Neutrino mass, Dark Matter and Baryon Asymmetry via TeV-Scale Physics without Fine-Tuning

Mayumi Aoki[1], Shinya Kanemura[2], and Osamu Seto[3]

[1]ICRR, University of Tokyo, Kashiwa 277-8582, Japan
[2]Department of Physics, University of Toyama, 3190 Gofuku, Toyama 930-8555, Japan
[3]Instituto de Física Teórica UAM/CSIC, Universidad Autónoma de Madrid, Cantoblanco, Madrid 28049, Spain

We propose an extended version of the standard model, in which neutrino oscillation, dark matter, and baryon asymmetry of the Universe can be simultaneously explained by the TeV-scale physics without assuming unnatural hierarchy among the mass scales. Tiny neutrino masses are generated at the three loop level due to the exact Z_2 symmetry, by which stability of the dark matter candidate is guaranteed. The extra Higgs doublet is required not only for the tiny neutrino masses but also for successful electroweak baryogenesis. The model provides discriminative predictions especially in Higgs phenomenology, so that it is testable at current and future collider experiments.

PACS numbers: 14.60.Pq, 14.60.St, 14.80.Cp, 12.60.Fr [July 2, 2008]

Although the standard model (SM) for elementary particles has been successful for over three decades, today we have definite reasons to consider a model beyond the SM. First of all, observed data for neutrino oscillation indicate that neutrinos have tiny masses and mix with each other[1]. Second, cosmological data have revealed that the density of dark energy and dark matter (DM) in the Universe dominates that of baryonic matter[2]. Essence of DM would be weakly interacting massive particles. Finally, asymmetry of matter and anti-matter in our Universe has been addressed for a long time as a serious problem regarding existence of ourselves[3]. They are all beyond the scope of the SM, so that a new model is required to explain these phenomena.

A simple scenario to generate tiny neutrino masses (m_ν) would be based on the see-saw mechanism with heavy right-handed (RH) neutrinos by $m_\nu \simeq m_D^2/M_R$, where M_R ($\sim 10^{13-16}$ GeV) is the Majorana mass of RH neutrinos and m_D is the Dirac mass of at most the electroweak scale[4]. This scenario would be compatible with the framework with large mass scales like grand unification. Introduction of such large scales, however, causes a problem of hierarchy. In addition, the decoupling theorem[5] makes it far from experimental tests.

Original idea of generating tiny neutrino masses radiatively due to TeV-scale physics has been proposed by Zee[6]. The extension with a TeV-scale RH neutrino has been discussed in Ref. [7], where the neutrino masses are generated at the three-loop due to the exact Z_2 symmetry, and the Z_2-odd RH neutrino is a candidates of DM. This has been extended with two RH neutrinos for the description of the neutrino data[8]. Several models with adding baryogenesis have been considered in Ref. [9].

In this letter, we propose a model which would simultaneously explain neutrino oscillation, origin of DM and baryon asymmetry by the TeV-scale physics. We do not impose unnatural hierarchy among the mass scales. We consider an extended Higgs sector with TeV-scale RH neutrinos. Tiny masses of left-handed (LH) neutrinos are generated at the three-loop level under an exact Z_2 symmetry. The lightest neutral Z_2-odd state is a candidate of DM. Baryon asymmetry can be generated at the electroweak phase transition (EWPT) by the non-decoupling property[10] and additional CP violating phases in the Higgs sector[11]. In this framework, a successful model can be built without contradiction of the current data. All the masses are between $\mathcal{O}(100)$ GeV and $\mathcal{O}(1)$ TeV. The model has discriminative features in Higgs phenomenology, lepton flavor physics and DM physics, so that it is testable.

In addition to the *known* SM fields, particle entries are

$$\Phi_1, \Phi_2, S^\pm, \eta, N_R^\alpha, \qquad (1)$$

where Φ_1, Φ_2 are scalar isospin doublet fields with hypercharge 1/2, $S^\pm$ are charged isospin singlet fields, η is a real scalar singlet field, and N_R^α is the α-th generation isospin-singlet RH neutrinos. It turns out that at least two generations are necessary for N_R^α to reproduce the neutrino data. In the following, we consider the minimum model with α = 1, 2. In order to generate tiny neutrino masses in the three-loop level, and at the same time in order to have a stable DM candidate, we impose a new (exact) Z_2 symmetry as in Ref. [7], which we refer as Z_2. We assign the Z_2 odd charge to N_R^α, $S^\pm$ and η, while ordinary gauge fields, quarks and leptons and Higgs doublets are Z_2 even. Introduction of two Higgs doublets would cause a dangerous flavor changing neutral current. To avoid this in a natural way, we impose another discrete symmetry that is softly broken[12], which we refer as $\tilde{Z}_2$. From a phenomenological reason discussed later, we assign $\tilde{Z}_2$ charges such that only Φ_1 couples to leptons

	Q^i	u_R^i	d_R^i	L^i	e_R^i	Φ_1	Φ_2	$S^\pm$	η	N_R^α
Z_2 (exact)	+	+	+	+	+	+	+	−	−	−
$\tilde{Z}_2$ (softly broken)	+	−	−	+	+	+	−	+	−	+

TABLE I: Particle properties under the discrete symmetries.

whereas Φ_2 does to quarks[13];

$$\mathcal{L}_Y = -y_{e_i}\overline{L}^i\Phi_1 e_R^i - y_{u_i}\overline{Q}^i\tilde{\Phi}_2 u_R^i - y_{d_i}\overline{Q}^i\Phi_2 d_R^i + \text{h.c.}, \quad (2)$$

where Q^i (L^i) is the ordinary i-th generation LH quark (lepton) doublet, and u_R^i and d_R^i (e_R^i) are RH-singlet up- and down-type quarks (charged leptons), respectively. We summarize the particle properties under Z_2 and $\tilde{Z}_2$ in TABLE I. Notice that the Yukawa coupling in Eq. (2) is different from that in the Type I or Type II two-Higgs-doublet model (THDM)[14]. Our $\tilde{Z}_2$ charge assignment for quarks is the same as that in Type I, but that for leptons is the same as of Type II. The charged Higgs boson couplings to leptons are multiplied by $\tan\beta$, while those to quarks are by $\cot\beta$ in an universal way, where $\tan\beta = \langle\Phi_2^0\rangle/\langle\Phi_1^0\rangle$. The scalar potential is then given by

$$\begin{aligned} V =& -\mu_1^2|\Phi_1|^2 - \mu_2^2|\Phi_2|^2 - (\mu_{12}^2\Phi_1^\dagger\Phi_2 + \text{h.c.}) + \lambda_1|\Phi_1|^4 \\ &+\lambda_2|\Phi_2|^4 + \lambda_3|\Phi_1|^2|\Phi_2|^2 + \lambda_4|\Phi_1^\dagger\Phi_2|^2 + \left\{\frac{\lambda_5}{2}(\Phi_1^\dagger\Phi_2)^2 \right. \\ &\left. + \text{h.c.}\right\} + \sum_{a=1}^{2}\left(\rho_a|\Phi_a|^2|S|^2 + \sigma_a|\Phi_a|^2\frac{\eta^2}{2}\right) \\ &+ \sum_{a,b=1}^{2}\left\{\kappa\,\epsilon_{ab}(\Phi_a^c)^\dagger\Phi_b S^-\eta + \text{h.c.}\right\} + V_{Z_2\text{odd}}(S^\pm,\eta), \end{aligned} \quad (3)$$

where ϵ_{ab} is the anti-symmetric tensor with $\epsilon_{12}=1$. The mass term and the interaction for N_R^α are given by

$$\mathcal{L}_Y = \sum_{\alpha=1}^{2}\left\{\frac{1}{2}m_{N_R^\alpha}\overline{N_R^{\alpha c}}N_R^\alpha - h_i^\alpha\overline{(e_R^i)^c}N_R^\alpha S^- + \text{h.c.}\right\}. \quad (4)$$

The parameters μ_{12}^2, λ_5 and κ (as well as h_i^α) are generally complex. The phases of λ_5 and κ can be eliminated by rephasing $S^\pm$ and Φ_1. The remaining phase of μ_{12}^2 is physical and causes CP violation in the Higgs sector, which is necessary for generating baryon asymmetry at the EWPT[11]. Although the CP violating phase is crucial for successful baryogenesis, it does not much affect in the following discussions on neutrino masses, DM and the strong first order EWPT required for electroweak baryogenesis. Thus, in the following, we neglect the phase of μ_{12}^2 (and h_i^α) for simplicity. We later give a comment on the case with the non-zero CP-violating phase.

Because Z_2 is exact, Z_2 even and odd fields cannot mix. Mass matrices for the Z_2 even scalars are diagonalized as in the usual THDM by introducing the mixing angles α and β, where α diagonalizes CP-even neutral states[14]. Consequently, the Z_2 even physical scalar states are two CP-even (h and H), a CP-odd (A) and charged ($H^\pm$) states. We define h and H such that h is the SM-like Higgs boson when $\sin(\beta-\alpha)=1$.

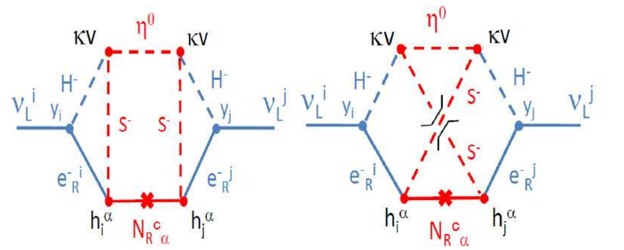

FIG. 1: The diagrams for generating tiny neutrino masses.

The LH neutrino mass matrix M_{ij} is generated by the three-loop diagrams in FIG. 1. The absence of lower order loop contributions is guaranteed by Z_2. $H^\pm$ and e_R^i play a crucial role to connect LH neutrinos with the one-loop sub-diagram by the Z_2-odd states. We obtain

$$M_{ij} = \sum_{\alpha=1}^{2} C_{ij}^\alpha F(m_{H^\pm}, m_{S^\pm}, m_{N_R^\alpha}, m_\eta), \quad (5)$$

where $C_{ij}^\alpha = 4\kappa^2\tan^2\beta(y_{e_i}^{\text{SM}}h_i^\alpha)(y_{e_j}^{\text{SM}}h_j^\alpha)$, and

$$\begin{aligned} F(m_{H^\pm}, m_{S^\pm}, m_{N_R}, m_\eta) =& \left(\frac{1}{16\pi^2}\right)^3\frac{(-m_{N_R}v^2)}{m_{N_R}^2 - m_\eta^2} \\ \times\int_0^\infty dx&\left[x\left\{\frac{B_1(-x,m_{H^\pm},m_{S^\pm}) - B_1(-x,0,m_{S^\pm})}{m_{H^\pm}^2}\right\}^2\right. \\ \times&\left.\left(\frac{m_{N_R}^2}{x+m_{N_R}^2} - \frac{m_\eta^2}{x+m_\eta^2}\right)\right], \quad (m_{S^\pm}^2 \gg m_{e_i}^2), \end{aligned} \quad (6)$$

with m_f representing the mass of the field f, $y_{e_i}^{\text{SM}} = \sqrt{2}m_{e_i}/v$, $v = 246$GeV and B_1 being the tensor coefficient function in the Passarino-Veltman formalism[15]. Magnitudes of $(\kappa\tan\beta)$ as well as $F(m_H, m_S, m_{N_R}, m_\eta)$ determine the universal scale of M_{ij}, whereas variation of h_i^α ($i = e, \mu, \tau$; α =1-2) reproduces the mixing pattern indicated by the neutrino data[1]. M_{ij} is related to these data by $M_{ij} = U_{is}(M_\nu^{\text{diag}})_{st}(U^T)_{tj}$, where $M_\nu^{\text{diag}} = \text{diag}(m_1, m_2, m_3)$ and U_{is} to be the Maki-Nakagawa-Sakata matrix[16]. Under the *natural* requirement that $h_e^\alpha \sim \mathcal{O}(1)$, and taking into account the $\mu \to e\gamma$ search results[17], we find that $m_{N_R^\alpha} \gtrsim \mathcal{O}(1)$ TeV, $m_{H^\pm} \lesssim \mathcal{O}(100)$ GeV and $\kappa\tan\beta \sim \mathcal{O}(10)$. In addition, with the LEP direct search and precision measurement data[1], possible values for the scalar masses uniquely turn out to be $m_{H^\pm} \simeq m_H \simeq 100$ GeV and $m_{S^\pm} \sim \mathcal{O}(100)$ GeV for $\sin(\beta-\alpha)\simeq 1$. In the Type II THDM, such light $H^\pm$ cannot be allowed because of the $b \to s\gamma$ data[18]. Consequently, the Yukawa coupling in Eq. (2) has been employed in our model. For values of h_i^α, we only require that $h_i^\alpha y_i \sim \mathcal{O}(y_e) \sim 10^{-5}$, since we cannot avoid to include the hierarchy among y_i^{SM}. Several sets for h_i^α are shown in TABLE II with the predictions

Set	h_e^1	h_e^2	h_μ^1	h_μ^2	h_τ^1	h_τ^2	$B(\mu \to e\gamma)$
A	1.2	1.2	-0.011	0.025	-0.0015	0.00070	5.3×10^{-12}
B	1.2	1.35	0.0037	0.022	-0.00075	0.0012	4.5×10^{-12}

TABLE II: Values of h_i^α for $m_{H^\pm} = m_{S^\pm} = 100$GeV $m_\eta = 50$ GeV, $m_{N_R^1} = m_{N_R^2}$ =3.5 TeV for the normal hierarchy. For Set A (B), $\kappa \tan\beta = 36(42)$ and $U_{e3} = 0(0.18)$. Predictions on the branching ratio of $\mu \to e\gamma$ are also shown.

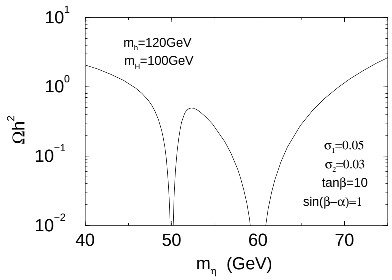

FIG. 2: The relic abundance of η.

on the branching ratio of $\mu \to e\gamma$, assuming the normal hierarchy; *i.e.*, $m_1 \simeq m_2 \ll m_3$ with $m_1 = 0$. We moreover find that larger values are required for $\kappa \tan\beta$ in the case of the inverted hierarchy, where $m_3 \ll m_1 \simeq m_2$ with $m_3 = 0$. Therefore, our model turns out to be better compatible with the normal hierarchy scenario[19].

Next, since Z_2 is exact, the lightest Z_2-odd particle is stable and can be a candidate of DM if it is neutral. In our model, N_R^α must be heavy, so that the DM candidate is identified as η. Since η is a singlet under the SM gauge group, the interactions with Z_2 even particles are only through the Higgs coupling. When η is lighter than the W boson, η predominantly annihilates into $b\bar{b}$ and $\tau^+\tau^-$ through s-channel Higgs (h and H) exchange diagrams. From their thermal averaged annihilation rate $\langle\sigma v\rangle$, the relic mass density $\Omega_\eta h^2$ is evaluated as

$$\Omega_\eta h^2 = 1.1 \times 10^9 \left. \frac{(m_\eta/T_d)}{\sqrt{g_*} M_P \langle\sigma v\rangle}\right|_{T_d} \text{GeV}^{-1}, \qquad (7)$$

where M_P is the Planck scale, g_* is the total number of relativistic degrees of freedom in the thermal bath, and T_d is the decoupling temperature[20]. FIG. 2 shows $\Omega_\eta h^2$ as the function of m_η for $\sin(\beta-\alpha) = 1$. The parameters are chosen as $(\sigma_1, \sigma_2, m_h, m_H) = (0.05, 0.03, 120\text{GeV}, 100\text{GeV})$. Strong annihilation can be seen near 50 GeV $\simeq m_H/2$ (60 GeV $\simeq m_h/2$) due to the resonance of H (h) mediation. The data ($\Omega_{\rm DM}h^2 \sim 0.11$[2]) indicate that m_η is around 48-62 GeV.

The model satisfies the Sakharov's conditions for baryogenesis[3]. In particular, departure from thermal equilibrium can be realized by the strong first order EWPT. The free energy with quantum corrections can be described at high temperatures T as[21]

$$V_{eff}[\varphi, T] = D(T^2 - T_0^2)\varphi^2 - ET\varphi^3 + \frac{\lambda_T}{4}\varphi^4 + ..., \qquad (8)$$

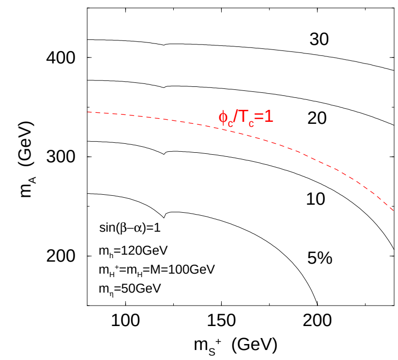

FIG. 3: The region of strong first order EWPT is shown in the m_A-$m_{S^\pm}$ plane. The contour plots represent deviations in the hhh coupling from the SM value.

where φ is the order parameter, and

$$E \simeq \frac{1}{12\pi v^3}(6m_W^3 + 3m_Z^3 + m_A^3 + 2m_{S^\pm}^3), \qquad (9)$$

with $D \simeq (6m_W^2 + 3m_Z^2 + 6m_t^2 + m_A^2 + 2m_{S^\pm}^2)/(24v^2)$, $T_0^2 \sim m_h^2/(4D)$ and $\lambda_T \sim m_h^2/(2v^2)$. A large value of E is crucial for the strong first order EWPT[10]. In Eq. (9), the quantum effect by h, H and $H^\pm$ is neglected since they are unimportant for $\sin(\beta-\alpha) \simeq 1$ and $m_{H^\pm}$ ($\simeq m_H$) being the same value as the soft $\tilde{Z}_2$ breaking scale $M(\equiv \sqrt{\mu_{12}^2/\sin 2\beta})$[26]. For sufficient sphaleron decoupling in the broken phase, it is required that[22]

$$\frac{\varphi_c}{T_c}\left(= \frac{2E}{\lambda_{T_c}}\right) \gtrsim 1, \qquad (10)$$

where φ_c ($\neq 0$) and T_c are the critical values of φ and T at the EWPT. In FIG. 3, the allowed region under the condition of Eq. (10) is shown. The condition is satisfied when $m_A \gtrsim 340$ GeV for $m_{S^\pm} \simeq 100$ GeV, $m_h \simeq 120$ GeV, $m_H \simeq m_{H^\pm}(\simeq M) \simeq 100$ GeV and $\sin(\beta-\alpha) \simeq 1$.

A successful scenario which can simultaneously solve the above three issues under the data[1, 17, 18] would be

$$\begin{array}{ll} \sin(\beta-\alpha) \simeq 1, & (\kappa\tan\beta) \simeq 30-40, \\ m_h = 120\text{GeV}, & m_H \simeq m_{H^\pm}(\simeq M) \simeq 100\text{GeV}, \\ m_A \gtrsim 350\text{GeV}, & m_{S^\pm} = 100-200\text{GeV}, \\ m_\eta \simeq 48-62\text{GeV}, & m_{N_R^1} \simeq m_{N_R^2} \simeq 3-4\text{TeV}. \end{array} \qquad (11)$$

This is realized without assuming unnatural hierarchy among the coupling constants. All the masses are between $\mathcal{O}(100)$ GeV and $\mathcal{O}(1)$ TeV. As they are determined by the data, the model has a predictive power.

We now outline phenomenological predictions in the scenario in (11) in order. The detailed analysis is shown elsewhere[23]. (I) h is the SM-like Higgs boson. Its production is the same as that in the SM, but it decays into a DM pair $\eta\eta$ when $m_\eta < m_h/2$. The branching ratio of $h \to \eta\eta$ amounts to about 50% (37%) for $m_h = 120$ GeV and $m_\eta \simeq 48$ (57) GeV. The decay rate is related to the DM abundance, so that our DM scenario can be testable at the Large Hadron Collider (LHC). (II) η is potentially detectable by direct DM searches at *e.g.* the CDMS and the XMASS[24], because η can scatter with nuclei through the scalar exchange[25]. (III) The quantum effect on the hhh coupling amounts to more than 10-20 % (see FIG. 3), due to the non-decoupling property[26] of A and $S^{\pm}$, which is required for successful baryogenesis[10]. Thus, it should be testable at the International Linear Collider (ILC)[27]. (IV) Because of the Yukawa coupling defined in Eq. (2), the phenomenology for the extra Higgs bosons is different from that in Type I or Type II especially for $\tan\beta \gtrsim 1$. H ($H^{\pm}$) can predominantly decay into $\tau^+\tau^-$ ($\tau^{\pm}\nu$) when $m_H \simeq m_{H^{\pm}} \sim 100$ GeV. As A is relatively heavy in our scenario, it decays into $H^{\pm}W^{\mp}$ and HZ in addition to the $\tau^+\tau^-$ mode. (V) Required large mass difference between A and $H^{\pm}$ and degeneracy between $H^{\pm}$ and H can be tested at the LHC via $pp \to W^* \to AH^{\pm}$ and $pp \to W^* \to HH^{\pm}$[28]. They would be main production channels for $H^{\pm}$, H and A, because their couplings to quarks are largely suppressed in our model. They would also be produced via τ-associated processes. (VI) $S^{\pm}$ would be produced via $q\bar{q}$ (e^+e^-) $\to S^+S^-$ at the LHC (the ILC)[29], and should decay into $\tau^{\pm}\nu\eta$ through the coupling κ. The signal would be a high-energy hadron pair[30] with a large missing transverse momentum. (VII) The couplings h_i^{α} cause lepton flavor violation such as $\mu \to e\gamma$, depending on $m_{N_R^{\alpha}}$. If such a phenomenon is observed at future experiments[31], we could obtain information on $m_{N_R^{\alpha}}$.

Finally, we comment on the case with the CP violating phases. Our model includes the THDM, so that the same generation mechanism can be applied in evaluation of baryon number at the EWPT unless $\tan\beta$ is too large[11]. The mass spectrum in the Higgs sector would be changed by including the CP violating phases, but most of the phenomenological features discussed above should be conserved with a little modification. In addition, a scenario with leptogenesis may also be considered.

We have discussed the model solving neutrino oscillation, DM and baryon asymmetry by the TeV scale physics without fine tuning, which shows specific features in Higgs phenomenology, flavor physics and DM physics, so that it is testable at current and future experiments.

This work was supported, in part, by Japan Society for the Promotion of Science, Nos. 1945210 and 18034004, and the MEC project FPA 2004-02015, the Comunidad de Madrid project HEPHACOS (P-ESP-00346), and the European Network of Theoretical Astroparticle Physics ILIAS/ENTApP (RII3-CT-2004-506222).

[1] W. M. Yao, et al., J. Phys. G **33** (2006) 1.
[2] E. Komatsu, et al., arXiv:0803.0547 [astro-ph].
[3] A. D. Sakharov, Pisma Zh. Eksp. Teor. Fiz. **5**, 32 (1967).
[4] T. Yanagida, In Proceedings of Workshop on *the Unified Theory and the Baryon Number in the Universe*, p.95 KEK Tsukuba, Japan (1979); M. Gell-Mann, P. Ramond and R. Slansky, in Proceedings of Workshop *Supergravity*, p.315, Stony Brook, New York, 1979.
[5] T. Appelquist and J. Carazzone, Phys. Rev. D **11**, 2856 (1975).
[6] A. Zee, Phys. Lett. B **93**, 389 (1980) [Erratum-ibid. B **95**, 461 (1980)]; A. Zee, Phys. Lett. B **161**, 141 (1985).
[7] L. M. Krauss, S. Nasri and M. Trodden, Phys. Rev. D **67**, 085002 (2003).
[8] K. Cheung and O. Seto, Phys. Rev. D **69**, 113009 (2004).
[9] E. Ma, Phys. Rev. D **73**, 077301 (2006); J. Kubo, E. Ma and D. Suematsu, Phys. Lett. B **642**, 18 (2006); T. Hambye, et al., Phys. Rev. D **75**, 095003 (2007). N. Sahu and U. Sarkar, Phys. Rev. D **76**, 045014 (2007); K. S. Babu and E. Ma, arXiv:0708.3790 [hep-ph]; N. Sahu and U. Sarkar, arXiv:0804.2072 [hep-ph].
[10] S. Kanemura, Y. Okada and E. Senaha, Phys. Lett. B **606**, 361 (2005).
[11] J. M. Cline, K. Kainulainen and A. P. Vischer, Phys. Rev. D **54**, 2451 (1996); L. Fromme, S. J. Huber and M. Seniuch, JHEP **0611**, 038 (2006).
[12] S. L. Glashow and S. Weinberg, Phys. Rev. D **15**, 1958 (1977).
[13] Y. Grossman, Nucl. Phys. B **426**, 355 (1994).
[14] J. F. Gunion, et al., "*The Higgs Hunters's Guide*" (Addison Wesley, 1990).
[15] G. Passarino and M. J. G. Veltman, Nucl. Phys. B **160**, 151 (1979).
[16] Z. Maki, M. Nakagawa and S. Sakata, Prog. Theor. Phys. **28**, 870 (1962).
[17] A. Baldini, Nucl. Phys. Proc. Suppl. **168**, 334 (2007).
[18] E. Barberio, et al., arXiv:hep-ex/0603003.
[19] The unitarity bound for the scattering process $h\eta \to H^+S^-$ turns out to give $\kappa \lesssim 25$, whereas $\tan\beta$ should not be too large for successful electroweak baryogenesis[11].
[20] E. W. Kolb and M. S. Turner, *The Early Universe* (Addison-Wesley, 1990).
[21] G. W. Anderson and L. J. Hall, Phys. Rev. D **45**, 2685 (1992); M. Dine, et al., Phys. Rev. D **46**, 550 (1992).
[22] G. D. Moore, Phys. Lett. B **439**, 357 (1998); Phys. Rev. D **59**, 014503 (1999).
[23] M. Aoki, S. Kanemura and O. Seto, in preparation.
[24] Y. D. Kim, Phys. Atom. Nucl. **69**, 1970 (2006); D. S. Akerib, et al., Phys. Rev. Lett. **96**, 011302 (2006).
[25] J. McDonald, Phys. Rev. D **50**, 3637 (1994); for a recent study, see *e.g.*, H. Sung Cheon, S. K. Kang and C. S. Kim, J. Cosmol. Astropart. Phys. 05 (2008) 004.
[26] S. Kanemura, et al., Phys. Lett. B **558**, 157 (2003); Phys. Rev. D **70**, 115002 (2004).
[27] M. Battaglia, E. Boos and W. M. Yao, arXiv:hep-ph/0111276; Y. Yasui, et al., arXiv:hep-ph/0211047.
[28] S. Kanemura and C. P. Yuan, Phys. Lett. B **530**, 188 (2002); Q. H. Cao, S. Kanemura and C. P. Yuan, Phys. Rev. D **69**, 075008 (2004).
[29] S. Kanemura, et al., Phys. Rev. D **64**, 053007 (2001).
[30] B. K. Bullock, K. Hagiwara and A. D. Martin, Phys. Rev. Lett. **67**, 3055 (1991).
[31] S. Ritt, Nucl. Phys. Proc. Suppl. **162**, 279 (2006).