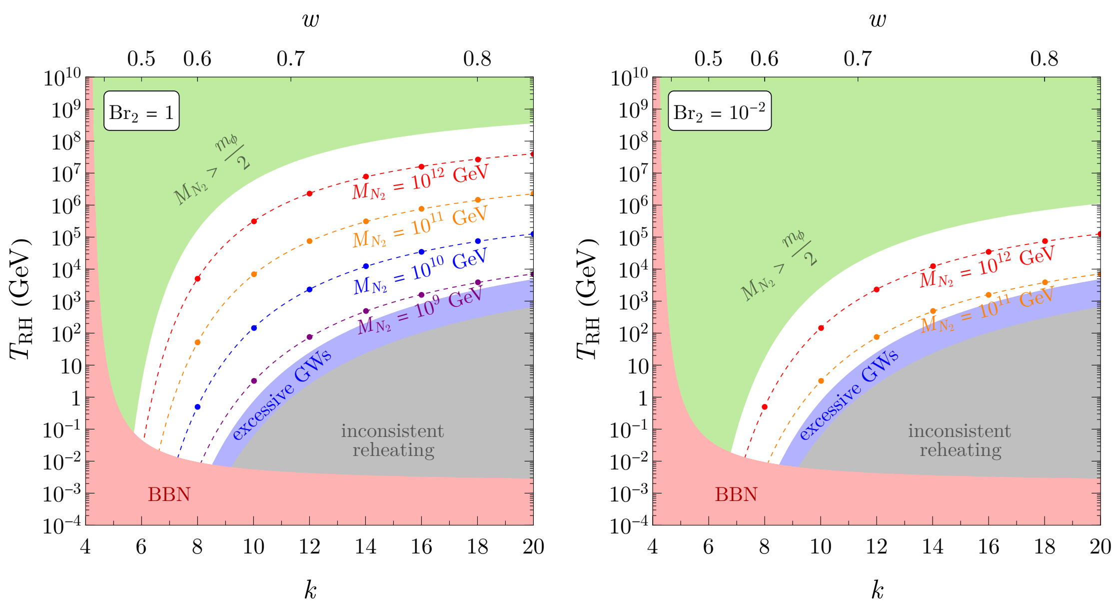

Figure 5: Contours of N_{N_2} corresponding to the observed baryon asymmetry for $\mathrm{Br}_2 = 1$ (left) and $\mathrm{Br}_2 = 10^{-2}$ (right) in the (k, T_{RH}) plane. Here only the contribution due to the intermediate scalar is included. The green-shaded region is kinematically inaccessible due to $M_{N_2} > m_\phi/2$.

non-minimal gravitational coupling $\xi_h R H^2$ where H the Standard Model Higgs field to enhance reheating, so that the ratio of gravitational wave energy density to the radiation is decreased. We also showed that N_1, if unstable, can explain the recent IceCube PeV events through its decay into SM neutrinos. In this case, if we want to accommodate simultaneously the correct DM relic

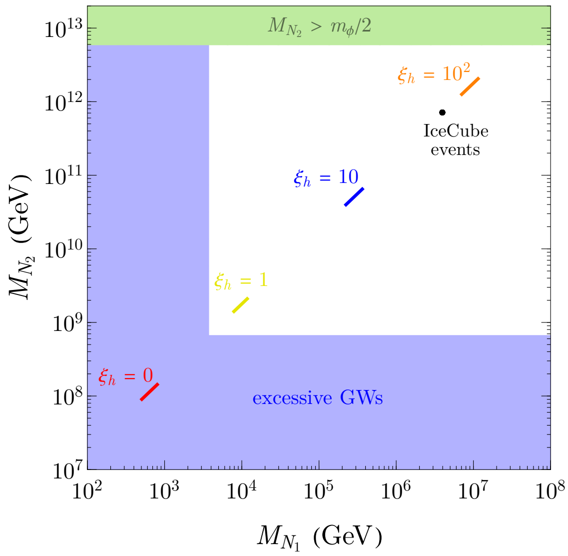

Figure 6: Parameter space satisfying the right dark matter relic abundance and baryon asymmetry, considering the production through S. The line colors correspond to different values of ξ_h, with $\xi_h = \{0, 1, 10, 10^2\}$ from bottom to top, and $\xi_h = 0$ corresponds to minimal graviton exchange. Each colored line segment shows the variation of the predicted masses with $k \in [6, 20]$. The black dot marks the parameter point that can also explain the IceCube high-energy neutrino excess.

abundance, the observed baryon asymmetry, gravitational reheating and the IceCube events, the value of ξ_h is fixed for a given k. Finally, we proposed a new scenario where the RHN and the dark matter are produced through an intermediate scalar state S, the CP-even partner of the Majoron. In this case, the gravitational production of the scalar is not helicity suppressed by the mass of the final state fermions. As a result, the mass ranges for N_1 and N_2 are increased.

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