Anomalous Higgs Yukawa couplings

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Abstract – In the standard model, the Higgs boson couples to the quarks and charged leptons according to the well-known formula \( \bar{\psi} h \psi \), where \( \psi \) is quark (q) or lepton (l) and \( v = 246 \text{GeV} \) is its vacuum expectation value. Suppose \( m_\phi \) is of radiative origin instead, then the effective \( h \bar{\psi} \psi \) Yukawa coupling will not be exactly \( m_\phi / v \). We show for the first time quantitatively how this may shift the observed branching fractions of \( h \to b \bar{b} \) and \( h \to \tau^+ \tau^- \) upward or downward. Thus, the precision measurements of the Higgs decay to fermions at the Large Hadron Collider, due to the resume operation in 2015, could be the key to possible new physics.

The 2012 discovery of the 125 GeV particle \([1,2]\) ushered in a new era of particle physics. The observed decay modes of this particle are consistent with it being the long-sought Higgs boson \( h \) \([3]\) of the standard model (SM) of particle interactions. In the near future, after the Large Hadron Collider (LHC) resume operation in 2015, more data will allow these determinations to be greatly improved. Of particular interest are the \( h \) branching fractions to fermions, such as \( b \bar{b} \) and \( \tau^+ \tau^- \). They are predicted by the SM to be proportional to \( 3m_t^2 \) and \( m_\tau^2 \), respectively, whereas current data are not definitive in this regard. In this paper, it is shown how these Yukawa couplings may be different from those of the SM if \( m_q \) or \( m_l \) is radiative in origin. Thus, the precision measurements of the Higgs decay to fermions could be the key to possible new physics.

The idea that quark and lepton masses may be radiative is of course not new. A review of such mechanisms already appeared 25 years ago \([4]\). In the context of the SM or its left-right extension, early work \([5–11]\) discussed how small radiative masses may be generated, but did not consider how the corresponding Higgs Yukawa couplings are affected. The implicit assumption is that the change is negligible. In supersymmetry, in the presence of soft breaking in the scalar sector, there are both tree-level and one-loop contributions to quark masses \([12–15]\). The resulting corrections to the Higgs Yukawa couplings have indeed been studied. Here we do something new. We consider radiative quark or lepton mass from dark matter in the context of the SM with only one Higgs boson. We show that the resulting Higgs Yukawa coupling may differ significantly from the SM prediction.

We follow the generic notion of a recent proposal which links radiative fermion mass with dark matter \([16,17]\). The first step is to forbid the usual Yukawa coupling \( \bar{\psi}_L \psi_R \phi^0 \) or \( \bar{\psi}_L \psi_R \phi^0 \) by some symmetry. The second step is to postulate new particles which allow this connection to be made in one loop with soft breaking of the assumed symmetry. A typical realization is shown below for \( m_\tau \). The new idea of this paper is the detailed analysis of the \( h \bar{\psi} \psi \) coupling which shows for the first time that it could be significantly different from the SM value of \( m_\phi / v \), where \( v = 246 \text{GeV} \). Note the important fact that this deviation comes entirely from a renormalizable theory. There are no hidden assumptions and all new particles and allowed renormalizable interactions are considered.

In fig. 1, \( \eta^+ \) is part of an electroweak doublet \( (\eta^+, \eta^0) \) and \( \chi^+ \) is a singlet. To realize this diagram in a simple specific model which also accommodates dark matter, consider the discrete symmetry \( Z_2 \times Z_2 \) under which \( N_L, \eta, \chi \) are (odd, even), \( N_L \) is (odd, odd), \( \tau_R \) is (even, odd), and all other fields are (even, even). As a result, the usual (hard) Yukawa term \( \bar{\tau}_L \tau_R \phi^0 \) is forbidden, but the (hard) Yukawa terms \( \bar{\eta}_L \eta_R \eta^+ \) and \( \bar{\chi}_R \tau_R \chi^+ \) are allowed, as well as the \( \mu (\eta^+ \phi^0 - \eta^0 \phi^+) \chi^- \) trilinear interaction which mixes \( \eta^\pm \) and \( \chi^\pm \). The first \( Z_2 \) is assumed to be exact, which accommodates dark matter \([18]\). The second \( Z_2 \) is assumed to be broken softly by the term \( m_N N_L N_R \). This allows the completion of the loop. The resulting radiative
$m_\tau$ is guaranteed to be finite and calculable as shown below.

The $2 \times 2$ mass-squared matrix spanning $(\eta^\pm, \chi^\pm)$ is given by

$$M_{\eta,\chi}^2 = \begin{pmatrix} m_\eta^2 & \mu v/\sqrt{2} \\ \mu v/\sqrt{2} & m_\chi^2 \end{pmatrix} = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} m_1^2 & 0 \\ 0 & m_2^2 \end{pmatrix} \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix}, \quad (1)$$

where $\langle \phi^0 \rangle = v/\sqrt{2}$ and $m_{\eta,\chi}^2$ already include the $\lambda_{\eta,\chi} v^2$ contributions from the quartic scalar terms $\lambda_{\eta}(\eta^2)(\Phi^\dagger\Phi)$ and $\lambda_{\chi}(\chi^2)(\Phi^\dagger\Phi)$, respectively. Now

$$\zeta_1 = \eta \cos \theta + \chi \sin \theta, \quad \zeta_2 = \chi \cos \theta - \eta \sin \theta \quad (2)$$

are the mass eigenstates, and the mixing angle $\theta$ is given by

$$\frac{\mu v}{\sqrt{2}} = \sin \theta \cos \theta (m_1^2 - m_2^2). \quad (3)$$

This imposes the constraint $\mu v/\sqrt{2}(m_1^2 - m_2^2) < 1/2$ on $\mu/(m_1^2 - m_2^2)$. Together with the requirement that $m_{1,2}^2 > 0$, this guarantees that there is no charge breaking minimum in the Higgs potential. The exact calculation of $m_\tau$ in terms of the exchange of $\tau$, $x$, $\tau R$, $N_\tau$, $N_R$ is given by

$$m_\tau = \frac{f_{\eta,\chi} \sin \theta \cos \theta}{16\pi^2} \sqrt{2} m_N \left[ \frac{m_1^2 - m_2^2}{m_1^2 m_2^2} \ln \frac{m_1^2}{m_2^2} - \frac{m_2^2}{m_1^2 m_2^2} \ln \frac{m_2^2}{m_1^2} \right] \quad (4)$$

where $x_{1,2} = m_{1,2}^2/m_N^2$ and

$$F(x, y) = \frac{1}{x - 1 - y} \ln \frac{x - 1}{x - 1 - y}$$

Previous calculations usually assume that $\mu v/\sqrt{2} \ll m_{\eta,\chi}^2$, so $\eta, \chi$ are kept as mass eigenstates and $\mu v/\sqrt{2}$ as a coupling or mass insertion. The presumed Higgs Yukawa coupling is then calculated using the same integral as that of the radiative mass, with their ratio unchanged from the SM prediction. This is an unjustified assumption because the correct comparison is $\mu v/\sqrt{2}$ against $m_1^2 - m_2^2$ and not $m_{1,2}^2$ as shown in eq. (3).

Let $\phi^0 = (v + h)/\sqrt{2}$ and consider the effective Yukawa coupling $h\phi^0 \bar{\tau} \tau$. In the SM, it is of course equal to $m_\tau/v$, but here it has three contributions, one from fig. 1 and two others which are new and have not been considered before. Assuming that $m_1^2$ is small compared to $m_{1,2}^2$ and $m_N^2$, fig. 1 yields the one-loop effective coupling $f_{\tau}^{(1)} h\phi^0 \bar{\tau} \tau$, where

$$f_{\tau}^{(1)} = \frac{f_{\eta,\chi} \mu}{16\sqrt{2} \pi^2 m_N} \left[ \cos^2 \theta + \sin^2 \theta F(x_1, x_2) \right]. \quad (5)$$

Comparing eq. (6) with eq. (4), we see that $f_{\tau}^{(1)} = m_\tau/v$ only in the limit $\sin 2\theta \to 0$. We see also that $F_+(x_1, x_2) + F_{-}(x_1, x_2)$ is always greater than $2F(x_1, x_2)$, so that $f_{\tau}^{(1)}$ is always greater than $m_\tau/v$. The correction due to nonzero $m_\tau$ is easily computed in the limit $m_1 = m_2 = m_N$, in which case it is $m_\tau^2/12m_N^2$. This shows that it should be generally negligible. Let

$$F_+(x_1, x_2) = \frac{F(x_1, x_1) + F(x_2, x_2)}{2F(x_1, x_2)} - 1, \quad (7)$$

$$F_-(x_1, x_2) = \frac{F(x_1, x_1) - F(x_2, x_2)}{2F(x_1, x_2)}$$

then $F_+ \geq 0$ and if $x_1 = x_2$, $F_+ = F_- = 0$. To get an idea of their behavior, we take for example $x_1 = x_2 + 2$ and plot $F$ and $F_\pm$ as functions of $x_2$ in fig. 2. Note that this choice allows all possible $\theta$ values as long as $\mu v/\sqrt{2} m_N^2 = \sin 2\theta$.

Two other one-loop contributions exist, i.e. $f_{\tau}^{(1,2)} h\phi^0 \bar{\tau} \tau$. They come from the quartic scalar couplings $(\lambda_{\eta}/2)(v + h)^2\eta^+\eta^-$ and $(\lambda_{\chi}/2)(v + h)^2\chi^+\chi^-$, respectively. Doing the corresponding integrals, we obtain

$$f_{\tau}^{(1)} = \frac{\lambda_{\eta} f_{\eta,\chi} \mu}{16\pi^2 m_N} \sin \theta \cos \theta \left[ \cos^2 \theta F(x_1, x_1) - \sin^2 \theta F(x_2, x_2) - \cos 2\theta F(x_1, x_2) \right], \quad (8)$$

![Fig. 2](image-url) (Color online) The functions $F, F_\pm$ plotted against $x_2$ for $x_1 = x_2 + 2$. 

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Fig. 3: (Color online) The ratio \((f_{r \nu} / m_{\tau})^2\) plotted against \(\theta\) for \(x_1 = 3\) and \(x_2 = 1\) with various \((r_{\eta}, r_{\chi})\).

\[
\begin{align*}
  f_{r \nu}^{(2)} &= \frac{\lambda_{\nu} v F_{r \nu}}{16 \pi^2 m_N} \sin \theta \cos \theta |F(x_1, x_1) - \cos \theta F(x_2, x_2) + \cos 2 \theta F(x_1, x_2)|, \\
  &\text{Combining all three contributions and using eq. (3), we have} \\
  \frac{f_{r \nu}}{m_{\tau}} &= \left[ \frac{f_{r \nu}^{(1)}}{m_{\tau}} + f_{r \nu}^{(2)} + f_{r \nu}^{(3)} \right]^2 \\
  &= 1 + \frac{1}{2} (\sin 2 \theta)^2 \left[ (F_+ + (x_1 - x_2) \cos 2 \theta (r_{\eta} - r_{\chi}) F_+ + (r_{\eta} + r_{\chi}) F_-) \right],
\end{align*}
\]

where \(r_{\eta, \chi} = \lambda_{\eta, \chi} (m_N / \mu)^2\). We plot \((f_{r \nu} / m_{\tau})^2\) or the analogous \((f_{r \nu} / m_b)^2\) as functions of \(\theta\) in fig. 3 for various \((r_{\eta}, r_{\chi})\) with \(x_1 = 3\) and \(x_2 = 1\). It shows that, in general, it is not equal to one as the SM predicts. The current LHC measurements of \(h \to \tau^+ \tau^-\) and \(h \to b \bar{b}\) provide the bounds

\[
\begin{align*}
  \left( \frac{f_{r \nu}}{m_{\tau}} \right)^2 &= 1.4 (0.5) (0.4), \\
  \left( \frac{f_{r \nu}}{m_b} \right)^2 &= 0.2 (0.7) (0.6) \text{ (ATLAS) [21]}, \\
  \left( \frac{f_{r \nu}}{m_{\tau}} \right)^2 &= 0.78 \pm 0.27, \\
  \left( \frac{f_{r \nu}}{m_b} \right)^2 &= 1.0 \pm 0.5 \text{ (CMS) [22]}. 
\end{align*}
\]

To get an idea of the numbers involved, we note that \(F(3, 1) = 0.324\). Thus, eq. (4) for \(m_{\tau}\) yields \(f_{r \nu} \tau / 4 \approx 0.4(m_{\tau} / \mu)\). This means that \(m_N / \mu < 1\) is preferred. On the other hand, eq. (3) yields \(\sin 2 \theta = \mu v / \sqrt{2} m_N^2\). This requires \(m_N > 174\) GeV for \(m_{\tau} / \mu < 1\). Hence a small \(\theta\) requires a large \(m_N\) as well as a large \(\mu\), whereas \(\theta = \pi / 4\) is perfectly allowed for \(\mu = m_N = v / \sqrt{2}\). We note also that in models with two Higgs doublets and tree-level fermion couplings, radiative contributions to Higgs Yukawa couplings [23] may be significant in the case of \(h \to b \bar{b}\) from \(t\) exchange.

Fig. 4: (Color online) The ratio \(\Gamma_{\gamma \gamma} / \Gamma_{SM}\) plotted against \(\theta\) for \(x_1 = 3\) and \(x_2 = 1\) with various \((r_{\eta}, r_{\chi})\) and \(\mu / m_N = 1\).

The charged scalars \(\zeta_{1,2}\) also contribute to \(h \to \gamma \gamma\) [24, 25]. Its decay rate is given by

\[
\Gamma_{\gamma \gamma} = \frac{G_F^2 a_2^2 m_h^4}{128 \pi^3} \left[ A_{1/2} \left( \frac{4 m_1^2}{m_h^2} \right) + A_1 \left( \frac{4 m_1^2}{m_h^2} \right) \right] + f_2 A_0 \left( \frac{4 m_1^2}{m_h^2} \right) + f_2 A_0 \left( \frac{4 m_2^2}{m_h^2} \right),
\]

where

\[
\begin{align*}
  f_1 &= \frac{1}{4 x_1} (\sin 2 \theta)^2 (x_1 - x_2) \\
  &\times \left\{ 1 + \frac{1}{2} (x_1 - x_2) [(r_{\eta} + r_{\chi}) + \cos 2 \theta (r_{\eta} - r_{\chi})] \right\}, \\
  f_2 &= \frac{1}{4 x_2} (\sin 2 \theta)^2 (x_1 - x_2) \\
  &\times \left\{ -1 + \frac{1}{2} (x_1 - x_2) [(r_{\eta} + r_{\chi}) - \cos 2 \theta (r_{\eta} - r_{\chi})] \right\}. 
\end{align*}
\]

The \(A\) functions are well known, \(i.e.\)

\[
\begin{align*}
  A_0(y) &= -y [1 - y f(y)], \\
  A_{1/2}(y) &= 2 y [1 + (1 - y) f(y)], \\
  A_1(y) &= -[2 + 3 y + 3 y (2 - y) f(y)], \\
  f(y) &= \arcsin^2 (y^{-1/2}) \text{ for } y \geq 1.
\end{align*}
\]

We plot in fig. 4 the ratio \(\Gamma_{\gamma \gamma} / \Gamma_{SM}\) as a function of \(\theta\) for various values of \((r_{\eta}, r_{\chi})\) with \(x_1 = 3\) and \(x_2 = 1\), assuming also \(\mu / m_N = 1\), \(i.e.\) \(m_N \sin 2 \theta = v / \sqrt{2}\). The current LHC measurements of \(h \to \gamma \gamma\) provide the bounds

\[
\begin{align*}
  \frac{\Gamma_{\gamma \gamma}}{\Gamma_{SM}} &= 1.57 \left( \begin{array}{c} +0.33 \\ -0.28 \end{array} \right) \text{ (ATLAS)}, \\
  \frac{\Gamma_{\gamma \gamma}}{\Gamma_{SM}} &= 0.78 \pm 0.27 \text{ (CMS)}.
\end{align*}
\]

Another important consequence of the radiative generation of fermion masses is the induced electromagnetic interaction which is now related to the fermion mass itself.
If we apply the above procedure to the muon (using a different $N$ and thus also different $f_{\eta,\chi}$), we find its anomalous magnetic moment to be given by

$$\Delta a_\mu = \frac{(g-2)\mu}{2} = \frac{m_\mu^2}{m_N^2} \left[ G(x_1) - G(x_2) \right],$$

where

$$G(x) = \frac{2x \ln x}{(x-1)^3} - \frac{x+1}{(x-1)^2}, \quad H(x) = \frac{x \ln x}{x-1}.$$  

The current discrepancy of the experimental measurement [26] vs. the most recently updated theoretical calculation [27] is

$$\Delta a_\mu = a_\mu^{\exp} - a_\mu^{SM} = 39.35 \pm 5.21_{th} \pm 6.3_{exp} \times 10^{-10}.$$  

We plot in fig. 5 this theoretical prediction for various $m_N$ as a function of $x_2$ with $x_1 = x_2 + 2$, i.e. $\sin 2\theta = \mu v/\sqrt{2m_N^2}$. Also shown is eq. (23) with the experimental and theoretical uncertainties combined in quadrature. We note that the maximum value of $\Delta a_\mu$ is obtained in the limit of $x_1 = x_2 = 1$ where $\Delta a_\mu = m_\mu^2/3m_N^2$. We also note that the subdominant contributions to $\Delta a_\mu$ from $f_0^2$ and $f_0^2$ are negative as expected [28], i.e.

$$\langle \Delta a_\mu \rangle' = \frac{-m_\mu^2}{16\pi^2 m_N^2} \left[ f_0^2 \cos 2\theta J(x_1) + \sin^2 2\theta J(x_2) \right]$$
$$+ f_0^2 \left[ \sin^2 2\theta J(x_1) + \cos^2 2\theta J(x_2) \right],$$

where

$$J(x) = \frac{x \ln x}{(x-1)^3} + \frac{5x - 2}{6(x-1)^3}.$$  

Using eq. (4) for $m_\mu$, we have checked that the value of $(f_0 f_0/4\pi)(\mu/m_N)$ varies between 0.01 and 0.1 in this range. If $m_\mu$ and $m_e$ are both radiative, then $\mu \to e\gamma$ would be severely constrained [29]. However, a flavor symmetry such as $Z_3$ may exist [30] to forbid it.

As for quarks, radiative $m_t$ would require very large corresponding Yukawa couplings in the loop. This is perhaps unrealistic, but such is not an issue with the other quarks. If the $b$ quark mass is radiative as proposed in ref. [16], the one-loop diagram is given by fig. 6. Here there are colored scalar triplets: $(\eta^{2/3}, \eta^{-1/3}, \chi^{-1/3})$, which are, respectively, doublet and singlet under $SU(2)$. The decay rate of $h \to gg$ is then modified.

$$\Gamma_{gg} = G_{\phi} \alpha_3^2 m_3^2 \left| A_{1/2} \right|^2 \left[ f_1^A_0 \left( \frac{4m_4^2}{m_h^2} \right) \right]^2 + f_2^A_0 \left( \frac{4m_4^2}{m_h^2} \right)^2.$$  

Here $m_{1,2}$ refer to the mass eigenvalues of the mixed $(\eta^{2/3}, \chi^{-1/3})$ system and $m_{\eta}^2$ is the mass of $\eta^{2/3}$ with $m_{\eta}^2 = m_4^2 \cos^2 \theta + m_2^2 \sin^2 \theta$. Using $r_{\eta,\chi} = \lambda_{\eta,\chi}^\prime (m_N^2/m')^2$ and $x_{1,2,\eta} = m_{1,2,\eta}^\prime / m_N^2$, we find

$$f_1^A = \frac{1}{4x_1^A} (\sin 2\theta)^2 (x_1^A - x_2^A)$$
$$\times \left\{ 1 + \frac{1}{2} (x_1^A - x_2^A) \left[ (r_\eta + r_\chi)' \cos 2\theta (r_\eta' - r_\chi') \right] \right\},$$

$$f_2^A = \frac{1}{4x_2^A} (\sin 2\theta)^2 (x_1^A - x_2^A)$$
$$\times \left\{ 1 - \frac{1}{2} (x_1^A - x_2^A) \left[ (r_\eta + r_\chi)' \cos 2\theta (r_\eta' - r_\chi') \right] \right\},$$

$$f_{\eta}^C = \frac{r_\eta}{4x_\eta} (\sin 2\theta)^2 (x_1^A - x_2^A)^2.$$
We plot in fig. 7 the ratio $\Gamma_{yy}/\Gamma_{SM}$ as a function of $\theta'$ for $x'_1 = 3$ and $x'_2 = 1$ with $\mu'/m_N = 1$ for various $(r'_y, r'_x)$. This shows that the production of $h$ via gluon fusion may be significantly affected. Currently the sample of $h \rightarrow bb$ decays at the LHC comes only from vector boson associated production, which is unchanged in this case.

In conclusion, we have shown in this paper in detail how the fermionic decay of the 125 GeV particle $h$ discovered at the LHC in 2012 may differ from the expectations of the standard model if a quark or lepton $\psi$ acquires its mass radiatively. The Yukawa coupling of $h$ to $\bar{\psi}\psi$ is then predicted to differ in general from the SM prediction of $m_\psi/v$ where $v = 246$ GeV. A large effect is possible with a modest mixing angle defined in eq. (3). This deviation is possibly observable in $h \rightarrow \tau^+\tau^-$ and $h \rightarrow bb$ with more data at the LHC, due to the resume operation in 2015. The new particles responsible for this deviation are different from other possible explanations of this effect, and if observed, could help to distinguish our proposal from others.

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