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Electroweak phase transition in the reduced minimal 3-3-1 model

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The electroweak phase transition is considered in the framework of the reduced minimal 3-3-1 model. The structure of the phase transition in this model is divided into two parts. The first part is the phase transition SU(3) → SU(2) at the TeV scale, and the second is SU(2) → U(1), which is like the standard model electroweak phase transition. When the mass of the neutral Higgs boson (h_1) is taken to be equal to the LHC value, $m_{h_1} = 125$ GeV, these phase transitions are first-order transitions; the mass of Z_2 is about 4.8 TeV, and we find the region of parameter space with the first-order phase transition at the $v_{\rho_0} = 246$ GeV scale, leading to an effective potential where the mass of the charged Higgs boson is in the range 3.258 TeV $< m_{h_{++}} <$ 19.549 TeV. Therefore, with this approach new bosons are the triggers of the first-order electroweak phase transition, which provides significant implications for the viability of electroweak baryogenesis scenarios.

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I. INTRODUCTION

An electroweak phase transition (EWPT) is a type of symmetry-breaking phase transition that plays an important role at the early stage of an expanding universe. Particularly, the EWPT is necessary to explain the baryon asymmetry of our Universe. As proposed by Sakharov [1], there are three necessary conditions that a baryon-generating interaction in a theoretical model must satisfy to produce an excess of baryons over antibaryons: baryon-number violation, C and CP violations, and a deviation from thermal equilibrium [1].

If baryon number (B) is conserved and is equal to zero, it will be equal to zero forever. In contrast, B will vanish in a state of thermal equilibrium. Therefore we need the third condition regarding the deviation from thermal equilibrium. The second condition is appropriate for ensuring a different decay rate for particles and antiparticles [1].

The baryon number and C and CP violations can be seen via the sphaleron rate and the Cabibbo-Kobayashi-Maskawa (CKM) matrix in various models [2]. The sphaleron rate tells us about the baryon-number violation, and the nonzero phases of the CKM matrix tell us about CP violation.

It is well known that in order to ensure that the third condition is satisfied, deviations from thermal equilibrium should be large enough, and therefore the EWPT should be a first-order phase transition. The EWPT is the transition from a symmetric phase to an asymmetric one that is needed to generate masses for particles. Therefore, the EWPT is related to the mass of the Higgs boson [1].

In the basic model of particles, the first and second conditions can be satisfied, but any conditions regarding thermal imbalance are difficult to satisfy. So at present an analysis of the third condition is the only approach that can explain the baryon asymmetry.

Why must the EWPT be a first-order phase transition? The effective potential is a function of temperature and vacuum expectation values (VEVs). For very large temperatures it only has one minimum at zero, and the symmetry is restored. As the temperature goes to T_0, a nonzero second minimum appears; this is the sign of symmetry breaking. When the temperature reaches the critical temperature ($T_c < T_0$), the values of the effective potential at the two minimums are equal, and the symmetry breaking is turned on. And at the critical temperature, if the two minimums are separated by a potential barrier, the phase transition will occur with bubble nucleations. Inside the bubbles, the scalar field stores a nonzero expectation value. If the bubble nucleation rate exceeds the universe's expansion rate, the bubbles collide and eventually fill all space [1]. Such a transition is called a first-order phase transition. It is very violent and one can expect large deviations from thermal equilibrium [1]. The other possible scenario takes place if the two minimums are never separated by a potential barrier. In this case, the phase transition is a smooth transition, rather than a violent or second-order phase transition.

To study the EWPT, ones consider the high-temperature effective potential as follows:

$$V_{\rm eff} = D(T^2 - {T'^2}_0)v^2 - ETv^3 + \frac{\lambda_T}{4}v^4,$$

where v is the VEV of the Higgs. In order to have a first-order phase transition, the strength of the phase transition should be larger than unity, i.e., $\frac{v_c}{T_c} \geq 1$.

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The EWPT has been investigated in the standard model (SM) [1,2] and in various extended models [3–10]. Also, very interesting research has shown that dark matter may trigger the electroweak phase transition [11]. For the SM, the strength of the EWPT is larger than unity at the electroweak scale, but it appears to be too weak for the experimentally allowed mass of the SM scalar Higgs boson [1,2]; therefore, it seems that electroweak baryogenesis requires new physics beyond the SM at the weak scale [3].

Before the neutral Higgs was found, the study of phase transitions in most models focused on two basic issues: determining the order of the phase transition, and the mass of the neutral Higgs. For the SM, a first-order phase transition problem has one variable, which is the mass of the neutral Higgs boson. However, for the extended models, this problem has at least two variables: the first one is the Higgs mass, and the others include the masses of heavy particles. Recently, the neutral Higgs was discovered at the LHC [12], so the electroweak phase transition problem has been reduced by one variable. This provides hope for the ability of the extended models to examine the EWPT. The remarkable successes of previous surveys include the following:

(i) In the SM, the sources of *CP* violation are smaller than the baryon asymmetry of the Universe and there is no first-order phase transition due to the large mass of the neutral Higgs, i.e., the SM does not have enough triggers for the first-order phase transition to be turned on [1].

(ii) The extended models—such as the two-Higgs-doublet model or minimal supersymmetric standard model—can explain the baryon asymmetry because the sources of *CP* violation in these models are stronger than in the SM and they have a first-order phase transition, with the mass of the neutral Higgs being about 120 GeV. Triggers for the first-order phase transition in these models are heavy bosons or dark matter candidates [7–9,11].

Among the extended models, those based on the $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ gauge group (called 3-3-1 for short) [13,14] have some intriguing features, such as the ability to account for the generation problem [13,14], the quantization of the electric charge [15], etc. Given these features, we hope that the 3-3-1 models can also answer the problem of the baryon asymmetry in our Universe.

The current 3-3-1 models have many different forms, but they are all based on the above-mentioned gauge group. The greatest disadvantage of the 3-3-1 models is the complication in the Higgs sector, namely, that these models need at least three Higgs triplets to generate masses for fermions. Recently there have been attempts to solve this problem, and some models with the simplest Higgs sector [with only two $SU(3)_L$ Higgs triplets] have been constructed.

With such a group structure, the 3-3-1 models must have at least two Higgs triplets [16,17]. Therefore, the number of bosons in the 3-3-1 models will be much greater than that in the SM, and the symmetry breaking structure will different.

In the present paper, we consider the EWPT in the reduced minimal 3-3-1 (RM331) model [17] because of its simplicity. This model consists of the minimal leptonic content (i.e., only the SM leptons) and bileptons: the singly and a doubly charged gauge bosons $V^{\pm}$ and $U^{\pm\pm}$, the heavy neutral boson Z_2, and exotic quarks. This model also has two Higgs triplets. Therefore, the physical scalar spectrum of the RM331 model is composed of a doubly charged scalar h^{++} and two neutral scalars h_1 and h_2 [17]. These new particles and exotic quarks can be triggers for the first-order phase transition.

The plan of the paper is as follows. In Sec. II we give a review of the RM331 model with regards to the boson, lepton and Higgs sectors. In Sec. III we find the effective potential in the RM331 model, which has a contribution from heavy bosons and a contribution similar to that in the SM. In Sec. IV we calculate in detail the structure of the phase transition in the RM331 model, find the first-order phase transition, and show the constraints on the mass of the charged Higgs boson. Finally, we summarize and describe outlooks in Sec. V.

II. A REVIEW OF THE RM331 MODEL

The fermion content of the RM331 model is the same as that of the minimal 3-3-1 model [13]. The difference is only in the Higgs sector.

A. Higgs potential

The Higgs potential in the RM331 model [17] is given by

$$V(\chi, \rho) = \mu_1^2 \rho^\dagger \rho + \mu_2^2 \chi^\dagger \chi + \lambda_1 (\rho^\dagger \rho)^2 + \lambda_2 (\chi^\dagger \chi)^2 + \lambda_3 (\rho^\dagger \rho)(\chi^\dagger \chi) + \lambda_4 (\rho^\dagger \chi)(\chi^\dagger \rho). \tag{1}$$

The scalar sector contains only two Higgs scalar triplets [17],

$$\rho = \begin{pmatrix} \rho^+ \\ \rho^0 \\ \rho^{++} \end{pmatrix} \sim (\mathbf{3}, 1), \qquad \chi = \begin{pmatrix} \chi^- \\ \chi^{--} \\ \chi^0 \end{pmatrix} \sim (\mathbf{3}, -1). \tag{2}$$

The expansion of ρ^0 and χ^0 around their VEVs is usually

$$\rho^0, \chi^0 \to \frac{1}{\sqrt{2}} (v_{\rho,\chi} + R_{\rho,\chi} + i I_{\rho,\chi}). \tag{3}$$

This potential immediately gives us two charged Goldstone bosons $\rho^{\pm}$ and $\chi^{\pm}$, which are eaten by the gauge bosons $W^{\pm}$ and $V^{\pm}$.

We return now to the content of the Higgs sector. The physical scalar spectrum of the RM331 model is composed of a doubly charged scalar h^{++} and two neutral scalars h_1

and h_2. Since the lightest neutral field h_1 is basically an $SU(2)_L$ component, we identify it as the SM Higgs boson. In the effective limit $v_\chi \gg v_\rho$, the Higgs content can be summarized as follows:

$$\rho = \begin{pmatrix} G_{W^+} \\ \frac{v_\rho}{\sqrt{2}} + \frac{1}{\sqrt{2}}(h_1 + iG_Z) \\ h^{++} \end{pmatrix}, \qquad \chi = \begin{pmatrix} G_{V^-} \\ G_{U^{--}} \\ \frac{v_\chi}{\sqrt{2}} + \frac{1}{\sqrt{2}}(h_2 + iG_{Z'}) \end{pmatrix}, \tag{4}$$

where the Higgs masses are given by

$$M_{h_1}^2 = \left(\lambda_1 - \frac{\lambda_3^2}{4\lambda_2}\right) v_\rho^2, \qquad M_{h_2}^2 = \lambda_2 v_\chi^2 + \frac{\lambda_3^2}{4\lambda_2} v_\rho^2, \tag{5}$$

$$M_{h^{--}}^2 = \frac{\lambda_4}{2}(v_\chi^2 + v_\rho^2). \tag{6}$$

B. Gauge-boson sector

The masses of the gauge bosons appear in the Lagrangian part

$$\mathcal{L} = (\mathcal{D}_\mu \chi)^\dagger (\mathcal{D}^\mu \chi) + (\mathcal{D}_\mu \rho)^\dagger (\mathcal{D}^\mu \rho), \tag{7}$$

where

$$\mathcal{D}_\mu = \partial_\mu - igA_\mu^a \frac{\lambda^a}{2} - ig_X X \frac{\lambda_9}{2} B_\mu, \tag{8}$$

with $\lambda_9 = \sqrt{\frac{2}{3}}\mathrm{diag}(1,1,1)$, so that $\mathrm{Tr}(\lambda_9\lambda_9) = 2$. The couplings of $SU(3)_L$ and $U(1)_X$ satisfy the relation

$$\frac{g_X^2}{g^2} = \frac{6s_W^2}{1 - 4s_W^2},$$

where $c_W = \cos\theta_W$, $s_W = \sin\theta_W$, and $t_W = \tan\theta_W$, where θ_W is the Weinberg angle.

Substituting the expansion in Eq. (3) into Eq. (7) leads to the following result:

$$\begin{aligned} W^\pm &= \frac{A^1 \mp iA^2}{\sqrt{2}} \to m_{W^\pm}^2 = \frac{g^2 v_\rho^2}{4}, \\ V^\pm &= \frac{A^4 \pm iA^5}{\sqrt{2}} \to m_{V^\pm}^2 = \frac{g^2 v_\chi^2}{4}, \\ U^{\pm\pm} &= \frac{A^6 \pm iA^7}{\sqrt{2}} \to m_{U^{\pm\pm}}^2 = \frac{g^2(v_\rho^2 + v_\chi^2)}{4}. \end{aligned} \tag{9}$$

From Eq. (9), it follows that $v_\rho = 246$ GeV, and we obtain the relation

$$m_U^2 - m_V^2 = m_W^2.$$

In the neutral gauge-boson sector, with the basis $(A_\mu^3, A_\mu^8, B_\mu)$, the mass matrix is given by

$$M^2 = \frac{g^2}{4}\begin{pmatrix} v_\rho^2 & -\frac{v_\rho^2}{\sqrt{3}} & -2\kappa v_\rho^2 \\ -\frac{v_\rho^2}{\sqrt{3}} & \frac{1}{3}(v_\rho^2 + 4v_\chi^2) & \frac{2}{\sqrt{3}}(v_\rho^2 + 2v_\chi^2) \\ -2\kappa v_\rho^2 & \frac{2}{\sqrt{3}}(v_\rho^2 + 2v_\chi^2) & 4\kappa^2(v_\rho^2 + v_\chi^2), \end{pmatrix},$$

where $\kappa = \frac{g_X}{g}$. We can easily identify the photon field A_μ as well as the massive neutral Z and Z' bosons [18],

$$A_\mu = s_W A_\mu^3 + c_W\left(\sqrt{3}t_W A_\mu^8 + \sqrt{1 - 3t_W^2}B_\mu\right),$$

$$Z_\mu = c_W A_\mu^3 - s_W\left(\sqrt{3}t_W A_\mu^8 + \sqrt{1 - 3t_W^2}B_\mu\right),$$

and

$$Z'_\mu = -\sqrt{1 - 3t_W^2}A_\mu^8 + \sqrt{3}t_W B_\mu,$$

where the mass-squared matrix for $\{Z, Z'\}$ is given by

$$\begin{pmatrix} m_Z^2 & m_{ZZ'}^2 \\ m_{ZZ'}^2 & m_{Z'}^2 \end{pmatrix},$$

with

$$m_Z^2 = \frac{1}{4}\frac{g^2}{\cos^2\theta_W} v_\rho^2,$$

$$m_{Z'}^2 = \frac{1}{3}g^2\left[\frac{\cos^2\theta_W}{1 - 4\sin^2\theta_W} v_\chi^2 + \frac{1 - 4\sin^2\theta_W}{4\cos^2\theta_W} v_\rho^2\right],$$

$$m_{ZZ'}^2 = \frac{1}{4\sqrt{3}}g^2\frac{\sqrt{1 - 4\sin^2\theta_W}}{\cos^2\theta_W} v_\rho^2.$$

Diagonalizing the mass matrix gives the mass eigenstates Z_1 and Z_2, which can be taken as mixtures,

$$Z_1 = Z\cos\phi - Z'\sin\phi, \qquad Z_2 = Z\sin\phi + Z'\cos\phi.$$

The mixing angle ϕ is given by

$$\tan 2\phi = \frac{m_Z^2 - m_{Z_1}^2}{m_{Z_2}^2 - m_Z^2},$$

where m_{Z_1} and m_{Z_2} are the *physical* mass eigenvalues,

$$m_{Z_1}^2 = \frac{1}{2}\{m_{Z'}^2 + m_Z^2 - [(m_{Z'}^2 - m_Z^2)^2 - 4(m_{ZZ'}^2)^2]^{1/2}\},$$

$$m_{Z_2}^2 = \frac{1}{2}\{m_{Z'}^2 + m_Z^2 + [(m_{Z'}^2 - m_Z^2)^2 - 4(m_{ZZ'}^2)^2]^{1/2}\}.$$

When diagonalized, the mass of Z_1 is approximately proportional to v_ρ (because $v_\chi \gg v_\rho$), so Z_1 is like the neutral gauge boson Z in the SM. The mass of the new heavy boson Z_2 depends on v_ρ and v_χ; in addition, $m_W = 80.39$ GeV and $v_{\rho_0} = 246$ GeV [19]. Choosing $v_{\chi_0} = 4$ TeV [17,20], we obtain $m_V = 1307.15$ GeV and $m_U = 1309.62$ GeV. If we choose $s_W^2 = 0.23116$ [19], we derive $m_{Z_1} \approx m_Z = 91.68$ GeV and $m_{Z_2} = 4.821$ TeV. Hence we can approximate $m_{Z_2} \approx 1.2 v_\chi$.

C. Fermion sector

The fermion sector in the model under consideration is the same as in the minimal 3-3-1 model [13]. The Yukawa couplings give the exotic quark masses [17],

$$\begin{aligned} L_{\mathrm{Yuk}}^{\mathrm{exot}} &= \lambda_{11}^{T}\bar{Q}_{1L}\chi T_R + \lambda_{ij}^{D}\bar{Q}_{iL}\chi^{*}D_{jR} + \mathrm{H.c.} \\ &= \lambda_{11}^{T}(\bar{u}_{1L}\chi^{-} + \bar{d}_{1L}\chi^{--} + \bar{T}_L\chi^{0})T_R + \lambda_{ij}^{D}(\bar{d}_{iL}\chi^{+} \\ &\quad - \bar{u}_{iL}\chi^{++} + \bar{D}_{iL}\chi^{0*})D_{jR} + \mathrm{H.c.} \end{aligned} \tag{10}$$

When the χ field develops its VEV, these couplings lead to the mass matrix in the basis (T, D_2, D_3),

$$M_J = \frac{v_\chi}{\sqrt{2}}\begin{pmatrix} \lambda_{11}^{T} & 0 & 0 \\ 0 & \lambda_{22}^{D} & \lambda_{23}^{D} \\ 0 & \lambda_{32}^{D} & \lambda_{33}^{D} \end{pmatrix}.$$

So the exotic quarks have masses around a few TeV because their masses are proportional to v_χ. Therefore we see that the masses of the exotic quarks are approximately equal to m_{Z_2}, but they are only involved in the transition phase SU(3) $\to$ SU(2).

The Yukawa couplings give the masses of the usual quarks through the triplet ρ. Therefore—as in the SM—the usual quarks are only involved in the transition phase SU(2) $\to$ U(1).

However, the charged lepton masses arise from the effective dimension-five operator through the couplings of both χ and ρ with the following Lagrangian [17]:

$$L_{\mathrm{Yuk}}^{l} = \frac{\kappa_l}{\Lambda}(\overline{f_L^c}\rho^{*})(\chi^{\dagger}f_L) + \mathrm{H.c.} \tag{11}$$

From the Lagrangian (11), we obtain $m_l = \frac{v_\chi}{\Lambda}\kappa_l v_\rho$, and the coupling constant $v_\chi/\Lambda \approx 1$, so that $m_l \approx \kappa_l v_\rho$. Finally, the masses of the charged leptons depend only on v_ρ. Therefore, they are only involved in the transition phase SU(2) $\to$ U(1). Taking into account $m_e = 0.5$ MeV, $m_\mu = 105$ MeV, and $m_\tau = 1.77$ GeV, ones get $k_e = 2 \times 10^{-5}$, $k_\mu = 4.3 \times 10^{-3}$, and $k_\tau = 7.2 \times 10^{-2}$.

III. EFFECTIVE POTENTIAL IN RM331

From the Higgs potential we obtain V_0 in a form that is dependent on the VEVs as follows:

$$V_0(v_\chi, v_\rho) = \mu_1^2 v_\chi^2 + \mu_2^2 v_\rho^2 + \lambda_1 v_\chi^4 + \lambda_2 v_\rho^4 + \lambda_3 v_\chi^2 v_\rho^2.$$

Here V_0 has a quartic form like in the SM, but it depends on two variables, v_χ and v_ρ, and has a mixing between v_χ and v_ρ. However, by developing the potential (1) we obtain two minimum equations. Therefore, we can transform the mixing between v_χ and v_ρ to a form that depends only on v_χ or v_ρ. Hence, we can write $V_0(v_\chi, v_\rho) = V_0(v_\chi) + V_0(v_\rho)$.

In order to derive an effective potential, we use the Higgs Lagrangian and the principle of least action to arrive at the equation of motion for fields. Expanding the Higgs fields around the VEVs and averaging over space for all fields, we obtain the one-loop effective potential (for details, see Ref. [1]).

The full Higgs Lagrangian in the RM331 model is given by

$$\mathcal{L} = (\mathcal{D}_\mu\chi)^{\dagger}(\mathcal{D}^{\mu}\chi) + (\mathcal{D}_\mu\rho)^{\dagger}(\mathcal{D}^{\mu}\rho) + V(\chi, \rho),$$

where

$$\begin{aligned} V(\chi, \rho) = {}& \mu_1^2\rho^{\dagger}\rho + \mu_2^2\chi^{\dagger}\chi + \lambda_1(\rho^{\dagger}\rho)^2 + \lambda_2(\chi^{\dagger}\chi)^2 \\ & + \lambda_3(\rho^{\dagger}\rho)(\chi^{\dagger}\chi) + \lambda_4(\rho^{\dagger}\chi)(\chi^{\dagger}\rho). \end{aligned}$$

Expanding ρ and χ around v_ρ and v_χ (which are considered as variables) [21], we obtain

$$\begin{aligned} \mathcal{L} = {}& \frac{1}{2}\partial^{\mu}v_\chi\partial_\mu v_\chi + \frac{1}{2}\partial^{\mu}v_\rho\partial_\mu v_\rho + V_0(v_\chi, v_\rho) \\ & + \sum_{\mathrm{boson}} m_{\mathrm{boson}}^2(v_\chi, v_\rho)W^{\mu}W_\mu, \end{aligned}$$

where W runs over all gauge fields and Higgs bosons. Through the boson-mass formulations (as in the above sections) we can split the masses of particles into two parts as follows:

$$m_{\mathrm{boson}}^2(v_\chi, v_\rho) = m_{\mathrm{boson}}^2(v_\chi) + m_{\mathrm{boson}}^2(v_\rho).$$

The effective potential is the function that depends on the VEVs and temperature. The masses of the particles depend on the VEV of the Higgs bosons. Therefore, when we consider the effective potential, we must consider contributions from fermions and bosons. However, for fermions, we have retained here only the top quark and exotic quarks, which dominate over the contributions from the other fermions [1]. And in the RM331 model there are two VEVs, so we have two motion equations according to v_χ and v_ρ,

$$\begin{aligned} & \partial^{\mu}v_\chi\partial_\mu v_\chi + \frac{\partial V_0(v_\chi)}{\partial v_\chi} + \sum \frac{\partial m_{\mathrm{bosons}}^2(v_\chi)}{\partial v_\chi}W^{\mu}W_\mu \\ & \quad + \sum \frac{\partial m_{\text{exotic-quarks}}(v_\chi)}{\partial v_\chi}Q\bar{Q} = 0, \end{aligned} \tag{12}$$

$$\begin{aligned} & \partial^{\mu}v_\rho\partial_\mu v_\rho + \frac{\partial V_0(v_\rho)}{\partial v_\rho} + \sum \frac{\partial m_{\mathrm{bosons}}^2(v_\rho)}{\partial v_\rho}W^{\mu}W_\mu \\ & \quad + \frac{\partial m_{\text{top-quark}}(v_\rho)}{\partial v_\rho}t\bar{t} = 0. \end{aligned} \tag{13}$$

The RM331 modle has the following gauge bosons: two massive bosons like the SM bosons Z_1 and $W^{\pm}$, the new heavy neutral boson Z_2, the singly and doubly charged gauge bosons $U^{\pm\pm}$ and $V^{\pm}$, two doubly charged Higgses h^{++} and h^{--}, one heavy neutral Higgs h_2, and one SM-like Higgs h_1. The masses of the gauge bosons and the Higgses in the RM331 model are presented in Table I

TABLE I. Mass formulations of bosons in the RM331 model.

Bosons	$m^2(v_\chi, v_\rho)$	$m^2(v_\chi)$	$m^2(v_\rho)$
$m^2_{W^\pm}$	$\frac{g^2 v_\rho^2}{4}$	0	80.39^2 (GeV)^2
$m^2_{V^\pm}$	$\frac{g^2 v_\chi^2}{4}$	1307.15^2 (GeV)^2	0
$m^2_{U^{\pm\pm}}$	$\frac{g^2(v_\rho^2+v_\chi^2)}{4}$	1307.15^2 (GeV)^2	80.39^2 (GeV)^2
$m^2_{Z_1} \sim m^2_Z$	$\frac{1}{4}\frac{g^2}{\cos^2\theta_W} v_\rho^2$	0	91.682^2 (GeV)^2
$m^2_{Z_2} \sim m^2_{Z'}$	$\frac{1}{3} g^2[\frac{\cos^2\theta_W}{1-4\sin^2\theta_W} v_\chi^2 + \frac{1-4\sin^2\theta_W}{4\cos^2\theta_W} v_\rho^2]$	4.8^2 (TeV)^2	14.53^2 (GeV)^2
$m^2_{h_1}$	$(\lambda_1 - \frac{\lambda_3^2}{4\lambda_2})v_\rho^2$	0	125^2 (GeV)^2
$m^2_{h^{--}}$	$\frac{\lambda_4}{2}(v_\chi^2 + v_\rho^2)$	$\frac{\lambda_4}{2} v_\chi^2$	$\frac{\lambda_4}{2} v_\rho^2$
$m^2_{h_2}$	$\lambda_2 v_\chi^2 + \frac{\lambda_3^2}{4\lambda_2} v_\rho^2$	$\lambda_2 v_\chi^2$	$\frac{\lambda_3^2}{4\lambda_2} v_\rho^2$

Using Eqs. (12) and (13) and averaging over space by using Bose-Einstein and Fermi-Dirac distributions for bosons and fermions, respectively, we obtain the following effective potentials:

$$\begin{aligned} V_{\rm eff}(v_\chi) = {} & V_0(v_\chi) + \frac{3}{64\pi^2}\left(m^4_{Z_2}(v_\chi)\ln\frac{m^2_{Z_2}(v_\chi)}{Q^2} - 12 m^4_Q(v_\chi)\ln\frac{m^2_Q(v_\chi)}{Q^2}\right) \\ & + \frac{1}{64\pi^2}\left(m^4_{h_2}(v_\chi)\ln\frac{m^2_{h_2}(v_\chi)}{Q^2} + 2m^4_{h^{++}}(v_\chi)\ln\frac{m^2_{h^{++}}(v_\chi)}{Q^2}\right) \\ & + \frac{3}{64\pi^2}\left(2m^4_U(v_\chi)\ln\frac{m^2_U(v_\chi)}{Q^2} + 2m^4_V(v_\chi)\ln\frac{m^2_V(v_\chi)}{Q^2}\right) \\ & + \frac{T^4}{4\pi^2}\left[F_-\left(\frac{m_{h_2}(v_\chi)}{T}\right) + 2F_-\left(\frac{m_{h^{++}}(v_\chi)}{T}\right) + 12F_+\left(\frac{m_Q(v_\chi)}{T}\right)\right] \\ & + \frac{3T^4}{4\pi^2}\left[F_-\left(\frac{m_{Z_2}(v_\chi)}{T}\right) + 2F_-\left(\frac{m_U(v_\chi)}{T}\right) + 2F_-\left(\frac{m_V(v_\chi)}{T}\right)\right], \end{aligned}$$

and

$$\begin{aligned} V_{\rm eff}(v_\rho) = {} & V_0(v_\rho) + \frac{3}{64\pi^2}\left(m^4_{Z_1}(v_\rho)\ln\frac{m^2_{Z_1}(v_\rho)}{Q^2} + m^4_{Z_2}(v_\rho)\ln\frac{m^2_{Z_2}(v_\rho)}{Q^2} + 2m^4_W(v_\rho)\ln\frac{m^2_W(v_\rho)}{Q^2} + 2m^4_U(v_\rho)\ln\frac{m^2_U(v_\rho)}{Q^2}\right. \\ & \left. - 4m^4_t(v_\rho)\ln\frac{m^2_t(v_\rho)}{Q^2}\right) + \frac{1}{64\pi^2}\left(m^4_{h_1}(v_\rho)\ln\frac{m^2_{h_1}(v_\rho)}{Q^2} + m^4_{h_2}(v_\rho)\ln\frac{m^2_{h_2}(v_\rho)}{Q^2} + 2m^4_{h^{++}}\ln\frac{m^2_{h^{++}}(v_\rho)}{Q^2}\right) + \frac{T^4}{4\pi^2} \\ & \times\left[F_-\left(\frac{m_{h_1}(v_\rho)}{T}\right) - F_-\left(\frac{m_{h_2}(v_\rho)}{T}\right) + 2F_-\left(\frac{m_{h^{++}}(v_\rho)}{T}\right)\right] + \frac{3T^4}{4\pi^2}\left[4F_+\left(\frac{m_t(v_\rho)}{T}\right) + F_-\left(\frac{m_{Z_1}(v_\rho)}{T}\right)\right. \\ & \left. + F_-\left(\frac{m_{Z_2}(v_\rho)}{T}\right) + 2F_-\left(\frac{m_W(v_\rho)}{T}\right) + 2F_-\left(\frac{m_U(v_\rho)}{T}\right)\right], \end{aligned}$$

where

$$F_\mp\left(\frac{m_\phi}{T}\right) = \int_0^{\frac{m_\phi}{T}} \alpha J^{(1)}_\mp(\alpha, 0) d\alpha,$$

$$J^{(1)}_\mp(\alpha, 0) = 2\int_\alpha^\infty \frac{(x^2 - \alpha^2)^{1/2}}{e^x \mp 1} dx.$$

The total effective potential in the RM331 model can be rewritten as follows:

$$V^{\rm RM331}_{\rm eff} = V_{\rm eff}(v_\chi) + V_{\rm eff}(v_\rho).$$

IV. ELECTROWEAK PHASE TRANSITION

The symmetry breaking in the RM331 model can take place sequentially due to the fact that the two scales of symmetry breaking are very different, $v_{\chi_0} \gg v_{\rho_0}$ ($v_{\chi_0} \sim$ 4–5 TeV [17,20], $v_{\rho_0} = 246$ GeV), and that the Universe is accelerating. The symmetry breaking SU(3) $\rightarrow$ SU(2) takes place before the symmetry breaking SU(2) $\rightarrow$ U(1).

The symmetry breaking SU(3) $\rightarrow$ SU(2) through χ_0 generates the masses of the heavy gauge bosons—such as $U^{\pm\pm}$, $V^\pm$, Z_2—and the exotic quarks. Therefore, the phase transition SU(3) $\rightarrow$ SU(2) only depends on v_χ.

When our Universe has been expanding and cooling due to the v_{ρ_0} scale, the symmetry breaking or phase transition SU(2) → U(1) is turned on through ρ_0, which generates the masses of the SM particles and the last part of the mass of $U^{\pm\pm}$. Therefore the phase transition SU(2) → U(1) only depends on v_ρ. The current baryon asymmetry of the Universe must have had its origins in the initial conditions of the Universe: if the early Universe does not possess the baryon asymmetry then the current Universe will not possess it as well [1]. In other words, this asymmetry exists throughout all periods of the Universe to date. In order to describe baryogenesis, models must satisfy three conditions given by Sakharov. In particular, the third condition—spontaneous symmetry breaking—must be associated with a first-order phase transition. In the RM331 model, the spontaneous symmetry breaking takes place in two different energy scales, i.e., the electroweak phase transition is a combination of two different phase transitions. Therefore, if the RM331 model is able to describe this phenomenon, both phase transitions must be first-order phase transitions. In contrast, if one of the phase transitions is not a first-order phase transition, the RM331 model will not fully describe this asymmetry, since this model does not ensure the continuity of baryogenesis in the Universe.

Through the boson-mass formulations in the above sections, we saw that the boson $V^\pm$ is only involved in the first phase transition, SU(3) → SU(2). The gauge bosons Z_1, $W^\pm$, and h_1 are only involved in the second phase transition, SU(2) → U(1). However, $U^{\pm\pm}$, Z_2, and h^{--} are involved in both phase transitions.

With this structure for the phase transition, we see that the mass of $U^{\pm\pm}$ is generated by both phase transitions. When the universe is at the v_{χ_0} scale, the symmetry breaking SU(3) → SU(2) generates masses for the exotic quarks and a part of $U^{\pm\pm}$, i.e., it is eaten by one of the Goldstone bosons $\chi^{\pm\pm}$ of the triplet χ. When the universe cools to the v_{ρ_0} scale, the symmetry breaking SU(2) → U(1) is turned on, which generates masses for the SM particles and the last part of $U^{\pm\pm}$, i.e., $U^{\pm\pm}$ is eaten by the other Goldstone boson $\rho^{\pm\pm}$ of the triplet ρ.

A. Phase transition SU(3) → SU(2)

This phase transition involves exotic quarks and heavy bosons, without the involvement of the SM particles: the effective potential of the EWPT SU(3) → SU(2) is

$$
\begin{aligned}
V_{\rm eff}(v_\chi) = {} & V_0(v_\chi) + \frac{3}{64\pi^2}\left(m^4_{Z_2}(v_\chi)\ln\frac{m^2_{Z_2}(v_\chi)}{Q^2} - 12 m^4_Q(v_\chi)\ln\frac{m^2_Q(v_\chi)}{Q^2}\right) \\
& + \frac{1}{64\pi^2}\left(m^4_{h_2}(v_\chi)\ln\frac{m^2_{h_2}(v_\chi)}{Q^2} + 2m^4_{h^{++}}(v_\chi)\ln\frac{m^2_{h^{++}}(v_\chi)}{Q^2}\right) \\
& + \frac{3}{64\pi^2}\left(2m^4_U(v_\chi)\ln\frac{m^2_U(v_\chi)}{Q^2} + 2m^4_V(v_\chi)\ln\frac{m^2_V(v_\chi)}{Q^2}\right) \\
& + \frac{T^4}{4\pi^2}\left[F_-\left(\frac{m_{h_2}(v_\chi)}{T}\right) + 2F_-\left(\frac{m_{h^{++}}(v_\chi)}{T}\right) + 12F_+\left(\frac{m_Q(v_\chi)}{T}\right)\right] \\
& + \frac{3T^4}{4\pi^2}\left[F_-\left(\frac{m_{Z_2}(v_\chi)}{T}\right) + 2F_-\left(\frac{m_U(v_\chi)}{T}\right) + 2F_-\left(\frac{m_V(v_\chi)}{T}\right)\right].
\end{aligned}
$$

The symmetry breaking scale is v_{χ_0}, which is chosen to be 4 TeV [17,20], and the masses of the three exotic quarks are m_Q. Therefore, the effective potential can be rewritten as

$$
V^{\rm eff}_{\rm SU(3)\to SU(2)} = D'(T^2 - T_0'^2)v_\chi^2 - E'Tv_\chi^3 + \frac{\lambda'_T}{4}v_\chi^4.
$$

The minimum conditions are

$$
V_{\rm eff}(v_{\chi_0}) = 0, \qquad \frac{\partial V_{\rm eff}(v_\chi)}{\partial v_\chi}(v_{\chi_0}) = 0, \qquad \frac{\partial^2 V_{\rm eff}(v_\chi)}{\partial v_\chi^2}(v_{\chi_0}) = m^2_{h_2}(v_\chi)|_{v_\chi = v_{\chi_0}},
$$

where

$$
D' = \frac{1}{24 v^2_{\chi_0}}\{6m^2_U(v_\chi) + 3m^2_{Z_2}(v_\chi) + 6m^2_V(v_\chi) + 18m^2_Q(v_\chi) + 2m^2_{h_2}(v_\chi) + 2m^2_{h^\pm}(v_\chi)\},
$$

$$
T_0'^2 = \frac{1}{D}\left\{\frac{1}{4}m^2_{h_2}(v_\chi) - \frac{1}{32\pi^2 v^2_{\chi_0}}(6m^4_U(v_\chi) + 3m^4_{Z_2}(v_\chi) + 6m^4_V(v_\chi) - 36m^4_Q(v_\chi) + m^4_{h_2}(v_\chi) + 2m^4_{h^\pm}(v_\chi))\right\},
$$

$$E' = \frac{1}{12\pi v_{\chi_0}^3}(6m_U^3(v_\chi) + 3m_{Z_2}^3(v_\chi) + 6m_V^3(v_\chi) + m_{h_2}^3(v_\chi) + 2m_{h^\pm}^3(v_\chi)),$$

$$\lambda'_T = \frac{m_{h_2}^2(v_\chi)}{2v_{\chi_0}^2}\Bigg\{1 - \frac{1}{8\pi^2 v_{\chi_0}^2 m_{h_2}^2(v_\chi)}\Bigg[6m_V^4(v_\chi)\ln\frac{m_V^2(v_\chi)}{bT^2} + 3m_{Z_2}^4(v_\chi)\ln\frac{m_{Z_2}^2(v_\chi)}{bT^2} + 6m_U^4(v_\chi)\ln\frac{m_U^2(v_\chi)}{bT^2} - 36m_Q^4(v_\chi)\ln\frac{m_Q^2(v_\chi)}{b_F T^2} + m_{h_2}^4(v_\chi)\ln\frac{m_{h_2}^2(v_\chi)}{bT^2} + 2m_{h^\pm}^4(v_\chi)\ln\frac{m_{h^\pm}^2(v_\chi)}{bT^2}\Bigg]\Bigg\}.$$

The values of $V_{\rm eff}(v_\chi)$ at the two minima become equal at the critical temperature,

$$T'_c = \frac{T'_0}{\sqrt{1 - E'^2/D'\lambda'_{T'_c}}}. \tag{14}$$

The problem here is that there are three variables: the masses of h_2, h^{--}, and Q. However, for simplicity, following the ansatz in Ref. [9], we assume $m_{h_2} = X$, $m_{h^{--}} = m_Q = K$, and $m_{Z2}(v_\chi) = 4.821$ TeV. Note that the contributions from h_2, h^{--}, and Z_2 in this phase transition are X or K, which is different from their contributions in the phase transition SU(2) → U(1). In order to have the first phase transition, the phase transition strength must be larger than unity, i.e., $\frac{v_{\chi c}}{T'_c} \geq 1$.

In Fig. 1 we plot K as a function of $m_{h_2}(v_\chi)$, with $m_{h_2}(v_\chi) > 1$ TeV.

According to Fig. 1, if X is larger than 1 TeV, the heavy particle masses must be in the range of a few TeV in order to have a first-order phase transition. In addition, this phase transition can be a strong first-order transition.

B. Phase transition SU(2) → U(1)

This phase transition dose not involve the exotic quarks or the boson $V^\pm$, and the contribution from $U^{\mp\mp}$ is equal to that from $W^\mp$. The effective potential of the EWPT SU(2) → U(1) is

$$V_{\rm eff}(v_\rho) = V_0(v_\rho) + \frac{3}{64\pi^2}\Bigg(m_{Z_1}^4(v_\rho)\ln\frac{m_{Z_1}^2(v_\rho)}{Q^2} + m_{Z_2}^4(v_\rho)\ln\frac{m_{Z_2}^2(v_\rho)}{Q^2} + 2m_W^4(v_\rho)\ln\frac{m_W^2(v_\rho)}{Q^2} + 2m_U^4(v_\rho)\ln\frac{m_U^2(v_\rho)}{Q^2} - 4m_t^4(v_\rho)\ln\frac{m_t^2(v_\rho)}{Q^2}\Bigg) + \frac{1}{64\pi^2}\Bigg(m_{h_1}^4(v_\rho)\ln\frac{m_{h_1}^2(v_\rho)}{Q^2} + m_{h_2}^4(v_\rho)\ln\frac{m_{h_2}^2(v_\rho)}{Q^2} + 2m_{h^{++}}^4\ln\frac{m_{h^{++}}^2(v_\rho)}{Q^2}\Bigg) + \frac{T^4}{4\pi^2}\Bigg[F_-\Bigg(\frac{m_{h_1}(v_\rho)}{T}\Bigg) + F_-\Bigg(\frac{m_{h_2}(v_\rho)}{T}\Bigg) + 2F_-\Bigg(\frac{m_{h^{++}}(v_\rho)}{T}\Bigg)\Bigg] + \frac{3T^4}{4\pi^2}\Bigg[4F_+\Bigg(\frac{m_t(v_\rho)}{T}\Bigg) + F_-\Bigg(\frac{m_{Z_1}(v_\rho)}{T}\Bigg) + F_-\Bigg(\frac{m_{Z_2}(v_\rho)}{T}\Bigg) + 2F_-\Bigg(\frac{m_W(v_\rho)}{T}\Bigg) + 2F_-\Bigg(\frac{m_U(v_\rho)}{T}\Bigg)\Bigg].$$

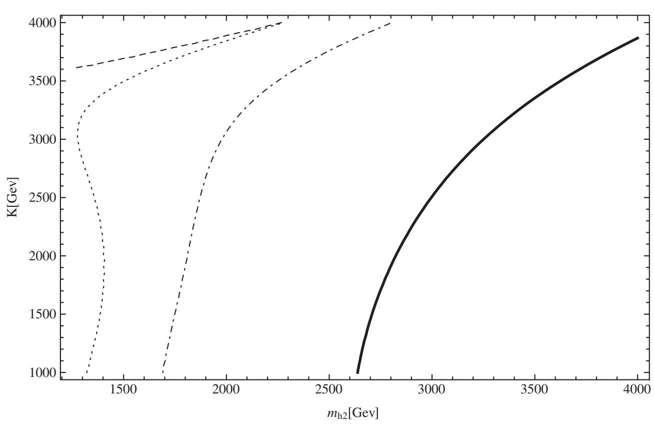

FIG. 1. Here, $X > 1$ TeV. Solid contour: $2E'/\lambda'_{T'_c} = 1$; dash-dotted contour: $2E'/\lambda'_{T'_c} = 2$; dotted contour: $2E'/\lambda'_{T'_c} = 3$; and dashed contour: $2E'/\lambda'_{T'_c} = 5$.

The minimum conditions are

$$V_{\rm eff}(v_{\rho_0}) = 0, \qquad \frac{\partial V_{\rm eff}}{\partial v_\rho}(v_{\rho_0}) = 0,$$

$$\frac{\partial^2 V_{\rm eff}}{\partial v_\rho^2}(v_{\rho_0}) = m_{h_1}^2 + m_{h_2}^2(v_\rho)|_{v_\rho = v_{\rho_0}}.$$

From the above minimum conditions, we see that in this EWPT $m_{h_2}^2(v_\rho)$ generates the masses of the last heavy particles and $m_{h_1}^2$ generates the masses of the SM particles.

With the symmetry-breaking scale equal to $Q \equiv v_{\rho_0} = v_0 = 246$ GeV, the high-temperature expansion of this potential has the form

$$V_{\rm eff}^{\rm RM331} = D(T^2 - T_0^2).v_\rho^2 - ET|v_\rho|^3 + \frac{\lambda_T}{4}v_\rho^4,$$

where

$$D = \frac{1}{24v_0^2}[6m_W^2(v_\rho) + 6m_U^2(v_\rho) + 3m_{Z_1}^2(v_\rho) + 3m_{Z_2}^2(v_\rho) + 6m_t^2(v_\rho) + m_{h_1}^2(v_\rho) + m_{h_2}^2(v_\rho) + 2m_{h^\pm}^2(v_\rho)],$$

$$T_0^2 = \frac{1}{D}\left\{\frac{1}{4}(m_{h_1}^2(v_\rho) + m_{h_2}^2(v_\rho)) - \frac{1}{32\pi^2 v_0^2}(6m_W^4(v_\rho) + 6m_U^4(v_\rho) + 3m_{Z_1}^4(v_\rho) + 3m_{Z_2}^4(v_\rho) - 12m_t^4(v_\rho) + m_{h_1}^4(v_\rho) + m_{h_2}^4(v_\rho) + 2m_{h^\pm}^4(v_\rho))\right\},$$

$$E = \frac{1}{12\pi v_0^3}(6m_W^3(v_\rho) + 6m_U^3(v_\rho) + 3m_{Z_1}^3(v_\rho) + 3m_{Z_2}^3(v_\rho) + m_{h_1}^3(v_\rho) + m_{h_2}^3(v_\rho) + 2m_{h^\pm}^3(v_\rho)),$$

$$\lambda_T = \frac{m_{h_1}^2(v_\rho) + m_{h_2}^2(v_\rho)}{2v_0^2}\left\{1 - \frac{1}{8\pi^2 v_0^2(m_{h_1}^2(v_\rho) + m_{h_2}^2(v_\rho))}\left[6m_W^4(v_\rho)\ln\frac{m_W^2(v_\rho)}{bT^2} + 3m_{Z_1}^4(v_\rho)\ln\frac{m_{Z_1}^2(v_\rho)}{bT^2} + 3m_{Z_2}^4(v_\rho)\ln\frac{m_{Z_2}^2(v_\rho)}{bT^2} + 6m_U^4(v_\rho)\ln\frac{m_U^2(v_\rho)}{bT^2} - 12m_t^4(v_\rho)\ln\frac{m_t^2(v_\rho)}{b_F T^2} + m_{h_1}^4(v_\rho)\ln\frac{m_{h_1}^2(v_\rho)}{bT^2} + m_{h_2}^4(v_\rho)\ln\frac{m_{h_2}^2(v_\rho)}{bT^2} + 2m_{h^\pm}^4(v_\rho)\ln\frac{m_{h^\pm}^2(v_\rho)}{bT^2}\right]\right\}. \tag{15}$$

The effective potential has two minimum points: the first minimum is at $v_\rho = 0$ and the second is at $v_{\rho_c} = \frac{2ET_c}{\lambda_{T_c}}$. In the limit $E \to 0$, we have a second-order phase transition. In order to have a first-order phase transition, the phase-transition strength has to be larger than unity, i.e., $\frac{v_{\rho_c}}{T_c} \geq 1$. The critical temperature T_c is given by

$$T_c = \frac{T_0}{\sqrt{1 - E^2/D\lambda_{T_c}}}. \tag{16}$$

Equation (16) is self-consistent with the critical temperature because λ_{T_c} is a function of T_c. According to the results at the LHC, we take $m_{h_1} = 125$ GeV, and put $m_{h_2}(v_\rho) = Z$, $m_{h^{--}}(v_\rho) = Y$ and $m_{Z2}(v_\rho) = 14.53$ GeV.

In Fig. 2 we show the mass regions of h_1 and h_{++} where the necessary condition for the first-order phase transition is imposed. According to Fig. 2 and the results of numerical evaluation, the strength of the EWPT is in the range $1 \leq 2E/\lambda_{T_c} < 5$. Therefore, in the RM331 model we always have a first-order phase transition, but it is weak (i.e., at the v_ρ scale).

The contributions from new particles (at the first symmetry breaking) generate the first-order phase transition, which is absent in the standard model. However, there is special feature: the heavy particles—such as $U^{\pm\pm}$, h_2, h^{--}, and Z_2—contribute only a small part from their total masses.

As can be seen in Fig. 3, when the temperature approaches T_c the second minimum slowly forms, i.e., the phase-transition nucleation appears. When the temperature goes to T_c the symmetry-breaking phase is turned on, and when temperature goes below T_c the system switches to the symmetry-breaking phase.

To finish this section, we conclude that the effective potential of this model is different from that of the SM, and that the contributions from the heavy bosons act as triggers for the first-order phase transition, with $m_{h_1} = 125$ GeV.

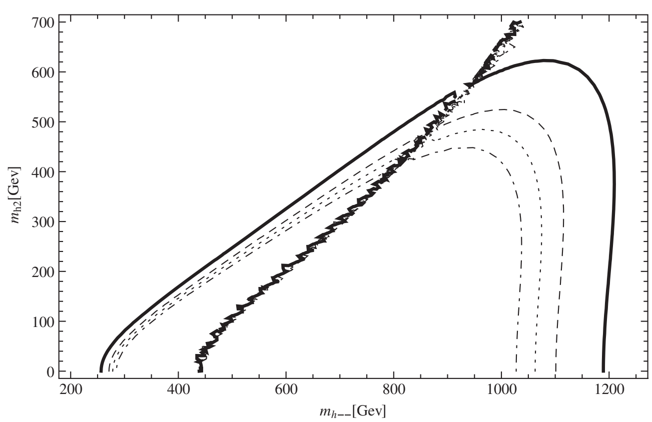

FIG. 2. Solid contour: $2E/\lambda_{T_c} = 1$; dashed contour: $2E/\lambda_{T_c} = 1.1$; dotted contour: $2E/\lambda_{T_c} = 1.15$; and dash-dotted contour: $2E/\lambda_{T_c'} = 1.2$.

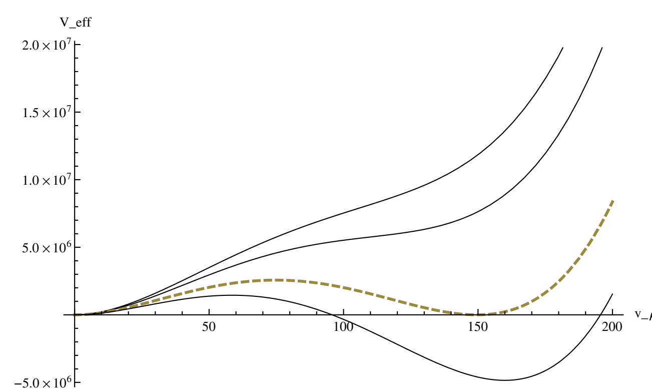

FIG. 3 (color online). The effective potential for $2E/\lambda_{T_c} = 1$, with $Y = 342$ GeV and $Z = 120.109$ GeV. The critical point is at $T_c = 149.549$ GeV. Dashed line: $T = T_c$; lines above the dashed line: $T > T_c$; and lines under the line: $T < T_c$.

C. Constraint on the mass of the charged Higgs boson

In this section, from the phase transition $SU(2) \longrightarrow U(1)$, we have derived

$$200\ \text{GeV} < Y = m_{h_{++}}(v_\rho) < 1200\ \text{GeV}$$

and

$$0 < Z = m_{h_2}(v_\rho) < 624\ \text{GeV}.$$

Therefore, we get

$$(200\ \text{GeV})^2 < \frac{\lambda_4}{2} v_{\rho_0}^2 < (1200\ \text{GeV})^2, \qquad (0\ \text{GeV})^2 < \frac{\lambda_3^2}{4\lambda_2} v_{\rho_0}^2 < (624\ \text{GeV})^2. \tag{17}$$

Taking into account the recent Higgs boson mass (125 GeV) and combining the above with Eq. (5), we also obtain

$$\left(\lambda_1 - \frac{\lambda_3^2}{4\lambda_2}\right) v_{\rho_0}^2 = (125\ \text{GeV})^2. \tag{18}$$

Combining Eqs. (17) and (18) leads to

$$0 < \lambda_1 < 6.692, \quad 0 < \frac{\lambda_3^2}{4\lambda_2} < 6.434, \quad 1.321 < \lambda_4 < 47.59.$$

From Eq. (19), we see that λ_1, λ_2, and λ_4 must be positive in order to satisfy the above boundary conditions and $3.258\ \text{TeV} < m_{h_{++}} < 19.549\ \text{TeV}$. Thus, since the mass of the heavy Higgs must be a few TeV, the electroweak phase transition in this model is a first-order phase transition. We hope that the heavy particles will uncover many more examples of new physics.

V. CONCLUSION AND OUTLOOKS

We have used the effective potential at finite temperature to study the structure of the electroweak phase transition in the RM331 model. This phase transition is separated into two phases. The first transition period is $SU(3) \longrightarrow SU(2)$, i.e., the symmetry breaking at the energy scale v_{χ_0} (in order to generate masses for the heavy particles and the exotic quarks). The second phase transition is $SU(2) \longrightarrow U(1)$ at v_{ρ_0}, which generates masses for all of the usual fermions and the SM gauge bosons. The electroweak phase transition in this model (at the scale v_ρ) may be a weak first-order electroweak phase transition with $m_{h_1} = 125$ GeV if the heavy boson masses are a few TeV. Therefore, this is strong enough to study the baryon asymmetry.

If Z_2 exists, its mass is a few TeV, so its contribution to the electroweak phase transition is very large. Therefore, the electroweak phase transition in this model is completely turned on. In other words, the baryon asymmetry problem in this model is directly related to the mass of Z_2.

The self-interactions of the Higgs in this model are more complicated than in the SM because heavy particles are involved in both phase transitions. Thus calculating the quantum corrections can reveal many new physical phenomena and open up new relations between cosmology and particle physics. In addition, from the phase transitions we can get some bounds on the Higgs self-couplings.

Although we only worked with the RM331 model, this calculation can still apply to other 3-3-1 models, such as the recent supersymmetric reduced minimal 3-3-1 model [22] and the 3-3-1-1 model [23]. Furthermore, the 3-3-1 models can have specific advantages over the SM in explaining the baryon asymmetry problem.

Our next works will calculate the sphaleron rate and *CP* violations in 3-3-1 models in order to analyze electroweak baryogenesis in greater detail. In addition, by using the electroweak phase transition or baryogenesis problem, we can predict the masses of the heavy particles beyond the SM with these 3-3-1 models.

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