Matter-antimatter asymmetry without loops

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We propose a new mechanism for generating matter-antimatter asymmetry via the interference of tree-level diagrams only. We first derive a general result that a nonzero CP-asymmetry can be generated via at least two sets of interfering tree-level diagrams involving either $2 \to 2$ or $1 \to \mathcal{N}$ (with $\mathcal{N} \geq 3$) processes. We illustrate this point in a simple TeV-scale extension of the Standard Model with an inert Higgs doublet and right-handed neutrinos, along with an electroweak-triplet scalar field. The imaginary part needed for the required CP-asymmetry comes from the trilinear coupling of the inert-doublet with the triplet scalar. Small Majorana neutrino masses are generated by both scotogenic and type-II seesaw mechanisms. The real part of the neutral component of the inert-doublet serves as a cold dark matter candidate. The evolutions of the dark matter relic density are considered by both scotogenic and type-II seesaw mechanisms. The real part of the neutral component of the inert-doublet serves as a cold dark matter candidate. The evolutions of the dark matter relic density are considered by both scotogenic and type-II seesaw mechanisms.

\textsuperscript{1}In principle, this condition can be somewhat relaxed if we consider flavor-dependent asymmetries, with zero net lepton or baryon number in the final state, as in flavored leptogenesis (for recent reviews, see Refs. [20, 21]). For simplicity, here we will not consider such flavor-dependent effects.
FIG. 1. Generic topologies for tree-level $2 \to 2$ subprocesses that can give rise to nonzero lepton or baryon asymmetry. Here $i_{1,2}$ and $f_{1,2}$ are respectively the initial and final states, and $m_{1,2}$ are the masses of two different mediators.

where $\text{Im}[C_i C_j^\ast]$ is the imaginary part coming from the couplings, which is required to be nonzero for CP violation, and $\text{Im}[M_1 M_2^\ast]$ incorporates the imaginary part from the sub-amplitudes $M_{1,2}$, which is reminiscent of the imaginary part coming from the interference of tree and loop diagrams in the $1 \to 2$ decay scenario. Note that for $1 \to 2$ decays, the tree-level decay rates for $f_1 f_2$ and $f_1^\ast f_2^\ast$ final states are both proportional to the modular square of the same coupling, which implies $\delta = 0$. Eq. (3) is a general result applicable to $2 \to 2$ scatterings, as well as $1 \to N$ decays (for $N \geq 3$).

In the tree-level $2 \to 2$ processes shown in Fig. 1, we have only one source for the complex sub-amplitudes, which is due to the finite widths of the mediators.\footnote{One may argue that the finite width is also a loop-induced effect for unstable particle decays, since it is related to the imaginary part of the self-energy [22, 23]. However, the crux of our discussion is that we only require a nonzero width, whereas the $1 \to 2$ decay case needs both nonzero width and interference between tree and loop (self-energy and/or vertex correction) diagrams.} In general, the sub-amplitudes $M_{1,2}$ can be written as

$$M_j = \frac{A_j}{x_j - m_j^2 + i m_j \Gamma_j}, \quad (4)$$

with $j = 1, 2$, $x_j = s, t, u$ the Mandelstam variables, $m_j$ and $\Gamma_j$ respectively the mediator masses and widths, and $A_j$ some arbitrary real parameters. The imaginary component of the product of amplitudes appearing in Eq. (3) can then be written as

$$\text{Im}[M_1 M_2^\ast] = \frac{A_1 A_2 [(x_1 - m_1^2) m_2 \Gamma_2 - (x_2 - m_2^2) m_1 \Gamma_1]}{[(x_1 - m_1^2)^2 + m_1^2 \Gamma_1^2] [(x_2 - m_2^2)^2 + m_2^2 \Gamma_2^2]}, \quad (5)$$

which is non-vanishing as long as $(x_1 - m_1^2) m_2 \Gamma_2 \neq (x_2 - m_2^2) m_1 \Gamma_1$. With the imaginary part of the couplings $\text{Im}[C_i C_j^\ast] \neq 0$, we can then produce a nonzero asymmetry [cf. Eq. (3)]. This general argument holds, irrespective of the specific subprocesses or the model details.

For the tree-level topologies shown in Fig. 1, we can have three distinct possibilities for the two subprocesses to realize $\text{Im}[M_1 M_2^\ast] \neq 0$ in Eq. (5):

i) If both subprocesses are in the $s$-channel [cf. Fig. 1 (a)+(b)], one just needs to replace $x_{1,2}$ by $s$ in Eq. (5). In this case, the $CP$-asymmetry factor $\delta$ in Eq. (3) can be largely enhanced in the vicinity of the resonance(s), with $s - m_i^2 \approx m_i \Gamma_i$ (with $i = 1, 2$), similar to the enhancement effect in resonant leptogenesis [24, 25].

ii) If one of the sub-amplitudes is in the $s$-channel and the other one in the $t$- or $u$-channel [cf. Fig. 1 (a)+(d) or (b)+(c)], one can safely neglect the imaginary part for the $t$- or $u$-channel propagator. For concreteness, we take $M_1$ as the $s$-channel and $M_2$ as the $x$-channel ($x = t$ or $u$) amplitude. In this case, Eq. (5) is simplified to

$$\text{Im}[M_1 M_2^\ast] \approx - \frac{A_1 A_2 m_1 \Gamma_1}{[(s - m_1^2)^2 + m_1^2 \Gamma_1^2] (x - m_2^2)}, \quad (6)$$

which is proportional to the $s$-channel mediator width $\Gamma_1$. Here also the $CP$-asymmetry could be largely enhanced at the $s$-channel resonance, i.e. $s - m_1^2 \approx m_1 \Gamma_1$.

iii) If both subprocesses are in the $t$- or $u$-channel [cf. Fig. 1 (c)+(d)], then the width terms in the denominator of Eq. (5) can be neglected, i.e.

$$\text{Im}[M_1 M_2^\ast] \approx \frac{A_1 A_2 [(x_1 - m_1^2) m_2 \Gamma_2 - (x_2 - m_2^2) m_1 \Gamma_1]}{(x_1 - m_1^2)^2 (x_2 - m_2^2)^2}. \quad (7)$$

As a result, the $CP$-asymmetry is suppressed by the ratio $m_1 \Gamma_1/(x_j - m_j^2)$ with $j = 1, 2$.

In what follows, we will consider a concrete model realization for the case ii) and illustrate the baryon asymmetry generation with a few benchmark points (BPs).

**The model.**—To illustrate our tree-level mechanism in a minimal realistic extension of the SM, we consider an amalgamation of the scotogenic model [26] and type-II seesaw [27–31]. For the purpose of scotogenic mechanism, an inert SU(2)$_L$-doublet scalar $\eta \equiv (\eta^+, \eta^0)$ and three right-handed neutrinos (RHs) $N_i$ (with $i = 1, 2, 3$) are introduced. To implement type-II seesaw, an SU(2)$_L$-triplet scalar $\Delta \equiv (\Delta^+, \Delta^+, \Delta^{0})$ is added. The inert doublet $\eta$ and the three RHs $N_i$ are odd under a discrete Z$_2$ symmetry, while all other particles are even. In this model, we assume the RHs are heavier than the $\eta$ scalars, thus the lightest neutral component $\eta^0$ plays the role of dark matter [26]. A nonminimal coupling of the inert doublet gravity can also successfully accommodate inflation [32].

The relevant Yukawa couplings are given by the Lagrangian

$$-\mathcal{L}_Y = Y_{\alpha i a}^N \bar{L}_\alpha N_i + Y_{\alpha \beta}^\Delta \bar{T}_\alpha \Delta L_\beta + \text{H.c.}, \quad (8)$$
with $L \equiv (\nu, \ell)$ being the SM lepton doublet, $C$ the charge conjugation operator, $\tilde{\eta} = i\sigma_2\eta^*$ ($\sigma_2$ being the second Pauli matrix), $\alpha, \beta = e, \mu, \tau$ the lepton flavor indices, and $i = 1, 2, 3$ the RHN indices. For simplicity, we assume there is no mixing nor $CP$ phase in the RHN sector. The most general scalar potential for the SM Higgs doublet $H \equiv (H^+, H^0)$, inert doublet $\eta$ and triplet $\Delta$ is given by

\[
V \supset -\frac{1}{2}(\mu_H H^\dagger H + \mu_{\eta} \Delta^\dagger \Delta + \mu_{\eta} H \eta + H.c.)
\]

where $\tilde{H} = i\sigma_2H^*$ and the mass parameters $\mu_{H, \eta, \Delta} > 0$ so that both $\tilde{H}$ and $\Delta$ obtain non-vanishing vacuum expectation values (VEVs), i.e. $\langle H^0 \rangle = v \simeq 246$ GeV and $\langle \Delta^0 \rangle = v_\Delta$. The mass parameter $\mu_{\eta} \Delta$ is chosen to be complex, which is crucial for the $CP$-asymmetry [cf. Eq. (3)]. All other parameters in Eq. (9) are assumed to be real. The physical masses for the neutral and charged scalars can be obtained from the minimization of the potential (9), which is detailed in the Appendix. Note that here the doublet $\eta$ is odd under the $Z_2$ symmetry and does not mix the SM Higgs and the triplet, which is necessary for the neutral real component $\eta_R$ to be a stable DM candidate.

In this setup the neutrino mass is generated from both loop-level scotogenic and tree-level type-II seesaw mechanisms, which are induced respectively by the Yukawa couplings $Y^N$ and $Y^\Delta$ given in Eq. (8):

\[
m_{\nu} = (Y^N)^\dagger \Lambda Y^N + Y^\Delta v_\Delta,
\]

where $\Delta$ is an effective loop-suppressed RHN mass scale, given by [26, 33]

\[
\Lambda_{ii} = \frac{m_{N_i}}{16\pi^2} \left[ \frac{m^2_{31}}{m^2_{3i}} - \frac{m^2_{32}}{m^2_{32}} \ln \left( \frac{m^2_{N_i}}{m^2_{N_3}} \right) \right].
\]

We have assumed the RHNs do not mix with each other, therefore $\Lambda$ is a diagonal matrix. The Yukawa couplings in Eq. (10) are related to the neutrino oscillation data, $\Delta$ and $v_\Delta$ as follows:

\[
Y^N_{i\alpha} = F^\frac{1}{2}_1 \left( \Lambda^{-1/2} O \tilde{m}_{\nu}^{1/2} U_{PMNS}^{\dagger} \right)_{i\alpha},
\]

\[
Y^\Delta_{\alpha\beta} = F^\frac{1}{2}_1 v^\dagger \Delta \left( U_{PMNS}^{\dagger} \tilde{m}_{\nu} U_{PMNS} \right)_{\alpha\beta},
\]

where $\tilde{m}_{\nu} = \{m_{\nu_1}, m_{\nu_2}, m_{\nu_3}\}$ the diagonal neutrino mass eigenvalues, and $U_{PMNS}$ the PMNS lepton mixing matrix. In Eq. (12) we have used the Casas-Ibarra parametrization [34] for the coupling $Y^N$, with $O$ an arbitrary orthogonal matrix. $F_1$ and $F_2$ are the fractions of contributions to neutrino mass matrix from the radiative scotogenic and tree-level type-II seesaw mechanisms respectively, with $F_1 + F_2 = 1$.

**Boltzmann equations.**– As stated above, the matter asymmetry is generated from the interference effects between two tree-level diagrams, which are shown in Fig. 2 for our scotogenic type-II seesaw model. In particular, we analyze the $2 \to 2$ $\Delta L = 2$ scattering processes

\[
\eta \to L_{\alpha} L_{\beta},
\]

which include $\eta^0 \eta^0 \to \ell_\alpha \ell_\beta$, $\eta^0 \eta^\pm \to \ell_\alpha \nu_\beta$ and $\eta^0 \eta^0 \to \nu_\alpha \nu_\beta$. These processes can be mediated by an $s$-channel triplet scalar $\Delta$, and also by RHNs $N_i$ in the $t$- and $u$-channels, as shown in Fig. 2. The effective $CP$-asymmetry factor [cf. Eq. (3)] is given by

\[
\delta = \sum_{\alpha\beta} 4 \text{Im} \left[ \mu_{\eta} \Delta Y_{i\alpha}^N \lambda_{\alpha\beta}^* Y_{i\beta}^N \right]
\times \frac{sm_{N_i} m_{\Delta} \Gamma_\Delta}{(s - m_{2i}^2)^2 + m_{\Delta}^2 \Gamma^2_\Delta} \left[ \frac{1}{t - m_{N_i}^2} + \frac{1}{u - m_{N_i}^2} \right],
\]

where $\Gamma_\Delta$ is the triplet scalar width. Here the imaginary part comes purely from the combinations of the Yukawa couplings $Y^N, Y^\Delta$ [cf. Eq. (8)] and the trilinear coupling $\mu_{\eta} \Delta$ [cf. Eq. (9)], which can be parametrized as

\[
\sum_{\alpha\beta} \text{Im} \left[ \mu_{\eta} \Delta Y_{i\alpha}^N \lambda_{\alpha\beta}^* Y_{i\beta}^N \right]
= F_{1} F_{2} v_{\Delta}^\dagger \text{Im} \left\{ \mu_{\eta} \Delta \left[ \Lambda^{-1/2} O \tilde{m}_{\nu}^* O^T \Lambda^{-1/2} \right] \right\}.
\]

In general the $O$ matrix might also be complex, thus contributing to the imaginary part in Eq. (16).

It is interesting that part of the same $2 \to 2$ process (14) containing $\eta^0$ also contributes to the (co)annihilation of DM particles. In this sense, the time evolutions of the DM relic density and the lepton asymmetry are related, as we will see below. The freeze-out mechanism for the DM is identical to the standard inert-doublet case [36, 37], where we can have $\eta \to \text{SM SM}$
with ‘SM’ standing for all SM fermions (quarks and leptons), gauge bosons $W,Z$ and Higgs boson $h$.

The genesis of DM relic density and leptonic asymmetry are both governed by the coupled Boltzmann equations

$$\frac{dY_\eta}{dz} = \frac{-s}{H(z)z} \left[ (Y_\eta^2 - (Y_\eta^{eq})^2)\langle \sigma v \rangle (\eta\eta \to SM SM) \right],$$

$$\frac{dY_{\Delta L}}{dz} = \frac{s}{H(z)z} \left[ (Y_{\Delta L}^2 - (Y_{\Delta L}^{eq})^2)\langle \sigma v \rangle (\eta\eta \to LL) - 2Y_{\Delta L}Y_{\eta}^{eq}\langle \sigma v \rangle (\eta\eta \to LL) - 2Y_{\Delta L}Y_{\eta}^{eq}\langle \sigma v \rangle (\eta\bar{L} \to \eta L) \right],$$

where $z = m^2_{\eta}/T$, $Y_i^{eq} = \langle \sigma v \rangle /s$ are the normalized number densities (in equilibrium) for the particles $i$ ($s$ being the entropy density), $Y_{\Delta L} = Y_L - Y_{\bar{L}}$, $r_\eta = Y_\eta^{eq}/Y_\eta^{eq}$, and $H(z) = \sqrt{8\pi^2g_*/90} m_\eta^2/(z^2M_{Pl})$ with $M_{Pl}$ the Planck scale and $g_*$ the number of relativistic degrees of freedom at temperature $T$. Here $\langle \sigma v \rangle$ are the thermally-averaged annihilation/scattering rates; $\langle \sigma v \rangle (\eta\eta \to SM SM)$ is the DM annihilation rate, $\langle \sigma v \rangle (\eta\eta \to LL)$ is the thermally-averaged scattering cross section for $\eta\eta \to LL$ that includes also the $\Delta L = 0$ processes such as $\eta^+\eta^- \to \ell^+\ell^-$, whereas $\langle \sigma v \rangle (\eta\eta \to LL)$ only includes the $\Delta L = 2$ processes listed below Eq. (14). The expressions for all thermal cross sections in Eq. (18) are collected in the Appendix.

Evaluating the Boltzmann equations above, one can obtain the lepton asymmetry $Y_{\Delta L}(z)$, which is then converted to baryon asymmetry $Y_{\Delta B} = -(28/51)Y_{\Delta L}$ [38] via the standard electroweak sphaleron processes [39] at the sphaleron transition temperature $T_{sph}$. In an analogous way, one can also calculate the evolution of the DM density $Y_\eta$ from Eq. (17) and get the final relic abundance $\Omega_{DM}h^2 = 2.755 \times 10^8 Y_\eta (m_\eta/\text{GeV})$ at DM freeze-out temperature $T_f \simeq m_\eta/20$.

**Numerical results.**– We solve the Boltzmann equations (17) and (18) numerically for three representative benchmark points (BP1, BP2, BP3) given in Table I in the Appendix, obtained by implementing our model in SARAH 4 [40] and after checking consistency with all existing experimental constraints. We assume $F_I = F_{II} = 1/2$ in Eqs. (12) and (13), i.e. equal contributions from scotogenic and tree-level type-II seesaw to neutrino masses. This choice maximizes the $C$P-asymmetry in Eq. (16), subject to keeping other factors the same. In addition, the $O$ matrix is taken to be identity, so that $O\tilde{m}_\nu^2O^T = \tilde{m}_\nu^2$, and the mass parameter $\mu_\Delta$ is assumed to be purely imaginary in Eq. (16). The RHNS are taken to be much heavier than the $\eta$ particles to avoid the wash-out of lepton asymmetry from the inverse decay processes $L\eta \to N_i$. For the benchmark points we take, the mass splitting $m_{\eta_1} - m_{\eta_2}$ is larger than 100 keV scale, such that the direct detection constraints for inelastic scattering of DM with nucleons [41–43] can be evaded.

The evolutions of the DM relic density $\Omega_{DM}h^2$ and the baryon asymmetry $Y_{\Delta B}$ are evaluated using micrOMEGAs 5.0 [44] and the results are presented in Fig. 3. In the left panel of Fig. 3, we have shown the time evolution of the DM relic density by red solid, dashed and dotted-dashed curves for BP1, BP2, and BP3 respectively. The
the right panel of Fig. 3. For each of the three di- 

tillator mass on the baryon asymmetry is further illus-

trates in the left panel of Fig. 3, we have fixed the ∆-mediator 

mass at the resonance point and have satisfied the re-

quired baryon asymmetry by appropriately fixing the tri-

linear coupling µη as shown in Table I. We find that the 

size of the trilinear coupling needed for the asymmetry 

decreases as the mass of the DM increases.

For each choice of the DM mass, the maximal contribu-

tion to the baryon asymmetry comes in the vicinity of the 

collision density goes as

\[ \sigma v = \frac{1}{m^2/\lambda_{H\eta}^2}. \]

Due to the Majorana nature of the heavy RHNs, we can 

generate baryogenesis, dark matter and neutrino mass 

in principle be distinguished from the pure scotogenic or 

monojet \cite{59, 60} searches at the LHC. Our model can 

be directly searched for at future high-energy colliders.

In the scotogenic model, the RHNs do not mix di-

mediate by \( \lambda_{H}\eta \). The inert doublet scalars can also 

be probed in the high-precision low-energy ex-

periments like MOLLER \cite{54}. The charged \( \eta^\pm \) scalars 

are produced in association with the neutral DM parti-

cle through the \( W \) boson, i.e. \( pp \rightarrow W^* \rightarrow \eta^\pm \eta^0 \rightarrow 

\eta^0\eta^0W^{(*)} \) \cite{55}. The inert doublet scalars can also be 

produced from their couplings to the SM Z boson via 

\( pp \rightarrow \eta_L\eta_Rj \) or the SM Higgs through \( pp \rightarrow \eta^0\eta^0j \) (with 

\( j \) being an energetic jet) \cite{56}. The inert-doublet sector 

can then be constrained by the mono-\( W \) \cite{57, 58} and 

monojet \cite{59, 60} searches at the LHC. Our model can 

in principle be distinguished from the pure scotogenic or 

pure type-II seesaw model at colliders using both inert-

doublet and triplet-scalar signatures.

In the scotogenic model, the RHNs do not mix di-

crty Fukui et al. \cite{54} for details. The off-shell decays 

are \( \eta^\pm \rightarrow \ell_N N_i \), followed by \( N_i \rightarrow \eta L \).

Due to the Majorana nature of the heavy RHNs, we can 

generate baryogenesis, dark matter and neutrino mass 

in principle be distinguished from the pure scotogenic or 

Type-II seesaw mechanism explicitly in 

a well-motivated scotogenic model with type-II seesaw, 

in which the asymmetry is generated in the \( \Delta L = 2 \) pro-

cesses \( \eta \eta \rightarrow L_\alpha L_\beta \) mediated by s-channel triplets and \( t \) 

or u-channel RHNs. The neutrino masses receive contribu-

tions from both scotogenic and type-II seesaw mecha-

nisms. The real part of the neutral component of the 

inert doublet \( \eta \) serves as a DM candidate. As shown in 

Fig. 3 the observed baryon asymmetry and DM relic 

density can be achieved for (sub)TeV inert-doublet and 

triplet masses.

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Appendix

Scalar masses.— All the scalar masses in the scoto-

genic plus type-II seesaw model with the SM Higgs \( H \), the 

inert doublet \( \eta \) and the triplet \( \Delta \) can be obtained from 

the scalar potential (9). When we take the first-order deriv-

ative of the potential with respect to the VEVs \( v \) and \( v_\Delta \), the 

solutions of the tadpole equations for \( \{\mu_H^2, \mu_\Delta^2\} \) are 

given by

\[ \mu_H^2 = \frac{1}{2}(\lambda_H v^2 - 2\sqrt{2}\mu_H v_\Delta + \lambda_H^2 v_\Delta^2), \]

\[ \mu_\Delta^2 = \frac{1}{2}(\lambda_\Delta v_\Delta^2 - 2\sqrt{2}\mu_H v_\Delta + \lambda_H^2 v^2). \]

After replacing \( \{\mu_H^2, \mu_\Delta^2\} \) in the scalar potential, the mass
matrix for the real scalars reads
\[ M^0 = \begin{pmatrix}
\lambda_H v^2 & \frac{\mu H \Delta v^2}{\sqrt{2} v}\n\frac{\mu H \Delta v^2}{\sqrt{2} v} & \frac{\mu^2 H^2 v^4}{4 v^2}
\end{pmatrix}, \tag{21}
\]
from which we can get the two mass eigenvalues for the Higgs and scalar \( \Delta^0 \), which is the real part of \( \Delta^0 \). In the case of \( \mu_H \Delta \sim O(100) \) keV, the two CP-even scalar masses turn out to be
\[ m_{h_R}^2 \approx \lambda_H v^2, \quad m_{\Delta^0}^2 \approx \frac{\mu H \Delta v^2}{\sqrt{2} v}, \tag{22}
\]
with the first one \((h)\) identified as the SM-like Higgs boson. The masses of the pseudo-scalar and the charged scalars from the triplet are
\[ m_{\Delta^+}^2 = m_{\Delta^0}^2 + \frac{1}{4} \lambda_H'^2 (v^2 + 2 v_\Delta^2), \tag{23}
\]
Finally the masses for real scalar \( \eta_R \), pseudo-scalar \( \eta_I \) and the charged scalars \( \eta^\pm \) from the \( Z_2 \)-odd doublet \( \eta \) are respectively
\[ m_{\eta_R}^2 = \frac{1}{2} (2 \mu_{\eta}^2 + (\lambda_H \eta + \lambda_H' \eta^\prime + \lambda_{H\eta} v^2) \].

\[ |M^{\text{tot}}(\eta \rightarrow LL)|^2 = \frac{\tilde{m}_\Delta^2}{v_\Delta^2} \frac{F^2_{\Delta^0} m_{\eta}^2 s}{(s - m_{\eta}^2)^2 + m_{\eta}^4}, \]
\[ + \sum_i F^2_{\Delta^0} \frac{\tilde{m}_\nu^2}{\tilde{m}_\nu^2} m_{\eta_i}^2 s \left[ \frac{1}{(t - m_{\eta_i}^2)} + \frac{1}{(u - m_{\eta_i}^2)} \right] \]
\[ + \frac{F_{\Delta^0}\rho_{\Delta^0} \mu_{\eta} m_{\eta_i} m_{\eta_j}^2}{(s - m_{\eta_i}^2)^2 + m_{\eta_i}^4} \sum_i \tilde{m}_\nu^2 \left[ \frac{1}{(t - m_{\eta_i}^2)} + \frac{1}{(u - m_{\eta_i}^2)} \right], \tag{31}
\]

The cross section \( \langle \sigma v \rangle(\eta \eta \rightarrow SM SM) \) in Eq. (17) can be found in [36, 37], with “SM SM” referring to all the possible channels involving the quarks, leptons, scalar and gauge bosons in the SM.

**Benchmark points.** – The three BP's used in our numerical analysis of the baryon asymmetry \( Y_{\Delta,B} \) and DM relic density \( \Omega_{\text{DM}} h^2 \) [cf. Fig. 3] are collected in Table I.

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TABLE I: Three benchmark points for the numerical analysis presented in Fig. 3. All the quartic couplings in Eq. (9) not listed in this table are set to be zero. Here $\Delta m_{\nu} = m_{\eta_R} - m_{\eta_L}$ is the mass splitting between the two scalars $\eta_R$ and $\eta_L$.

| BP1 | BP2 | BP3 |
|-----|-----|-----|
| $m_{\Delta}$ | 1 keV | 1 keV | 1 keV |
| $m_{\eta}$ | 600 GeV | 1 TeV | 1.5 TeV |
| $\mu_{\Delta}$ | 33.6 keV | 93.5 keV | 210 keV |
| $\mu_{\eta}$ | 15i GeV | 7.1i GeV | 6i GeV |
| $m_{\eta_1}$ | 6 TeV | 10 TeV | 15 TeV |
| $m_{\eta_2}$ | 6.6 TeV | 11 TeV | 16.5 TeV |
| $m_{\eta_3}$ | 7.2 TeV | 12 TeV | 18 TeV |
| $m_{\eta_0}$ | 600 GeV | 1 TeV | 1.5 TeV |
| $\Delta m_{\eta} = m_{\eta_R} - m_{\eta_L}$ | 506 keV | 300 keV | 200 keV |
| $m_{\eta_+}$ | 606 GeV | 1 TeV | 1.5 TeV |
| $m_{\eta_\pm}$ | 1.2 TeV | 2 TeV | 3 TeV |
| $m_{\Delta \pm}$ | 1.6 TeV | 2 TeV | 3 TeV |
| $\lambda_{H}$ | 0.253 | 0.253 | 0.253 |
| $\lambda_{H \eta}$ | 0.24 | 0.59 | 0.93 |
| $\lambda_{\eta \eta}$ | -0.24 | -0.59 | -0.93 |
| $\lambda_{\eta \eta}^2$ | $1 \times 10^{-5}$ | $1 \times 10^{-5}$ | $1 \times 10^{-5}$ |

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