(Re)visiting (Non)decoupling

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based in part on our work in progress with J. Brehmer (ITP Heidelberg), A. Freitas (Pittsburg U.) & T. Plehn (ITP Heidelberg)

CP3 - Université catholique de Louvain

Lunch seminars @ CP3 - May 27th 2015

Outline

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Motivation (I) - Historical Background

- TIMEFRAME ['80s] CONTEXT **GUT model building**

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Kazama & Yao Phys. Rev. D 25 (1982) 1605

♠ *A central idea behind the unification of forces of vastly different strength is that such an apparent hierarchy arises not from the difference of the fundamental coupling constants of the theory but rather from that of the masses of the exchanged particles.*

♠ *In constructing a viable unified theory, these heavy particles must be incorporated into the structure with due caution. Among the most important requirements are (i) that <u>superheavy particles must effectively decouple at low energies</u>, (ii) that <u>correct effective light–particle theory must emerge at low energies</u>, and (iii) that the mass hierarchy, arranged at the tree level, should be stable against radiative corrections. As we shall see, these requirements are deeply interrelated. None of them are trivial to satisfy.*

Motivation (II) - Phenomenological interest

♠ **DECOUPLING** – a somehow implicit notion

- Collider physics itself
- Effective field theory & simplified model approaches to new physics
- Model–independent new physics parametrization (S, T, U parameters)

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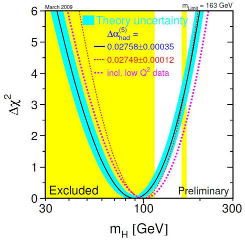
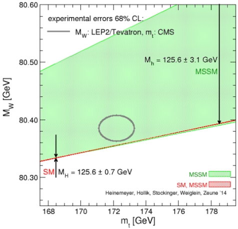
♠ **NONDECOUPLING** – – a somehow implicit notion (as well)

- Precision EW and heavy flavor physics as probes of missing SM building blocks
  $m_t - m_W$ correlations, m_H blueband
- Precision EW & heavy flavor physics constraints on new physics





Keywords

High scale Λ_{UV} **Heavy fields** Φ **Large masses** M $\hat{\mathcal{O}}(\Lambda_{\mathsf{UV}})$ $\mathcal{L}_{\mathsf{UV}}(\Phi,\varphi;g,g^*,m,M)$

for $\boxed{M \to \infty}$

Low scale Λ_{IR} **Light fields** φ **low masses** m $\hat{\mathcal{O}}(\Lambda_{\mathsf{IR}})$ $\mathcal{L}_{\mathsf{IR}}(\varphi;g,m)$

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- ♠ $\hat{\mathcal{O}}(\Lambda_{\text{IR}})$ UV–insensitive for $M \to \infty$
- ♠ heavy field effects undetectable for $M \to \infty$

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A more formal definition

The property by which field theory amplitudes computed from $\mathcal{L}_{\text{IR}}$ and $\mathcal{L}_{\text{UV}}$ are coincident in the in the asymptotic heavy mass limit – up to at most a redefinition of the renormalizable parameters in $\mathcal{L}_{\text{IR}}$

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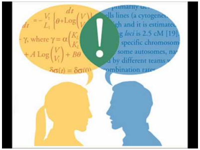

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♠ **NONDECOUPLING** $\left\{ \right.$

- ♠ $\mathcal{O}(\Lambda_{\text{IR}})$ are UV–sensitive for $M \to \infty$
- ♠ heavy field effects remain for $M \to \infty$

A more formal definition

Remainder of the UV–scale dynamics in low–energy observables for asymptotically large heavy field masses

Questions ahead

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- What fundamental field theory mechanisms are responsible for decoupling?
- What field theory mechanisms lead to the breakdown of decoupling conditions?
- What are the implications of non–decoupling for BSM Higgs physics?
- What is the interplay between non–decoupling effects and effective field theories?

Outline

The Appelquist–Carazzone theorem

♠ The notion of heavy field decoupling - a first QFT formalization:

♠ K. Symanzik (1973), *Infrared singularities and small distance behavior analysis*

♠ T. Appelquist, J. Carazzone, (1975), *Infrared Singularities and Massive Fields*

Phys.Rev. D11 (1975) 2856

We examine some problems associated with the low-momentum behavior of gauge theories and other renormalizable field theories. Our main interest is in the infrared structure of unbroken non-Abelian gauge theories and how this is affected by the presence of other heavy fields coupled to the massless gauge fields. It is shown in the context of a simple model of gauge mesons coupled to massive fermions that the heavy fields decouple at low momenta except for their contribution to renormalization effects.

The Appelquist–Carazzone theorem

♠, Setup: Simplified model of massless gauge bosons coupled to massive fermions:

$$\mathcal{L}_{\textsf{full}} = \underbrace{-\frac{1}{4} F^{\mu\nu} F_{\mu\nu}}_{\text{light sector}} + \underbrace{i\overline{\psi}\not{D}\,\psi + M\overline{\psi}\psi}_{\text{heavy sector}} + \underbrace{\mathcal{L}_{\varphi_\mu - \psi}}_{\text{light–heavy cross–talk}}$$

$$\mathcal{L}_{\textsf{full}} = \mathcal{L}_{\textsf{IR}(\varphi,\{g\})} + \mathcal{L}_{\textsf{UV}(\varphi,\{g,g^*\,M\})} + \mathcal{L}_{\varphi_\mu - \psi}$$

$\boxed{G^{(n)}(p_1,\ldots,p_n;\mu)}$ generic n–point light field Green's function computed from $\mathcal{L}_{\textsf{full}}$

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$G^{(n)}(p_1,\ldots,p_n;\mu)$ generic n–point light field Green's function computed from $\mathcal{L}_{\textbf{full}}$

♠ Statement: In the limit $\{p_i\} \ll \{M\}$,

$$\underbrace{G^{(n)}(p_1,\ldots,p_n;\mu)}_{\text{from } \mathcal{L}_{\textbf{full}}} = Z^{n/2}\,\underbrace{\tilde{G}^{(n)}(p_1,\ldots,p_n;\mu)}_{\text{from } \mathcal{L}_{\textbf{eff}}} + \mathcal{O}(1/M),$$

$\mathcal{L}_{\textbf{eff}} \equiv \mathcal{L}_{\textbf{IR}}(\{\tilde{g}_i\})$, with $\tilde{g}_i \equiv \tilde{g}_i(\{g_i, M\})$ and $Z \equiv Z(\{g_i, M\})$.

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Proof keypoints:

♠ SCALING ♠ RENORMALIZABILITY

Generalizing Appelquist–Carazzone

♠ Puzzling!! – while GUT fields decouple, heavy SM fields do not

Examples

- Large charm mass effect in the $\mathcal{O}(\alpha_s)$ corrections to the axial isoscalar weak current: Collins, Wilczek, Zee ['78]
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Generalized conditions for decoupling

1. **A separably renormalizable light sector**, which remains so at each possible stage of symmetry breaking
2. **Asymptotically weak couplings**

Decoupling breakdown: categories

A Nondecoupling taxonomy

1. **Intrinsic**, in situations beyond the validity range of the AC theorem
2. **Accidental**, in situations for which AC holds though is apparently violated

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♠ **Intrinsic**

- One VEV + hierarchical Yukawas: c or t quarks in the SM
- Strong coupling: Higgsless SM

♠ **Accidental**

- Renormalization scheme – induced: On–shell Z' sector renormalization
- UV–scale induced: linked gauge couplings from GUTs
- Delayed decoupling GUT–induced $\not{B}$–violating interactions

Sample case: $b \to s\gamma$

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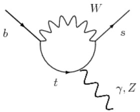

EW penguin contributing to the $\boxed{B \to X_S\, \gamma \text{ FCNC decay}}$

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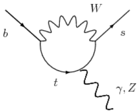

EW penguin contributing to the $B \to X_S\,\gamma$ FCNC decay

$$\boxed{\mathcal{M}(b \to s\gamma) = \sum_{q=u,c,t} \lambda_q\, F(x_q)} \quad \text{where} \quad \lambda_q = V^*_{qb} V_{qs}$$

♠ CKM unitarity implies $\boxed{\lambda_u + \lambda_c + \lambda_t = 0}$ (tree–level GIM)

♠ $\boxed{F(x_q) \sim \frac{m_q^2}{m_W^2}}$ resulting from a chirality flip $\Rightarrow$ $\boxed{\mathcal{O}\left(\frac{m_t^2}{m_W^2}\right)}$

- If all quarks where close–by in mass and hence separably renormalizable, GIM suppression would be stronger – non–decoupling would be tamed
- Example: FCNC top–quark decays – non–decoupling is at most $\mathcal{O}(m_b^2/m_W^2)$

The role of symmetries (I)

♠ Symmetries play a key role in understanding the field theory mechanisms responsible for protecting or breaking the decoupling conditions

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We have seen how superheavy particles in grand unified theories decouple at low momenta. The reason, we believe is that the ***gauge invariance*** $SU(2)_L \otimes U(1)_Y$ *of the SM which remains unbroken at the energies of the order of the superheavy particle masses. This explanation was suggested to us by Frank Wilczek. We thank him for a very useful discussion on that point.*

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- renormalizability protects Green's functions from receiving divergent contributions not absorvable into CTs
- While guaranteed, UV finiteness may rely on delicate cancellations between light and heavy virtual particle contributions
- Heavy mass dependence as the remainder of the would–be UV poles if the heavy fields were removed.

The role of symmetries (II)

Heavy fields decouple if:

- respect $SU(2)_L \otimes U(1)_Y$ and hence its breaking pattern
- **Example:** See–saw neutrino mass

$$\mathcal{L} \sim M (N_R^c)^T\, C^{-1}\, N_R + \overline{L}_L\, Y_N\, N_R + \text{h.c.} \quad \Rightarrow \quad M_{\nu,\text{eff}} \sim -\frac{Y_N^T\, Y_N}{M}\, v^2$$

- break $SU(2)_L \otimes U(1)_Y$ via a singlet VEV whilst remain separably renormalizable
- **Example:** Degenerate $SU(2)_L$ doublets – contributing to the T parameter as

$$\boxed{T \sim G_F\, \frac{M_1^2\, M_2^2}{M_1^2 - M_2^2}\, \log\left(\frac{M_1^2}{M_2^2}\right) - \frac{M_1^2 + M_2^2}{2} \sim \frac{(M_1^2 - M_2^2)}{M_1^2}}$$

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Heavy fields do not decouple if:

- Strongly coupled – mass hierarchy $=$ single scale(v^2) $+$ coupling hierarchy)
- Not separably renormalizable
- **Example:** m_t–imprints on low–energy observables

$$\boxed{\alpha T = \frac{3 G_F}{8\sqrt{2}\pi^2} F_0(m_t^2, m_b^2)} \quad \text{where} \quad \boxed{F_0(m_t^2, m_b^2) = m_t^2 + m_b^2 - \frac{m_t^2\, m_b^2}{m_t^2 - m_b^2} \log\left(\frac{m_t^2}{m_b^2}\right)}$$

The role of symmetries (III)

♠ Underlying/Accidental symmetries can temper non–decoupling effects

♠ Examples:

- <u>**Custodial symmetry**</u> leads to Veltman screening
- <u>**GIM mechanism**</u> keeps FCNC meson decay rates small
- <u>**Supersymmetry**</u> prevents mass splittings (viz. large T) from SUSY heavy Higgs bosons

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- **delayed decoupling** [$\not{B}$–violating interactions in GUTs]
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♠ **Non–linearly realized symmetries** are intrinsically linked to **non–decoupling**

(Non)decoupling and Effective Field Theories

♠ Back to History: Kazama & Yao Phys. Rev. D 25 (1982) 1605

♠ *It is clear in the case of theories without spontaneously broken symmetry one cannot talk about decoupling without the existence of an effective light-particle theory, because in its absence we cannot absorb the large mass effects by redefinition of the parameters of the light theory. These two concepts are, therefore, two sides of one and the same subject.*

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$$\boxed{\mathcal{L} = \mathcal{L}_{\mathsf{SM}} + \sum_{i_D} \sum_{D>4} \frac{C_{i_D}^{(D)}}{\Lambda^{D-2}} \mathcal{O}_{i_D}^{(D)}}$$

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WEAK COUPLING: Passarino ['12]

- **Decoupling**: $\boxed{C^{(D)}_{i_D} \sim \mathcal{O}(1/\Lambda)}$
- **Screening**: $\boxed{C^{(D)}_{i_D} \sim \log(\Lambda)}$
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Non–decoupling **modifies** the naive correspondence between mass dimension and perturbative expansions

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STRONG COUPLING: effective operators built from gauge–invariant operators for which the gauge symmetry is realized non–linearly CCWZ formalism - EWχ Lagrangian

Outline

Loop–induced Higgs couplings: H $\to gg$

♠ We consider a Higgs scalar partner $\mathcal{L} \supset -\mu_S^2\, |S|^2 - \lambda_s\, \Phi^\dagger\, \Phi\, |S|^2$,

$$\boxed{m_S^2(v) = \frac{\lambda_s\, v^2}{2} + \mu_S^2} \qquad \boxed{\mathcal{L}_{hSS} \supset \lambda_s\, v\, |S|^2\, \mathrm{H}} \qquad \boxed{\lambda_s = \frac{2(m_S^2 - \mu_S^2)}{v^2}}$$

$$\boxed{\mathcal{L} \supset -\frac{1}{4}\, g_{hGG}\, G^{\mu\nu\, A}\, G^A_{\mu\nu} \mathrm{H}} \quad \text{with} \quad \boxed{g_{hGG} = -\frac{g_s^2}{4\pi^2} \left[\sum_S \mathrm{dim}(r_L)\, C_{2,S}(r_C)\, \frac{\lambda_s\, v}{2m_S^2}\, A_s(\tau_S) \right]}$$

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$$\boxed{g_{hGG}} \sim \frac{\lambda\,v^2}{m_S^2}\begin{cases} \boxed{\to 0} \text{ if } m_s \to \infty \text{ as } \mu_S \to \infty \text{ while } \lambda \sim \mathcal{O}(1) \\ \\ \boxed{\to \text{constant}} \text{ if } m_S \to \infty \text{ while } \mu_S \ll m_s \end{cases}$$

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$$\boxed{g_{hGG}} \sim \frac{\lambda\,v^2}{m_S^2} \begin{cases} \boxed{\to 0} \text{ if } m_s \to \infty \text{ as } \mu_S \to \infty \text{ while } \lambda \sim \mathcal{O}(1) \\ \boxed{\to \text{constant}} \text{ if } m_S \to \infty \text{ while } \mu_S \ll m_s \end{cases}$$

♠ Non–decoupling contributions are connected to **LETs**

$$\boxed{\mathcal{L} \supset \frac{\alpha_s}{12\pi\,v}\left[s\,C_2(r_C)\,\frac{\partial\,\log(m(v))}{\partial\,\log(v)}\right]\mathrm{H}\,G^{\mu\nu\,A}\,G_{\mu\nu\,A}}.$$

$\boxed{\Rightarrow}$ Φ as a background field to which the (dynamical) heavy fields couple as $M_\Phi = M_\Phi(v)$

2HDM mass spectrum

- **Scalar potential**

$$
\begin{aligned}
V(\Phi_1,\Phi_2) = {} & m_{11}^2\, \Phi_1^\dagger \Phi_1 + m_{22}^2\, \Phi_2^\dagger \Phi_2 - \left[m_{12}^2\, \Phi_1^\dagger \Phi_2 + \text{h.c.} \right] \\
& + \frac{\lambda_1}{2}\, (\Phi_1^\dagger \Phi_1)^2 + \frac{\lambda_2}{2}\, (\Phi_2^\dagger \Phi_2)^2 + \lambda_3\, (\Phi_1^\dagger \Phi_1)\, (\Phi_2^\dagger \Phi_2) + \lambda_4\, |\Phi_1^\dagger\, \Phi_2|^2 \\
& + \left[\frac{\lambda_5}{2}\, (\Phi_1^\dagger \Phi_2)^2 + \lambda_6\, (\Phi_1^\dagger \Phi_1)\, (\Phi_1^\dagger \Phi_2) + \lambda_7\, (\Phi_2^\dagger \Phi_2)\, (\Phi_1^\dagger \Phi_2) + \text{h.c.} \right] \; ,
\end{aligned}
$$

2HDM mass spectrum

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$$+\left[\frac{\lambda_5}{2}\,(\Phi_1^\dagger\Phi_2)^2 + \lambda_6\,(\Phi_1^\dagger\Phi_1)\,(\Phi_1^\dagger\Phi_2) + \lambda_7\,(\Phi_2^\dagger\Phi_2)\,(\Phi_1^\dagger\Phi_2) + \text{h.c.}\right]\,,$$

♠ Mass eigenvalues

$$m_{\mathrm{H}^\pm}^2 = \underbrace{\frac{m_{12}^2}{\sin\beta\cos\beta}}_{\textbf{Decoupling}} - \underbrace{\frac{v^2}{2}\,[\lambda_4+\lambda_5+\lambda_6\,\cot\beta+\lambda_7\tan\beta]}_{\textbf{non–decoupling}}$$

$$m_{\mathrm{A}^0}^2 = \frac{m_{12}^2}{\sin\beta\cos\beta} - \frac{v^2}{2}\,[2\lambda_5+\lambda_6\,\cot\beta+\lambda_7\,\tan\beta]$$

2HDM mass spectrum

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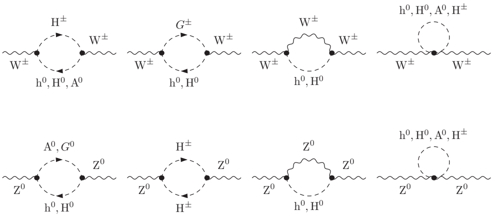
- **Decoupling limit**

$$\Phi_1 \supset \begin{pmatrix}\mathrm{H}^+\\ \mathrm{A}^0\\ \mathrm{H}^-\end{pmatrix} \oplus \mathrm{H}^0 \quad \text{or} \quad \Phi_1 \supset \begin{pmatrix}\mathrm{H}^+\\ \mathrm{H}^0\\ \mathrm{H}^-\end{pmatrix} \oplus \mathrm{A}^0 \qquad \Phi_2 \supset \begin{pmatrix}G^+\\ G^0\\ G^-\end{pmatrix} \oplus \frac{v+\mathrm{h}^0}{\sqrt{2}}\,.$$

all Higgs states fall into the singlet and triplet representations of $SU(2)_{L+R}$

2HDM contributions to the T parameter



$$
\begin{aligned}
(\alpha_{em}\,T)_{\mathbf{2HDM}} \supset \frac{G_F}{8\sqrt{2}\,\pi^2}\Bigg\{ & m^2_{H\pm}\left[1 - \frac{m^2_{A^0}}{m^2_{H\pm} - m^2_{A^0}}\ln\frac{m^2_{H\pm}}{m^2_{A^0}}\right] \\
& + \cos^2(\beta-\alpha)\, m^2_{h^0}\left[\frac{m^2_{A^0}}{m^2_{A^0} - m^2_{h^0}}\ln\frac{m^2_{A^0}}{m^2_{h^0}} - \frac{m^2_{H\pm}}{m^2_{H\pm} - m^2_{h^0}}\ln\frac{m^2_{H\pm}}{m^2_{h^0}}\right] \\
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\end{aligned}
$$

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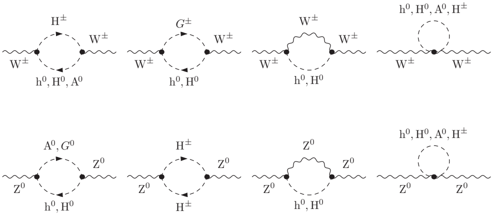

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$$+\sin^2(\beta-\alpha)\,m^2_{H^0}\left[\frac{m^2_{A^0}}{m^2_{A^0}-m^2_{H^0}}\ln\frac{m^2_{A^0}}{m^2_{H^0}} - \frac{m^2_{H\pm}}{m^2_{H\pm}-m^2_{H^0}}\ln\frac{m^2_{H\pm}}{m^2_{H^0}}\right]\Big\}.$$

♠ Non–decoupling beyond screening as power–like prefactors to Veltman logs

♠ $\boxed{(\alpha_{em}\,T)_{[\mathbf{2HDM}]} \simeq -\frac{3\,G_F\,m_Z^2\,s_W^2}{8\sqrt{2}\pi^2}\,\xi^2\,\log\left(\frac{m^2_\Phi}{m^2_W}\right)}$ for degenerate $m_{\Phi=\mathrm{H}^0,\mathrm{A}^0,\mathrm{H}^\pm}$

Loop–corrected Higgs couplings in the 2HDM

♠ **Coupling shift parametrization:** $\hat{\kappa}_f \equiv \frac{\Gamma^S_{hff}}{\Gamma^S_{hff,\mathsf{SM}}} \equiv \boxed{(1+\Delta_f)}$

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- **Leading quantum effects on Higgs couplings** [Kanemura et. al. arXiv:1502.07716]

  - **Loop Form factor**: $\Gamma_{hff}(p, k_1, k_2) = \Gamma^S_{hff} + \gamma_5\, \Gamma^P_{hff} + \not{k}_1\, \Gamma_{hff}{}^{V1} + \ldots$

  - **Coupling shift:**

$$1 + \Delta_t = \xi\, \cot\beta - \frac{1}{2}\xi^2 - \frac{1}{96\pi^2} \sum c_\Phi \frac{m_\phi^2}{v^2} \left(1 - \frac{M^2}{m_{\phi^2}^2}\right)^2$$

$$- \frac{1}{12\,\pi^2} \cot^2\beta \sum c_\Phi \left[\frac{m_t^4}{v^2\, m_\phi^2} + \tan^2\beta\, \frac{m_b^2\, m_t^2}{v^2\, m_{\mathsf{H}^\pm}^2}\right]$$

with $\Phi = \mathsf{A}^0, \mathsf{H}^0, \mathsf{H}^\pm$ and $\xi = \cos(\beta - \alpha) \simeq \dfrac{\tilde{\lambda} v^2}{M^2}$

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♠ Non–decoupling contributions — NLO EW in the full theory

♠ Non–decoupling contributions — $\mathcal{O}^{\mathsf{D=8}}$ in the EFT — $\boxed{\frac{G^{\mu\nu A}G_{\mu\nu A}[\Phi^\dagger\Phi]^2}{\Lambda^4}}$

Outline

Closing remarks

Take–home ideas

♠ **General conditions for decoupling**

- A separably renormalizable heavy sector
- Asymptotically weak couplings
- Light and heavy sectors not linked by renormalization conditions

♠ **Categories of non–decoupling:**

- **Accidental** UV–scale induced; renormalization–induced; delayed decoupling
- **Intrinsic:** strong coupling; one single scale for light and heavy masses

♠ **Symmetries** (gauge, spontaneous breaking, accidental) are the **key QFT ingredient**

♠ Non–decoupling effects are **ubiquitous in Higgs physics**

♠ **EFT perspective**: **screening** (*polynomial*) non–decoupling contributions lead to Wilson coefficients with **logarithmic** (*power–like*) dependence in Λ_{UV}.

Additional material

ADDITIONAL MATERIAL

Non–decoupling: a broaden perspective

♠ Non–decoupling cross–talk between distant scales

Non–decoupling: a broaden perspective

♠ Non–decoupling ⟺ cross–talk between distant scales

Many phenomena can be adscribed into this category

♠ Low–energy manifestation of **relevant** or **marginal** operators

- $\hat{\mathcal{O}}^{D=0}$: the Cosmological Constant
- $\hat{\mathcal{O}}^{D=2}$: the SM Hierarchy problem – m_{H} radiative instability

♠ Long–range interactions

- **Gravity**
- **Temperature**

♠ Non–local correlations

- In space – quantum entanglement
- In time – secular effects

♠ Multi–scale & criticallity

NLO EW corrections to $h \to \gamma\gamma$

Leading $\mathcal{O}(\alpha_{ew})$ corrections to $h \to \gamma\gamma$: Degrassi et al., ['94]

$$\boxed{\mathcal{L}_{gg\mathsf{H}} \supset (\sqrt{2}\,G_F)^{1/2}\,\frac{\alpha_s}{12\pi}\mathsf{H}\,G^{\mu\nu}G_{\mu\nu}\left[1 + \frac{G_F\,\sqrt{2}}{32\,\pi^2}m_t^2\right]}$$

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♠ Non–decoupling **non–trivially relates** coupling constant & operator mass dimension expansions

SUSY D–term contributions to the BMSSM Higgs potential

♠ Toy model: $\boxed{\text{MSSM}}$ $\longrightarrow$ $\boxed{SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U_x(1)}$

[Batra, Delgado, Kaplan, Tait hep-ph/0309149]

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- Additional (super)field content:
  - $U(1)_X$–breaking complex scalars φ, φ^C $\quad \boxed{\langle\varphi\rangle^2 \sim m_\varphi^2}$
  - Heavy Z′ boson $\quad \boxed{M_{Z'} = 2 q_X g_X \langle\varphi\rangle}$

- Extended Lagrangian:

- D–terms: $\mathcal{L}^{\text{D–term}}_{MSSM \otimes U(1)_X} \supset \mathcal{L}^{\text{D–term}}_{MSSM} + \frac{g_X^2}{2}\left[\frac{1}{2}|\tilde{\Phi}_u| + \frac{1}{2}|\tilde{\Phi}_d|^2 + q_X|\varphi|^2 - q_X|\varphi^C|^2\right]$

- Soft SUSY breaking: $\mathcal{L}^{\text{soft}}_{\varphi} \supset m_\varphi^2 \left[|\varphi|^2 + |\varphi^C|^2\right]$

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- If $m_\varphi^2 > v^2$ SUSY breaks down for φ, Z′ before EWSB

$\Rightarrow$ Naturalness remains protected

$\Rightarrow$ $U(1)_X$ D–terms do not decouple and contribute to the light Higgs potential

$$\boxed{V_{\mathsf{H}} \supset \frac{g_X^2}{2}\left[\frac{1}{2}|\Phi_u|^2 + |\tilde{\Phi}_d|^2\right] \times \left(1 + \frac{m_{Z'}^2}{2m_\Phi^2}\right)^{-1}}$$

Higgs triplet contributions to the T parameter

Heavy scalar triplet $T \equiv T^a \, \frac{\sigma^a}{2} = \begin{pmatrix} \frac{T^0 + v_T}{2} & -\frac{T^+}{\sqrt{2}} \\ -\frac{T^-}{\sqrt{2}} & -\frac{T^0 + v_T}{\sqrt{2}} \end{pmatrix}$ [Chen, Dawson ['05]

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♠ Tree–level analysis:

$$\boxed{m_Z^2 = \frac{1}{4}(g^2 + g'^2)\, v_D^2} \qquad \boxed{m_W^2 = \frac{1}{4} g^2\, v_D^2 + g^2\, v_T^2} \qquad \boxed{\alpha T = \frac{1}{1 - 4\sqrt{2}\, G_F\, v_T^2} - 1}.$$

♠ Tree–level $SU(2)_V$ breaking $\Rightarrow$ δv_T **must be added at one–loop**

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♠ One–loop analysis

$$\boxed{T^{\textsf{1–loop}} \sim -T^{\textsf{tree}} \left[\frac{8\sqrt{2}\, G_F\, v_T}{32\pi^2 (1 - 4\sqrt{2}\, G_F\, v_T)} \times \frac{3\mu\, v_T (\kappa\, v_D^2 + \mu v_T)}{2\, m_\Phi^2} \left(\log \frac{m_\Phi^2}{Q^2} - 1 \right) \right]}$$

♠ Once $v_D \neq 0$ breaks $SU(2)_L \otimes U(1)_Y$, gauge symmetry can no longer keep $v_T = 0$ radiatively stable (barring fine–tuned Q^2) ⇒ Nondecoupling appears for $\mu \sim m_\Phi$

Sample case (II): QCD corrections to J_5^μ

Collins, Wilczek & Zee [PRD 18 (1978)] *Low–energy manifestations of heavy particles*

♠ consider a sequential four–quark Weinberg–Salam model of weak interactions.

- The states (u, d, s) are treated as **light** and c as **massive** separately
- The full theory, as well as the **light** (u, d) and **heavy** (c, s) subsectors are **separably gauge invariant** and hence **renormalizable**

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Screening–like non–decoupling contribution induced by the charm mass

- Gauge invariant theory $\boxed{\Rightarrow}$ all quarks (regardless of their masses) contribute
- Gauge symmetry spontaneously broken $\boxed{\Rightarrow}$ quarks become massive
- One single mass scale & hierarchical Yukawas $\boxed{\Rightarrow}$ $m_{u,d,s} \ll m_c$
- Heavy quark sector alone **Not separably renormalizable**