Physics Overview

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Abstract

Recent developments of physics at the TeV energy scale, especially physics related to the $e^+e^-$ linear colliders are briefly reviewed. The topics include the present status of the standard model, Higgs physics, supersymmetry, strongly interacting electroweak symmetry breaking mechanism, and top quark physics.

I. Introduction

One of the most remarkable results in recent particle physics is that the Standard Model (SM) has successfully passed the recent precision experimental tests at the CERN LEP, the SLAC SLC and the Fermilab Tevatron with the LEP and SLC tests of the precision of one-loop electroweak radiative corrections [1–4]. Theoretically, this means that the present available tests support the renormalizable $SU(2) \times U(1)$ gauge theory of the SM as the theory of electroweak interactions. An important ingredient of such a theory is the spontaneous breakdown of the $SU(2) \times U(1)$ gauge symmetry.

Despite the present success of the SM, the mechanism of the electroweak symmetry breaking (EWSB) is not clearly known yet. In the $SU(2) \times U(1)$ electroweak theory, all particle masses come from the EWSB sector. Thus probing the EWSB mechanism concerns

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the understanding of the origin of all particle masses, which is a very deep and fundamental problem in physics. In the SM, it is assumed that an elementary Higgs field is responsible for the EWSB. After careful experimental searches at LEP, the SM Higgs boson has not been found, and the present data are not so sensitive to the Higgs mass $m_H$ due to Veltman’s screening theorem \[5\]. From the theoretical point of view, there are several unsatisfactory features in the Higgs sector of the SM, e.g. there are so many free parameters related to the Higgs sector, and there are the well-known problems of triviality and unnaturalness \[6,7\]. Usually, people take the point of view that the present theory of the SM is only valid up to a certain energy scale $\Lambda$, and new physics will become important above $\Lambda$. Possible new physics are supersymmetry (SUSY) and dynamical EWSB mechanism concerning new strong interactions. So that probing the mechanism of EWSB also concerns the discovery of new physics beyond the SM. Such an important problem can be experimentally studied at LEP and the future TeV energy colliders such as the upgraded Fermilab Tevatron, the CERN LHC, and the future $e^+e^-$ linear colliders.

In this paper, we give a short theoretical review on the recent developments of physics at the TeV energy scale, especially on physics related to the $e^+e^-$ linear colliders. We shall first briefly review the present status of the SM in Sec.II as the basis of the other parts of this review. In Sec. III, we shall deal with the physics of the Higgs boson. Sec. IV is a brief review on SUSY, and Sec. V concerns some recent developments of the study of strongly interacting EWSB mechanisms. In Sec. VI, we shall discuss some topics related to the top quark. The conclusions will be given in Sec. VII.

II. Status of the Standard Model

Recent data from LEP, SLC and the Tevatron supports the electroweak SM as the theory of electroweak interactions to the precision of one-loop radiative corrections. There are recent reports on this subject \[1\]. Here we quote a table of the global fit results of the Z pole precision observables given in Ref. \[1\] to show the situation (Table I).
Table I. $Z$ pole precision observables from LEP and the SLC. Shown are the experimental results, the SM predictions, and the pulls. The SM errors are from the uncertainties in $M_Z$, $\ln M_H$, $m_t$, $\alpha(M_Z)$, and $\alpha_s$. They have been treated as Gaussian and their correlations have been taken into account. Quoted from Ref. [1].

| Quantity | Group(s) | Value      | Standard Model | pull |
|----------|----------|------------|----------------|------|
| $M_Z$ [GeV] | LEP | 91.1867 ± 0.0021 | 91.1865 ± 0.0021 | 0.1 |
| $\Gamma_Z$ [GeV] | LEP | 2.4939 ± 0.0024 | 2.4957 ± 0.0017 | −0.8 |
| $\Gamma(\text{had})$ [GeV] | LEP | 1.7423 ± 0.0023 | 1.7424 ± 0.0016 | — |
| $\Gamma(\text{inv})$ [MeV] | LEP | 500.1 ± 1.9 | 501.6 ± 0.2 | — |
| $\Gamma(\ell^+\ell^-)$ [MeV] | LEP | 83.90 ± 0.10 | 83.98 ± 0.03 | — |
| $\sigma_{\text{had}}$ [nb] | LEP | 41.491 ± 0.058 | 41.473 ± 0.015 | 0.3 |
| $R_e$ | LEP | 20.783 ± 0.052 | 20.748 ± 0.019 | 0.7 |
| $R_\mu$ | LEP | 20.789 ± 0.034 | 20.749 ± 0.019 | 1.2 |
| $R_\tau$ | LEP | 20.764 ± 0.045 | 20.794 ± 0.019 | −0.7 |
| $A_{FB}(e)$ | LEP | 0.0153 ± 0.0025 | 0.0161 ± 0.0003 | −0.3 |
| $A_{FB}(\mu)$ | LEP | 0.0164 ± 0.0013 | 0.2 |
| $A_{FB}(\tau)$ | LEP | 0.0183 ± 0.0017 | 1.3 |
| $R_b$ | LEP + SLD | 0.21656 ± 0.00074 | 0.2158 ± 0.0002 | 1.0 |
| $R_c$ | LEP + SLD | 0.1735 ± 0.0044 | 0.1723 ± 0.0001 | 0.3 |
| $R_{s,d}/R_{(d+s+u)}$ | OPAL | 0.371 ± 0.023 | 0.3592 ± 0.0001 | 0.5 |
| $A_{FB}(b)$ | LEP | 0.0990 ± 0.0021 | 0.1028 ± 0.0010 | −1.8 |
| $A_{FB}(c)$ | LEP | 0.0709 ± 0.0044 | 0.0734 ± 0.0008 | −0.6 |
| $A_{FB}(s)$ | DELPHI + OPAL | 0.101 ± 0.015 | 0.1029 ± 0.0010 | −0.1 |
| $A_b$ | SLD | 0.867 ± 0.035 | 0.9347 ± 0.0001 | −1.9 |
| $A_c$ | SLD | 0.647 ± 0.040 | 0.6676 ± 0.0006 | −0.5 |
| $A_s$ | SLD | 0.82 ± 0.12 | 0.9356 ± 0.0001 | −1.0 |
| $A_{LR}$ (hadrons) | SLD | 0.1510 ± 0.0025 | 0.1466 ± 0.0015 | 1.8 |
| $A_{LR}$ (leptons) | SLD | 0.1504 ± 0.0072 | 0.5 |
| $A_\mu$ | SLD | 0.120 ± 0.019 | −1.4 |
| $A_\tau$ | SLD | 0.142 ± 0.019 | −0.2 |
| $A_c(Q_{LR})$ | SLD | 0.162 ± 0.043 | 0.4 |
| $A_\tau(P_{LR})$ | LEP | 0.1431 ± 0.0045 | −0.8 |
| $A_\tau(P_{\tau})$ | LEP | 0.1479 ± 0.0051 | 0.3 |
| $\delta^2_{Q_{FB}}$ | LEP | 0.2321 ± 0.0010 | 0.2316 ± 0.0002 | 0.5 |

We see that the agreement of the SM with the recent data is quite good. The deviations are all within 2$\sigma$. In the SM fit analysis, the top quark mass $m_t$ is taken as one of the fitting parameters. The best fit requires $m_t = 171.4 \pm 4.8$ GeV [1] which is very close to the directly measured value at the Tevatron, $m_t = 173.8 \pm 3.2(stat.) \pm 3.9(syst.)$ GeV [18].
This is a remarkable success of the SM. Such a good agreement supports the $SU(2) \times U(1)$ gauge interactions in the SM.

Despite of the above success of the SM, its EWSB sector is still not clear. The assumed Higgs boson in the EWSB sector of the SM has not been found yet. The experimental lower bound of the Higgs mass from the recent LEP experiments is $m_H > 94.1 \text{ GeV}$ \cite{9}. Although the global fit analysis favors the SM with a light Higgs boson \cite{1,2}, there exist examples of models with new physics without a light Higgs boson, which can also be consistent with the present precision data \cite{10,11}. So that whether there is a light Higgs boson or not should be tested by future experiments, especially at the LHC and the LC.

Furthermore, as has been pointed out in Refs. \cite{12,13} that the global fit includes a large number of ‘raw’ observables with large errors, which may dilute some deviations of more precise observables. In Refs. \cite{12,13}, a sharper test of the SM is presented from a model-independent analyses of the forward-backward asymmetry data in $Z$ decays. Instead of assuming the SM, they extract the left-handed and right-handed effective coupling constants of leptons and quarks from the forward-backward asymmetry data by imposing a weaker assumption of lepton-quark universality. The obtained results show that the SM fits the extracted left-handed and right-handed couplings of the leptons, and the $u$, $d$ and $c$ quarks very well. However, for the $b$ quark, the extracted left-handed and right-handed couplings are \cite{12}

$$\bar{g}_b^L = -0.4159(24), \quad \bar{g}_b^R = 0.1050(90).$$

These are to be compared with the corresponding SM predictions up to two-loop corrections: $\bar{g}_b^L = -0.4208$ and $\bar{g}_b^R = 0.0774$. We see that the agreement of $\bar{g}_b^L$ is good, while the SM predicted $\bar{g}_b^R$ is smaller than the experimental value by $3.1\sigma$ which is more significant than what we have seen in the global fit.

A similar deviation has also been mentioned in Ref. \cite{1}. The experimentally measurable forward-backward asymmetry $A_{FB}(b)$ and the mixed forward-backward and left-right asymmetry $A_{LR}^{FB}(b)$ for the $b$ quark can be expressed as $A_{FB}(b) = \frac{3}{4}A_eA_b$ and $A_{LR}^{FB}(b) = \frac{3}{4}A_b$, respectively. One can thus determine the asymmetry parameter $A_b$ model-independently both by using the SLC measured value of $A_{LR}^{FB}(b) = \frac{3}{4}(0.867 \pm 0.035)$ and by using the
LEP measured $A_{FB}(b) = 0.0990 \pm 0.0021$ together with the averaged value of the LEP and SLC measured lepton asymmetry parameter $A_l = 0.1489 \pm 0.0018$ for $A_e$. The former gives $A_b = 0.867 \pm 0.035$ and the latter gives $A_b = 0.887 \pm 0.022$. The averaged value is then $A_b = 0.881 \pm 0.019$ which is almost $3\sigma$ below the SM prediction listed in Table I [1].

From the above analyses, we see that there are still considerable deviations between the experiments and the SM predictions provided one concentrates to certain well measured data and analyzes the data model-independently. Of course, the above discrepancies still cannot lead to a definite conclusion of needing new physics beyond the SM. Definite conclusion can only be made together with more future experiments, e.g. experiments at the LC. It is interesting to notice that the discrepancies are related to the $b$ quark in the third family. If the discrepancies really reflect new physics beyond the SM, it will not be surprising since the top quark, as the $SU(2)$ partner of the $b$ quark in the third family, is the heaviest particle yet discovered whose mass is close to the EWSB scale $v = 246$ GeV so that the third family is more sensitive to new physics in the EWSB sector than the first two families do.

III. The Higgs Boson

1. Theoretical Aspect of the SM Higgs Field

Although the SM is successful at present energies, it contains several theoretically unsatisfactory features, especially its Higgs sector.

- There are so many free parameters in the SM and most of them are related to the Higgs filed.
- The self-energy of the elementary Higgs field theory is quadratically divergent which makes the theory unnatural [3,4] and will lead to the hierarchy problem when grand unification is concerned.
- The renormalized coupling constant $\lambda$ of the Higgs self-interaction vanishes in the
continuum limit (i.e. when the regularization momentum cut-off goes to infinity). This is the so-called *triviality* problem [6].

To avoid *triviality*, people usually take the point of view that the SM may not be a fundamental theory but is a low energy effective theory of a more fundamental theory below a certain physical scale $\Lambda$. Since the Higgs mass $m_H$ is proportional to $\lambda$, there is an upper bound on the Higgs mass for a given value of $\Lambda$ [14]. Such a triviality bound on the Higgs mass is shown as the upper curve in Fig. 1 [14]. By definition, $m_H$ cannot exceed $\Lambda$. This determines the maximal value of $m_H$ which is of the order of 1 TeV [6,14].

![Fig. 1. The triviality bound (upper curve) and the vacuum stability bound (lower curve) on $m_H$ in the SM. The solid areas as well as the cross-hatched area indicate theoretical uncertainties. Quoted from Ref. [14].](image)

On the other hand, the Higgs boson self-interaction makes the physical vacuum a stable ground state with nonvanishing vacuum expectation value (VEV), while the fermion-loop contribution to the effective potential tends to *violate the vacuum stability*. The heavier the fermion is, the stronger the violation of vacuum stability will be. The $t$ quark is the heaviest fermion yet found which gives a strong violation of vacuum stability. Therefore we need a strong enough Higgs self-interaction (large enough $m_H$) to overcome the violation effect from the $t$ quark loop and maintain vacuum stability. This requirement gives a lower bound on the Higgs mass $m_H$. The vacuum stability bound on $m_H$ is shown as the lower curve in Fig. 1 [14].
We see from Fig. 1 that the scale of new physics $\Lambda$ can be of the order of the Planck mass only if $m_H \sim 160$ GeV. Of special interest is that if a very light Higgs boson with $m_H \sim 100$ GeV or a heavy Higgs boson with $m_H \gtrsim 500$ GeV is found, the scale of new physics will be of the order of TeV. If the Higgs boson is not found below 1 TeV, we should find new physics beyond the SM in this region.

2. Searching for the Higgs Boson at High Energy Colliders

In this subsection, we briefly review the searches for the Higgs boson at high energy colliders. For technical details, we refer to the talk by S. Yamashita at this workshop.

2.1 The CERN LEP2

At the LEP2 energy, the dominant production mechanism for the SM Higgs boson is the Higgs-strahlung process

\[ e^+e^- \rightarrow Z^* \rightarrow Z \ H, \]  

in which the Higgs boson is emitted from a virtual Z boson. The cross section of this channel is nearly two orders of magnitude larger than that of the W fusion of Higgs except at the edge of the phase space for Higgs-strahlung where both cross sections are very small. In the range of $50 \text{ GeV} < m_H < 120 \text{ GeV}$, the main decay mode of the SM Higgs boson is the $b\bar{b}$ mode [$B(H \rightarrow b\bar{b}) \sim (80 - 90)\%$]. Other branching ratios are smaller by an order of magnitude or more. Experimental detections are in the following four modes

(a) $Z \rightarrow e^+e^-; \mu^+\mu^-$, \quad $H \rightarrow \text{anything}$
(b) $Z \rightarrow \tau^+\tau^-$, \quad $H \rightarrow \text{hadrons}$, and vice versa
(c) $Z \rightarrow \nu\bar{\nu}$, \quad $H \rightarrow \text{hadrons}$
(d) $Z \rightarrow \text{hadrons}$, \quad $H \rightarrow b\bar{b}$.

Main detection backgrounds are fermion-pair productions, single W and Z productions and W, Z pair productions. For eliminating the background $e^+e^- \rightarrow ZZ$ when $m_H \sim M_Z$, the $b$-tagging technique is helpful since $B(Z \rightarrow b\bar{b})$ is only $(15.46\pm0.14)\%$ \cite{1}. The averaged signal
to background ratio at the ALEPH, DELPHI, L3 and OPAL detectors is $(302 \pm 7) : (217 \pm 11)$ \cite{15}. At present no evidence is found. The lower bound on $m_H$ at LEP2 is now \cite{9}.

\[
m_H > 94.1 \text{ GeV}. \tag{3}
\]

2.2 The Upgraded Fermilab Tevatron

In a recent paper \cite{16}, it is shown that the SM Higgs boson in the mass region $135 \text{ GeV} < m_H < 180 \text{ GeV}$ is able to be detected at the upgraded Tevtron with the center-of-mass energy of 2 TeV and an integrated luminosity of 30 fb$^{-1}$ via the process

\[
pp \rightarrow gg \rightarrow H \rightarrow W^+W^- \rightarrow l\bar{\nu}jj, \ l\bar{\nu}\ell\nu. \tag{4}
\]

For the details of the signal and background analysis, see Ref. \cite{16}.

2.3 The CERN LHC

- $m_H > 140$ GeV

  The clearest search is for $m_H > 2M_Z$. In this case the following channel is available

  \[
  pp \rightarrow HX \rightarrow ZZX \rightarrow l^+l^-l^+l^-X \text{ (or } l^+l^-\nu\bar{\nu}X), \tag{5}
  \]

  in which the four-lepton final state is very clear with rather small backgrounds. This channel is usually called the golden channel. Theoretical study shows that the resonance behavior can be clearly seen when $m_H < 800$ GeV \cite{17}.

  If $140 \text{ GeV} < m_H < 2M_Z$, the four-lepton final state can still be detected with one of the $Z$'s virtual \cite{18}.

- $M_Z < m_H < 140$ GeV

  When the Higgs mass is in the intermediate range $M_Z < m_H < 140$ GeV, the above detection is not possible since the branching ratio of the four-lepton channel drops very rapidly as $m_H < 140$ GeV. Detection of such an intermediate-mass SM Higgs boson is much more difficult. Fortunately, the $H \rightarrow \gamma\gamma$ branching ratio has its
maximal value in this $m_H$ range. Thus the best way is to detect the $\gamma\gamma$ final state for the Higgs boson. Recently, it is shown that the SM Higgs boson in the mass range of 100 GeV – 150 GeV can be detected at the LHC via

$$pp \rightarrow H(\gamma\gamma) + jet$$

if a transverse-momentum cut of 2 GeV on the tracks is made for reducing the background \[19\]. The statistical significance $S = N_S / \sqrt{N_B}$ ($N_S$ and $N_B$ stand for the number of signal and number of background, respectively) can be as large as $S > 8$ for $100 \text{ GeV} < m_H < 150 \text{ GeV} \[19\].

To find channels with larger rates and smaller backgrounds, people suggested the following associate productions of $H$ \[20\].

$$pp \rightarrow WHX \rightarrow l\bar{\nu}\gamma\gamma X, \hspace{1cm} (7)$$

$$pp \rightarrow t\bar{t}HX \rightarrow l\bar{\nu}\gamma\gamma X. \hspace{1cm} (8)$$

In the $t\bar{t}H$ associate production channels, the leptons $l\bar{\nu}$ in the final state come from the $W$ decay in $t \rightarrow Wb$, thus it is an inclusive detection (without detecting $b$ and the decay products of $\bar{t}$). The signal and backgrounds of the $WH$ associate production channel have been calculated in Ref. \[20\] which shows that the backgrounds are smaller than the signal even for a mild photon detector with a 3% $\gamma\gamma$ resolution. The inclusive search for the $t\bar{t}H$ associate production suffers from a further large background from $pp \rightarrow W(\rightarrow l\bar{\nu})\gamma\gamma(n-jet)$, ($n = 1, \cdots, 4$) \[21\], and the search is possible only when the $\gamma\gamma$ resolution of the photon detector is of the level of 1% \[21\]. The photon detectors of the CMS and ATLAS collaborations at the LHC are just of this level. Actually, if the jets are also detected, the background can be effectively reduced with certain choices of the jets, and such a detection is possible even for the mild photon detector with 3% $\gamma\gamma$ resolution \[21\].
2.4 The LC

At the LC, the Higgs boson can be produced either by the Higgs-strahlung process (2) or by $WW$ and $ZZ$ fusions

$$e^+e^- \rightarrow \nu\bar{\nu}(WW) \rightarrow \nu\bar{\nu}H,$$  \hspace{1cm} (9)

$$e^+e^- \rightarrow e^+e^-(ZZ) \rightarrow e^+e^-H.$$  \hspace{1cm} (10)

The cross sections for the Higgs-strahlung and $WW$ fusion processes are $\sigma \sim 1/s$ and $\sigma \sim \frac{\ln s}{M_W^2}$, respectively. So that the Higgs-strahlung process is important at $\sqrt{s} \lesssim 500$ GeV (e.g. the KEK JLC), while the $WW$ fusion process is important at $\sqrt{s} > 500$ GeV (e.g. the DESY TESLA and the SLAC NLC). Since there are less hadronic backgrounds at the LC, the Higgs boson can be tagged via the $H \rightarrow b\bar{b}$ channel. Several thousands of events can be produced for the envisaged luminosities [22].

By means of laser back-scattering, $\gamma\gamma$ and $e\gamma$ colliders can be constructed based on the LC. It has been shown recently that the $s$-channel Higgs production rate at the photon collider will be about an order of magnitude larger than the production rate in the Higgs-strahlung process at the LC [23].

It is shown that in the minimal SUSY extension of the SM, the cross section of

$$\gamma\gamma \rightarrow hh$$  \hspace{1cm} (11)

can be significantly enhanced if there is large $\tilde{t}$-mixing [24], so that the lightest SUSY Higgs boson $h$ can be distinguished from that in the usual two-Higgs-doublet model through this process.

It has also been shown that at the $e\gamma$ collider, the process

$$e\gamma \rightarrow eh$$  \hspace{1cm} (12)

is enhanced by an order of magnitude by the chargino-loop contributions. So that this is a plausible process for detecting the lightest SUSY Higgs boson $h$ [25].
The study of the $f \bar{f} \gamma$ decay mode of the Higgs bosons shows that, for large $\tan \beta$ (say $\tan \beta \sim 30$), the decay rates of the SUSY Higgs bosons $h, H, A \rightarrow f \bar{f} \gamma$ are significantly smaller than that of the SM Higgs boson [26], so that the SUSY Higgs bosons can be distinguished from the SM Higgs boson via this decay mode.

2.5 The Muon Collider

In recent years, there is an increasing attention to the construction of muon colliders which is specially advantageous in detecting the SM and the SUSY Higgs bosons in the $s$-channel [27].

$$\mu^+ \mu^- \rightarrow H \rightarrow b(t) \bar{b}(\bar{t}).$$

(13)

If a light Higgs boson can be discovered at other colliders, e.g. the LHC or LC, a muon collider can be set up with energy around $m_H$ which will be specially useful for a better determination of $m_H$ and a measurement of the total width $\Gamma^{\text{tot}}_H$ by scanning around $\sqrt{s} = m_H$ if the energy spread is small enough. Otherwise, one can make a wider scan at the muon colliders to search for the Higgs boson. For $m_H \leq 2M_W$, the branching ratio $B(H \rightarrow \mu^+ \mu^-)$ is not small due to the fact that the important decay channel $H \rightarrow WW, ZZ$ are not open, so that this $s$-channel detection is very useful. In this case, the important decay channels can be $b\bar{b}, WW^*,$ and $ZZ^*$. The $WW^*$ channel can be measured via the final states $l\nu2j$, and the $ZZ^*$ channel can be measured via $2l2j, 2\nu2j, 4l$ and $2l2\nu$. Then without knowing the precise value of $\Gamma(H \rightarrow \mu^+ \mu^-)$, one can measure the ratios of the $H \rightarrow b\bar{b}, H \rightarrow WW^*$ and $H \rightarrow ZZ^*$ rates, and from which one can obtain the $(Hb\bar{b})^2, (HWW^*)^2$ and $(HZZ^*)^2$ coupling-squared ratios to the precision of a few percent [27], which can test the Higgs boson interactions. Higher energy muon colliders can also be used for the discovery of heavier SM Higgs boson and the SUSY Higgs bosons $H$ and $A$ [27]. More on the muon collider physics can be found in Ref. [27].
3. Probing the Higgs Interactions

If a light Higgs resonance is found from the above searches, it is not the end of the story. It is needed to test whether it is the SM Higgs or something else (for instance, a composite Higgs boson in certain strongly interacting EWSB models). This can be done by examining its interactions. We know the self-interactions of the SM Higgs boson contain the following trilinear and quartic terms

\[ \frac{1}{8} \frac{m_H^2}{v^2} \left[ 4v H^3 + H^4 \right], \]  

where \( v \) is the VEV of the Higgs field, \( v = 246 \) GeV. For detecting the trilinear interaction, it is possible to look at the double Higgs-boson productions, i.e. \( pp \rightarrow HHX \) at the LHC and \( e^+e^- \rightarrow HHZ, \) \( HH\bar{\nu}_e\nu_e \) at the LC. It has been shown that the detection at the LHC is almost not possible due to the large background \[22,19\], while the detection at the LC is possible at the C.M. energy \( E = 1.6 \) TeV requiring a very large integrated luminosity, \( \int L dt = 1000 \text{ fb}^{-1} \). Therefore the detection is not easy. The signals of the quartic interactions are so small that it is hard to detect. So that other methods of distinguishing the SM Higgs boson from other Higgs resonances are needed.

Since the top quark has the largest Yukawa coupling to the SM Higgs boson, it is possible to detect the Higgs Yukawa coupling via the process

\[ e^+e^- \rightarrow t\bar{t}H. \]  

This detection has been studied in Refs. \[28,29\].

IV. Supersymmetry

The way of avoiding the unnaturalness problem without giving up the elementary Higgs fields is to introduce SUSY in which the quadratic divergence of the Higgs self-energy causing unnaturalness is canceled by the contribution of its SUSY partner. Furthermore,
inclusion of the SUSY partners of known particles also weakened the problem of triviality. In most SUSY models motivated by avoiding the unnaturalness problem, the mass of the lightest Higgs boson $h$ is rather low, say below 130 GeV [30]. Thus it is weakly interacting and is perturbatively calculable \(^1\). SUSY theory is one of the promising theories of new physics which receives most attention at present. However, it contains even more free parameters than the SM does. The number of free parameters may be reduced when going up to grand unification or even to superstring theory. Anyway, the most convincing test of SUSY models is to find SUSY partners experimentally. In this talk, we only briefly review some processes of searching for SUSY particles. For details, we refer to the talk given by K. Fujii at this workshop.

In the minimal SUSY extension of the SM (MSSM) [31], both SUSY and anomaly-free require that the MSSM should contain two Higgs doublets $H_1$ and $H_2$. Every known particle and its SUSY partner (denoted by the corresponding symbol with a tilde) belong to a SUSY multiplet described by a superfield (denoted by a symbol with a hat). To enforce lepton and baryon number conservation, a discrete symmetry called $R$-parity is usually imposed

$$R = (-1)^{2s+3B+L},$$

(16)

where $s$ is the spin quantum number. Then ordinary particles have $R = 1$ and SUSY particles have $R = -1$. The $R$-parity conserving superpotential related to the Higgs sector can be written as [31]

$$W = \epsilon_{ij}[(h_L)_{mn}\hat{H}_1^i\hat{L}_m^j\hat{E}_n + (h_D)_{mn}\hat{H}_1^i\hat{Q}_m^j\hat{D}_n_{-} - (h_U)_{mn}\hat{H}_2^i\hat{Q}_m^j\hat{U}_n - \mu \hat{H}_1^i\hat{H}_2^j],$$

(17)

where $\hat{H}_1$, $\hat{H}_2$, $\hat{L}_m$, $\hat{E}_n$, $\hat{Q}_m$, $\hat{D}_n$, and $\hat{U}_n$ are, respectively, superfields for $H_1$, $H_2$, the left-handed lepton, the right-handed lepton, the left-handed quark, the right-handed $D$- and $U$-quark, $h$’s are Yukawa couplings, and $\mu$ is a Higgs superfield mass parameter. Since the real world does not respect SUSY exactly, SUSY should be broken. In order to keep the solution of the unnaturalness problem, the SUSY breaking mechanism should not cause 

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\(^1\)It is still possible that there are other non-minimal heavy Higgs bosons which are strongly interacting
uncanceled quadratic divergences so that it should be soft, and the scale of SUSY breaking should not be much higher than the TeV scale. With $R$-parity conservation, the general soft-SUSY-breaking terms are \[ V_{\text{soft}} = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 - m_{12}^2 (\epsilon_{ij} H_1^i H_2^j + \text{h.c.}) + (M_Q^2)_{mn} \tilde{Q}_m^i \tilde{Q}_n^i + (M_U^2)_{mn} \tilde{U}_m^i \tilde{U}_n^i + (M_D^2)_{mn} \tilde{D}_m^i \tilde{D}_n^i + (M_L^2)_{mn} \tilde{L}_m^i \tilde{L}_n^i + (M_E^2)_{mn} \tilde{E}_m^i \tilde{E}_n + \epsilon_{ij} [(h_{L\Lambda} A_{L})_{mn} \tilde{H}^i_1 \tilde{L}^j_m \tilde{E}_n + (h_{D\Lambda} A_{D})_{mn} \tilde{H}^i_1 \tilde{Q}^j_m \tilde{D}_n - (h_{U\Lambda} A_{U})_{mn} \tilde{H}^i_2 \tilde{Q}^j_m \tilde{U}_n + \text{h.c.}]) + \frac{1}{2} [M_3 \tilde{g} \tilde{g} + M_2 \tilde{W}^a \tilde{W}^a + M_1 \tilde{B} \tilde{B} + \text{h.c.}], \tag{18} \]

where the coefficients are complex matrices. The total number of free parameters in the MSSM is 124 (105 new parameters in addition to the 19 SM parameters) \[.\] It is hard to make phenomenological predictions with so many free parameters. Actually, requirement on separate conservation of the lepton numbers $L_e$, $L_\mu$, $L_\tau$ and the experimental bounds on the flavor changing neutral current and the electric dipole moment of the electron and neutron restrict the MSSM parameters to a limited subspace of the total parameter space. There are two general approaches which make the MSSM phenomenologically viable.

- **The minimal supergravity-inspired MSSM (constrained MSSM).**

  In this approach the scalar squared masses and the $A$’s are flavor diagonal and universal at the Planck scale $M_P$

  \[
  M_Q^2(M_P) = M_U^2(M_P) = M_D^2(M_P) = M_L^2(M_P) = M_E^2(M_P) = m^2_0 \mathbf{1},
  \]

  \[
  m^2_1(M_P) = m^2_2(M_P) = m^2_0,
  \]

  \[
  A_U(M_P) = A_D(M_P) = A_L(M_P) = A_0 \mathbf{1}, \tag{19}
  \]

  and, furthermore, the gauge couplings and the gaugino mass parameters are assumed to unify at the grand unification scale $M_X$

  \[
  \sqrt{5/3} g_1(M_X) = g_2(M_X) = g_3(M_X) = g_U,
  \]

  \[
  M_1(M_X) = M_2(M_X) = M_3(M_X) = m_{1/2}. \tag{20}
  \]

  The total number of free parameters in this approach reduces to 23, i.e. the 18 SM parameters (not including the Higgs mass) and 5 additional new parameters
\(m_0, \ m_{1/2}, \ A_0, \ \tan \beta\) (the ratio of the VEV’s of the two Higgs bosons) and \(\text{sgn}(\mu)\). This approach is also called the \textit{constrained MSSM} (CMSSM).

**Models of gauge-mediated SUSY breaking.**

This kind of approach consists of a hidden sector in which SUSY is broken, a "messenger" sector containing messenger fields with \(SU(3) \times SU(2) \times U(1)\) quantum numbers, and a sector containing the fields of the MSSM. The coupling of the messengers to the hidden sector generates a SUSY-breaking spectrum in the messenger sector, and the SUSY breaking is transformed to the MSSM via the virtual exchange of the messengers. In the simplest case of the models, there is one effective mass scale \(\Lambda\) which determines all the low-energy scalar and gaugino masses though loop-effects (no \(A\)-parameters are generated). One must have \(\Lambda \sim 100\) TeV in order to result the SUSY partner masses of \(O(1)\) TeV or less. The generation of \(\mu\) and \(m_{12}\) lies outside the ansatz of this SUSY breaking. There is a messenger scale \(M\) characterizing the average mass of the messenger particles which can lie between \(\Lambda\) and \(10^{16}\) GeV. The total number of free-parameters in this approach is 22, i.e. the 18 SM parameters (not including the Higgs mass), and the four additional new parameters \(\Lambda, \ \tan \beta, \ \text{sgn}(\mu), \ \text{and} \ M\). This simplest approach is called the \textit{minimal gauge-mediated model} (MGM).

Different approaches give rise to different SUSY particle spectra. Particles with the same \(SU(3) \times U(1)_{\text{em}}\) quantum numbers may mix to form mass eigen-states. For example, the charginos \(\tilde{\chi}_j^\pm\), \((j = 1, 2)\) are linear combinations of the charged winos and higgsinos, and the neutralinos \(\tilde{\chi}_k^0\) \((k = 1, \cdots, 4)\) are linear combinations of the neutral wino, bino and neutral higgsinos. With \(R\)-parity conservation, SUSY particles \((R = -1)\) can only be produced in pairs, and heavy SUSY particles \((R = -1)\) cannot decay into ordinary particles \((R = 1)\) but will decay eventually to the \textit{lightest SUSY particle} (LSP) which behaves as a missing energy \(E_T\) in the experiments. These characterize the phenomenology of SUSY searches. In the following, we briefly review some of the SUSY searches at high energy colliders.

**LEP2**

Searching for SUSY particles has been carried out at LEP for many years, and no evidence is found. We quote here some recent gives the following lower bounds of the
SUSY particles given in the $\sqrt{s} = 183$ GeV run.

The DELPHI data on $e^+e^- \rightarrow \tilde{\chi}_1^0\tilde{\chi}_1^0 \rightarrow \tilde{G}\gamma\tilde{G}\gamma$ leads to the lower bound of $m_{\tilde{\chi}_1^0}$ in the MGM assuming $\tilde{\chi}_1^0 \rightarrow \tilde{G}\gamma$ to have 100% probability and negligible lifetime \[32\]

$$m_{\tilde{\chi}_1^0} > 83 \text{ GeV}$$

with $m_{\tilde{\varepsilon}_R} = 1.1m_{\tilde{\chi}_1^0}$ and $\tilde{\chi}_1^0 \approx \tilde{B}$.

The OPAL search for sleptons leads to the following bounds. For $\mu < -100$ GeV and $\tan \beta = 1.5$, the 95% C.L. bounds are \[33\]

$$m_{\tilde{\varepsilon}_R} < 77 \text{ GeV} \quad \text{for} \quad m_{\tilde{\varepsilon}_R} - m_{\tilde{\chi}_1^0} > 5 \text{ GeV,}$$

$$m_{\tilde{\mu}_R} < 65 \text{ GeV} \quad \text{for} \quad m_{\tilde{\mu}_R} - m_{\tilde{\chi}_1^0} > 2 \text{ GeV,}$$

$$m_{\tilde{\tau}_R} > 60 \text{ GeV} \quad \text{for} \quad m_{\tilde{\tau}_R} - m_{\tilde{\chi}_1^0} > 9 \text{ GeV.} \quad (22)$$

It has been shown that the LEP precision measurements at the Z-pole and low-energy electroweak experiments give constraints on the MSSM parameters, especially the $b \rightarrow s\gamma$ data gives severe constraint on the parameter space with $\mu > 0$ and large $\tan \beta$ \[34\].

**LHC**

At the LHC, SUSY particles pairs $\tilde{g}\tilde{g}$, $\tilde{g}\tilde{q}$, $\tilde{q}\tilde{q}$, $\tilde{\ell}\tilde{\ell}$, $\tilde{\chi}_1^0\tilde{\chi}_1^\pm$ can be searched for via the following channels \[35\]

$$l^\pm + \text{jets} + E_T : \quad \tilde{g}\tilde{g}, \tilde{q}\tilde{q} \quad (23)$$

$$l^\pm l^\pm + \text{jets} + E_T : \quad \tilde{g}\tilde{g}, \tilde{q}\tilde{q}, \tilde{q}\tilde{q}$$

$$l^\pm l^\pm l^\pm + E_T : \quad \tilde{\chi}_1^0\tilde{\chi}_1^\pm$$

$$l^\pm l^\pm + E_T : \quad \tilde{\ell}\tilde{\ell}.$$ 

The LHC is advantageous for detecting colored SUSY particles. For example, in the first channel in \[24\], $\tilde{g}$ and $\tilde{q}$ can be detected up the mass values of 3.6 TeV \[37\].

**LC**

The LC is advantageous due to the small backgrounds. It provides high precision detections of all SUSY particles up to the mass values of 1 TeV \[22\].
Neutralinos and charginos are easy to detect and study with high accuracy at the LC via \[22\]

\[
e^+e^- \rightarrow \tilde{\chi}_i^+\tilde{\chi}_j^- \quad i, j = 1, 2 \tag{24}
\]

\[
e^+e^- \rightarrow \tilde{\chi}_i^0\tilde{\chi}_j^0 \quad i, j = 1, 2, 3, 4
\]

with \(\tilde{\chi}_1\tilde{\chi}_1 \rightarrow \tilde{G}\tilde{G}\gamma\gamma \rightarrow \gamma\gamma E_T\). The neutralino and chargino masses can be measured to the accuracy of \(\delta m_{\tilde{\chi}_1^\pm} \approx 100\) MeV, \(\delta m_{\tilde{\chi}_1^0} \approx 600\) MeV.

The sleptons can be detected and studied via \[22\]

\[
e^+e^- \rightarrow \tilde{l}_i^+\tilde{l}_j^-, \tilde{l}_R^+\tilde{l}_R^-, \text{ etc.} \tag{25}
\]

The smuon mass can be measured to the accuracy of \(\delta m_{\tilde{\mu}} \approx 1.8\) GeV.

It has been shown that the cross sections of \(e^-e^- \rightarrow \tilde{e}_{L,R}^-\tilde{e}_{L,R}^-\) are much larger than those of \(e^+e^- \rightarrow \tilde{e}_{L,R}^+\tilde{e}_{L,R}^-\) \[30\], so that the \(e^-e^-\) linear collider is specially advantageous for the search for selectron.

We see that, with the LHC and the LC, it is possible to search for all the SUSY particles with the mass values up to 1 TeV which covers the range required by solving the unnaturalness problem by SUSY. If SUSY particles are found below 1 TeV, SUSY will be confirmed and one should study the properties of the SUSY particles seriously to see if the model is just the simple MSSM or certain complicated SUSY model beyond the MSSM. If SUSY particles are not found in this energy region, SUSY will be irrelevant to the solution of the unnaturalness problem, and we should consider other possible solutions, e.g. strongly interacting EWSB mechanism (without introducing elementary Higgs fields) which will be reviewed in the next section.

Another interesting feature is that the careful theoretical study on the MSSM Higgs mass up to two-loop calculations show that the mass of the lightest CP even Higgs boson \(h\) in the MSSM cannot exceed a bound \(m_h |_{max}\approx 130\) GeV \[30\]. Thus the Higgs boson \(h\) can be searched for at all the designed LC’s (NLC, TESLA and JLC). If a light scalar resonance is found below 130 GeV, it will be a good support to the MSSM, and further study of its
properties will be needed to see if it is really the MSSM Higgs boson. If no light scalar resonance is found below 130 GeV, MSSM will be in a bad shape and SUSY models beyond the MSSM should be seriously considered.

V. Strongly Interacting Electroweak Symmetry Breaking Mechanism

1. Strongly Interacting Electroweak Symmetry Breaking Models

Introducing elementary Higgs field is the simplest but not unique way of breaking the electroweak gauge symmetry spontaneously. The way of completely avoiding the problems of triviality and unnaturalness is to abandon elementary scalar fields and introducing new strong interactions causing certain fermion condensates to break the electroweak gauge symmetry. This idea is similar to those in the theory of superconductivity and chiral symmetry breaking in QCD. The simplest model realizing this idea is the original QCD-like technicolor model \[37,4\]. However, such a simple model predicts a too large oblique correction parameter \(S\) \[38\] and is already ruled out by the LEP data. A series of improved models have been proposed to overcome the shortcomings of the simplest model \[39–44,11\]. In the following, we briefly review two of the recently proposed models.

- Topcolor-Assisted Technicolor Models

This model combines the technicolor idea and the top-condensate idea, in which a large enough \(m_t\) and \(m_b\) mass difference can be obtained without causing large oblique correction parameters \(S, T\) and \(U\) \[12\]. It is assumed in this model that at the energy scale \(\Lambda \gtrsim 1\) TeV, there is a topcolor theory with the gauge group \(SU(3)_1 \times U(1)_{Y1} \times SU(3)_2 \times U(1)_{Y2} \times SU(2)_{L}\) in which \(SU(3)_1 \times U(1)_{Y1}\) preferentially couples to the third-family fermions and \(SU(3)_2 \times U(1)_{Y2}\) preferentially couples to the first- and second-family fermions, so that the third-family fermions are different from those in the first two families. It is assumed that there is also a technicolor sector mainly in charge of the EWSB and will break the topcolor gauge group into \(SU(3)_{QCD}\)
and $U(1)_Y$ at the scale $\Lambda$:

$$SU(3)_1 \times U(1)_{Y_1} \times SU(3)_2 \times U(1)_{Y_2} \times SU(2)_L \to SU(3)_{QCD} \times U(1)_Y \times SU(2)_L. \quad (26)$$

The $SU(3)_1 \times U(1)_{Y_1}$ couplings are assumed to be much stronger than those of $SU(3)_2 \times U(1)_{Y_2}$. The strong $SU(3)_1 \times U(1)_{Y_1}$ interactions will form top quark condensate $\langle t\bar{t} \rangle$ but not bottom quark condensate from the simultaneous effects of the $SU(3)_1$ and $U(1)_{Y_1}$ interactions. The technicolor dynamics gives rise to the masses of the $u$, $d$, $s$, $c$, and $b$ quarks and a small portion of the top quark mass, while the main part of the top quark mass comes from the topcolor dynamics causing the top quark condensate just like the constituent quarks acquiring their large dynamical masses from the dynamics causing the quark condensates in QCD. In this prescription, the technicolor dynamics does not cause a large oblique correction parameter $T$. Improvement of this kind of model is still in progress [15]. It has been shown that this kind of model gives rise to a positive enhancement to $R_b$ with the right order of magnitude, which depends on the parameters in the model [16].

This kind of model contains the usual pseudo-Goldstone bosons (PGB’s) in the technicolor sector with masses of few hundred GeV and an isospin triplet top-pions with masses around 200 GeV. These light particles characterizing the phenomenology of the model.

- **Top Quark Seesaw Theory**

Very recently, a new promising theory of strongly interacting EWSB related to the top quark condenstate called top quark seesaw theory was proposed [11]. The gauge group in this theory is $SU(3)_1 \times SU(3)_2 \times SU(2)_W \times U(1)_Y$ which breaks into $SU(3)_{QCD} \times U(1)_Y$ at a scale $\Lambda$. Instead of introducing techniquarks, certain $SU(2)_W$-singlet quarks, $\chi, \cdots$, with topcolor interactions and specially assigned $U(1)_Y$ quantum numbers are introduced in this theory. Topcolor causes the following $t$ ($b$) and $\chi$ bound state scalar field

$$\varphi = \begin{pmatrix} \chi_R & t_L \\ \chi_R & b_L \end{pmatrix} \quad (27)$$

which behaves like a Higgs doublet whose VEV breaks the electroweak symmetry.
Furthermore, the VEV of $\varphi$ will cause a dynamical mass $m_{t\chi} \sim 600$ GeV, and the dynamics in this theory will cause the following mass terms in the $\chi - t$ sector

$$- \mu_{\chi\chi} \chi L \chi_R - \mu_{\chi t} \chi t L R + h.c.$$

(28)

Then the mass matrix of the heavy charge $2/3$ quarks takes the form

$$\begin{pmatrix} t_L & \chi_L \end{pmatrix} \begin{pmatrix} 0 & m_{t\chi} \\ \mu_{\chi t} & \mu_{\chi\chi} \end{pmatrix} \begin{pmatrix} t_R \\ \chi_R \end{pmatrix}$$

(29)

which gives the seesaw mechanism for the top quark mass. Diagonalizing the mass matrix in (29) for $\mu_{\chi\chi} \gg m_{t\chi}$ leads to the physical top quark mass

$$m_t \approx m_{t\chi} \frac{\mu_{\chi t}}{\mu_{\chi\chi}}.$$

(30)

Eq. (30) can yield the desired top quark mass for appropriate dynamical value of $\mu_{\chi t}/\mu_{\chi\chi}$.

The composite scalar spectrum in this theory is:

- $h^0$: neutral Higgs boson of mass $m_{t\chi}$ times a factor of order one (or smaller);
- $H^0, H^\pm, A^0$: heavy states of a two Higgs-doublet sector (roughly degenerate);
- $H^0_{\chi t}, A^0_{\chi t}$: a CP-even and a CP-odd state (can be light);
- $A^0_{\chi\chi}$: a CP-odd state (can be light);
- $\phi_{bb}$: neutral complex scalar (with arbitrary mass);
- $H^\pm_{tb}$: charged scalar which can be as light as 250 GeV;
- $H^0_{\chi\chi}, H^\pm_{\chi b}$: CP-even neutral state and a charged scalar with large masses.

This theory has several advantages. (a) In this model, one of the particles responsible for the EWSB is just the known top quark, and the $SU(2)_W$-doublet nature of the Higgs field just comes from the $SU(2)_W$-doublet nature of the third family quarks. (b) The new quark $\chi$ introduced in this theory is $SU(2)_W$-singlet so that there is no large custodial symmetry violation, and the experimental constraint on the oblique correction parameter $T$ can be easily satisfied. (c) The problem of predicting a too
large oblique correction parameter $S$ in technicolor theories due to introducing many
technifermion-doublets does not exist in the present theory since there is only one top
quark condensate. (d) Unlike the original top quark condensate model which leads to
a too large top quark mass, the present theory can give rise to the desired top quark
mass via the seesaw mechanism.

Various possible ways of building models in this theory are given in Ref. [11].

2. Model-Dependent Tests

Due to the nonperturbative nature of the strong interaction dynamics, it is hard to
make precision predictions from the strongly interacting electroweak symmetry breaking
models. However, most of the models contain certain PGB’s with masses in the region of
few hundred GeV. The properties of the PGB’s are different from model to model, therefore
the PGB’s are characteristics of the models, and their effects can be experimentally tested.
Productions of PGB’s, especially the PGB’s in the technicolor sector, at the existing high
energy colliders have been extensively studied in the literatures [47,48].

Since the top quark couples to the EWSB sector strongly due to its large mass, a feasible
way of testing the strongly interacting EWSB models is to test the PGB effects in top
quark productions at high energy colliders. It is shown that at the Tevatron [49], the LHC
[50] and the photon collider [51], the s-channel PGB effects in top quark pair productions
are experimentally detectable, and with which models with and without topcolor can be
experimentally distinguished and can be distinguished from the MSSM [51].

3. Model-Independent Probe of the EWSB Mechanism

We have seen that there are various kinds of EWSB models proposed. Some of them
contain light Higgs boson(s) (elementary or composite) and some of them do not. We still do
not know whether the actual EWSB mechanism in the nature looks like one of the proposed
models or not. Therefore, only testing the proposed models seems to be not enough, and
certain model-independent probe of the EWSB mechanism is needed. Since the scale of
new physics is likely to be several TeV, electroweak physics at energy $E \lesssim 1$ TeV can be effectively described by the *electroweak effective Lagrangian* in which composite fields are approximately described by effective local fields. The electroweak effective Lagrangian is a *general description* (including all kinds of models) which contains certain yet unknown coefficients whose values are, in principle, determined by the underlying dynamics. Different EWSB models give rise to different sets of coefficients. The model-independent probe (the first step) is to investigate through what processes and to what accuracy we can measure these coefficients in the experiments. Once this is done, the second step is to examine what kind of model can give rise to a set of coefficients fitting the experimental values. The second step concerns the difficult strong interaction dynamics, and we shall not discuss it in this short review article. We shall mainly consider the first step.

From the experimental point of view, the most challenging case of probing the EWSB mechanism is that there is no light scalar resonance found below 1 TeV. We shall take this case as the example in this review. Effective Lagrangian including a light Higgs boson has also been studied in the literature [53]. In the case we are considering, the effective Lagrangian is the so called electroweak chiral Lagrangian (EWCL) which is an effective Lagrangian for the would-be Goldstone bosons $\pi^a$ in the nonlinear realization $U = e^{i\tau^a \pi^a / f_\pi}$ with electroweak interactions. The bosonic sector of which, up to the $p^4$-order, reads

$$L_{\text{eff}} = L_G + L_S,$$

with

$$L_G = -\frac{1}{2}\text{Tr}(W_{\mu\nu}W^{\mu\nu}) - \frac{1}{4}B_{\mu\nu}B^{\mu\nu},$$

$$L_S = L^{(2)} + L^{(2)'} + \sum_{n=1}^{14} L_n,$$

where $W_{\mu} \equiv W_{\mu}^{a} \tau^{a} / 2$, $B_{\mu} \equiv B_{\mu} \tau^{3} / 2$, $W_{\mu\nu} = \partial_{\mu} W_{\nu} - \partial_{\nu} W_{\mu} + ig[W_{\mu}, W_{\nu}]$, $B_{\mu\nu} = \partial_{\mu} B_{\nu} - \partial_{\nu} B_{\mu}$, and

---

2This is in a similar spirit as the study of the chiral Lagrangian in QCD by Gasser and Leutwyler [52].
\[ \mathcal{L}^{(2)} = \frac{f^2}{4} \text{Tr}[(D_\mu U)^\dagger (D^\mu U)] \quad \text{and} \quad \mathcal{L}^{(2)\nu} = \frac{\ell_0}{4} (\frac{f}{\Lambda})^2 \frac{f^2}{4} \text{Tr}(\mathcal{T} \mathcal{V}_\mu)]^2 \]

\[ \mathcal{L}_1 = \ell_1 (\frac{f}{\Lambda})^2 \frac{g g'}{2} B_{\mu\nu} \text{Tr}(\mathcal{T} \mathcal{W}^{\mu\nu}) \]
\[ \mathcal{L}_2 = \ell_2 (\frac{f}{\Lambda})^2 \frac{g g'}{2} \mathcal{B}_{\mu\nu} \text{Tr}(\mathcal{T} [\mathcal{V}_\mu, \mathcal{V}_\nu]) \]
\[ \mathcal{L}_3 = \ell_3 (\frac{f}{\Lambda})^2 ig \text{Tr}(\mathcal{W}_{\mu\nu} [\mathcal{V}_\mu, \mathcal{V}_\nu]) \]
\[ \mathcal{L}_4 = \ell_4 (\frac{f}{\Lambda})^2 [\text{Tr}(\mathcal{V}_\mu \mathcal{V}_\nu)]^2 \]
\[ \mathcal{L}_5 = \ell_5 (\frac{f}{\Lambda})^2 [\text{Tr}(\mathcal{V}_\mu \mathcal{V}_\mu)]^2 \]
\[ \mathcal{L}_6 = \ell_6 (\frac{f}{\Lambda})^2 [\text{Tr}(\mathcal{V}_\mu \mathcal{V}_\nu)] [\text{Tr}(\mathcal{T} \mathcal{V}_\mu) [\text{Tr}(\mathcal{T} \mathcal{V}_\nu)] \]
\[ \mathcal{L}_7 = \ell_7 (\frac{f}{\Lambda})^2 [\text{Tr}(\mathcal{V}_\mu \mathcal{V}_\nu)] [\text{Tr}(\mathcal{V}_\nu \mathcal{V}_\nu) [\text{Tr}(\mathcal{T} \mathcal{V}_\mu)]^2 \]
\[ \mathcal{L}_8 = \ell_8 (\frac{f}{\Lambda})^2 \frac{2}{4} [\text{Tr}(\mathcal{T} \mathcal{W}_{\mu\nu})]^2 \]
\[ \mathcal{L}_9 = \ell_9 (\frac{f}{\Lambda})^2 \frac{ig}{2} \text{Tr}(\mathcal{T} \mathcal{W}_{\mu\nu}) [\text{Tr}(\mathcal{T} [\mathcal{V}_\nu, \mathcal{V}_\nu]) \]
\[ \mathcal{L}_{10} = \ell_{10} (\frac{f}{\Lambda})^2 \frac{1}{2} [\text{Tr}(\mathcal{T} \mathcal{V}_\mu) [\text{Tr}(\mathcal{T} \mathcal{V}_\nu)]^2 \]
\[ \mathcal{L}_{11} = \ell_{11} (\frac{f}{\Lambda})^2 g e^{\mu\nu\rho\lambda} \text{Tr}(\mathcal{T} \mathcal{V}_\mu) [\text{Tr}(\mathcal{V}_\nu \mathcal{W}_{\rho\lambda}) \]
\[ \mathcal{L}_{12} = \ell_{12} (\frac{f}{\Lambda})^2 \frac{2g}{2} [\text{Tr}(\mathcal{T} \mathcal{V}_\mu) [\text{Tr}(\mathcal{V}_\nu \mathcal{W}_{\mu})^2 \]
\[ \mathcal{L}_{13} = \ell_{13} (\frac{f}{\Lambda})^2 \frac{g g'}{4} e^{\mu\nu\rho\lambda} B_{\mu\nu} \text{Tr}(\mathcal{T} \mathcal{W}_{\rho\lambda}) \]
\[ \mathcal{L}_{14} = \ell_{14} (\frac{f}{\Lambda})^2 \frac{g^2}{8} e^{\mu\nu\rho\lambda} \text{Tr}(\mathcal{T} \mathcal{W}_{\mu\nu}) [\text{Tr}(\mathcal{T} \mathcal{W}_{\rho\lambda}) \]

in which \( D_\mu U = \partial_\mu U + ig \mathcal{W}_\mu U - ig' U \mathcal{B}_\mu \), \( \mathcal{V}_\mu \equiv (D_\mu U)^\dagger \), and \( \mathcal{T} \equiv U T_3 U^\dagger \).

The coefficients \( \ell \)'s reflect the strengths of the \( \pi^a \) interactions, i.e. the EWSB mechanism. \( \ell_1 \), \( \ell_6 \) and \( \ell_8 \) are related to the oblique correction parameters \( S \), \( T \) and \( U \), respectively; \( \ell_2 \), \( \ell_3 \), \( \ell_9 \) are related to the triple-gauge-couplings; \( \ell_{12} \), \( \ell_{13} \) and \( \ell_{14} \) are CP-violating. The task now is to find out experimental processes to measure the yet undetermined \( \ell \)'s.

Note that the would-be-Goldstone bosons \( \pi^a \) are not physical particles, so that they are not experimentally observable. However, due to the Higgs mechanism, the degrees of freedom of \( \pi^a \) are related to the longitudinal components of the weak bosons \( V_L^a \) (\( W_L^\pm \), \( Z_L^0 \)) which are experimentally observable. Thus the \( \ell \)'s are able to be measured via \( V_L^a \)-processes. So that we need to know the qualitative relation between the \( V_L^a \)-amplitude (related to the experimental data) and the GB-amplitude (reflecting the EWSB mechanism), which is the

\(^3\text{Note that with} \ell_1, \ell_6, \ell_8 \text{ taking the values from the experimentally values of} S, T \text{ and} U, \ell_2, \ell_3, \ell_9 \text{ consistent with the experimental bounds on the triple-gauge-couplings which are still rather weak, and a suitable} Zb\bar{b} \text{-coupling, the theory described by this chiral Lagrangian (without a light Higgs boson) can fit the LEP precision electroweak data.}\)
The precise formulation of the ET is

$$T[V_L^{a_1}, V_L^{a_2}, \ldots] = C \cdot T[-i\pi^{a_1}, 1\pi^{a_2}, \ldots] + B$$

(34)

with

$$E_j \sim k_j \gg M_W, \quad (j = 1, 2, \ldots, n),$$

$$C \cdot T[-i\pi^{a_1}, -i\pi^{a_2}, \ldots] \gg B$$

(35)

where $T[V_L^{a_1}, V_L^{a_2}, \ldots]$ and $T[-i\pi^{a_1}, -i\pi^{a_2}, \ldots]$ are, respectively, the $V_L^a$-amplitude and the $\pi^a$-amplitude, $E_j$ is the energy of the $j$-th external line, $C$ is a gauge and renormalization scheme dependent constant factor, and $B$ is a process-dependent function of the energy $E$. This precise formulation has been proved both in the SM and in the EWCL formalism [58]. By taking special convenient renormalization scheme, the constant $C$ can be simplified to $C = 1$ [58–60]. In the EWCL theory, the $B$-term may not be small even when the center-of-mass energy $E \gg M_W$, and it is not sensitive to the EWSB mechanism. Therefore the $B$-term serves as an intrinsic background when probing $T[-i\pi^{a_1}, -i\pi^{a_2}, \ldots]$ via $T[V_L^{a_1}, V_L^{a_2}, \ldots]$ in (34). Only when $|B| \ll |C \cdot T[-i\pi^{a_1}, -i\pi^{a_2}, \ldots]|$ the probe can be sensitive. In Ref. [61], a new power counting rule for semi-quantitatively estimating the amplitudes in the EWCL theory was proposed, and with which a systematic analysis on the sensitivities of probing the EWSB mechanism via the $V_L^a$ processes were given. The results are summarized in Table II.
Table II. Probing the EWSB Sector at High Energy Colliders: A Global Classification for the NLO Bosonic Operators. Quoted from Ref. [1].

(Notations: √ = Leading contributions, △ = Sub-leading contributions, and ⊥ = Low-energy contributions. Notes:  Here, \( L_{13} \) or \( L_{14} \) does not contribute at \( O(1/\Lambda^2) \).  At LHC(14), \( W^+W^+ \rightarrow W^+W^+ \) should also be included.)

| Operators       | \( L^{(2)} \) | \( L_{1,13} \) | \( L_{2} \) | \( L_{3} \) | \( L_{4,5} \) | \( L_{6,7} \) | \( L_{8,14} \) | \( L_{9} \) | \( L_{10} \) | \( L_{11,12} \) | T1 || B | Processes                  |
|-----------------|---------------|----------------|------------|------------|------------|------------|------------|----------|----------|------------|-----|----------------|---------------------------------|
| LEP-I (S,T,U)   | ⊥             | ⊥ †            | ⊥ †        |            |            |            |            |          |          |            |     |               | \( g^4 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( e^+e^- \rightarrow Z \rightarrow f\bar{f} \) |
| LEP-II          | ⊥             | ⊥             | ⊥ †        | ⊥          | ⊥          | ⊥          |            |          |          |            |     |               | \( g^4 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( e^+e^- \rightarrow W^-W^+ \) |
| LC(0.5)/LHC(14) | √             | √             | √          | √          | ⊥          | ⊥          |            |          |          |            |     |               | \( g^2 E^2 \Lambda^2 \leftrightarrow g^3 M_W^2 \Lambda^2 \) | \( f\bar{f} \rightarrow W^-W^+ Z/(LL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^3 E^2 \Lambda^2 \leftrightarrow g^3 M_W^2 \Lambda^2 \) | \( f\bar{f} \rightarrow W^-W^+ Z/(LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^3 E^2 \Lambda^2 \leftrightarrow g^3 M_W^2 \Lambda^2 \) | \( f\bar{f} \rightarrow W^-W^+ Z/(LLL) \) |
|                 | √             | √             | √          | √          | ⊥          | ⊥          |            |          |          |            |     |               | \( g^2 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( f\bar{f} \rightarrow ZZ Z/(LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^3 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( f\bar{f} \rightarrow ZZ Z/(LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^2 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | ZZ \( \rightarrow ZZ (LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^3 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | ZZ \( \rightarrow ZZ (LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^2 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | ZZ \( \rightarrow ZZ (LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^3 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | ZZ \( \rightarrow ZZ (LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^2 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | ZZ \( \rightarrow ZZ (LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^3 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( q\bar{q} \rightarrow W^+Z/(LL) \) |
| LHC(14)         | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^3 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( q\bar{q} \rightarrow W^+Z/(LT) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^2 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( q\bar{q} \rightarrow W^+W^+ Z/(LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^3 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( q\bar{q} \rightarrow W^+W^+ Z/(LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^2 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( q\bar{q} \rightarrow W^+Z/(LL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^3 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( q\bar{q} \rightarrow W^+Z/(LT) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^2 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( q\bar{q} \rightarrow ZZ Z/(LLL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( g^3 E^2 \Lambda^2 \leftrightarrow g^2 M_W^2 \Lambda^2 \) | \( q\bar{q} \rightarrow ZZ Z/(LLL) \) |
| LC(\( e^-\gamma \)) | √             | √             | √          | √          | √          | √          |            |          |          |            |     |               | \( e^2 E^2 \Lambda^2 \leftrightarrow e^2 M_W^2 \Lambda^2 \) | \( e^-\gamma \rightarrow \nu e^-W^-Z, e^-WW/(LL) \) |
| LC(\( \gamma\gamma \)) | √             | √             | √          | √          | √          | √          |            |          |          |            |     |               | \( e^2 E^2 \Lambda^2 \leftrightarrow e^2 M_W^2 \Lambda^2 \) | \( \gamma\gamma \rightarrow W^+Z/(LL) \) |
|                 | ∆             | ∆             | ∆          | ∆          | ∆          | ∆          |            |          |          |            |     |               | \( e^2 E^2 \Lambda^2 \leftrightarrow e^2 M_W^2 \Lambda^2 \) | \( \gamma\gamma \rightarrow W^-W^+ /(LT) \) |
We see that the coefficients $\ell$'s can be experimentally determined via various $V_L^a$ processes at various phases of the LHC and the LC (including the $e\gamma$ collider) complementarily. Without the LC, the LHC itself is not enough for determining all the coefficients. Quantitative calculations on the determination of the quartic-$V_L^a$-couplings $\ell_4$ and $\ell_5$ at the 1.6 TeV LC has been carried out in Ref. [62]. The results are shown in Fig. 2 which shows that with polarized electron beams, $\ell_4$ and $\ell_5$ can be determined at a higher accuracy. Determination of custodial-symmetry-violating-term coefficients $\ell_6$ and $\ell_7$ via the interplay between the $V_LV_L$ fusion and $VVV$ production has been studied in Ref. [63].

![Fig. 2. Determining the coefficients $\ell_4$, $\ell_5$ at the 1.6 TeV $e^+e^-/e^-e^-$ LC’s. The $\pm 1\sigma$ exclusion contours are displayed. (a) unpolarized case; (b) the case of 90% (65%) polarized $e^- (e^+)$ beam. The thick solid lines are contributions from certain simple theoretical models. Quoted from Ref. [62].](image)

As we have mentioned, the top quark couples strongly to the EWSB sector, so that studying effective anomalous couplings of the top quark will be very helpful for model-independent probe of the EWSB mechanism. This kind of study has been carried out in Refs. [64–67].
VI. The Top Quark

The details of top quark physics are given in the talk by G.P. Yeh at this workshop. Here we only mention a few topics related to top quark physics at the LC.

The precise measurements of the properties of the top quark are important for testing the SM and probing new physics. For $m_t \sim 175$ GeV, the maximal cross section of

$$e^+ e^- \xrightarrow{\gamma,Z} t\bar{t}$$  \hspace{1cm} (36)

at the LC is

$$\sigma(t\bar{t}) \sim 800 \text{ fb}$$  \hspace{1cm} (37)

at about 30 GeV above the threshold. With an integrated luminosity of $\int \mathcal{L} dt = 50 \text{ fb}^{-1}$, detailed experimental simulations predict that the top quark mass can be determined to the accuracy of

$$\delta m_t = 200 \text{ MeV}$$  \hspace{1cm} (38)

at $m_t \sim 180$ GeV \cite{68}. This is better than the claimed $\delta m_t = 3$ GeV at the proton colliders.

If the electron in (36) is left-handed polarized, the produced top quark is preferentially left-handed in the forward direction while only a small fraction is produced as right-handed particle in the backward direction. Thus the backward direction is very sensitive to a small anomalous magnetic moment of the top quark which can be measured to a precision of a few percent \cite{68}. Electric dipole moment of the top quark can also be well measured \cite{68}.

Recent Fermilab CDF data on branching ratio of $t \rightarrow b + W^+ \rightarrow b + l^+ + \nu_l$ is

$$B(t \rightarrow bl^+\nu_l) = 0.188 \pm 0.048,$$  \hspace{1cm} (39)

which is consistent with the tree-level SM prediction $2/9$ \cite{69}. Further precise measurements can give constraints on new physics. For instance, it can give constraint on $\tan \beta$ in the
MSSM since $t \rightarrow b + H^+$ and $t \rightarrow \tilde{b} + \tilde{\chi}_1^+$ increase rapidly with \( \tan \beta \); it can give constraint on the technipion and top-pion masses since $t \rightarrow \pi^+, \pi_i^+$ will be large for light $\pi^+$ and $\pi_i^+$.

It has been expected that the threshold measurement of $t\bar{t}$ production at the LC can provide precision measurements of $m_t$ and $\alpha_s$ due to the large cross section near the threshold. However, there are subtleties in the theoretical calculation. Near the threshold, both $\alpha_s$ and the relative velocity $v$ are small, and $\alpha_s/v$ is of the order of $1$ and have to be summed up to all orders even in the leading order approximation. To next-to-leading order (NLO) and next-to-next-to-leading order (NNLO), one has to resum all $(\alpha_s/v)^N(1, \alpha_s, v, \alpha_s^2, \alpha_s v, v^2)$ terms. Recent calculations of the resummation [71] show that the NNLO terms are of the same order as the NLO terms ($\sim 10\%$) (cf. Fig. 3). So that the convergence is actually not good, and thus the theoretical uncertainty cannot be so small as expected.

![Fig. 3. The LO (dashed-dotted lines), NLO (dashed lines) and NNLO (solid lines) contributions to the rate of $e^+e^- \rightarrow t\bar{t}$. Quoted from Ref. [71].](image)

In view of the large $t\bar{t}$ production cross section at the photoh-photon collider, another recent investigation suggested to do the precision measurements via $\gamma\gamma \rightarrow t\bar{t}$ off the threshold region to avoid the difficult threshold calculation [72]. The calculation in Ref. [72] shows that the optimal energy is $\sqrt{s_{\gamma\gamma}} = 420$ GeV which should be accessible at a $\sqrt{s_{e^+e^-}} \gtrsim 500$ GeV LC, and at this collider with a typical $\gamma\gamma$ integrated luminosity of $20$ fb$^{-1}$, $\alpha_s$ can be
determined at the accuracy of 3% statistically [72]. A better accuracy can be achieved with a larger luminosity.

VII. Conclusions

The SM has successfully passed a lot of precision experimental tests, but there are still certain notable deviation when doing the model-independent analyses, namely the $3\sigma$ deviation of the right-handed effective coupling of the $b$-quark. At present, we cannot draw conclusion on the need of new physics beyond the SM only from this deviation. Clearer tests of new physics will be done at future high energy collider experiments.

Despite of the success of the SM, its EWSB sector is still not clear. The assumed Higgs boson in the SM has not been found, and the SM Higgs sector suffers from the well-known problems of triviality and unnaturalness, so that the EWSB sector may concern new physics. Although the global fit of the SM to the precision electroweak data favors a light Higgs boson, there are possible new physics models without a light Higgs boson that can also fit the precision data. So that the search for the Higgs boson should be carried out in the whole possible energy range up to 1 TeV.

Since the EWSB mechanism concerns the understanding of the origin of particle masses, the probe of it is a very interesting and important topic in current particle physics. If a light Higgs boson (elementary or composite) exists, it can be found, as we have seen, at the future high energy colliders such as the LHC, the LC (including the $\gamma\gamma$ and $e\gamma$ colliders), etc. The LC has the advantage of low hadronic backgrounds. After finding the Higgs boson, we have to further study its properties to see if it is just the SM Higgs boson, or a Higgs boson in a more complicated new physics model (e.g. the MSSM), or it is composite. If there is no light Higgs boson, a feasible way of probing the EWSB mechanism is to study the longitudinal weak boson reactions and $t\bar{t}$ productions. We have seen that, for this purpose, the LHC alone is not enough and the LC (including the $\gamma\gamma$ and $e\gamma$ colliders) are needed.
SUSY and strongly interacting EWSB mechanism are two possible candidates of new physics which can avoid the shortcomings of the SM Higgs sector. The most convincing way of testing SUSY is to search for SUSY particles at high energy colliders. If a SUSY particle is found, we should study its properties to see if the SUSY model is just the simplest MSSM or more complicated ones. An interesting conclusion is that if a light Higgs boson is not found below 130 GeV, the MSSM will be in a bad shape and non-minimal SUSY models will be favored. After finding the SUSY particles, our direction may be lead to SUSY GUT and superstring to find out the origin of the large number of free parameters in the SUSY model. If no SUSY particles are found below 1 TeV, SUSY may not be relevant to the solution of the unnaturalness problem. Then the possible solution may be the strongly interacting EWSB mechanism in which both the unnaturalness and triviality problems no long exist. In this case, our direction may be lead to the study of new strong interactions and new composite particles above 1 TeV. In all these studies, the LC will play an important role.

Because of its large mass, the top quark couples strongly to the EWSM sector, so that it plays an important role in probing the EWSB mechanism. In addition, high precision study of the properties of the top quark is important in testing the SM and studying new physics. This can be effectively done at the LC (including the $\gamma\gamma$ and $e\gamma$ colliders).

In summary, particle physics will be in a crucial status of clarifying the choice of different directions of new physics when we go to the TeV energy scale. The LC will be an important equipment for studying TeV physics and will help us to know to which direction we should further go. Further theoretical and experimental studies of LC physics are needed.

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