Phenomenology Of A Non-Standard Top Quark Yukawa Coupling

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ABSTRACT

There are theoretical speculations that the top quark may have different properties from that predicted by the standard model. We use an effective Lagrangian technique to model such a non-standard top quark scenario. We parametrize the CP violating interactions of the top quark with the bubble wall in terms of an effective top quark Yukawa coupling, then study its effects on the electroweak baryogenesis. We also discuss the phenomenology of such an effective Yukawa coupling in low and high energy regions.
I. Introduction

The standard model (SM) is perfectly consistent with all evidence gathered to date, however, it raises as many questions as it answers. One important issue is that of mass generation. In the standard model, the fermion and gauge boson masses come from interactions which couple them to the symmetry breaking sector. Since the top quark is heavier than all other observed fermions and gauge bosons and has a mass of the order the electroweak symmetry breaking scale, it couples to the symmetry breaking sector strongly.

The symmetry breaking sector of the standard model consists of a complex fundamental scalar \( \Phi \). However, there are the theoretical arguments of “Triviality”[1] and “Naturalness”[2] against such a scalar sector. So one believes that the Higgs sector of the standard model is an effective theory and expects that the new physics will manifest itself in the effective interactions of the top quark. In this paper, we will study the phenomenology of a non-standard top quark Yukawa coupling in electroweak baryogenesis and in low and high energy processes.

One can propose various models in which the top quark has a non-standard Yukawa coupling. For example, consider the top quark condensate models[3]. In these models, the electroweak symmetry breaking is realized by the top-antitop condensate caused by a strong interaction involving top quarks above the electroweak symmetry breaking scale. However, in the original formulation of this idea, a new strong interaction takes place at a very high energy and the effective theory in the TeV energy region corresponds to the SM. There are also models where the new physics which triggers the top quark condensation appears around a few TeV. Since the new physics scale is of the order of the Fermi scale, there would not be the hierarchy problem. And, more importantly, one would expect many interesting phenomena at the TeV energy scale. One of them can be that the
effective Yukawa coupling, $\Gamma_{t}^{\text{eff}}$, differs from the SM. Models with weak interactions above the electroweak symmetry breaking scale can also produce an effective Yukawa coupling $\Gamma_{t}^{\text{eff}}$ different from the SM. In the Appendix, we will give an example of this latter approach and derive $\Gamma_{t}^{\text{eff}}$ in a class of the Left-Right (L-R) symmetric models.

Given the fact that the underlying theory beyond the SM is presently unknown and so many disparate models exist in the literature, it is worthwhile to extract phenomenologically the essence of the predictions in a “fundamental” symmetry breaking and mass generation theory by means of an effective Lagrangian technique[4]. In this paper we will show that if $\Gamma_{t}^{\text{eff}}$ violates CP symmetry it will help to produce the observed baryon asymmetry at the weak scale.

In the context of electroweak baryogenesis, an effective Lagrangian approach has been taken in Ref.[5]. Following Ref.[5], we parametrize the new physics effects by a set of higher dimension operators $O^{i}$:

$$L^{\text{new}} = \Sigma_{i} \frac{c_{i}}{\Lambda d_{i}} O^{i}, \quad (1)$$

where $d_{i}$ are integers greater than 4. The $O^{i}$, which have dimension $[\text{mass}]^{d_{i}}$, are invariant under the SM gauge symmetry and contain only the SM fields. The parameters $c_{i}$, which determine the strength of the contribution of $O_{i}$, can in principal be calculated by matching the effective theory with the underlying theory. However, they are taken as free parameters here since we do not know the underlying theory. $\Lambda$ is a mass parameter which is the scale at which the effective theory breaks down. Here we shall take $\Lambda$ to be the order of a few TeV. Barring the presence of anomalously large coefficient $c_{i}$, we expect that $L^{\text{new}}$ will be dominated by low dimension operators.

The lowest dimension $O^{i}$, which modifies the Yukawa coupling and satisfies
the above requirements, is given by

\[ \mathcal{O}^t = c_t \ e^{i\xi} \left( \frac{\phi^2 - \frac{v^2}{2}}{\Lambda^2} \right) \Gamma_t \ \overline{\Psi_L} \Phi R, \quad (2.a) \]

where \( \overline{\Psi_L} \equiv (\overline{t}_L, \overline{b}_L) \), \( \Gamma_t \) is the SM Yukawa coupling, \( \xi \) is a CP violation phase and \( c_t \) is a real parameter. Note that in our notation the \( \Phi \) is an \( SU_L(2) \times U_Y(1) \) doublet scalar and \( \phi \equiv (\Phi^\dagger \Phi)^{1/2} \). By definition, the operator in Eq.\( (2.a) \) does not renormalize the top quark mass \( m_t = \frac{\Gamma_t}{\sqrt{2}} v \), but it modifies the top-Higgs interaction (see Eq.\( (7) \) below). Combining \( (2.a) \) with the SM Yukawa coupling, we have an effective Yukawa coupling for the top quark,

\[ \Gamma_{eff}^t = \Gamma_t \left\{ 1 + c_t \ e^{i\xi} \left( \frac{\phi^2 - \frac{v^2}{2}}{\Lambda^2} \right) \right\}. \quad (2.b) \]

In the following sections we will examine in detail the phenomenology of the effective Yukawa coupling \( \Gamma_{eff}^t \) in electroweak baryogenesis and in relevant low and high energy processes. The paper is organized as follows. In section II we estimate the coefficients in \( (2) \) by considering their effect on electroweak baryogenesis. In section III, we calculate the electric dipole moments of fermions induced by \( \Gamma_{eff}^t \). In section IV, we consider the phenomenological implication of the \( \Gamma_{eff}^t \) in a future Linear Collider. Section V includes a summary and discussion.

II. Baryogenesis With An Effective Top Quark Yukawa Coupling

The first order electroweak phase transition proceeds by the formation of bubbles of the true vacuum which grow and fill the universe. The baryon asymmetry can be produced in the bubble wall and in front of it. In the presence of a CP violating effective top quark Yukawa coupling \( \Gamma_{eff}^t \) (Eq.\( (2.b) \), the most efficient mechanism for baryogenesis is the charge transport mechanism considered by Cohen, Kaplan and Nelson[6]. For this mechanism to work, the thickness of
the bubble wall is required to be smaller than the typical particle mean free path \( \sim 4/T \) for quarks[7], where \( T \) is the phase transition temperature. In general, this requirement can be satisfied in electroweak models with \( SU_L(2) \) singlet scalar fields[7]. One would also expect that the thin bubble wall scenario will be applicable in the effective theory. As argued in Ref.[8], an important effect of the new physics on the effective potential is to shift the quartic coupling of the Higgs field from \( \lambda_T \sim m_H^2/2v^2 \) to \( \lambda_T \sim (m_H^2/2v^2) - (4v^2/\Lambda^2) \). Thus if \( \Lambda \) is not too far from the Fermi scale, one can make \( \lambda_T \) small enough to have a thin wall without conflicting with the experimental lower bound on the Higgs boson mass \( m_H > 60 \) GeV[9].

The asymmetry arising from the different number of the left-handed and right-handed top quarks which scatter off the bubble wall serves as a force to bias the anomalous baryon number violating reaction. However, to compute the net baryon density produced, one needs to know the distribution of the asymmetry in front of the bubble wall by solving numerically appropriate Boltzmann equations[10]. If the system is partially in thermal equilibrium, however, the procedure could be simplified.

Approximately, the final ratio of baryon number density to entropy can be expressed as

\[
\frac{n_B}{s} \sim \kappa \alpha_W^4 \delta_{CP} F ,
\]

where \( F \) is a factor which depends on the properties associated with the phase transition and the \( \kappa \) subsumes our ignorance about the sphaleron rate in the unbroken phase. Quantitatively, \( \kappa > 0.5 \) as derived in lattice simulation[11] and \( \kappa \sim 20 \) as estimated in a different method given in Ref.[12]. From Ref.[6], \( F \) could be as large as of the order of 0.1, which gives \( \frac{n_B}{s} \sim 10^{-7} \) for \( \kappa \sim 1 \) and a maximal
CP violation in the two-Higgs model. Here, $\delta_{CP} \sim c_t \sin \xi v^2/2\Lambda^2$, we would expect that for $\Lambda = 1$ TeV,

$$\frac{n_B}{s} \sim \kappa c_t \sin \xi \ 10^{-9} \ .$$  \hspace{1cm} (4)

In order to explain the observed asymmetry $\frac{n_B}{s} \sim (0.4 - 1.4) \times 10^{-10}$, it is required that,

$$\kappa c_t \sin \xi \geq 4 \times 10^{-2} \ .$$  \hspace{1cm} (5)

We would like to point out that since the electroweak baryogenesis calculations available so far are qualitative, the result in (5) is only accurate to within a couple of orders of magnitude. For instance, if the effect of the QCD sphaleron[13] is taken into consideration, the asymmetry in (4) would be suppressed by a factor $10^{-2}$[14][15]. As a result, the CP violation required should be larger than that shown in (5), namely,

$$\kappa c_t \sin \xi \geq 4 \ .$$  \hspace{1cm} (6)

III. Electric Dipole Moments Of The Electron And Neutron

In the unitary gauge, $\Phi = \frac{(v+H)}{\sqrt{2}} \left( \begin{array}{c} 0 \\ 1 \end{array} \right)$ with $H$ being the physical Higgs particle, the effective Yukawa coupling $\Gamma^e_{eff}$ generates a non-standard CP violating top-Higgs interaction,

$$\mathcal{L}^e_{eff} \sim \frac{m_t}{v} \left[ (1 + \left( \frac{c_t}{16} \right) \cos \xi \right] + i \left( \frac{c_t}{16} \right) \sin \xi \gamma_5 \right] t \ H \ ,$$  \hspace{1cm} (7)

where we have fixed $\Lambda = 1$ TeV.

In the last section, we have put a lower bound on $\kappa c_t \sin \xi$ from baryogenesis. In this section, we will calculate the contribution of $\mathcal{L}^e_{eff}$ to the electric
dipole moment of the fermion. A recent discussion of this kind of calculation in a
gauge theory is given in Ref.[16]. For the electric dipole moment of the electron,
\(d_e\), the dominant contribution is from the two-loop diagram[17] in Fig.1. In the
calculation, we first evaluate the top quark triangle to get an effective photon-
photon-Higgs operator, then we evaluate the one-loop diagram in Fig.2, and cut
off the logarithmic divergence by the top quark mass. We have

\[
\frac{d_e}{e} \sim \frac{m_e}{v} \left(\frac{1}{27}\right) c_t \sin \xi \frac{\alpha_{em}}{\pi v} \frac{1}{16\pi^2} \ln \frac{m_t^2}{m_H^2} .
\] (8)

The value of the Higgs boson mass is constrained in electroweak baryogenesis. Its
upper limit in the standard model is about 40 GeV[18]. In an effective theory, it can
be relaxed to the experimentally allowed region[8]. In the numerical calculation of
d_e, we take \(m_H \sim 80\) GeV and \(m_t \sim 174\) GeV[19]. Given the constraints on the
quantity \(\kappa c_t \sin \xi\) as shown in (5) and (6), we have the prediction on the \(d_e\) from
the electroweak baryogenesis,

\[
\frac{d_e}{e} \sim \frac{6}{\kappa} \times (10^{-28} \sim 10^{-30}) \text{ cm} .
\] (9)

Clearly, for \(\kappa\) in the range mentioned above, the electric dipole moment of the
electron is just below the current experimental upper bound. This opens the
possibility of detecting \(d_e\) in the near future.

The estimate of the electric dipole moment of the neutron \(d_n\) contains large
uncertainties. So we will not go into this calculation in detail. As an estimate, we
use the nonrelativistic quark model to calculate the electric dipole moments of the
up and down quarks, \(d_u\) and \(d_d\), using a diagram similar to that in Fig.1. We have
\(d_n = \frac{4}{3} d_d - \frac{1}{3} d_u \sim \frac{m_u}{m_e} d_e \sim \frac{1}{\kappa} \times (2 \times 10^{-26} \sim 8 \times 10^{-29}) \text{ e.cm}\). Thus the electric
dipole moment of the neutron is also within the experimental limit.

Before concluding this section, we would like to examine the contribution of
$L_{eff}$ to the QCD $\theta$, which is given by

$$L^\theta = \overline{\theta} \frac{g_s^2}{32\pi^2} G^{\mu\nu} \tilde{G}^{\mu\nu},$$  \hspace{1cm} (10)

where $\overline{\theta} = \theta + \text{arg det } M$, with $M$ being the mass matrix of the quark fields. Strong interactions, through the $\overline{\theta}$ dependence, lead to an estimate for $d_n$. The experimental bound on $d_n$ implies that

$$\overline{\theta} < 10^{-9}. \hspace{1cm} (11)$$

The relevant diagram we should evaluate is the top quark self energy correction shown in Fig.3. In our calculation, the dimensional regularization is used and the infinity associated with the phase of the top quark mass ( $\text{arg } m_t$ ) induced by loop effects can be absorbed in the counterterms appearing in the effective Lagrangian. The coefficient of this counterterm is a new independent coupling in the effective theory and can not be computed. So in what follows, we shall focus on the ”log-enhanced” term and compute the leading correction by setting the counterterm to zero for the renormalization scale to be $\Lambda (= 1\text{TeV})$, the cut-off of the theory. Thus we have

$$\text{arg } m_t \sim c_t \sin \xi \frac{v^2}{\Lambda^2} \frac{2}{16\pi^2} \frac{m_t^2}{v^2} \ln \frac{\Lambda^2}{m_t^2} \sim \frac{3}{\kappa} \times \left(10^{-4} \sim 10^{-6}\right).$$ \hspace{1cm} (12)

Clearly, if no cancellation occurs between the original $\theta$ and the induced $\theta = \text{arg } m_t$, an axion[20] must be included in the effective lagrangian.

IV. Phenomenology Of The Top Quark Yukawa Coupling At Linear Collider

There is a direct way to determine the Yukawa coupling from top and Higgs production at a high energy collider[21]. In this section, we consider the possibility
of testing the top quark Yukawa coupling at future $e^+e^-$ colliders. For a light Higgs boson, $m_H < 120$ GeV, as required by electroweak baryogenesis, the dominant process for top pair plus Higgs boson production is the bremsstrahlung process in Fig.4. For the standard model coupling of the Higgs boson to top quark, the cross section has been calculated in Ref.[22]. They concluded that for an integrated luminosity of $\int \mathcal{L} = 20 \text{ fb}^{-1}$, some 100 events can be expected at Higgs masses of order 60 GeV, falling to 20 events at 120 GeV. Even though the production rate is low, as argued in Ref.[22], the signature of the process $e^+e^- \rightarrow t\bar{t}H \rightarrow W^+W^-b\bar{b}b\bar{b}$ is spectacular, so that it is not impossible to isolate the events experimentally.

Given the effective lagrangian (7) for the Top-Higgs coupling, the Dalitz plot density for $e^+e^- \rightarrow t\bar{t}H$ can be written as

$$
d\sigma(e^+e^- \rightarrow t\bar{t}H) = \left\{ \frac{d\sigma}{dx_1 dx_2} \right\}_{SM} (1 + 2\delta \cos \xi + \delta^2) - \left\{ \frac{d\sigma}{dx_1 dx_2} \right\}_{BSM} (\delta \sin \xi)^2, \tag{13}
$$

where $\delta = c_t/16$. The $\left\{ \frac{d\sigma}{dx_1 dx_2} \right\}_{SM}$ is given in Eq.(7) of Ref.[22]. And our calculation gives

$$
\left\{ \frac{d\sigma}{dx_1 dx_2} \right\}_{BSM} = 4N_c \sigma_0 g_{tH}^2 \left\{ G_+ \frac{f}{x_{12}} [h - x + \frac{x^2}{x_{12}} (1 - f)] + G_- \frac{3f}{x_{12}} [x - h + \frac{x^2}{x_{12}} f] \right\}, \tag{14}
$$

where $g_{tH} = \frac{m_t}{v}$, $\sigma_0 = \frac{4\pi\alpha^2}{3s}$, $h = \frac{m_H^2}{s}$, $f = \frac{m_t^2}{s}$, $x_{12} = (1 - x_1)(1 - x_2)$, $x_1 = 2E_t/\sqrt{s}$, $x_2 = 2E_H/\sqrt{s}$, $x = 2E_H/\sqrt{s}$, $N_c = 3$ is the number of colors, and

$$G_\pm = Q_e^2 Q_t^2 + \frac{2Q_e Q_t \hat{v}_t \hat{a}_t}{1 - M_Z^2/s} + \frac{(\hat{v}_t^2 + \hat{a}_t^2)(\hat{v}_t^2 \pm \hat{a}_t^2)}{(1 - M_Z^2/s)^2}, \tag{15.a}
$$

with

$$
\hat{v}_t = \frac{1 - 4Q_t x_W}{4\sqrt{x_W(1 - x_W)}}, \quad \hat{a}_t = \frac{1}{4\sqrt{x_W(1 - x_W)}}, \tag{15.b}
$$
and \( x_W = \sin^2 \theta_W \simeq 0.23 \).

To illustrate the effects of the non-standard top quark Yukawa coupling in Eq.(2.b) on the production of the top and Higgs, in Fig.5 we plot in the \( \delta - \xi \) parameter space contours of \( \sigma/\sigma_{SM} \), the ratio of the integrated cross section to the SM cross section, for \( m_t = 174 \text{ GeV} \) and \( m_H = 80 \text{ GeV} \). The \( \sigma/\sigma_{SM} = 0.60 \) and 1.40 contours represent the parameters for which the integrated cross section is 3 standard deviations (considering statistical uncertainty only) from the SM value, assuming an integrated luminosity of 20 \( fb^{-1} \). Regions above these contours should therefore be readily distinguishable from the SM. The dotted line gives the lower bound on \( \delta \) from electroweak baryogenesis from Eq.(6) with \( \kappa = 0.5 \), which is the most restrictive limit using the indicated ranges of parameters in Sec. II. At the other extreme, using Eq.(5) with \( \kappa = 20 \) gives no substantial bounds on \( \delta \) and \( \xi \).

V. Conclusion And Discussions

In this paper we have examined the impact of a non-standard top quark on electroweak baryogenesis. We parametrize the CP violating interaction of the top quark with the bubble wall in terms of an effective Yukawa coupling[23] (Eq.(2.b)). We found that to explain the observed baryon number asymmetry, the electric dipole moments of the electron and neutron must be very close to the present experimental limits. We have explored the possibility of testing the effective Yukawa coupling in the future linear collider and concluded that its effect on the cross section of the top and Higgs production is sizable (see Fig.5). However, we would like to point out that the electroweak baryogenesis calculations available so far are qualitative and the quantitative results obtained are probably only accurate to within a couple of orders of magnitude. For instance, the uncertainty of \( \kappa \)
in the sphaleron rate generates an uncertainty in the predictions on $d_e$ and $d_n$. Thus future experimental inputs will test the present knowledge about electroweak baryogenesis.

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Appendix

In this Appendix we illustrate how to generate the operator $O_t$ in a weakly interacting Higgs model. Let us examine the effective Lagrangian of the left-right symmetric (L-R) models at an energy below the $SU_R(2)$ breaking scale. Specifically we consider a class of L-R models described in Ref.[24], where the Higgs sector consists of a complex bidoublet $\phi(2, 2, 0)$ and one set of L-R symmetric lepton-number-carrying triplets $\Delta_L(3, 1, 2) \oplus \Delta_R(1, 3, 2)$. In the fermion sector, there are L-R symmetric doublets of three generations of the quarks and leptons.

After the $SU_R(2)$ is broken at the scale $V_R$, the right-handed neutrinos receive Majorana masses $\sim V_R$ and decouple at low energy. However, they do have effects on low energy physics, such as making the left-handed neutrinos light via a see-saw mechanism. In terms of an effective Lagrangian below $V_R$, the neutrino majorana mass terms can be parametrized by an $SU_L(2) \times U_Y(1)$ gauge invariant higher dimension operators. Now we will show that an higher dimension operator $O^t$ can be generated by heavy Higgs scalars in this model.

The bidoublet $\phi$ will be split into two $SU_L(2) \times U_Y(1)$ doublets below the scale $V_R$, which we denote by $\phi_1$ and $\phi_2$. The top quark Yukawa sector can be rewritten in terms of Higgs fields $\phi_1$ and $\phi_2$ by

$$L^{top} = h_1 \bar{\Psi}_L \phi_1 t_R + h_2 \bar{\Psi}_L \tilde{\phi}_2 t_R ,$$

where $h_1$ and $h_2$ are Yukawa couplings in the L-R symmetric Lagrangian. The vacuum expectation values of $\phi_1$ and $\phi_2$ are related to that of the bidoublet $\phi$.

For a general CP violating Higgs potential, one has,

$$< \phi > = e^{i\alpha} \begin{pmatrix} \kappa & 0 \\ 0 & \kappa' \end{pmatrix} ,$$

where $\alpha$ is a CP phase. So $< \phi_1 > = e^{i\alpha} \kappa$, $< \phi_2 > = e^{i\alpha} \kappa'$. Thus the mass of the
top quark is given by
\[ m_t = h_1 \kappa e^{i\alpha} + h_2 \kappa' e^{-i\alpha} \quad . \]

The detailed Higgs potential involving \( \phi_1 \) and \( \phi_2 \) can be obtained from that of \( \phi, \Delta_R \) and \( \Delta_L \). Here we give a few terms which are relevant to our discussion,
\[ V(\phi_1, \phi_2) = |\lambda_1| e^{-i\delta} V_R^2 \phi_1 \phi_2 + \lambda_2 (\phi_1 \phi_2)^2 + ... \quad , \]

where \( \lambda_1 \) and \( \lambda_2 \) are two coupling constants, and \( \delta \) is a CP phase.

After \( SU_L(2) \times U_Y(1) \) is broken, there are two “extra” neutral scalars as well as a charged scalar in addition to the neutral scalar corresponding to the Higgs boson of the SM. The experimental constraints[25] on the flavor-changing neutral currents require these “extra” scalars be heavier than a few TeV. However, this in general does not need any fine tuning, for \( V_R > 1 \) TeV, in accordance with the extended survival hypothesis[26]: Higgs bosons acquire the maximum mass compatible with the pattern of the symmetry breaking. So the effective theory below \( V_R \) contains only a doublet scalar. At this point, we would like to point out that Haber and Nir[27] have made a detailed study on the low energy structure of the multi-Higgs models with a cut-off \( \Lambda \). They concluded that if performing minimally the required fine-tuning in order to set the electroweak scale \( v \), the low energy scalar spectrum is identical to that of the SM, up to corrections of order \( \frac{v^2}{\Lambda^2} \). Here, \( \Lambda \sim V_R \).

To examine the effects of heavy scalars on the low energy physics in term of \( SU_L(2) \times U_Y(1) \) invariant effective operators, we follow the procedure in Ref.[27], and derive the effective Lagrangian by using a “rotated” basis \( \{ \Phi, \Phi' \} \), so that \( \langle \Phi \rangle = v = \sqrt{\kappa^2 + \kappa'^2} \) and \( \langle \Phi' \rangle = 0 \). From eq.(17), one has
\[ \begin{pmatrix} \Phi \\ \Phi' \end{pmatrix} = e^{-i\alpha} \begin{pmatrix} \cos \zeta & \sin \zeta \\ -\sin \zeta & \cos \zeta \end{pmatrix} \begin{pmatrix} \phi_1 \\ \phi_2 \end{pmatrix} \quad , \]
\[ (18) \]
where \( \tan \zeta = \kappa' / \kappa \). In this basis, \( \Phi \) serves as the SM doublet and the \( \Phi' \) is just an massive scalar with mass \( \sim V_R \). Integrating out the heavy field \( \Phi' \) will generate many higher dimension operators[28]. The contribution of a Feynman diagram in Fig.6 will give an operator similar to \( O' \). In this model, \( m_{\Phi'}^2 \sim V_R^2 \sim \Lambda^2 \), the parameter \( c_t \) and CP phase \( \xi \) in (2), as functions of the couplings \( h_1, h_2, \lambda_2 \), the mixing angle \( \zeta \) and the CP phases \( \alpha, \delta \), are calculable. However, a detailed examination of the whole structure for the effective Lagrangian and the calculation of the various coefficients in the effective Lagrangian of the L-R models are beyond the scope of this paper.
Figure Captions

[Fig.1] Dominant contribution to $d_e$, the electric dipole moment of the electron.

[Fig.2] Same as Fig.1 with the fermion loop replaced by an effective photon-photon-Higgs coupling.

[Fig.3] Higgs boson contribution to the top quark self energy.

[Fig.4] Dominant contribution to $e^+e^- \rightarrow Ht\bar{t}$. Another graph where H couples to $\bar{t}$ is also present.

[Fig.5] Contours of $\sigma/\sigma_{SM}$ in the $\delta - \xi$ plane for $m_t = 174$ GeV and $m_H = 80$ GeV. The region with $1 \leq \xi/\pi \leq 2$ is a mirror image of the region shown. The dotted line indicates the lower bounds on $\delta$ from electroweak baryogenesis from Eq.(5) with $\kappa = 0.5$.

[Fig.6] A possible diagram for the generation of $\mathcal{O}'$ in eq.(2) in left-right symmetric models.
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