Topics in Electroweak Physics

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We briefly discuss five topics in Precision Electroweak Physics: i) the recently proposed Effective Scheme of Renormalization, ii) evidence for electroweak bosonic corrections derived from the radiative correction \( \Delta r_{\text{eff}} \), iii) an approach to estimate the scale of new physics in a hypothetical Higgs-less scenario, iv) simple and accurate formulae for \( s_{\text{eff}}^2, M_W, \Gamma_2 \), and their physical applications, v) a recent proposal concerning the field renormalization constant for unstable particles.

1. Effective Scheme of Renormalization

Precise calculations in the Standard Model (SM) are based on a number of renormalization frameworks. Two of the most frequently employed are: 1) the On-Shell Scheme (OS) [1–3] 2) the \( \overline{\text{MS}} \) approach [4]. The OS scheme is “very physical” in the sense that the renormalized parameters are identified with physical, scale-independent observables, such as \( \alpha, G_F, M_Z, M_W, \ldots \). The \( \overline{\text{MS}} \) approach is frequently applied in a hybrid version, with couplings defined by \( \overline{\text{MS}} \) subtractions, but retaining physical masses. It employs scale-dependent parameters such as \( s^2 = \sin^2 \theta(\mu) \), \( c^2(\mu) \) (usually evaluated at \( \mu = M_Z \)) and exhibits very good convergence properties [5]. It plays an important role in the analysis of Grand Unified Theories. However, it leads to a residual scale dependence in finite orders of perturbation theory (PT). Very recently, a novel approach, called the Effective Scheme of Renormalization (EFF), was proposed [6,7]. It shares the good convergence properties of the \( \overline{\text{MS}} \) scheme, but it eliminates the residual scale dependence in finite orders of PT. A distinguishing feature is that the basic electroweak mixing parameter (EWMP) is directly identified with \( s_{\text{eff}}^2 \equiv \sin^2 \theta_{\text{eff}} \), employed by the Electroweak Working Group (EWWG) to describe the on-resonance asymmetries measured at LEP and SLC. It may be evaluated by means of the basic relation [6–8]

\[
s_{\text{eff}}^2 c_{\text{eff}}^2 = \frac{\pi \alpha}{\sqrt{2} G_F M_Z^2 (1 - \Delta r_{\text{eff}})} ,
\]

where \( \Delta r_{\text{eff}} \) is the relevant radiative correction. In order to calculate \( \Delta r_{\text{eff}} \) the following strategy was followed:

i) Since current calculations of \( s_{\text{eff}}^2 \) incorporate two-loop effects enhanced by powers \( (M_Z^2/M_W^2)^n \) (with \( n = 1, 2 \)), we first express \( \Delta r_{\text{eff}} \) in terms of corrections \( \Delta \hat{r}_W, \Delta \hat{\rho}, \Delta \hat{k} \) and \( \hat{f} \), for which the irreducible contributions of this order have been evaluated [9,10].

ii) To ensure the absence of a residual scale dependence, we use scale-independent couplings, such as \( c^2, s_{\text{eff}}^2, G_F, M_Z^2 \), retain only two-loop effects enhanced by factors \( (M_Z^2/M_W^2)^n \) (\( n = 1, 2 \)), and employ a simple definition of the EWMP, identified with \( s_{\text{eff}}^2 \). In particular, \( s^2 \) can be expressed in terms of \( s_{\text{eff}}^2 \) by means of the relation

\[
s_{\text{eff}}^2 = \left[ 1 + \frac{\hat{c}^2}{s_{\text{eff}}^2} \Delta \hat{k} (M_Z^2, \mu) \right] s^2(\mu) ,
\]

where \( \Delta \hat{k}(q^2, \mu) \) is an electroweak form factor [11]. The analysis leads to the expression [6]:

\[
\Delta r_{\text{eff}} = \Delta \hat{r}_W - \frac{\hat{c}^2}{s_{\text{eff}}^2} \left[ \Delta \hat{\rho} - \Delta \hat{k} \left( 1 - \frac{s_{\text{eff}}^2}{c_{\text{eff}}^2} \right) \right] + \frac{\hat{c}^2}{s_{\text{eff}}^2} \left[ 2 \Delta \hat{\rho} - (\Delta \hat{\rho})_{\text{lead}} - \hat{f} + \Delta \hat{k} \frac{s_{\text{eff}}^2}{c_{\text{eff}}^2} \right] ,
\]
where $\Delta \hat{\rho} \equiv \text{Re} [A_{WW}(M_W^2) - c^2 A_{ZZ}(M_Z^2)] / M_W^2$,

\[ x_t = 3G_F M_t^2 / (8\sqrt{2}\pi^2), \quad \Delta \hat{r}_W = -2\delta e / e + (c^2/s^2) f, \quad (\Delta \hat{\rho})_{\text{lead}} = (3/64\pi^2) M_t^2 / M_W^2, \]

\[ \hat{f} \equiv (ReA_{WW}(M_W^2) - A_{WW}(0)) / M_W^2 + V_W + M_W^2 B_W, \quad A_{WW} \text{ and } A_{ZZ} \text{ are the } W \text{ and } Z \text{ self-energies modulo a factor } c^2/s^2, \quad V_W \text{ and } B_W \text{ are vertex and box corrections contributing to } \mu \text{-decay, and } \delta e \text{ stands for the charge renormalization counterterm.} \]

It is understood that, in Eq. (3), $s^2$ is replaced everywhere by $s^2_{\text{eff}}$. The corrections $\Delta \hat{r}_W$, $\Delta \hat{k}$, and $\hat{f}$ depend also on $c^2 = M_W^2 / M_Z^2$. In order to obtain an expression that depends solely on $s^2_{\text{eff}} = 1 - s^2_{\text{eff}}$, $M_W^2$ is replaced by $c^2 M_Z^2$; in two-loop contributions, $c^2$ is replaced by $c^2_{\text{eff}}$, since the difference is of third order; in one-loop corrections, a Taylor expansion about $c^2 = c^2_{\text{eff}}$ is made, in conjunction with the one-loop expression for $c^2 - c^2_{\text{eff}}$. The corresponding expression for $M_W$ is given in Ref. [6]. An interesting feature is that the calculation of $s^2_{\text{eff}}$ is completely decoupled from that of $M_W$, while the $s^2_{\text{eff}}$ results are employed to calculate $M_W$. The results for the leptonic partial widths $\Gamma_l$ of the $Z$ have been recently obtained [12]. A detailed comparison shows that, for $M_H = 100 \text{ GeV}$, the difference $|s^2_{\text{eff}}(\overline{MS}) - s^2_{\text{eff}}(\text{EFF})|$ of the $\overline{MS}$ and EFF calculation of $s^2_{\text{eff}}$ is $\lesssim 10^{-5}$ over the range $30 \text{ GeV} \leq \mu \leq 200 \text{ GeV}$ and exhibits a maximum at $\mu \approx 70 \text{ GeV}$. In the $M_W$ case, one finds $|M_W(\overline{MS}) - M_W(\text{EFF})| \leq 1 \text{ MeV}$ over the range $50 \text{ GeV} \leq \mu \leq 205 \text{ GeV}$ and a maximum at $\mu \approx 100 \text{ GeV}$. At $\mu = 300 \text{ GeV}$, the differences amount to $\approx 3 \times 10^{-5}$ and 3 MeV, respectively. These findings give support to the choice $\mu = M_Z$ in the $\overline{MS}$ calculations of observables in the resonance region. It should be pointed out, however, that this satisfactory state of affairs holds when the corrections of $\mathcal{O}(\alpha^2(M_t/M_W)^2)$ are included. If these contributions are excluded, $M_W(\overline{MS})$ is a monotonically decreasing function of $\mu$ over the range $30 \text{ GeV} \leq \mu \leq 500 \text{ GeV}$ [13], and the choice of scale is very ambiguous. In summary, the EFF approach has the virtue of eliminating the scale ambiguity which, in some cases, may create a significant theoretical uncertainty.

2. Evidence for Electroweak Bosonic Corrections

It turns out that $\Delta r_{\text{eff}}$ is very sensitive to electroweak bosonic contributions (EWBC), i.e. corrections involving virtual bosons: $W$, $Z$, $H$, $\phi's$. They are subleading numerically, but very important conceptually! One way to obtain sharp evidence for these corrections is to measure $\Delta r_{\text{eff}}$. Using the current experimental value $(s^2_{\text{eff}})_{\exp} = 0.23149 \pm 0.00017$ and Eq.(1), we find $(\Delta r_{\exp})_{\text{eff}} = 0.06047 \pm 0.00048$. On the other hand, subtracting the EWBC, the theoretical evaluation leads to: $(\Delta r_{\text{eff}})_{\text{subtr}} = 0.05106 \pm 0.00083$. The difference is $0.00941 \pm 0.00096$, thus providing evidence for the presence of EWBC at the 9.8 $\sigma$ level [5,8,14]!

3. The Higgs-less Scenario

The corrections $\Delta r_{\text{eff}}$ and $\Delta r$ have been also employed to discuss the scale of new physics in a hypothetical scenario in which the Higgs boson is absent [15]. At the one-loop level, the Higgs boson contribution to $\Delta r_{\text{eff}}$ is a complicated function of $\xi = M_H^2 / M_Z^2$, given in Ref. [15]. It may be written in the form

\[ (\Delta r_{\text{eff}})_{H} = \frac{\alpha}{4\pi} \frac{3}{2} (\frac{5}{3} - \frac{3}{2} c^2) \left( \frac{1}{n - 4} + C + \ln \frac{M_Z}{\mu} \right), \]

where the first term is the divergent part and the second one is the $\overline{MS}$-renormalized contribution evaluated at $\mu = M_Z$ ($C = [\gamma - \ln 4\pi]/2$, $\mu = 't$ Hooft scale). Subtracting $(\Delta r_{\text{eff}})_{H}$ from $\Delta r_{\text{eff}}$ we have

\[ \Delta r_{\text{eff}} - (\Delta r_{\text{eff}})_{H} = \Delta r_{\text{eff}} - (\Delta r_{\text{eff}})_{\overline{MS}} \]

\[ - \frac{\alpha}{4\pi} \frac{3}{2} (\frac{5}{3} - \frac{3}{2} c^2) \left( \frac{1}{n - 4} + C + \ln \frac{M_Z}{\mu} \right) \]

Clearly, Eq.(5) is divergent and scale dependent. We now conjecture that contributions from unknown new physics (NP) cancel the divergence and scale dependence of Eq.(5). Thus, the NP contribution to $\Delta r_{\text{eff}}$ must be of the form:

\[ X = \frac{\alpha}{4\pi} \frac{3}{2} \left( \frac{5}{3} - \frac{3}{2} c^2 \right) \left( \frac{1}{n - 4} + C + \ln \frac{M_Z}{\mu} \right), \]
We note that in the $\overline{MS}$ renormalization approach, the term proportional to $\ln M/\mu$ represents the NP contribution to $\Delta r_{\text{eff}}$ at scale $\mu$. If the NP is characterized by a scale $\Lambda$, we may decompose

$$\ln \frac{M}{\mu} = \ln \frac{\Lambda}{\mu} + K,$$

where the term involving $K \equiv \ln \frac{M}{\mu}$ represents the NP contribution to $\Delta r_{\text{eff}}$ at scale $\Lambda$. Adding $X$ to $\Delta r_{\text{eff}} - (\Delta r_{\text{eff}})_{\text{H}}$ we find the expression for $\Delta r_{\text{eff}}$ in the new scenario (NS) in which the Higgs boson contribution has been replaced by new physics:

$$\Delta r_{\text{eff}} \left( \Delta \right)_{\text{NS}} = \left( \Delta r_{\text{eff}} \right)_{\text{H}} + \frac{\alpha}{4\pi} \frac{\Lambda^5}{c^2} \frac{3}{2} \ln \frac{M}{M_Z}.$$  

(8)

The last term represents the NP contribution to $\Delta r_{\text{eff}}$ at scale $M_Z$. Calculating $\Delta r_{\text{eff}} - (\Delta r_{\text{eff}})_{\text{H}}$ and equating $(\Delta r_{\text{eff}})_{\text{NS}} = (\Delta r_{\text{eff}})_{\text{exp}}$, we can determine $\ln \frac{M}{M_Z}$. Employing $\Delta \alpha_h^{(5)} = 0.02761 \pm 0.00036$ and the other experimental inputs, one finds

$$\ln \frac{M}{M_Z} = 0.307 \pm 0.485,$$

(9)

which corresponds to a central value $M_c = 124$ GeV and a 95% CL upper bound $M^\text{95} = 275$ GeV. If the model-dependent constant $K$ is positive, we see from Eq.(7) that $\Lambda$ is sharply bounded: $\Lambda \leq 275$ GeV @ 95% CL. Instead, if $K < 0$, $\Lambda$ is not bounded by these considerations. Thus, we can group the NP models into two classes, according to the sign of $K$. Furthermore, if for instance $\Lambda = 1$ TeV, we have

$$\ln \frac{M}{M_Z} = 2.395$$

and we find from Eqs.(7,9) that $K = -2.088 \pm 0.435$. Thus, for such $\Lambda$ values, a substantial cancellation of logarithmic and constant terms is required [15,16]. Similar results are obtained from the corresponding analysis of $\Delta r$ [15].

4. Simple formulae for $s_{\text{eff}}$, $M_W$, $\Gamma_l$

Simple formulae that reproduce accurately the numerical results of the codes in the range $20$ GeV $\leq M_H \leq 300$ GeV, probed by recent experiments, have been presented [12]. They are of the form:

$$s_{\text{eff}}^2 = (s_{\text{eff}}^0) + c_1 A_1 + c_2 A_2$$

and

$$M_W = M_W^0 - d_1 A_1 - d_2 A_2$$

where $\Gamma_l$ is given by

$$\Gamma_l = \Gamma_l^0 - g_1 A_1 - g_2 A_2$$

and

$$\Gamma_l = \frac{\alpha}{4\pi} \frac{\Lambda^5}{c^2} \frac{3}{2} \ln \frac{M}{M_Z}.$$  

(10)

where $A_1, A_2$ are the model-dependent constant terms. Similar results are obtained from recent calculations. Using Eq. (10) in the EFF scheme and $(s_{\text{eff}}^2)_{\text{exp}} = 0.23419 \pm 0.00017$, one finds $M_H = 124.82 \pm 0.52$ GeV and a 95% CL upper bound $M^\text{95} = 280$ GeV. Instead, Eq. (11) and $(M_W)_{\text{exp}} = 80.451 \pm 0.033$ GeV lead to $M_H = 23 \pm 23$ GeV, $M^\text{95} = 122$ GeV. Thus, $M_W$ constrains $M_H$ much more sharply than $s_{\text{eff}}$. It is important to note that the $M^\text{95}$ value derived from $M_W$, and the direct exclusion bound $M_H > 140$ GeV @ 95% CL, suggest a very narrow window for $M_H$! One may also extract $A_1$ from $(s_{\text{eff}}^2)_{\text{exp}}$ and Eq. (10), to predict $M_W$ via Eq. (11): $(M_W)_{\text{indir}} = 80.374 \pm 0.025$ GeV, which is close to the corresponding value $(M_W)_{\text{indir}} = 80.379 \pm 0.023$ GeV obtained in the global analysis [17], and differs from $(M_W)_{\text{exp}}$ by 1.86 $\sigma$. Finally, we may use simultaneously Eq. (10-12) in conjunction with $(s_{\text{eff}}^2)_{\text{exp}}, (M_W)_{\text{exp}}$, and $(\Gamma_l)_{\text{exp}}$ to obtain $M_H = 97^{+66}_{-41}$ GeV, $M^\text{95} = 223$ GeV, to be compared with $M_H = 85^{+54}_{-34}$ GeV, $M^\text{95} = 196$ GeV in the recent EWWG fit.

The current determination of $(s_{\text{eff}}^2)_{\text{exp}}$ has $\chi^2$/d.o.f. $= 10.6/5$, corresponding to a CL of only 6%, and shows an intriguing dichotomy: the leptonic observables ($A_l(SLD), A_l(P_-), A_l^{(0,1)})$ lead to $(s_{\text{eff}}^2)_l = 0.23113 \pm 0.00021$, while the value from the hadronic ones ($A_l^{(0,1)} , A_l^{(0,1)} , < Q^2 >$) is $(s_{\text{eff}}^2)_h = 0.23220 \pm 0.00029$. Thus, there
is a $3\sigma$ difference between the two determinations! Furthermore, from $(s_{eff}^2)_{t}$ one finds $M_H = 59^{+50}_{-29}$ GeV, $M_H^{0} = 158$ GeV, closer to the result from $(M_W)_{\text{exp}}$. If $(s_{eff}^2)_{t} - (s_{eff}^2)_{h}$ reflects a statistical fluctuation, one possibility is to enlarge the error by $[\chi^2/d.o.f.]^{1/2}$ (PDG prescription), leading to $s_{eff}^2 = 0.23149 \pm 0.00025$. Interestingly, increasing the error in $s_{eff}^2$ leads to smaller $M_H^{0}$ in the combined $s_{eff}^2$-$M_W$-$\Gamma_t$ analysis: 223 GeV $\rightarrow$ 201 GeV! The reason is that this procedure gives enhanced weight to the $M_W$ input, which prefers a smaller $M_H$. If $(s_{eff}^2)_{t} - (s_{eff}^2)_{h}$ is due to new physics involving the $(t, b)$ generation, a substantial, tree-level change in the $Zb_R\bar{b}_R$ coupling is required [18]. If the discrepancy were to settle on the leptonic side, a scenario with light $\tilde{e}_L$ would improve the agreement with the electroweak data and the direct lower bound on $M_H$ [19].

It has been pointed out by several people that, if the central values of $M_t$ and $M_W$ remain as they are now, but the errors shrink sharply as expected at Tevatron/LHC or even much better at LC + GigaZ, a discrepancy would be established with the SM, that can be accommodated in the MSSM!

The comparison of the calculations of $s_{eff}^2$, $M_W$, and $\Gamma_t$ in the EFF, $\overline{MS}$, and OS frameworks has been applied to study the scheme dependence and to estimate the theoretical error arising from the truncation of the perturbative series [12]. Including QCD uncertainties, the theoretical errors have been estimated to be $\delta s_{eff}^2 \approx 6 \times 10^{-5}$ and $\delta M_W \approx 7$ MeV.

5. Field Renormalization Constant for Unstable Particles

In Ref. [20], it was proposed that the first problem can be solved by considering the complex valued position $\bar{\sigma}$ of the propagator’s pole, which is gauge invariant. We have: $\bar{\sigma} = M_0^2 + A(\bar{\sigma})$, where $M_0$ is the bare mass and $A(s)$ the self-energy. Decomposing $\bar{\sigma} = m_2^2 - i m_2 \Gamma_2$, where $m_2$ and $\Gamma_2$ are real, one identifies $m_2$ and $\Gamma_2$ with the mass and width of the particle:

$$m_2^2 = M_0^2 + \text{Re} A(\bar{\sigma}),$$  

$$m_2 \Gamma_2 = -\text{Im} A(\bar{\sigma}).$$

In Ref. [21], it was proposed that the second problem can be solved by defining the field renormalization constant $\bar{Z}$ by means of the normalization condition

$$m_2 \Gamma_2 = -\bar{Z} \text{Im} A(m_2^2),$$

which, in conjunction with Eq. (14), leads to

$$\bar{Z}^{-1} = 1 + \frac{\text{Im} (A(\bar{\sigma}) - A(m_2^2))}{m_2 \Gamma_2}.$$  

In the narrow width approximation, the r.h.s. of Eq. (16) becomes $1 - \text{Re} A'(m_2^2) \approx 1 - \text{Re} A'(M^2)$, where $M^2$ is the on-shell mass, and $\bar{Z}$ reduces to the conventional expression. We note that: i) $\bar{Z}$, defined by Eq. (16), involves a finite difference, rather than a derivative, thus avoiding the threshold problem; ii) using Eq. (16) we see that the r.h.s. of Eq. (15) is gauge invariant, since it equals $m_2 \Gamma_2$ as a mathematical identity. It was also shown that the use of Eq. (16) removes unphysical threshold singularities in the relation between on-shell and pole widths [22].}

This approach has been recently discussed in the framework of renormalization theory [23, 24]. Dividing the unrenormalized transverse propagator $-i Q_{\mu\nu} [s - M_0^2 - A(s)]^{-1}$ ($Q_{\mu\nu} = g_{\mu\nu} - q_\mu q_\nu/q^2$) by $\bar{Z}$, and introducing $S(s) \equiv \bar{Z} A(s)$, $\delta M^2 \equiv \text{Re} S(\bar{\sigma})$, $\bar{Z} \equiv 1 - \delta \bar{Z}$, we obtain the renormalized propagator:

$$D = -i Q_{\mu\nu} / (s - m_2^2 - S(\bar{\sigma})), $$  

where

$$S(\bar{\sigma}) = S(s) - \delta M^2 + \delta \bar{Z}(s - m_2^2)$$

stands for the renormalized self-energy. Since $\delta M^2$ and $\delta \bar{Z}$ are real, they should be chosen so
that \( \text{Re} \, S^{(r)}(s) \) is ultraviolet convergent to all orders. Once this is done, \( \text{Im} \, S^{(r)}(s) = \text{Im} \, S(s) = \hat{Z} \, \text{Im} \, A(s) \) must also be convergent, since there are no additional counterterms available. This means that \( \hat{Z} \) may be defined by imposing an appropriate normalization condition on \( \text{Im} \, S(s) \). A particularly simple one is

\[
\text{Im} S(m_2^2) = -m_2 \Gamma_2, \tag{19}
\]

which coincides with Eq. (15)! It was proposed independently in Ref. [21] to solve the threshold and gauge-dependence problems, and in Ref. [24] to implement a systematic order by order removal of the ultraviolet divergences in \( S^{(r)}(s) \). In Ref. [23] it was also emphasized that in this formulation one can derive closed and exact expressions for the mass and field-renormalization counterterms, to wit

\[
\delta M^2 = \text{Re} \, S(\bar{s}) ; \quad \delta \hat{Z} = \frac{\text{Im} \left[ S(\bar{s}) - S(m_2^2) \right]}{m_2 \Gamma_2}. \tag{20}
\]

In many cases, \( \Gamma_2 = \mathcal{O}(g^2) \), where \( g \) is a generic gauge coupling. If \( \delta M^2 \) and \( \delta \hat{Z} \) admit expansions in powers of \( \Gamma_2 \), they can be expressed as series involving \( R \equiv \text{Re} \, S(s) \), \( I \equiv \text{Im} \, S(s) \), and their powers and derivatives evaluated at \( s = m_2^2 \). These expansions of Eq. (20) coincide with the order by order analysis in Ref. [24]. However, in other important instances, such as the photonic corrections to the \( W \) self-energy, such expansions are ill-defined and lead to power-like infrared divergences! In such cases one should employ the exact formulae in Eq. (20), which lead to sensible expressions for \( \delta M^2 \), \( \delta \hat{Z} \), and the renormalized propagator [23].

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