Temperature effects on the Z_2 symmetry breaking in the scotogenic model

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It is well known that the scotogenic model for neutrino mass generation can explain correctly the relic abundance of cold dark matter. There have been claims in the literature that an important part of the parameter space of the simplest scotogenic model can be constrained by the requirement that no Z_2 -breaking must occur in the early universe. Here we show that this requirement does not give any constraints on the underlying parameter space at least in those parts, where we can trust perturbation theory. To demonstrate this, we have taken into account the proper decoupling of heavy degrees of freedom in both the thermal potential and in the RGE evolution.

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I. INTRODUCTION

Among the known shortcomings of the standard model (SM), probably dark matter (DM) and neutrino masses are the most important ones. While many extensions of the SM have been proposed that could either explain DM or neutrino masses, the “scotogenic” model [1] can be considered the prototype for a class of beyond SM (BSM) models that can potentially explain both phenomena at the same time.

The setup of the scotogenic model is very simple: Add to the SM (two or) three right-handed neutrinos, N_i , and a new scalar doublet, usually denoted by η . In order to avoid a tree-level type-I seesaw, a global Z_2 is added by hand. This symmetry eliminates a certain term from the scalar potential and avoids a vacuum expectation value for η , as long as certain conditions on the scalar parameters are fulfilled, see Sec. II. Under this symmetry, N_i and η are odd, while all SM particles are even. The lightest Z_2 odd particle, either N_1 or η , is then absolutely stable, providing potentially a good cold dark matter candidate. At the same time, the model is able to explain correctly the observed neutrino masses and mixing angles (for a recent global fit of all neutrino data see, for example [2]) via the one-loop diagram shown in Fig. 1.

Due to its economic structure the scotogenic model has been studied in many publications in the literature. Both DM candidates, the lightest N_1 [3] and the neutral component of η [4], can give a relic density consistent with DM abundance in viable parts of parameter space, while respecting bounds from direct [5,6] and indirect [7,8] detection.

On the other hand, for any model to generate the correct dark matter abundance observed today requires the DM candidate to be stable from the time of its decoupling from the thermal bath in the early universe onwards. The scotogenic model is no exception: Breaking the Z_2 at this—or any later—time would lead to a rapid decay of the DM candidate and consequently, both the DM and the motivation to study the model would be lost. The authors of [9] have pointed out that when evolving the parameters of the scotogenic model to larger energies using RGEs, very often points which had a conserved Z_2 at low energy develop a deeper minimum at a nonzero value of the vev of η at high energies and, thus, Z_2 is spontaneously broken. This is induced by a negative contribution arising from a one-loop correction to the mass of the Z_2 -odd scalar doublet mediated by the Z_2 -odd Majorana right-handed heavy neutrinos, N . It was then argued in [9], that all such points should be excluded from the parameter space for consistency.

However, the early universe is a hot thermal bath and temperature effects, while briefly discussed, were not taken into account in a consistent manner in [9]. It is well-known that the addition of thermal effects is responsible of the electroweak symmetry restoration at temperatures around 165 GeV [10,11]. One then would expect that thermal contributions should mitigate or even avoid these Z_2 breaking minima at high scales. In this paper, we

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FIG. 1. One-loop neutrino mass diagram in the scotogenic model.

discuss how to build the effective potential of the scotogenic model [1] including one-loop thermal and nonthermal effects, as well as the corresponding resummation terms. We then *improved* the potential by taking into account the RGE evolution of the parameters, in order to consider the possibility that the Z_2 is spontaneously broken at a high energy scale, which is characterized by the temperature T . As a result, we find that the inclusion of thermal effects change the conclusion about the stability of the DM in the early universe. With the possible exception of parameter points close to nonperturbativity, where we cannot trust our calculation, we found no points in our scans in which the Z_2 is broken spontaneously. In particular, our conclusions hold in the parameter region studied in [9].

The rest of this paper is organized as follows. In Sec. II we give a description of the scotogenic model and its parameters. We also discuss briefly the tree-level stability conditions and the constraints these put on the parameter choices for the scalar potential. Section III is devoted to a discussion of the effective potential, temperature effects and the correct decoupling of the heavy degrees of freedom. In Sec. IV we discuss our numerical results. The paper then closes with a short summary.

II. THE SCOTOGENIC MODEL

This section gives a brief description of the scotogenic model and a discussion of its scalar potential. The model is a simple extension of the SM, to which three copies of right-handed neutrinos, $N \propto (1, 1, 0)$, and a new scalar doublet, $\eta \propto (1, 2, 1/2)$, are added. Here, (x, y, z) indicate quantum numbers under the SM group in the usual order $SU(3)_C \times SU(2)_L \times U(1)_Y$. N_i and η are assumed odd under a new Z_2 symmetry. The Lagrangian of the model then contains the terms:

$$\mathcal{L} \supset -(Y_N)_{\alpha\beta} L_\alpha \eta^\dagger N_\beta - \frac{1}{2} M_{N_\alpha} N_\alpha N_\alpha - \mathcal{V} \quad (1)$$

$$\begin{aligned} \mathcal{V} = & \mu_1^2 H^\dagger H + \frac{1}{2} \lambda_1 (H^\dagger H)^2 + \mu_2^2 \eta^\dagger \eta + \frac{1}{2} \lambda_2 (\eta^\dagger \eta)^2 \\ & + \lambda_3 (H^\dagger H)(\eta^\dagger \eta) + \lambda_4 (H^\dagger \eta)(\eta^\dagger H) \\ & + \frac{1}{2} [\lambda_5 (H^\dagger \eta)^2 + \text{H.c.}] \end{aligned} \quad (2)$$

Here, the fermions are described in terms of 2-component spinors and H is the usual SM Higgs doublet. Note that the Z_2 symmetry eliminates terms such as $m_{12}^2 H^\dagger \eta$. We will assume in the following that λ_5 is real, as this does not affect any of our findings.

Neutrino masses arise at one-loop order as shown in Fig. 1. The 3×3 Majorana neutrino mass matrix from this diagram is given by

$$\begin{aligned} (m_{\text{SC}})_{\alpha\beta} = & \sum_{i=1}^3 \frac{Y_{Ni\alpha} Y_{Ni\beta}}{2(4\pi)^2} M_{N_i} \left[\frac{m_R^2}{m_R^2 - M_{N_i}^2} \log \left(\frac{m_R^2}{M_{N_i}^2} \right) \right. \\ & \left. - \frac{m_I^2}{m_I^2 - M_{N_i}^2} \log \left(\frac{m_I^2}{M_{N_i}^2} \right) \right] \\ \equiv & \frac{1}{32\pi^2} (Y_N^T \hat{M} Y_N)_{\alpha\beta}, \end{aligned} \quad (3)$$

where $m_{R/I}^2$ are the masses of the neutral scalar and pseudoscalar components of η :

$$m_R^2 = \mu_2^2 + \frac{1}{2} \lambda_L v_1^2, \quad m_I^2 = \mu_2^2 + \frac{1}{2} \lambda_S v_1^2, \quad (4)$$

with $v_1 = \sqrt{2} \langle H \rangle$, $\lambda_{L,S} = \lambda_3 + \lambda_4 \pm \lambda_5$. Equation (3) is of a form which allows to use a minimally modified version [12] of the Casas-Ibarra parametrization [13] to fit all neutrino data. Since such fits have been discussed many times in the literature, we will not repeat any details here. We stress the importance of the term proportional to λ_5 in (2). In the limit $\lambda_5 \rightarrow 0$ the model conserves lepton number and the Majorana neutrino mass diagram vanishes identically.¹ Since this corresponds to an enhanced symmetry, a small value of λ_5 is technically natural and somewhat small values of λ_5 are usually used in the neutrino mass fit, allowing the Yukawas to be even order $\mathcal{O}(1)$.

We are interested in the question whether Z_2 breaking points can appear during the thermal evolution of the universe, which for temperature $T = 0$ would nevertheless appear to be Z_2 symmetric. For this we have to take into account the thermal contributions to the scalar potential as discussed in the subsequent section. These corrections can be expressed in terms of tree-level masses for which we give here the corresponding formulas allowing for a potential spontaneous Z_2 breaking. We decompose the scalar doublets as follows

¹In more technical terms, one sees from (4) that $\lambda_5 = 0$ implies $m_R^2 = m_I^2$, which in turn gives $\hat{M} = 0$.

$$\begin{aligned} H^T &= \left(G^+, \frac{1}{\sqrt{2}}(v_1 + h + iG_0) \right), \\ \eta^T &= \left(\eta^+, \frac{1}{\sqrt{2}}(v_2 + \eta_R + i\eta_I) \right), \end{aligned} \quad (5)$$

where we explicitly allow for a possible nonzero vacuum expectation value for the Z_2 -odd field η . We have slightly abused notation here, as in case of a Z_2 breaking the would-be Goldstone fields are an admixture of G^+ and η^+ , as well as an admixture of G^0 and η_I .

We get the background potential by expanding (2):

$$V_0(v_1, v_2) = \frac{1}{2}\mu_1^2 v_1^2 + \frac{1}{2}\mu_2^2 v_2^2 + \frac{1}{8}\lambda_1 v_1^4 + \frac{1}{8}\lambda_2 v_2^4 + \frac{1}{4}\lambda_L v_1^2 v_2^2. \quad (6)$$

We briefly summarize here the tree-level conditions leading to the correct electroweak symmetry breaking while respecting the Z_2 symmetry, for details see, e.g., [14] and references therein. There are in principle four possible minima²:

- (1) $v_1^2 = v_2^2 = 0$;
- (2) $v_1^2 = -\frac{2\mu_1^2}{\lambda_1}$, $v_2^2 = 0$;
- (3) $v_1^2 = 0$, $v_2^2 = -\frac{2\mu_2^2}{\lambda_2}$;
- (4) $v_1^2 = -2\frac{\mu_1^2\lambda_2 - \mu_2^2\lambda_L}{\lambda_1\lambda_2 - \lambda_L^2}$, $v_2^2 = -2\frac{\mu_2^2\lambda_1 - \mu_1^2\lambda_L}{\lambda_1\lambda_2 - \lambda_L^2}$.

The corresponding minimum should fulfill the conditions: $v_{1,2}^2 \geq 0$, the Hessian should be positive, i.e., $\det[\frac{\partial^2 V_0}{\partial v_1 \partial v_2}] > 0$. The potential at minimum 2 should be the lowest of the four: Taking $\lambda_1, \lambda_2 > 0$, as required to avoid unbounded from below directions, one finds beside $\mu_1^2 < 0$ the conditions,

$$\lambda_1 < \lambda_L \frac{\mu_1^2}{\mu_2^2} \quad \text{if } \mu_2^2 < 0 \text{ and } \lambda_L > 0, \quad (7a)$$

$$\lambda_1 > \lambda_L \frac{\mu_1^2}{\mu_2^2} \quad \text{if } \mu_2^2 > 0 \text{ and } \lambda_L < 0. \quad (7b)$$

Moreover, minimum 2 is always the global one if $\lambda_L > 0$ and $\mu_2^2 > 0$. The requirement, that the potential is bounded from below also implies [16],³

$$0 < \sqrt{\lambda_1\lambda_2} + \lambda_3 + \lambda_4 - |\lambda_5| \quad (8a)$$

$$0 < \lambda_3 + \sqrt{\lambda_1\lambda_2} \quad (8b)$$

²There is also the possibility of a charge breaking local minimum which, however, has a value larger than the correct global electroweak symmetry breaking one [15].

³The conditions here listed are only the necessary conditions, but sufficient for our purposes. For a more complete set see, e.g., [17,18].

It has been shown in [9] that due to the evolution of the renormalization group equations (RGEs) the parameters get changed such that the desired Z_2 gets spontaneously broken at high energy scales. The most relevant RGE in this context is the one for μ_2^2 :

$$\begin{aligned} Q \frac{d\mu_2^2}{dQ} &= \frac{1}{16\pi^2} \left[\left(6\lambda_2 - \frac{3}{2}(g_Y^2 + 3g_L^2) + 2\text{tr}(Y_N^\dagger Y_N) \right) \mu_2^2 \right. \\ &\quad \left. + 2(2\lambda_3 + \lambda_4)\mu_1^2 - 4 \sum_{i=1}^3 M_{N_i}^2 (Y_N Y_N^\dagger)_{ii} \right] \end{aligned} \quad (9)$$

where Q denotes the renormalization scale. In particular, the last contribution proportional to $M_{N_i}^2$ can drive μ_2^2 to negative values if the corresponding Yukawa couplings are sizable and the larger the ratios $M_{N_i}^2/\mu_2^2$ are.

We will investigate to which extent this implies that the Z_2 symmetry could be broken at high energies during the thermal evolution of the Universe. For this, one has to take into account the thermal contributions to the potential and check if there are indeed parameter points where the Z_2 breaking minimum could potentially arise at a high temperature T . In such a minimum, also v_2 would be nonzero and we give here the corresponding tree-level mass matrices from which we calculated the input to the one-loop contributions discussed in the next section. The mass matrices for the CP -odd, -even, and charge scalars in the basis (H, η) are given by:

$$\mathcal{M}_R^2 = \begin{pmatrix} \mu_1^2 + \frac{3}{2}\lambda_1 v_1^2 + \frac{1}{2}\lambda_L v_2^2 & \lambda_L v_1 v_2 \\ \lambda_L v_1 v_2 & \mu_2^2 + \frac{3}{2}\lambda_2 v_2^2 + \frac{1}{2}\lambda_L v_1^2 \end{pmatrix}, \quad (10)$$

$$\mathcal{M}_I^2 = \begin{pmatrix} \mu_1^2 + \frac{1}{2}\lambda_1 v_1^2 + \frac{1}{2}\lambda_S v_2^2 & \lambda_5 v_1 v_2 \\ \lambda_5 v_1 v_2 & \mu_2^2 + \frac{1}{2}\lambda_2 v_2^2 + \frac{1}{2}\lambda_S v_1^2 \end{pmatrix}, \quad (11)$$

$$\mathcal{M}_\pm^2 = \begin{pmatrix} \mu_1^2 + \frac{1}{2}\lambda_1 v_1^2 + \frac{1}{2}\lambda_3 v_2^2 & \frac{1}{2}(\lambda_4 + \lambda_5)v_1 v_2 \\ \frac{1}{2}(\lambda_4 + \lambda_5)v_1 v_2 & \mu_2^2 + \frac{1}{2}\lambda_2 v_2^2 + \frac{1}{2}\lambda_3 v_1^2 \end{pmatrix}, \quad (12)$$

where the lightest eigenvalues of $\mathcal{M}_I^2$ and $\mathcal{M}_\pm^2$ need to be zero when evaluated at any of the minima, since they correspond to the Goldstones of electroweak symmetry breaking. It is worth noticing that in the unbroken Z_2 phase $\lambda_4 - |\lambda_5| < 0$ when N is heavier than all scalars to avoid an electrically charged stable dark matter candidate. The masses of the electroweak gauge bosons are given by

$$m_W^2 = \frac{1}{4}g_L^2(v_1^2 + v_2^2), \quad m_Z^2 = \frac{1}{4}(g_L^2 + g_Y^2)(v_1^2 + v_2^2), \quad m_\gamma^2 = 0.$$

The standard model charged fermion masses are simply calculated from the SM Higgs Yukawa couplings

as, $m_f^2 = \frac{1}{2} Y_f^2 v_1^2$ whereas for $v_2 \neq 0$ a tree-level type-I seesaw contribution to the neutrino masses appears, given by:

$$M_\nu = \begin{pmatrix} 0 & m_D \\ m_D & M_N \end{pmatrix} \quad \text{with} \quad m_D = \frac{1}{\sqrt{2}} Y_N v_2.$$

III. EFFECTIVE POTENTIAL, TEMPERATURE EFFECTS AND DECOUPLING OF HEAVY DEGREES OF FREEDOM

We will combine the one-loop RGEs with the effective potential and include also the thermal contributions. The region of parameter space, where potentially a Z_2 breaking could occur, is where the right-handed neutrinos are heavier than the additional scalars and have sizable neutrino Yukawa couplings [9]. The mass hierarchy between N_i and η can be quite large. This implies that one has to decouple the heavy states properly as outlined for example in [19]. We will briefly discuss this after presenting the thermal contributions. In the following we will neglect the contributions from SM fermions except for the top quark, as they do not play any role here.

The effective potential at one-loop and zero temperature is given by $V_{\text{eff}}^{(0)} = V_0 + V_1$, where the temperature-independent one-loop correction to the potential is the Coleman-Weinberg potential is given by,

$$V_1 = \frac{1}{64\pi^2} \sum_i n_i m_i^4 \left(\ln \frac{m_i^2}{Q^2} - c_i \right). \quad (13)$$

Here, m_i^2 are the field dependent mass squares of the various particles and the n_i count the degrees of freedom and the sign is positive for bosons and negative for fermions: $n_Z = 3$, $n_W = 6$, $n_\gamma = 2$, $n_{S^0} = 1$, $n_{S^\pm} = 2$, $n_t = -12$ and $n_{N_i} = -2$. S^0 and $S^\pm$ refers to any neutral or charged scalar. We take the Landau gauge ($\xi = 0$) and work in the $\overline{\text{MS}}$ scheme [20,21]. Thus, one has $c_i = 3/2$ in case of scalars and fermions and $c_i = 5/6$ for the vector bosons.

We note for completeness that (13) depends on the choice of the renormalization scale Q . To minimize its effect one should use the improved effective potential by using the renormalization group, i.e.,

$$\left[Q \frac{\partial}{\partial Q} + \beta_i \frac{\partial}{\partial \lambda_i} - \gamma_{m_i} m_i^2 \frac{\partial}{\partial m_i^2} - \gamma_{v_i} v_i \frac{\partial}{\partial v_i} \right] V_{\text{eff}}^{(0)}(Q, \lambda_i, m_i^2, v_i) = 0. \quad (14)$$

The solution is given by $V_{\text{eff}}^{(0)}(Q, \lambda_i, m_i^2, v_i) = V_{\text{eff}}^{(0)}(Q, \lambda_i(Q), m_i^2(Q), v_i(Q))$, with $\lambda_i(Q)$, $m_i^2(Q)$ and $v_i(Q)$ the running parameters [22,23].

The effective potential including finite temperature effects is defined as $V_{\text{eff}}^{(T)} = V_0 + V_1 + V_T$ at one-loop level. Here, the leading order temperature-dependent correction is given by,

$$V_T = \frac{T^4}{2\pi^2} \left[\sum_{i \in \text{bosons}} n_i J_B(m_i^2/T^2) + \sum_{i \in \text{fermions}} n_i J_F(m_i^2/T^2) \right], \quad (15)$$

where the functions $J_{F,B}(y^2)$ come from the evaluation of the thermal loop with the corresponding statistic for bosons or fermions and they are defined as,

$$J_{F,B}(y^2) = \int_0^\infty dx x^2 \ln [1 \pm e^{-\sqrt{x^2+y^2}}], \quad (16)$$

see, e.g., [21] and references therein.

One should take the thermal mass $m(v_i, T)$ to avoid infrared divergences for $T \gg m$. The thermal mass is only important for the zero Matsubara mode [21], so only boson masses have to be modified according to $m_i^2(v_i) \rightarrow \Pi_i^2(v_i, T) \equiv m_i^2 + k_i T^2$.

$$\Pi_{W_L}^2 = m_{W_L}^2 + 2g_L^2 T^2, \quad (17)$$

$$\begin{aligned} \Pi_{\mu_1}^2 = & \mu_1^2 + \left[\frac{1}{8} g_L^2 + \frac{1}{16} (g_L^2 + g_Y^2) + \frac{1}{4} \lambda_1 + \frac{1}{6} \lambda_3 \right. \\ & \left. + \frac{1}{12} \lambda_4 + \frac{1}{4} Y_t^2 \right] T^2, \end{aligned} \quad (18)$$

$$\begin{aligned} \Pi_{\mu_2}^2 = & \mu_2^2 + \left[\frac{1}{8} g_L^2 + \frac{1}{16} (g_L^2 + g_Y^2) + \frac{1}{4} \lambda_2 + \frac{1}{6} \lambda_3 + \frac{1}{12} \lambda_4 \right. \\ & \left. + \frac{1}{24} \text{Tr}(Y_N^\dagger Y_N) \right] T^2, \end{aligned} \quad (19)$$

$$\Pi_{Z_L}^2 = \frac{1}{2} [m_Z^2 + \Delta(m_Z^2, T^2)] + (g_L^2 + g_Y^2) T^2, \quad (20)$$

$$\Pi_{\gamma_L}^2 = \frac{1}{2} [m_Z^2 - \Delta(m_Z^2, T^2)] + (g_L^2 + g_Y^2) T^2, \quad (21)$$

with $\Delta^2 = m_Z^4 + (g_L^2 - g_Y^2)^2 (4T^2 + v_1^2 + v_2^2) T^2$. Note, that in case of the gauge bosons only the longitudinal components get a thermal contribution from the zero Matsubara mode. The method described so far is known as the Parwani method [24]. Alternatively, one could use the so-called Arnold-Espinosa method [25] which is based on a resummation of daisy-diagrams [25]. Both approaches have been compared for example in [26], where it has been found that they lead to very similar numerical results in the context of singlet extensions of the SM, whereas the results of [27] indicate that the Arnold-Espinosa method leads to stronger

constraints on the parameter space in the context of two Higgs doublets models.

The renormalization scale Q should be chosen such, that large logarithms are avoided in the calculation of the loop corrected effective potential in (13). We use here the $\overline{\text{MS}}$ scheme where one has to decouple heavy states by hand. In view of the fact that a potential Z_2 breaking requires $M_{N_i}^2/\mu_2^2 \gtrsim 1$, we have chosen the following procedure: assuming that the additional scalars have masses in the range of a few hundred GeV we decouple the right-handed neutrinos if their mass is above max (1 TeV, T). The idea here is that right-handed neutrinos will only contribute if the temperature T is sufficiently high. The renormalization scale is then chosen to be:

$$Q = \sqrt[2n]{\prod_{k=1}^n m_k^2(v_i, T)}, \quad (22)$$

where $3 \leq n \leq 6$ depending on how many neutrinos decoupled, for instance $n = 3$ if all of them are decoupled and 6 if none is decoupled. In the RGE evolution, as well as in the calculation of the effective potential including the thermal effects, the right-handed neutrinos are included if their masses $m_{N_i} \leq Q$, otherwise they are decoupled. We have checked that the results presented in the next section remain the same if we replace this geometric mean for Q by the maximum of the masses which are taken into account.

IV. NUMERICAL RESULTS

As already explained in the previous sections, the improvement of the RGEs to the effective potential implies that the potential only depends on four scales that can be related: the renormalization scale Q , the temperature T and the background fields values v_1 and v_2 . To keep the logarithmic corrections under control, we considered the renormalization scale as a function of the thermal mass (22). Then, for each set of input values, we numerically minimized the effective potential in terms of v_1 and v_2 for different values of the temperature ranging from 0 to 10^{16} GeV, building our way back in time in the evolution of the universe. The mass spectrum and the one-loop RGEs are numerically computed by a modified version of SPHENO [28,29] generated by SARA [30,31]. The minimization of the thermal effective potential is done using MINUIT [32]. Certain checks and cuts are applied during the procedure regarding unboundedness from below and nonperturbativity. The initial values should generate a valid electroweak minimum, i.e., fulfill the conditions (7). Also, we only considered points that satisfy the bounded from below conditions (8) at every step of the RGE running. One should be especially cautious about this point, as the well-known bounded from below conditions of the scotogenic or the IDM model are just necessary conditions, but not

sufficient.⁴ Then, we checked possible nonperturbative couplings, as depending on the initial values at low scale, the RGE running can drive the couplings outside the perturbative regime. Our computation of the effective potential cannot be trusted in such cases, so we omitted these points.

In Fig. 2 we show the behavior of the v_1 and v_2 minima for a benchmark point as a function of the energy scale. We consider three different RGE-improved potentials: tree-level potential only, one-loop zero-temperature potential and thermal effective potential. For this point, μ_2^2 runs into negative values around 10^5 – 10^6 GeV (9). We see that when considering only the tree-level potential the vev of η becomes nonzero around this scale, when the condition (7a) fails. This contrasts with the behavior of the one-loop potential or the thermal potential, where no Z_2 -breaking minimum is found. Once the correct treatment of the potential with the effective one-loop improved potential is taken into account, the spontaneous breaking of Z_2 arising from the RGEs disappear. For completeness, in the last panel we show a *zoom* of the region of interest for the last potential just to compare with the previous case. As stated in [10,11] the symmetry restoration should occur at temperatures around 150–170 GeV, which is what we found here. This leads to a merging of the global and local minima at $v_{1,2} = 0$.

We scanned over the input BSM parameters $(Y_N)_{ij}$, M_{N_k} , μ_2 and $\lambda_{2,3,4}$ searching for a Z_2 breaking minima at any step in the thermal evolution. For simplicity, we considered the mass hierarchy $M_{N_k} = (1, 1.5, 2) \times M_n$, as well as all the entries of the diagonal Yukawa matrix Y_N to be the same and equal to Y_n . We randomly generate 5×10^4 points in logarithmic scale in the ranges,

$$\begin{cases} |\lambda_{2,3,4}|: [10^{-5}, 1] \\ |Y_n|: [10^{-5}, 1] \\ \mu_2: [10^2, 10^3] \text{ GeV} \\ M_n: [10^2, 10^4] \text{ GeV} \end{cases} \quad (23)$$

Note that λ_5 is not among the input parameters, as the neutrino mass (3) can be used to fix its value. All input parameters are given at the scale $Q = 160$ GeV [$\simeq m_t(m_t)$]. As λ_3 and λ_4 can be negative, the sign is also randomized. In the RGE evolution the top-Yukawa coupling Y_t and the strong coupling g_3 are numerically important, in particular for λ_1 . We have set their values at $Q = 160$ GeV to $Y_t = 0.93761$ and $g_3 = 1.16787$. Dimensionless parameters are conservatively taken up to 1, as even for this value the running will drive the parameter to the nonperturbative

⁴We indeed found that from the set of randomly generated input values, around a 0.05% led to an unbounded from below potential even passing the conditions (8). Of course, such points were excluded from our result.

FIG. 2. Comparison of the behavior of the vevs at the minimum for a benchmark point for different RGE-improved potentials. From top to bottom: tree-level potential (top), one-loop corrected potential, and the thermal effective potential. The last panel zooms into the region where the symmetry is restored as a function of the energy scale (left) and temperature (right). The values for v_1 and v_2 are given in GeV. Values of the parameters: $(\lambda_2, \lambda_3, \lambda_4, Y_n, \mu_2^2, M_n) = (0.2, 0.2, -0.15, 0.2, 3 \text{ TeV}^2, 1 \text{ TeV})$.

regime, $\sqrt{4\pi}$ in our case, at a very low scale. μ_2 should not be very large to maximize the possibility of getting a Z_2 breaking minimum, but cannot be very low, as there are already constraints to the mass of η around 100 GeV [33]. Regarding M_n , naively one would expect that its effect gets larger the larger it is. However, a large mass means that it gets effective at higher temperatures implying that it contributes only at a later stage to the RGE evolution. We are thus interested in making this value large to drive μ_2 faster toward negative values when doing the running, see (9), but not too large so right-handed neutrinos couple at sufficiently low energies before the thermal contribution becomes dominant. For each set of input values we computed the RGE-improved effective thermal potential and minimized it

with respect to v_1 and v_2 for different temperatures. Finally, from the list of minima for each temperature, we search for Z_2 symmetry breaking minima, i.e., $|v_2| > 0$ (actually, larger than the numerical uncertainty of MINUIT). We have considered here temperatures T up to 10^{16} GeV in the cases where all couplings remain perturbative. We have stopped at lower temperatures if the modulus of one of the quartic couplings exceeds $\sqrt{4\pi}$ as latest at this stage perturbation theory breaks down. As can be seen in Fig. 3, we found no Z_2 breaking points in the parameter space considered.

Figure 3 also shows, for comparison, the points which break Z_2 symmetry according to [9]. The black dots are points within our scan that would be considered to be Z_2 symmetry breaking in [9], but not in our approach, using

FIG. 3. Results of the scan over the parameter space of the scotogenic model. No Z_2 -breaking minimum was found at any step of the thermal evolution for the parameter space we studied. Green dots represent points which do not break Z_2 symmetry when the effective thermal potential is considered. Black dots are Z_2 -breaking according to [9], but not in our approach.

the whole thermal effective potential (see Fig. 2 and explanations therein). Note that in [9] the authors evaluated and minimized only the tree-level potential at different steps of the RGE running, finding that the RGEs can drive μ_2^2 toward large negative values generating a Z_2 breaking minimum. This behavior is clear in Fig. 3, the negative contribution to the running of μ_2^2 is proportional to M_n^2 and Y_n^2 , so the black points appear toward larger values of M_n and Y_n , and of course lower initial values of μ_2 .

V. CONCLUSIONS

In this work we have studied the effects of temperature on the Z_2 symmetry breaking in the scotogenic model. It had previously been shown [9] that in the running of the parameters of the scotogenic model to high energy, the mass squared parameter for the inert doublet turns negative in sizable parts of parameter space. It was concluded then in [9] that such points should be excluded from the parameter space, since dark matter would be lost due to the spontaneous breaking of the Z_2 symmetry by the vacuum expectation value of η .

Here, however, we have shown that this conclusion was premature. For a more realistic calculation, we have constructed the effective potential of the scotogenic model,

taking special care on the inclusion of thermal effects. We then take into account the RGE evolution of parameters, including the correct decoupling of the right-handed neutrinos in case that they are heavy. Our numerical scans over the parameter space then show, that no spontaneous Z_2 breaking occurs in the parts of parameter space studied in [9]. Our scans are quite general, but since our numerical programs fail for points very close to nonperturbativity, we cannot make any claims for that part of parameter space. Models with such large quartic couplings have been investigated for example in [34,35] in the context of two-step electroweak phase transitions. With this possible caveat, however, our main result is that considering the complete one-loop contribution to the effective potential with the thermal corrections changes the conclusion about the stability of dark matter in the early universe.

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- [1] E. Ma, *Phys. Rev. D* **73**, 077301 (2006).
- [2] P. F. de Salas, D. V. Forero, S. Gariazzo, P. Martínez-Miravé, O. Mena, C. A. Ternes, M. Tórtola, and J. W. F. Valle, *J. High Energy Phys.* **02** (2021) 071.
- [3] D. Suematsu, T. Toma, and T. Yoshida, *Phys. Rev. D* **79**, 093004 (2009).
- [4] L. Lopez Honorez, E. Nezri, J. F. Oliver, and M. H. G. Tytgat, *J. Cosmol. Astropart. Phys.* **02** (2007) 028.
- [5] C. Guo, S.-Y. Guo, and Y. Liao, *Chin. Phys. C* **43**, 103102 (2019).
- [6] D. Schmidt, T. Schwetz, and T. Toma, *Phys. Rev. D* **85**, 073009 (2012).
- [7] T. de Boer, R. Busse, A. Kappes, M. Klasen, and S. Zeinstra, *J. Cosmol. Astropart. Phys.* **08** (2021) 038.
- [8] T. de Boer, R. Busse, A. Kappes, M. Klasen, and S. Zeinstra, *Phys. Rev. D* **103**, 123006 (2021).
- [9] A. Merle and M. Platscher, *Phys. Rev. D* **92**, 095002 (2015).
- [10] M. Dine, R. G. Leigh, P. Y. Huet, A. D. Linde, and D. A. Linde, *Phys. Rev. D* **46**, 550 (1992).
- [11] A. Katz and M. Perelstein, *J. High Energy Phys.* **07** (2014) 108.
- [12] I. Cordero-Carrión, M. Hirsch, and A. Vicente, *Phys. Rev. D* **99**, 075019 (2019).
- [13] J. A. Casas and A. Ibarra, *Nucl. Phys.* **B618**, 171 (2001).
- [14] P. M. Ferreira, L. A. Morrison, and S. Profumo, *J. High Energy Phys.* **04** (2020) 125.
- [15] P. M. Ferreira, R. Santos, and A. Barroso, *Phys. Lett. B* **603**, 219 (2004); **629**, 114(E) (2005).
- [16] K. G. Klimenko, *Theor. Math. Phys.* **62**, 58 (1985).
- [17] K. Kannike, *Eur. Phys. J. C* **72**, 2093 (2012).
- [18] K. Kannike, *Eur. Phys. J. C* **76**, 324 (2016); **78**, 355(E) (2018).
- [19] A. V. Manohar and E. Nardoni, *J. High Energy Phys.* **04** (2021) 093.
- [20] S. P. Martin, *Phys. Rev. D* **65**, 116003 (2002).
- [21] M. Laine and A. Vuorinen, *Lect. Notes Phys.* **925**, 1 (2016).
- [22] M. Bando, T. Kugo, N. Maekawa, and H. Nakano, *Phys. Lett. B* **301**, 83 (1993).
- [23] M. Quiros, in *ICTP Summer School in High-Energy Physics and Cosmology, Trieste, Italy, 1998*, edited by A. Masiero, G. Senjanovic, and A. Smirnov (World Scientific, Singapore, 1999), pp. 187–259.
- [24] R. R. Parwani, *Phys. Rev. D* **45**, 4695 (1992); **48**, 5965(E) (1993).
- [25] M. E. Carrington, *Phys. Rev. D* **45**, 2933 (1992).
- [26] G.-C. Cho, C. Idegawa, and E. Senaha, *Phys. Lett. B* **823**, 136787 (2021).
- [27] P. Basler, M. Krause, M. Muhlleitner, J. Wittbrodt, and A. Wlotzka, *J. High Energy Phys.* **02** (2017) 121.
- [28] W. Porod, *Comput. Phys. Commun.* **153**, 275 (2003).
- [29] W. Porod and F. Staub, *Comput. Phys. Commun.* **183**, 2458 (2012).
- [30] F. Staub, *Comput. Phys. Commun.* **185**, 1773 (2014).
- [31] F. Staub, *Adv. High Energy Phys.* **2015**, 840780 (2015).
- [32] F. James and M. Roos, *Comput. Phys. Commun.* **10**, 343 (1975).
- [33] P. A. Zyla *et al.* (Particle Data Group), *Prog. Theor. Exp. Phys.* **(2020)**, 083C01.
- [34] D. Curtin, P. Meade, and C.-T. Yu, *J. High Energy Phys.* **11** (2014) 127.
- [35] S. Fabian, F. Goertz, and Y. Jiang, *J. Cosmol. Astropart. Phys.* **09** (2021) 011.