The Oblique Corrections from Heavy Scalars in Irreducible Representations

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The contributions to $S$, $T$, and $U$ from heavy scalars in any irreducible representation of the electroweak gauge group $SU(2)_L \times U(1)_Y$ are obtained. We find that in the case of a heavy scalar doublet there is a slight difference between the $S$ parameter we have obtained and that in previous works.

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I. INTRODUCTION

In spite of its tremendous success, the standard model (SM) has several drawbacks. The Higgs particle has not yet been found in experiments; on the other hand, the SM suffers unnaturalness and triviality from a theoretical point of view. Thus, the SM may not be correct, or at least it is just an effective theory at the electroweak scale. There are many new physics possibilities beyond the SM. Although we do not know whether nature really behaves like one of them or not, we can estimate their effects on the current electroweak precision measurements. Peskin and Takeuchi's $S$-$T$-$U$-formalism is a practical way to do this job [1]. Since the current SM parameter fits indicate that $S$ and $T$ are small negative numbers, and $U$ is also close to zero [2], those new physics models which give large positive contributions to $S$ and $T$ are presumably excluded. Thus, the oblique correction parameters $S$, $T$, and $U$ are often used to judge whether a new model is compatible with experiments or not. If the SM is not a full theory, there will be new heavy particles above the electroweak scale. Provided the new particles feel the electroweak interactions, they should give corrections to $S$, $T$, and $U$ no matter whether they are fermions, scalars, or gauge bosons. The contributions of heavy fermions and scalars to the oblique correction parameters has been studied extensively in the literature [1, 3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14]. And in a previous work we have obtained the oblique corrections from fermions in higher-dimensional representations of the electroweak gauge group $SU(2)_L \times U(1)_Y$ [15].

In this paper, we obtain the oblique corrections from scalars in any irreducible representation of $SU(2)_L \times U(1)_Y$. We begin by computing the contributions to $S$, $T$, and $U$ from a heavy complex scalar doublet which has no vacuum expectation value (VEV). The $T$ parameter from this scalar doublet is exactly the same as the left-handed fermion doublet contribution, which has been shown first in Ref.[7]. However, there is a slight difference between the $S$ parameter we have obtained and that in Ref.[7]. This difference may come from the convention inconsistency of Ref.[7]. We then proceed to obtain the contributions to $S$, $T$, and $U$ from higher-dimensional scalar representations. Throughout this paper, we assume that the heavy scalar multiplet we consider has never acquired a VEV, and there is a slight mass non-degeneracy in the scalar multiplet which may be caused by unknown ultraviolet completions.

This paper is organized as follows. In Sec. II we compute the contributions to $S$, $T$, and $U$ from a heavy complex scalar doublet. In Sec. III we obtain the explicit expressions of the oblique corrections from a $(2j + 1)$-dimensional scalar multiplet. Conclusions and discussions are given in Sec. IV. In Appendices A and B some basic definitions and relations are listed, and general scalar vacuum polarization amplitudes that help to calculate vector-boson self energies are given.

II. A HEAVY SCALAR DOUBLET

To begin with, let us consider a heavy complex scalar doublet

$$\Phi = \left( \phi_1 \phi_2 \right)$$

which belongs to the representation $(2, \ Y)$ of $SU(2)_L \times U(1)_Y$. The heavy scalar fields $\phi_1$ and $\phi_2$ can acquire their masses from either bare mass terms or some other fields' VEVs much larger than the electroweak scale. In general, the doublet $\Phi$ is non-degenerate, and the masses of $\phi_1$ and $\phi_2$ are denoted by, respectively, $m_1$ and $m_2$. These scalars
Using these Feynman rules, the contributions of the scalar doublet to the vertices
\[ \equiv \mu_{\nu} q_{\mu} q_{\nu} + 2 e e = \frac{i e}{\sqrt{2}} \left( \mu_{\nu} q_{\mu} q_{\nu} + 2 e e \right), \]
\[ \equiv \mu_{\nu} q_{\mu} q_{\nu} + 2 e = \frac{i e}{\sqrt{2}} \left( \mu_{\nu} q_{\mu} q_{\nu} + 2 e e \right), \]
\[ \equiv \mu_{\nu} q_{\mu} q_{\nu} + 2 e = \frac{i e}{\sqrt{2}} \left( \mu_{\nu} q_{\mu} q_{\nu} + 2 e e \right), \]

where \( c \equiv \cos \theta_W, \ s \equiv \sin \theta_W, \ W^{\pm} = \frac{1}{\sqrt{2}} (W^{1} \pm i W^{2}), \ T^{\pm} = T^{1} \pm i T^{2}, \ Q = T^{3} + Y, \) and \( T^{a} = \tau^{a} / 2 \) with \( \tau^{a} \) (\( a = 1, 2, 3 \)) being Pauli matrices. In order to calculate the oblique corrections from this heavy scalar doublet, it is sufficient to calculate its contributions to the self energies of \( W \) and \( Z \) bosons. The relevant Feynman rules can be easily read off from Eq. (1) as follows:

- **\( V\phi\phi \) vertices**

\[ W^{+}_{\mu} (-p-q) \phi_{2}(p) \phi_{1}^{\dagger}(q) : \frac{i e}{\sqrt{2} s} (p-q)^{\mu} \]
\[ W^{-}_{\mu} (-p-q) \phi_{1}(p) \phi_{2}^{\dagger}(q) : \frac{i e}{\sqrt{2} s} (p-q)^{\mu} \]
\[ Z_{\mu} (-p-q) \phi_{1}(p) \phi_{1}^{\dagger}(q) : \frac{i e}{s c} \left[ 1 - (1 + 2 Y) s^{2} \right] (p-q)^{\mu} \]
\[ Z_{\mu} (-p-q) \phi_{2}(p) \phi_{2}^{\dagger}(q) : \frac{i e}{s c} \left[ -1 + (1 - 2 Y) s^{2} \right] (p-q)^{\mu} \]

with all momenta incoming.

- **\( VV\phi\phi \) vertices**

\[ W^{+}_{\mu} W^{-}_{\nu} \phi_{1}^{\dagger} \phi_{1}^{\dagger} : \frac{i e^{2}}{2 s^{2}} g^{\mu \nu} \]
\[ W^{+}_{\mu} W^{-}_{\nu} \phi_{2} \phi_{2}^{\dagger} : \frac{i e^{2}}{2 s^{2}} g^{\mu \nu} \]
\[ Z_{\mu} Z_{\nu} \phi_{1}^{\dagger} \phi_{1}^{\dagger} : \frac{i e^{2}}{s^{2} c^{2}} \left[ 1 - (1 + 2 Y) s^{2} \right]^{2} g^{\mu \nu} \]
\[ Z_{\mu} Z_{\nu} \phi_{2} \phi_{2}^{\dagger} : \frac{i e^{2}}{s^{2} c^{2}} \left[ 1 - (1 - 2 Y) s^{2} \right]^{2} g^{\mu \nu} \]

Using these Feynman rules, the contributions of the scalar doublet to the \( W \)-boson self energy can be written in terms of the building blocks \( \imath \Pi^{\mu \nu}_{1}(q^{2}, m^{2}) \) and \( \imath \Pi^{\mu \nu}_{2}(q^{2}, m_{1}^{2}, m_{2}^{2}) \) defined in Appendix B as follows

\[ \imath \Pi^{\mu \nu}_{WW}(q^{2}) = \frac{e^{2}}{2 s^{2}} \imath \Pi^{\mu \nu}_{1}(q^{2}, m^{2}) + \frac{e^{2}}{2 s^{2}} \imath \Pi^{\mu \nu}_{1}(q^{2}, m_{1}^{2}) + \left( \frac{e}{\sqrt{2} s} \right)^{2} \imath \Pi^{\mu \nu}_{2}(q^{2}, m_{1}^{2}, m_{2}^{2}) \]

\[ = \frac{e^{2}}{s^{2} c^{2}} g^{\mu \nu} \left[ m_{1}^{2} \log \frac{m_{2}^{2}}{\mu^{2}} + m_{2}^{2} \log \frac{m_{1}^{2}}{\mu^{2}} - 2 f_{2}(m_{1}^{2}, m_{2}^{2}) + q^{2} \left[ \frac{1}{3} Y + 2 f_{1}(m_{1}^{2}, m_{2}^{2}) \right] \right] \]

\[ + (q^{4} \imath \Pi^{\mu \nu}_{W} terms) + \mathcal{O}(q^{4}) \]

\[ \imath \Pi^{\mu \nu}_{ZZ}(q^{2}) = \frac{e^{2}}{s^{2} c^{2}} \left[ \frac{1 - (1 + 2 Y) s^{2}}{2} \right]^{2} \imath \Pi^{\mu \nu}_{1}(q^{2}, m^{2}) + \frac{e^{2}}{s^{2} c^{2}} \left[ \frac{1 - (1 - 2 Y) s^{2}}{2} \right]^{2} \imath \Pi^{\mu \nu}_{1}(q^{2}, m_{2}^{2}) \]

\[ = \frac{e^{2}}{s^{2} c^{2}} \left( \frac{1}{4 \pi^{2}} \right) g^{\mu \nu} \left[ \frac{1}{2} \left( -2 Y + \log \frac{m_{2}^{2}}{\mu^{2}} + \log \frac{m_{1}^{2}}{\mu^{2}} \right) - 2 s^{2}[ -2 Y + (1 + 2 Y) \log \frac{m_{1}^{2}}{\mu^{2}} ] \right] \]

\[ + \mathcal{O}(q^{4}) \]
where μ, T, f_1 and f_2 have been defined in Appendix B. Now, combining Eqs. (2) and (3) with Eqs. (A1), (A2) and (A3) in Appendix A we have

\[
\Pi_{11}(0) = \frac{1}{2(4\pi)^2} \left[ m_1^2 \log \frac{m_2^2}{\mu^2} + m_2^2 \log \frac{m_2^2}{\mu^2} - 2 f_2(m_1^2, m_2^2) \right],
\]

\[
\Pi'_{11}(0) = \frac{1}{2(4\pi)^2} \left[ - \frac{1}{3} \Upsilon + 2 f_1(m_1^2, m_2^2) \right],
\]

\[
\Pi_{33}(0) = 0, \quad \Pi_{3Q}(0) = 0 ,
\]

\[
\Pi'_{33}(0) = \frac{1}{12(4\pi)^2} \left[ - 2 \Upsilon + \log \frac{m_1^2}{\mu^2} + \log \frac{m_2^2}{\mu^2} \right],
\]

\[
\Pi'_{3Q}(0) = \frac{1}{12(4\pi)^2} \left[ - 2 \Upsilon + + (1 + 2 \Upsilon) \log \frac{m_1^2}{\mu^2} + (1 - 2 \Upsilon) \log \frac{m_2^2}{\mu^2} \right].
\]

Thus, from Eqs. (4) and (A4), we obtain

\[
S = - \frac{Y}{6\pi} \log \frac{m_1^2}{m_2^2},
\]

\[
T = \frac{1}{8\pi s^2c^2m_Z^2} \left[ m_1^2 \log \frac{m_2^2}{\mu^2} + m_2^2 \log \frac{m_2^2}{\mu^2} - 2 f_2(m_1^2, m_2^2) \right],
\]

\[
U = \frac{1}{\pi} \left[ f_1(m_1^2, m_2^2) - \frac{1}{12} \left( \log \frac{m_1^2}{\mu^2} + \log \frac{m_2^2}{\mu^2} \right) \right],
\]

from which, the explicit expressions for T and U can be obtained by figuring out the functions f_1 and f_2 as follows:

\[
T = \frac{1}{16\pi s^2c^2m_Z^2} \left( m_1^2 + m_2^2 - \frac{2 m_1^2 m_2^2}{m_1^2 - m_2^2} \log \frac{m_1^2}{m_2^2} \right),
\]

\[
U = \frac{1}{12\pi} \left[ \frac{5 m_1^4 - 22 m_1^2 m_2^2 + 5 m_2^4}{3(m_1^2 - m_2^2)^2} + \frac{m_1^6 - 3 m_1^4 m_2^2 - 3 m_1^2 m_2^4 + m_2^6}{(m_1^2 - m_2^2)^3} \log \frac{m_1^2}{m_2^2} \right].
\]

Comparing them with the results of Ref. [1], it is interesting to note that the contribution to the T parameter from the heavy scalar doublet is exactly the same as that from one heavy SM-like fermion doublet. And the contribution of the scalar doublet to the U parameter differs only by a factor of 1/2 from that of one SM-like fermion doublet. Moreover, except for an additional positive piece 1/(6\pi) in the SM-like fermion case, the difference of the scalar and SM-like fermion doublets' contributions to S is also a factor of 1/2.\(^1\)

In Ref. [8], the contributions to S and T from a heavy scalar doublet with hypercharge Y = −1/2 are given. The expression for the T parameter in our paper, i.e. Eq. (5), coincides with that in Ref. [8]. However, for the case of Y = −1/2 in Eq. (6), we have

\[
S = \frac{1}{12\pi} \log \frac{m_1^2}{m_2^2},
\]

which differs by a factor of 2 from the result quoted in Eq. (18) of Ref. [8]. We reexamine Ref. [8] and find that this difference stems from the inconsistency of the convention of Ref. [8]. Obviously, Eq. (5) of Ref. [8] describes one of the self energy diagrams for the charged W\(^\pm\) while Eq. (6) of that paper corresponds to W\(^a\) (a = 1, 2, 3). If one adds up these two diagrams, then the convention is inconsistent. Were Eq. (18) of Ref. [8] imprecise, some subsequent literature directly quoting the result should be modified as well. For instance, Eq. (8) of Ref. [8] has also an over-multiplying factor of 1/2.

\(^1\) Note that in this paper the hypercharge Y is defined by Q = T_3 + Y, while in Ref. [8] Q = T_3 + Y/2. Thus, if taken this paper’s notation, the Y appearing in the original expressions in the Ref. [8] must be replaced by 2Y.
III. HIGHER-DIMENSIONAL REPRESENTATIONS

In this section we proceed to give the contributions to $S$, $T$ and $U$ from scalars in more general representations. Consider a heavy complex scalar multiplet $\Phi$ with quantum numbers of $SU(2)_L \times U(1)_Y$ as follows

$$\Phi = \left( \begin{array}{c} \phi_j \\ \phi_{j-1} \\ \vdots \\ \phi_{-j} \end{array} \right) \sim (2j + 1, \ Y) .$$  \hfill (11)

The mass of $\phi_l$ is denoted by $m_l$ for $l$ running from $-j$ to $j$. The interactions between these scalars and the electroweak gauge bosons is of the same form as Eq.

$$(T^\pm)_{kl} = \sqrt{(j \mp l)(j \pm l + 1)} \delta_{k \pm 1,l} ,$$

$$(T^3)_{kl} = l \delta_{kl} , \quad Q_{kl} = (Y + l) \delta_{kl} , \quad -j \leq k, l \leq j .$$  \hfill (12)

Likewise, we compute their contributions to $S$, $T$, and $U$ resulting in

$$S = -Y \frac{3}{2 \pi} \sum_{l=-j}^{j} l \log \frac{m_l^2}{\mu^2} ,$$  \hfill (13)

$$T = \frac{1}{4 \pi S^2 \epsilon^2 m_Z^2} \left[ \sum_{l=-j}^{j} (j^2 + j - l^2) m_l^2 \log \frac{m_l^2}{\mu^2} - \sum_{l=-j}^{j-1} (j - l)(j + l + 1) f_2(m_l^2, m_{l+1}^2) \right] ,$$  \hfill (14)

$$U = \frac{1}{\pi} \left[ \sum_{l=-j}^{j-1} (j - l)(j + l + 1) f_1(m_l^2, m_{l+1}^2) - \sum_{l=-j}^{j} \frac{t^2}{3} \log \frac{m_l^2}{\mu^2} \right] .$$  \hfill (15)

It is easily checked that, by taking $j = 1/2$, Eqs.\(14\), \(15\) and \(16\) agree with Eqs.\(5\), \(6\) and \(7\). Comparing the above results with those of Ref.\(13\), we can see that the equality of the contributions to the $T$ parameter from a scalar $(2j + 1)$-plet and a SM-like fermion $(2j + 1)$-plet is just a coincidence for $j = 1/2$. It is interesting to find that the contribution of a scalar $(2j + 1)$-plet to the $S$ parameter is always one half of that of a vectorlike fermion $(2j + 1)$-plet, and the contribution of a scalar $(2j + 1)$-plet to the $U$ parameter is always one half of that of a SM-like fermion $(2j + 1)$-plet. Note that in the degenerate mass limit these expressions for $S$, $T$ and $U$ all vanish, and thus the current precision electroweak measurements cannot rule out a heavy degenerate scalar multiplet. Moreover, the small experimental values of $S$, $T$ and $U$ could impose rather stringent constraints on the mass non-degeneracy of the scalar multiplet.

IV. CONCLUSIONS

In this paper, we have obtained the oblique corrections from heavy scalars in any irreducible representations of the electroweak gauge group $SU(2)_L \times U(1)_Y$. Our expression for the $S$ parameter in the case of a heavy scalar doublet with $Y = -1/2$ is slightly different from that in Ref.\(1\). We have pointed out that the convention inconsistency of that paper most likely causes the disagreement. We have used dimensional regularization to perform scalar loop calculations with an implicit assumption that possible quadratic divergent terms are exactly canceled in the expressions of $S$, $T$ and $U$, which must be true for the purpose of the oblique correction formalism. The scalar multiplet we considered above has no VEV and no mixing with the SM Higgs doublet. We have shown that such a heavy degenerate scalar multiplet is not excluded by the current $S$-$T$-$U$ fits.

APPENDIX A: BASIC DEFINITIONS AND RELATIONS

In this appendix we briefly recapitulate some basic definitions and relations given in Ref.\(1\). The vacuum-polarization amplitudes are defined as follow:

$$i \Pi_{\mu \nu}^{XY}(q^2) = ig^{\mu \nu} \Pi_{XY}(q^2) + (q^\mu q^\nu \text{terms}) = \int d^4 xe^{-iqx} \left( J_K^\mu(x) J_Y^\nu(0) \right) ,$$  \hfill (A1)
where \((XY) = (11), (22), (33), (3Q), \) and \((QQ)\). And \(\Pi_{XY}(q^2)\) is defined by
\[
\Pi_{XY}(q^2) \equiv \Pi_{XY}(0) + q^2\Pi'_{XY}(q^2) .
\] (A2)

The relations between the one-particle irreducible (1PI) self-energies of the gauge bosons and the vacuum-polarization amplitudes are given by
\[
\begin{align*}
\Pi_{AA} &= e^2\Pi_{QQ}, \\
\Pi_{ZA} &= \frac{e^2}{s}\left(\Pi_{3Q} - s^2\Pi_{QQ}\right), \\
\Pi_{ZZ} &= \frac{e^2}{s^2c^2}\left(\Pi_{33} - 2s^2\Pi_{3Q} + s^4\Pi_{QQ}\right), \\
\Pi_{WW} &= \frac{e^2}{s^2}\Pi_{11} ,
\end{align*}
\] (A3)

where \(c \equiv \cos\theta_W\), \(s \equiv \sin\theta_W\), and \(e\) is the coupling constant of the electromagnetic interaction. The three oblique correction parameters \(S, T\) and \(U\) are defined respectively by
\[
\begin{align*}
\alpha S &\equiv 4e^2[\Pi'_{33}(0) - \Pi_{33}(0)] , \\
\alpha T &\equiv \frac{e^2}{s^2c^2m_Z^2}[\Pi_{11}(0) - \Pi_{33}(0)] , \\
\alpha U &\equiv 4e^2[\Pi_{11}(0) - \Pi_{33}(0)] ,
\end{align*}
\] (A4)

where \(\alpha \equiv e^2/(4\pi)\) is the fine-structure constant.

**APPENDIX B: GENERIC COMPLEX SCALAR VACUUM POLARIZATION AMPLITUDES**

\[
i\Pi_{\Omega}^{\mu \nu}(q^2, m^2) \equiv (\mu) \phi(q) \phi(\nu) \Phi(q)
\]

\[
i\Pi_{\Phi}^{\mu \nu}(q^2, m_1^2, m_2^2) \equiv (\mu) \phi(q) \phi(\nu) \Phi(q)
\]

**FIG. 1:** Two types of generic complex scalar vacuum polarization amplitudes, \(i\Pi_{\Omega}^{\mu \nu}(q^2, m^2)\) and \(i\Pi_{\Phi}^{\mu \nu}(q^2, m_1^2, m_2^2)\), which are related to vector-boson self energies. The subscripts \(\Omega\) and \(\Phi\) are chosen due to the shape of the respective Feynman diagram.

In this appendix we present some notations to simplify the calculations of vector-boson self energies. Rather than computing the specific complex scalar vacuum polarization amplitudes one by one, it is convenient to compute, once and for all, the most general ones which may give contributions to vector-boson self energies. There are two types of generic complex scalar vacuum polarization graphs as shown in Fig 1 with the Feynman rules defined in Fig 2. Then, performing dimensional regularization gives

\[
i\Pi_{\Omega}^{\mu \nu}(q^2, m^2) = ig^{\mu \nu} \int \frac{d^d k}{(2\pi)^d} \frac{i}{k^2 - m^2} \rightarrow ig^{\mu \nu} \mu^{-d} \int \frac{d^d k}{(2\pi)^d} \frac{i}{k^2 - m^2} \]
\[
= -ig^{\mu \nu} \frac{m^2}{(4\pi)^2} (\Upsilon + 1 - \log \frac{m^2}{\mu^2}) ,
\]

\[
i\Pi_{\Phi}^{\mu \nu}(q^2, m_1^2, m_2^2) = \int \frac{d^d k}{(2\pi)^d} \frac{(2k + q)^\mu (2k + q)^\nu}{(k^2 - m_1^2)(k^2 - m_2^2)} \rightarrow \mu^{-d} \int \frac{d^d k}{(2\pi)^d} \frac{(2k + q)^\mu (2k + q)^\nu}{(k^2 - m_1^2)(k^2 - m_2^2)} \]
\[
= ig^{\mu \nu} \frac{1}{(4\pi)^2} \left[ (m_1^2 + m_2^2)(\Upsilon + 1) - 2f_2(m_1^2, m_2^2) + q^2 \left[ -\frac{1}{3} \Upsilon + 2f_1(m_1^2, m_2^2) \right] \right] + (q^4 q^6) \text{ terms} + \mathcal{O}(q^8) ,
\] (B2)
\[
\begin{align*}
\text{Scalar Propagator} & \quad = \frac{i}{k^2 - m^2} \\
\phi & \quad \rightarrow \quad \phi^\dagger
\end{align*}
\]

3-line Vertex
\[
\begin{array}{c}
\mu \\
\phi_1
\end{array} \quad = \quad i(p - q)^\mu
\]

4-line Vertex
\[
\begin{array}{c}
\mu \\
\nu
\end{array} \quad = \quad ig^{\mu\nu}
\]

FIG. 2: The defining Feynman Rules for the generic complex scalar vacuum polarization amplitudes \(i\Pi_\mu_\nu^{(\phi)}(q^2, m^2)\) and \(i\Pi_\mu_\nu^{(\Phi)}(q^2, m_1^2, m_2^2)\).

where \(\mu\) is an arbitrary mass scale parameter and the infinity \(\Upsilon \equiv 2/(4 - d) - \gamma + \log(4\pi)\), and where the functions \(f_1(m_1^2, m_2^2)\) and \(f_2(m_1^2, m_2^2)\) are defined respectively by

\[
\begin{align*}
f_1(m_1^2, m_2^2) & \equiv \int_0^1 dx \, x(1 - x) \log \left( \frac{xm_1^2 + (1 - x)m_2^2}{\mu^2} \right), \\
f_2(m_1^2, m_2^2) & \equiv \int_0^1 dx \, xm_1^2 + (1 - x)m_2^2) \log \left( \frac{xm_1^2 + (1 - x)m_2^2}{\mu^2} \right).
\end{align*}
\]

In particular, if \(m_1 = m_2 = m\) in Eq. (B2), we have

\[
i \Pi_\mu_\nu^{(\Phi)}(q^2, m^2, m^2) = \frac{i g^{\mu\nu}}{(4\pi)^2} \left[ 2m^2(\Upsilon + 1 - \log \frac{m^2}{\mu^2}) + \frac{q^2}{3}(-\Upsilon + \log \frac{m^2}{\mu^2}) \right] + O(q^4).
\]

Eqs. (B1), (B2) and (B5) are very practical in calculating the contributions of scalar loops to vector boson self energies.

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