Corrigendum

Corrigendum to: “Complete electroweak chiral Lagrangian with a light Higgs at NLO” [Nucl. Phys. B 880 (2014) 552–573]

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The matching of the Higgs-portal example to the chiral Lagrangian in Section 7.2 is corrected to properly account for the systematics of the relevant strong-coupling limit. The text of Section 7.2 given below supersedes the previous version. The main conclusions of Section 7.2 and the rest of the paper are not affected by this correction.

The Higgs-portal model discussed here is equivalent to the Standard Model extended by a heavy scalar singlet. The chiral Lagrangian as the low-energy effective field theory of this model in the strong-coupling regime is further elaborated on in [1].

7.2. Higgs portal

As a specific model for a UV completion we consider the Higgs portal (see [2–5] and references therein). This model postulates the existence of a new, Standard-Model singlet scalar particle, which has allowed dimension-4 couplings to the Higgs field. This interaction modifies the scalar potential of Eq. (4) to

\[ V = -\frac{\mu_s^2}{2} |\phi_s|^2 + \frac{\lambda_s}{4} |\phi_s|^4 - \frac{\mu_h^2}{2} |\phi_h|^2 + \frac{\lambda_h}{4} |\phi_h|^4 + \frac{\eta}{2} |\phi_s|^2 |\phi_h|^2, \] (1)

where \( \phi_s \) refers to the standard scalar doublet and \( \phi_h \) denotes the hidden scalar. Both of them acquire a vacuum expectation value, which can be written as

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\[
\frac{v_s}{\sqrt{2}} = \sqrt{\frac{\lambda_s \mu_s^2 - \eta_s^2}{\lambda_s \lambda_{12} - \eta^2}}, \quad \frac{v_h}{\sqrt{2}} = \sqrt{\frac{\lambda_h \mu_h^2 - \eta_h^2}{\lambda_s \lambda_{12} - \eta^2}}
\]

Expanding both scalars around their vacuum expectation value, i.e. \(|\phi_i| = \frac{1}{\sqrt{2}}(v_i + h_i)\), leads to a potential of the form

\[
V = \frac{\lambda_s v_s^2}{4} h_s^2 + \frac{\lambda_h v_h^2}{4} h_h^2 + \frac{\eta}{2} v_s v_h h_s h_h + \mathcal{O}(h_i^3)
\]

The transformation

\[
\begin{pmatrix}
H_1 \\
H_2
\end{pmatrix} = \begin{pmatrix}
\cos \chi & -\sin \chi \\
\sin \chi & \cos \chi
\end{pmatrix} \begin{pmatrix} h_s \\ h_h \end{pmatrix}
\]

diagonalizes the mass matrix. The rotation angle \(\chi\) is defined as

\[
\tan(2\chi) = \frac{2\eta v_s v_h}{\lambda_h v_h^2 - \lambda_s v_s^2}
\]

The masses of the physical states \(H_1\) and \(H_2\) are given by

\[
M_{1,2}^2 = \frac{1}{4} (\lambda_h v_h^2 + \lambda_s v_s^2) \mp \frac{\lambda_h v_h^2 - \lambda_s v_s^2}{4 \cos(2\chi)}
\]

The Lagrangian relevant for the two scalars then reads

\[
\mathcal{L}_H = \frac{1}{2} \partial_{\mu} H_1 \partial^{\mu} H_1 + \frac{1}{2} \partial_{\mu} H_2 \partial^{\mu} H_2 - V(H_1, H_2)
+ \frac{v^2}{4} (D_{\mu} U^+ D^{\mu} U) \left(1 + \frac{2a_1}{v} H_1 + \frac{2a_2}{v} H_2 + \frac{b_1}{v^2} H_1^2 + \frac{b_2}{v^2} H_2^2 + \frac{b_3}{v^2} H_1 H_2 + \frac{b_4}{v^2} H_1^2 H_2 + \frac{b_5}{v^2} H_1 H_2^2 + \frac{b_6}{v^2} H_1^2 H_2^2\right) - v \left(\bar{q} Y_u U P_+ r + \bar{q} Y_d U P_- r + \bar{Y}_e U P_\eta + h.c.\right) \left(1 + \frac{c_1}{v} H_1 + \frac{c_2}{v} H_2\right),
\]

where

\[
V(H_1, H_2) = \frac{1}{2} M_1^2 H_1^2 + \frac{1}{2} M_2^2 H_2^2 - \lambda_1 H_1^4 - \lambda_2 H_1^2 H_2^2 - \lambda_3 H_1 H_2^3 - \lambda_4 H_2^4
- \lambda_5 H_1^3 H_2 - \lambda_6 H_1^2 H_2^3 - \lambda_7 H_1 H_2^4 - \lambda_8 H_2^5
\]

The couplings \(\lambda_i\) and \(z_i\) depend on \(\mu_s, \mu_h, \lambda_s, \lambda_h\) and \(\eta\). With the parameters of the Higgs-portal model

\[
a_1 = \sqrt{b_1} = c_1 = \cos \chi, \quad a_2 = \sqrt{b_2} = c_2 = \sin \chi, \quad b_{12} = 2 \sin \chi \cos \chi,
\]

the theory is renormalizable and unitary. The scalar \(H_1\) is now identified with the light scalar \(h\) that was found at the LHC. \(H_2\) is assumed to be heavy such that it can be integrated out. In doing so, we take its mass \(M_2\) to be larger than all other energy scales in the model, \(M_2 \gg v_h, v_s, M_1\). In this limit the couplings \(\lambda_s, \lambda_h\) and \(\eta\) become large. We will assume, however, that they still remain in a regime where perturbation theory is a sufficiently reliable approximation. \(H_2\) can then be integrated out at tree level by solving its equation of motion and inserting the solution into the Lagrangian (7). The \(H_2\)-part of this Lagrangian can be written as

\[
\mathcal{L}_{H_2} = \frac{1}{2} H_2 (-\partial^2 - M_2^2) H_2 + J_1 H_2 + J_2 H_2^2 + J_3 H_2^3 + J_4 H_2^4
\]
where the $J_l$ can be read off from (7) and (8). Making the dependence on $M_2$ explicit, the $J_l$ take the form

$$J_l \equiv M_2^2 J_l^0 + \tilde{J}_l$$

(11)

where $J_l^0$ is a pure polynomial in $H_1 \equiv h$.

The equation of motion for $H_2$ reads

$$(- \partial^2 - M_2^2 + 2J_2) H_2 + J_1 + 3J_3 H_2^2 + 4J_4 H_2^3 = 0$$

(12)

It can be solved order by order in powers of $1/M_2^2$ by expanding

$$H_2 = H_2^{(0)} + H_2^{(1)} + H_2^{(2)} + \ldots, \quad H_2^{(l)} = O(1/M_2^{2l})$$

(13)

$H_2^{(0)}$ can be determined from the $O(M_2^2)$ piece of (12) as an infinite series in powers of $h$

$$H_2^{(0)} = \sum_{k=2}^{\infty} r_k h^k$$

(14)

$H_2^{(1)}$ can then be obtained in terms of $H_2^{(0)}$, etc. Inserting the solution (13) of (12) back into (10), (7), and expanding in $1/M_2^2$, one arrives at the effective Lagrangian of the model with $H_2$ integrated out in the limit described above. At leading order, $O(1/M_2^0)$, the result has the form of the electroweak chiral Lagrangian in (4), with the functions $F_U(h)$, $V(h)$ and $\sum_n \hat{Y}_f^{(n)} (h/v)^n$ given as infinite series in $h$. For example,

$$F_U(h) = 2a_1 \frac{h}{v} + (b_1 - a_1 a_2^2 (a_1 + a_2 v/v_h)) \frac{h^2}{v^2} + O(h^3)$$

(15)

Extending the derivation to the NLO terms of $O(1/M_2^2)$ one finds

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{LO}} + \Delta \mathcal{L}_{\text{NLO}} + O\left(\frac{1}{M_2^4}\right),$$

(16)

where ($H_2^{(0)} \equiv H_0$)

$$\Delta \mathcal{L}_{\text{NLO}} = \frac{\left[(- \partial^2 + 2\tilde{J}_2) H_0 + \tilde{J}_1 + 3\tilde{J}_3 H_0^2 + 4\tilde{J}_4 H_0^3\right]^2}{2M_2^2 (1 - 2J_2^0 - 6J_0^0 H_0 - 12J_0^0 H_0^2)}$$

(17)

The effective Lagrangian $\Delta \mathcal{L}_{\text{NLO}}$ contains operators that modify the leading-order Lagrangian (4) as well as a subset of the next-to-leading operators of Section 4. In particular, we have

$$O_{D1}, O_{D7}, O_{D11}, O_{\psi S1}, O_{\psi S2}, O_{\psi S7}, O_{\psi S14}, O_{\psi S15}, O_{\psi S18},$$

(18)

the hermitian conjugates of the $O_{\psi S1}$ in (18), and 4-fermion operators coming from the square of the Yukawa bilinears contained in $\tilde{J}_1$. The 4-fermion operators that are generated have the same structure as those in the heavy-Higgs model discussed in [6], which are

$$O_{FY1}, O_{FY3}, O_{FY5}, O_{FY7}, O_{FY9}, O_{FY10}, O_{ST5}, O_{ST9},$$

(19)

$$O_{LR1}, O_{LR2}, O_{LR3}, O_{LR4}, O_{LR8}, O_{LR9}, O_{LR10}, O_{LR11}, O_{LR12}, O_{LR13}, O_{LR17}, O_{LR18}$$

and their hermitian conjugates, but they are now dressed with functions $F_1(h/v)$.

\footnote{The terms $O_{LR2}, O_{LR4}, O_{LR11}$ and $O_{LR13}$ had been missed in the discussion of the heavy-Higgs models in [6,7].}
This discussion shows explicitly how a subset of our NLO operators is generated in the Higgs-portal scenario. After integrating out the heavy scalar $H_2$ in the non-decoupling limit $M_2 \gg v_h, v_s, M_1$ the effective theory takes the form of a chiral Lagrangian. In particular, even for $F_i (h/v) \rightarrow 1$, it is seen that operators of canonical dimension 4 ($\mathcal{O}_{D1}$), 5 ($\mathcal{O}_{qSi}$) and 6 (4-fermion terms) contribute at the same (next-to-leading) order $1/M_2^2$. This shows that the effective Lagrangian is not simply organized in terms of canonical dimension.

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References

[1] G. Buchalla, O. Catà, A. Celis, C. Krause, arXiv:1608.03564 [hep-ph].
[2] R. Schabinger, J.D. Wells, Phys. Rev. D 72 (2005) 093007, arXiv:hep-ph/0509209.
[3] B. Patt, F. Wilczek, arXiv:hep-ph/0605188.
[4] M. Bowen, Y. Cui, J.D. Wells, J. High Energy Phys. 0703 (2007) 036, arXiv:hep-ph/0701035.
[5] C. Englert, T. Plehn, D. Zerwas, P.M. Zerwas, Phys. Lett. B 703 (2011) 298, arXiv:1106.3097 [hep-ph].
[6] G. Buchalla, O. Catà, J. High Energy Phys. 1207 (2012) 101, arXiv:1203.6510 [hep-ph].
[7] G. Buchalla, O. Catà, C. Krause, Nucl. Phys. B 880 (2014) 552, arXiv:1307.5017 [hep-ph].