Effective field theories, like the Standard Model Effective Field Theory (SMEFT), are defined by a chosen field content and a set of symmetries, up to a cut off scale $\Lambda$. Usually, in order to perform calculations, gauge independent field re-definitions consistent with the symmetries of the theory are then used to redefine the fields. This procedure results in a fixed (non-redundant) operator basis, that is not itself field re-definition invariant. Recently, an alternative approach of identifying and calculating with field space geometry has been developed. Field redefinition invariants, characterising field space geometry, appear in observables in amplitude perturbations, and have an expansion in terms of local operators. In the case of the SMEFT, calculating via the geometric approach is known as the geoSMEFT. This approach makes it much easier to calculate at high orders in $1/\Lambda$ in the SMEFT, and can directly result in a complete characterisation of an amplitude perturbation in the $1/\Lambda$ expansion. Using the geoSMEFT, several consistent and complete $O(1/\Lambda^4)$ results are now known. We define the geoSMEFT and demonstrate its use in some examples.

1 Introduction

At this time, there is no statistically significant evidence in particle physics experiments pointing to an explicit new state to add to Standard Model (SM). Simultaneously, several interesting hints of deviations from the SM expectation in the global pattern of experimental results exist. This argues that a mass gap is present between any putative scale of new physics $\sim \Lambda$ and the electroweak scale $v \simeq 246$ GeV; possibly leading to deviations in the predicted SM pattern for some experimental results. Such a mass gap, if limited to $v/\Lambda < 1/4\pi$ can also be consistent with expectations of UV physics motivated by naturalness concerns for the Higgs mass ($m_h$). Broad classes of new physics scenarios consistent with this assumption can be studied efficiently using Effective Field Theory (EFT) methods to analyse data sets gathered at energies $\sqrt{s} \sim v << \Lambda$.

It is of interest to examine current, and possible future deviations in measurements, in a consistent theoretical EFT framework, that also incorporates lower energy experimental data gathered on the $Z, W^\pm$ and $h$ particle phase space poles. This theoretical framework has come to be known as the linear Standard Model Effective Field Theory (SMEFT). Here “linear SMEFT” refers to the assumption that the particle spectrum contains a SU(2) scalar doublet $H$, before the Higgs takes on a vacuum expectation value. Developing predictions for particle physics experiments in the linear SMEFT to $O(1/\Lambda^2)$ is essentially a solved problem$^{1,2}$, and fully automated when using the SMEFTsim package$^{3,4}$. However, the SMEFT is an expansion in $1/\Lambda$. When interpreting the data at $O(1/\Lambda^2)$ in the SMEFT, the sub-leading $O(1/\Lambda^4)$ terms in the expansion exist and cannot be ignored if the relative suppression between these corrections is not numerically negligible. This is likely to be the case if deviations from the SM are found. As such, the effects of these sub-leading terms should be considered, not naively ignored. In addition, loop corrections in the SMEFT also exist, which also need consideration, and at times should be incorporated into global SMEFT fits$^{5,6,7,8}$. 
The development of SMEFT loop calculations has advanced steadily for years, since the foundation of developing these corrections was laid down by the systematic dimension six (gauge independent) renormalization results\textsuperscript{1,9,10}. In recent years, a dimension eight operator basis has also been defined and published\textsuperscript{11}. However, despite initial studies and partial results\textsuperscript{12}, systematically advancing the predictions of the SMEFT to dimension eight was a daunting prospect. Such theoretical calculations are of value to have an informed understanding of neglected dimension eight terms in SMEFT studies at $O(1/\Lambda^2)$ to avoid over interpreting the results of global SMEFT fits. This point has long been made in the literature\textsuperscript{13,14,15,16} when considering the SMEFT as a model independent (bottom up) EFT. To define such error estimates in a bottom up fashion within the EFT here we focus on how to develop theoretical calculations consistently to $O(1/\Lambda^4)$ rapidly, completely and efficiently\textsuperscript{a}.

The key to enabling the systematic advance to $O(1/\Lambda^4)$ was to reconsider a basic feature in the SMEFT. Field redefinitions are used to define an operator basis in an effective field theory. But the resulting operator basis is not field redefinition invariant. This is well known and not an intrinsic problem at any mass dimension. An interesting question is: “Is there a way to calculate that is more field redefinition invariant?” It is well known that operator bases are unphysical, in that they combine up in a contribution to a specific observable in a manner that a consistent global pattern of deviations are present – irrespective of the operator basis chosen. This implies that combinations of operators can be considered more “physical”, than an individual operator. The immediate follow up question then becomes: - “Are particular combinations of operators more useful to consider?” And if so - what defines such combinations of operators systematically?

A useful set of intermediate geometric quantities that the operators combine into in the SMEFT are now known. The key point is that abstract interaction field spaces are present in the EFT and define tensors - metrics and connections on such spaces. These tensors are useful as they then project geometric invariants of these spaces onto amplitudes perturbations in the SMEFT. It is these geometric invariants the the operators combine up into in amplitude perturbations (as observables are field redefinition invariant). When considering low n-point interactions critical for phenomenology, the number of geometric quantities that so project is limited and tractable to identify. Defining and exploiting these geometric quantities leads to the geometric formulation of the SMEFT (geoSMEFT).

1.1 geoSMEFT

Building on many key steps in the previous literature\textsuperscript{19,20,21,22,23,24} The geoSMEFT\textsuperscript{25} is constructed in the following fashion. For the sake of constructing CP even two, three point functions, a set of field space metrics and connections are defined as follows. The scalar potential is

$$V(\phi) = -\mathcal{L}_{\text{SMEFT}}\big|_{\mathcal{L}(\alpha,\beta\longrightarrow 0)}.$$  

(1)

The field space metric for the scalar field bilinear, dependent on the SM field coordinates, is defined via

$$h_{IJ}(\phi) = \frac{g^{\mu\nu}}{d} \frac{\delta^2 \mathcal{L}_{\text{SMEFT}}}{\delta (D_\mu \phi)^I \delta (D_\nu \phi)^J} \bigg|_{\mathcal{L}(\alpha,\beta\longrightarrow 0)}.$$  

(2)

The notation $\mathcal{L}(\alpha, \beta \cdots)$ corresponds to non-trivial Lorentz-index-carrying Lagrangian terms and spin connections, i.e. \{${W^A_{\mu\nu}}$, $(D^\mu \Phi)^K$, $\bar{\psi} \sigma^\mu \psi$, $\bar{\psi} \psi \cdots$\}. Here $d = 4$ is the number of space

\textsuperscript{a}Note that an alternative approach of assuming a specific UV completion, or a set of UV completions (i.e. insisting on a top down understanding), to fix matching patterns is doomed to fail in characterising the full space of the SMEFT as a bottom up EFT. As then one is assuming a different theory then the SMEFT to define the effect of sub-leading terms.Asserting that only a top down EFT approach is possible for error estimates in EFT studies\textsuperscript{3}, or EFT in general, would be as absurd as asserting only a bottom up EFT approach is possible for such studies. The approach pursued for an error estimate is always a freely chosen convention. Here, we report and summarise results in some recent published literature\textsuperscript{18} developed (in part) to enable robust error estimates in a bottom up approach to the SMEFT.
time dimensions and \( g^{\mu\nu} \) is the usual Minkowskian space time metric. Note that this definition reduces the connection \( h_{IJ} \) to a function of \( SU(2)_L \times U(1)_Y \) generators, scalar fields coordinates \( \phi_i \) and \( \bar{v}_T \). Here \( \sqrt{2} \langle H^H H \rangle_{SM} \equiv v_T \) is the vev, including the tower of higher order corrections in the SMEFT. The CP even gauge field scalar manifolds, for the \( SU(2)_L \times U(1)_Y \) fields interacting with the scalar fields, give

\[
g_{AB}(\phi) = -\frac{2 g^{\mu\nu} g^{\sigma\rho}}{d^2} \frac{\delta^2 \mathcal{L}_{\text{SMEFT}}}{\delta W^A_{\mu\sigma} \delta W^B_{\nu\rho}} \bigg|_{\mathcal{L}(\alpha,\beta\ldots) \to 0, \text{CP-even}},
\]

and (here \( A, B \) run over \( 1 \cdots 8 \))

\[
k_{AB}(\phi) = -\frac{2 g^{\mu\nu} g^{\sigma\rho}}{d^2} \frac{\delta^2 \mathcal{L}_{\text{SMEFT}}}{\delta G^A_{\mu\sigma} \delta G^B_{\nu\rho}} \bigg|_{\mathcal{L}(\alpha,\beta\ldots) \to 0, \text{CP-even}}.
\]

We also have

\[
k_{I,J}(\phi) = \frac{g^{\mu\rho} g^{\nu\sigma}}{2d^2} \frac{\delta^3 \mathcal{L}_{\text{SMEFT}}}{\delta (D_\mu \phi)^I \delta (D_\nu \phi)^J W^A_{\rho\sigma}} \bigg|_{\mathcal{L}(\alpha,\beta\ldots) \to 0},
\]

and

\[
f_{ABC}(\phi) = \frac{g^{\mu\rho} g^{\sigma\alpha} g^{\beta\mu}}{3d^3} \frac{\delta^3 \mathcal{L}_{\text{SMEFT}}}{\delta W^A_{\mu\rho} \delta W^B_{\nu\sigma} \delta W^C_{\alpha\beta}} \bigg|_{\mathcal{L}(\alpha,\beta\ldots) \to 0, \text{CP-even}},
\]

\[
k_{ABC}(\phi) = \frac{g^{\mu\rho} g^{\sigma\alpha} g^{\beta\mu}}{3d^3} \frac{\delta^3 \mathcal{L}_{\text{SMEFT}}}{\delta G^A_{\mu\rho} \delta G^B_{\nu\sigma} \delta G^C_{\alpha\beta}} \bigg|_{\mathcal{L}(\alpha,\beta\ldots) \to 0, \text{CP-even}}.
\]

We also define the fermionic connections

\[
Y_{\psi^1}^{\psi^1}(\phi_I) = \frac{\delta \mathcal{L}_{\text{SMEFT}}}{\delta (\psi^1_{\mu\nu} \psi^1_{\nu\mu})} \bigg|_{\mathcal{L}(\alpha,\beta\ldots) \to 0}, \quad L_{\psi,pr}^{\psi,pr} = \frac{\delta^2 \mathcal{L}_{\text{SMEFT}}}{\delta (D_\mu \phi)^I \delta (\bar{\psi}^I_{\mu\nu} \tau_A \mu^A \psi^I_{\nu\mu})} \bigg|_{\mathcal{L}(\alpha,\beta\ldots) \to 0},
\]

and

\[
d_{\psi^1,pr}^{\psi,pr} = \frac{\delta^2 \mathcal{L}_{\text{SMEFT}}}{\delta (\psi^1_{\mu\nu} \sigma_{\mu\nu} \psi^1_{\nu\mu})} \bigg|_{\mathcal{L}(\alpha,\beta\ldots) \to 0}.
\]

The explicit form of these field space metrics and connects are given in the core geoSMEFT paper\(^\text{25}\) to all orders in \( \bar{v}_T / \Lambda \). The remarkable compactness of the all orders form of these objects in a key expansion parameter for precision SMEFT phenomenology, is due to the fact that they are functions of arbitrary numbers of Higgs field emissions, and symmetry generators acting on these fields. The completeness relations of the \( SU(2) \times U(1) \) algebras reduce the complexity of these objects to a compact form whose complexity usually saturates around mass dimension eight – when considering operator expansions of the metrics and tensors. This can also be seen in Table 1.

### 1.2 Mass eigenstate transformations at all orders in \( \bar{v}_T / \Lambda \)

The utility of understanding the SMEFT through the field space geometric quantities present is multi-fold. For example, one can immediately confirm that the transformation at all orders in \( \bar{v}_T / \Lambda \) between weak and mass eigenstates can be compactly expressed as

\[
\alpha^A = U^A_C \beta^C, \quad \psi^{A,\mu} = U^A_C A^{C,\mu}, \quad \phi^J = V^K_J \Phi^K,
\]
Table 1: Counting of operators contributing to two- and three-point functions from Hilbert series. $N_f$ corresponds to the number of fermion generations.

| Field space connection | Mass Dimension |
|------------------------|----------------|
| $h_{1J}(\phi) (D_\mu \phi)^J (D^\mu \phi)^J$ | 6  | 8  | 10 | 12 | 14 |
| $g_{AB}(\phi) W^A_{\mu \nu} V^{B,\mu \nu}$ | 2  | 2  | 2  | 2  | 2  |
| $k_{1J}(\phi) (D_\mu \phi)^J (D^\nu \phi)^J W^{A,\mu \nu}$ | 3  | 4  | 4  | 4  | 4  |
| $f_{ABC}(\phi) W^A_\mu W^{B,\sigma \rho} W^{C,\mu \rho}_{\sigma \nu}$ | 1  | 2  | 2  | 2  | 2  |
| $Y^{u}_{pr}(\phi) Q u + $ h.c. | $2 N_f^2$ | $2 N_f^2$ | $2 N_f^2$ | $2 N_f^2$ | $2 N_f^2$ |
| $Y^{d}_{pr}(\phi) Q d + $ h.c. | $2 N_f^2$ | $2 N_f^2$ | $2 N_f^2$ | $2 N_f^2$ | $2 N_f^2$ |
| $Y^{\phi}_{pr}(\phi) L e + $ h.c. | $2 N_f^2$ | $2 N_f^2$ | $2 N_f^2$ | $2 N_f^2$ | $2 N_f^2$ |
| $d^u_{pr}(\phi) L \sigma_{\mu \nu} e W^A_{\mu \nu} + $ h.c. | $4 N_f^2$ | $6 N_f^2$ | $6 N_f^2$ | $6 N_f^2$ | $6 N_f^2$ |
| $d^d_{pr}(\phi) Q \sigma_{\mu \nu} u W^A_{\mu \nu} + $ h.c. | $4 N_f^2$ | $6 N_f^2$ | $6 N_f^2$ | $6 N_f^2$ | $6 N_f^2$ |
| $d^\phi_{pr}(\phi) Q \sigma_{\mu \nu} d W^A_{\mu \nu} + $ h.c. | $4 N_f^2$ | $6 N_f^2$ | $6 N_f^2$ | $6 N_f^2$ | $6 N_f^2$ |
| $L^{\nu}_{pr,A}(\phi) (D_\mu \phi)^J (\tilde{\psi}_{p,R} \gamma_\mu \tau_A \psi_{r,R})$ | $N_f^2$ | $N_f^2$ | $N_f^2$ | $N_f^2$ | $N_f^2$ |
| $L^{\nu}_{pr,A}(\phi) (D_\mu \phi)^J (\tilde{\psi}_{p,L} \gamma_\mu \tau_A \psi_{r,L})$ | $2 N_f^2$ | $4 N_f^2$ | $4 N_f^2$ | $4 N_f^2$ | $4 N_f^2$ |

where in the SM limit $\alpha^A = \{g_2, g_2, g_1\}$ and $W^A = \{W_1, W_2, W_3, B\}$ and

\[
\beta^C = \left\{ \frac{g_2 (1 - i)}{\sqrt{2}}, \frac{g_2 (1 + i)}{\sqrt{2}}, \sqrt{2g_1^2 - 2g_2^2 + s_\theta^2}, \frac{2g_1 g_2}{\sqrt{g_1^2 + g_2^2}} \right\}, \quad (10)
\]

\[
A^C = (W^+, W^-, Z, A). \quad (11)
\]

Here $\phi^J = \{\phi_1, \phi_2, \phi_3, \phi_4\}$, $\Phi^K = \{\Phi^-, \Phi^+, \chi, h\}$ for the scalar fields with normalisation

\[
H(\phi_1) = \frac{1}{\sqrt{2}} \left[ \phi_2 + i \phi_3 \right], \quad \tilde{H}(\phi_1) = \frac{1}{\sqrt{2}} \left[ -\phi_2 + i \phi_3 \right]. \quad (12)
\]

Note $\phi_4$ is expanded around the vacuum expectation value with the replacement $\phi_4 \rightarrow \phi_4 + \tilde{\nu}_T$ and

\[
U_{BC} = \begin{bmatrix}
\frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} & 0 & 0 \\
\frac{1}{\sqrt{2}} & -\frac{1}{\sqrt{2}} & 0 & 0 \\
0 & 0 & c_\theta & s_\theta \\
0 & 0 & -s_\theta & c_\theta
\end{bmatrix}, \quad V_{JK} = \begin{bmatrix}
-\frac{i}{\sqrt{2}} & \frac{i}{\sqrt{2}} & 0 & 0 \\
\frac{i}{\sqrt{2}} & \frac{i}{\sqrt{2}} & 0 & 0 \\
0 & 0 & 1 & 0 \\
0 & 0 & 0 & 1
\end{bmatrix}.
\]

We are using a short-hand notation for the transformation matrices that lead to the canonically normalized mass eigenstate fields

\[
U^A_C = \sqrt{g}^{AB} U_{BC}, \quad Y^I_K = \sqrt{h}^{IJ} V_{JK}.
\]

Here $\sqrt{g}^{AB}$ and $\sqrt{h}^{IJ}$ are square-root metrics, which are understood to be matrix square roots of the expectation value $\langle \cdot \rangle$ of the field space connections for the bilinear terms in the SMEFT. Also note $\sqrt{h}^{IJ} \sqrt{h}_{JK} = \delta^I_K$ and $\sqrt{g}^{AB} \sqrt{g}_{BC} = \delta^A_B$. The rotation angles $c_\theta, s_\theta$ are functions of $\alpha_A$ and $\langle g^{AB} \rangle$ and are defined geometrically in the core geoSMEFT paper.

1.3 The $h\bar{A}A$ amplitude perturbation

One can also express amplitude perturbations at all orders in the $\tilde{\nu}_T/\Lambda$ expansion of the theory. For example, the effective coupling of $h-\gamma-\gamma$, including the tower of $\tilde{\nu}_T^2/\Lambda^2$ corrections, is given...
by

\[ \langle h | A(p_1) A(p_2) | \rangle = -\langle h | A^{\mu\nu} A_{\mu\nu} | \rangle \frac{\sqrt{\hbar}}{4} \left[ \left( \frac{\delta g_{33}(\phi)}{\delta \phi_4} \right) \frac{\bar{e}^2}{g_2^2} + 2 \left( \frac{\delta g_{34}(\phi)}{\delta \phi_4} \right) \frac{\bar{e}^2}{g_1 g_2} + \left( \frac{\delta g_{44}(\phi)}{\delta \phi_4} \right) \frac{\bar{e}^2}{g_1^2} \right], \]

where \( \bar{e} = g_2 (s_\theta \sqrt{g}^3 + c_\theta \sqrt{g}^4) = g_1 (c_\theta \sqrt{g}^4 + s_\theta \sqrt{g}^3) \).

### 1.4 \( W, Z \) couplings to \( \bar{\psi} \psi \)

The mass eigenstate coupling of the \( Z \) and \( W \) to \( \bar{\psi} \psi \) are obtained by summing over more than one field space connection. For couplings to fermion fields of the same chirality, the sum is over \( L^{\psi, pr}_{1, A} \) and the modified \( \bar{\psi} \gamma_{\mu} \gamma_5 \psi \), that includes the tower of SMEFT corrections in \( U^A \). A compact expression for the mass eigenstate connection is

\[ -A^A_{\mu}(\bar{\psi}_{1} \gamma_{\mu} \gamma_5 \psi_{1}) \delta_{pr} + A_{C, \mu}(\bar{\psi}_{1} \gamma_{\mu} \sigma_{A} \psi_{1}) \langle L^{\psi, pr}_{1, A} \rangle (-\gamma_{\mu} T_{C, A} \bar{\psi}_{1}), \]  

(13)

where the fermions are in the weak eigenstate basis. Rotating the fermions to the mass eigenstate basis is straightforward, where the \( V_{CKM} \) and \( U_{PMNS} \) matrices are introduced as usual. Expanding out to make the couplings explicit, the Lagrangian effective couplings for \( \{ Z, A, W^\pm \} \) are

\[ \langle Z | \bar{\psi}_p \psi_r \rangle = \frac{g_2}{2} \bar{\psi}_p f_Z \left[ (2 s_{\theta}^2 Q_{\psi} - \sigma_3) \delta_{pr} + \sigma_3 \bar{v}_T \langle L^{\psi, pr}_{3, 3} \rangle + \bar{v}_T \langle L^{\psi, pr}_{3, 4} \rangle \right] \psi_r, \]  

(14)

\[ \langle A | \bar{\psi}_p \psi_r \rangle = -\bar{\psi}_p f_A (Q_{\psi} \delta_{pr} \psi_r), \]  

(15)

\[ \langle W^\pm | \bar{\psi}_p \psi_r \rangle = -\frac{g_2}{2} \bar{\psi}_p (f_{W^\pm}) T^\pm \left[ (\delta_{pr} \bar{v}_T \langle L^{\psi, pr}_{1, 1} \rangle \pm i \bar{v}_T \langle L^{\psi, pr}_{1, 2} \rangle) \right] \psi_r. \]  

(16)

The last expressions simplify due to SU(2)\(_L\) gauge invariance. Similarly the SMEFT has the right-handed \( W^\pm \) couplings to (weak eigenstate) quark fields.

\[ \langle W^\mu_+ | d_r \bar{u} \rangle = \bar{v}_T \langle L^{ud, pr}_{1} \rangle \frac{g_2}{2} \bar{u}_p f_{W^+} d_r, \quad \langle W^\mu_- | d_r \bar{u} \rangle = \bar{v}_T \langle L^{ud, pr}_{1} \rangle \frac{g_2}{2} \bar{u}_p f_{W^-} d_r. \]

Using these expressions, EWPD was analysed for the first time to dimension eight. Dirac masses, mixings, and CP-violation parameter were directly also derived in the geoSMEFT. The geoSMEFT has also been used to calculate consistent expressions for \( \sigma(G G \rightarrow h) \) and \( \Gamma(h \rightarrow G G) \) to \( O(1/A^4) \) between simultaneously developing the corresponding amplitude results to \( O(1/16\pi^2 A^2) \) using the background field method and it was noted that these two sub-leading terms in the SMEFT are not independent.

It is interesting that the loop and operator mass dimension series expansions are cross correlated in order to maintain gauge invariance in the SMEFT, i.e. satisfy the corresponding Ward identity relations between n-point functions order by order in the 1/A power counting in the theory. This fact is explicitly demonstrated in some examples at one loop in the literature. It should also be intuitive that these expansions are not independent in general, as gauge fixing acts on fields that themselves are redefined order by order in the operator expansion of the theory when the Higgs takes on a vacuum expectation value. For this basic reason, strong general claims on gauge independence of results in the SMEFT at \( O(1/A^4) \) between n-point functions, simply due to squaring gauge (parameter) independent results at \( O(1/A^2) \) clearly require detailed mathematical demonstrations, and cannot be established only by assertion. The interested reader is encouraged to examine some recent debate on this issue present in the literature, from dramatically different perspectives.
Conclusions

The geoSMEFT is a consequence of the field redefinition invariance defining the SMEFT. It is a manifestly useful construction, allowing all order results in the $\tilde{\nu}_T/\Lambda$ expansion of the theory to be directly and immediately developed. In the literature, these theoretical techniques have already pushed calculations in the SMEFT to unprecedented theoretical precision in multiple processes, and many interesting formal questions have been investigated using field space geometry in recent months\textsuperscript{34,35,36}. Using field space geometries to ones advantage can be the key to efficient sub-leading order calculation in the SMEFT, and other field theories defined by field re-definitions. These results critically inform efforts to perform consistent fits to global data sets including a full set of $O(1/\Lambda^2)$ corrections, by defining the corresponding SMEFT corrections to the SM to $O(1/\Lambda^4)$ rapidly, completely and efficiently. This allows us to better understand the effect of neglected higher order terms when studying the global data set in the SMEFT. Fitting and constraining parameters introduced as corrections to the SM at $O(1/\Lambda^2)$ should be maximally informed by these results for robust bottom up SMEFT studies, which are expected to be a core legacy of the Large Hadron Collider physics program.

Acknowledgments

The author thanks the Villum Fonden project number 00010102. The Danish National Research Foundation (DNRF91) through the Discovery center, and DFF for support.

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