We study the contribution from dipole operators to one-loop Fierz identities and provide the resulting QCD and QED shifts to the tree-level relations for all four-fermion operators. The results simplify one-loop basis changes as well as matching computations and allow one to consistently eliminate operators from an operator basis which give rise to complications, e.g. traces involving $\gamma_5$.

I. INTRODUCTION

Fierz identities relate tensor products of Dirac structures in $D = 4$ space-time dimensions \(^1\). When using dimensional regularization, space-time is continued to $D = 4 - 2\varepsilon$ dimensions and the breaking of Fierz identities is handled by introducing evanescent operators \(^2\). \(^3\). Such evanescent structures, when inserted into divergent loop diagrams lead to finite contributions, which can be interpreted as one-loop shifts to the regular Fierz transformations. In Ref. \(^4\), such shifts are avoided by fixing the renormalization scheme (and therefore operator basis) to preserve tree-level Fierz relations; the Regularization Independent (RI) scheme. While this scheme avoids the issue of shifted Fierz identities, it has the downside that the now-fixed operator basis may not be the most convenient choice for the calculation at hand. For example, the basis may include operators that result in traces featuring $\gamma_5$ when computing loop-level diagrams that may be avoided in a different basis.

A different approach is to fix the operator basis which is most convenient to the calculation at hand and directly account for the modification of Fierz identities. This strategy has been successfully applied in the case of $\Delta F = 1$ processes \(^5\)\(^6\) and $\Delta F = 2$ \(^7\), as well as in matching computations involving Leptoquarks (LQs) \(^8\). Furthermore, when performing two-loop computations, such shifts are necessary to take into account since they describe the mixing of evanescent operators into the physical sector \(^9\). The shifts for all four-fermion operators, together with the corresponding renormalization constants, have been computed recently in \(^10\) in the generalized BMU scheme \(^11\). The effects of dipole operators on the one-loop QCD and QED shifts were not discussed in \(^11\), since mass effects were neglected. In Ref. \(^12\), dipole contributions to SMEFT basis changes were calculated using path-integral techniques. In this letter we give all one-loop shifts that result from electric and chromomagnetic dipole operators and therefore complete the findings in \(^11\) up to one-loop order in QCD and QED for effective four-fermion operators up to mass-dimension-six.

II. PROCEDURE

In order to obtain one-loop shifts to tree-level Fierz relations, one considers insertions of four-fermion operators into one-loop diagrams. The tree-level transformation from an operator basis, $\{O_i\}$ to its Fierz-conjugated counterpart, $\{\tilde{O}_i\}$ is given by

$$
\langle \tilde{O}_i \rangle^{(0)} = \mathcal{F}_{ij} \langle O_j \rangle^{(0)}. \tag{1}
$$

Denoting the one-loop amplitudes with $O$ and $\tilde{O}$ operator insertions

$$
\langle O_i \rangle^{(1)} = (\delta_{ij} + \tau_{ij}^{(1)}) \langle O_j \rangle^{(0)},
\langle \tilde{O}_i \rangle^{(1)} = (\delta_{ij} + \tilde{\tau}_{ij}^{(1)}) \langle \tilde{O}_j \rangle^{(0)}, \tag{2}
$$

respectively, the corresponding one-loop-corrected transformation is given by

$$
(\delta_{ij} + \tau_{ij}^{(1)}) \langle \tilde{O}_j \rangle^{(0)} = (\mathcal{F}_{ij} + \mathcal{F}_{ik} \tau_{kj}^{(1)} + \Delta_{ij}) \langle O_j \rangle^{(0)}, \tag{3}
$$

where $\Delta$ is the one-loop correction to the tree-level Fierz transformation. Using Eq. \(^1\) yields

$$
\Delta_{ij} = \tau_{ik}^{(1)} \mathcal{F}_{kj} - \mathcal{F}_{ik} \tau_{kj}^{(1)}. \tag{4}
$$

In the case of dipole operators, only penguin diagrams need to be considered in the computation of $\tau^{(1)}$ and $\tilde{\tau}^{(1)}$, since genuine vertex corrections do not produce dipole structures. Furthermore, one-loop shifts arise from $O(\epsilon)$

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{penguin_diagrams}
\caption{Two topologies of penguin diagrams with four-fermion operator insertions: open (left) and closed (right). Here, the fermion lines follow spinor indices.}
\end{figure}
terms in $D$-dimensional Dirac algebra multiplied by $1/\epsilon$ poles from loops, so only the divergent parts of the loop integrals need to be calculated. Therefore, considering the general operators

\begin{align}
Q_o &= \overline{\Gamma}_1^\alpha g_3 (\overline{\Gamma}_2^\alpha g_2), \\
Q_e &= \overline{\Gamma}_1^\alpha g_2 (\overline{\Gamma}_2^\alpha g_2),
\end{align}

together with our convention for the covariant derivative of a quark field

\[ D_\mu q = [\partial_\mu + ieQ_A \gamma_\mu + ig_T A^A \gamma_\mu] q, \]

we consider the following master formulae for the pole structure of the open and closed penguin diagrams shown in Fig. 1.

The resulting shifts are then obtained by considering all possible operator insertions into the above formulae and projecting the difference of the two onto the dipole operators. In the next section we will present the results for all possible four-fermion operators. We report the results in the basis used in Fig. 1, for which we adopt the following notation for four-quark operators

\[ V^{AB}_{q_1q_2q_3q_4} = (\overline{\Gamma}_1^\alpha \gamma_\mu P_A q_2^\alpha) (\overline{\Gamma}_3^\alpha \gamma_\mu P_B q_4^\alpha), \]

the semi-leptonic operators

\[ V^{AB}_{q_1q_2\ell_1\ell_2} = (\overline{\ell}_1^\gamma \gamma_\mu P_A q_2^\alpha) (\overline{\ell}_1^\gamma \gamma_\mu P_B \ell_2^\alpha), \]

and the four-lepton operators

\[ V^{AB}_{\ell_1\ell_2\ell_3\ell_4} = (\overline{\ell}_3^\gamma \gamma_\mu P_A \ell_4^\alpha) (\overline{\ell}_3^\gamma \gamma_\mu P_B \ell_4^\alpha), \]

where \( P_A, P_B = P_R/L = \frac{1}{2}(1 \pm \gamma_5) \) are the right/left fermion projection operators and \( \alpha, \beta \) are color indices.

In addition, we introduce the chromo- and electromagnetic dipole operators

\[ D^B_{q_1q_2G} = \frac{1}{g_5} m_q (\overline{\Gamma}_1^\alpha \gamma_\mu P_B T^A q_2^\alpha) G^A_{\mu\nu}, \]

\[ D^B_{f_1f_2\gamma} = \frac{1}{e} m_f (\overline{\ell}_1^\gamma \gamma_\mu P_B f_2^\alpha) F_{\mu\nu}. \]

Using this notation we find for the one-loop amplitudes with \( D^L_{q_1q_2G} \) and \( D^L_{f_1f_2\gamma} \) insertions with incoming momentum \( q \), depicted in Fig. 2,

\[ -i \frac{2m^2_q (\overline{\Gamma}_1^\alpha \gamma_\mu P_L T^A q_2) \equiv D^L_{q_1q_2G}, \]

\[ -i \frac{2m^2_f (\overline{\ell}_1^\gamma \gamma_\mu P_L f_2) \equiv D^L_{f_1f_2\gamma}, \]

and analogous expressions for different flavours and chiralities. We will report our results for the one-loop QCD and QED shifts in terms of the amplitudes \( D^L_{q_1q_2G} \) and \( D^L_{f_1f_2\gamma} \), respectively.

For the computation we will adopt the naive dimensional regularization (NDR) scheme. One complication arises however when tensor operators are considered, since closed-topology diagrams with tensor-operator insertions will include traces involving \( \gamma_5 \) which cannot be evaluated in NDR (see the last term of Eq. (7)). To treat these, we use the 't Hooft Veltman (HV) scheme.
TABLE I: One-loop QCD and QED shifts for VLR four-fermion operators. The shifts are given in units of $(\frac{e^2}{m^2})$ and $(\frac{e^2}{m^2})$ for QCD and QED respectively.

| Operator         | QCD shift | QED shift |
|------------------|-----------|-----------|
| $V_{LR}^{q_1\bar{q}_1 q_2 \bar{q}_2}$ | $\frac{m_g}{m_q} D_{q_1 q_2}^R$ | $A_{q_1} D_{q_1 q_2}^R$ |
| $\tilde{V}_{LR}^{q_1\bar{q}_1 q_2 \bar{q}_2}$ | 0 | $N_c A_{q_1} D_{q_1 q_2}^R$ |
| $V_{LR}^{q_1\bar{q}_1 q_2 \bar{q}_2}$ | $\frac{m_g}{m_q} D_{q_1 q_2}^R + \frac{m_g}{m_q} D_{q_2 q_1}^L$ | $A_{q_1} D_{q_1 q_2}^R + A_{q_1} D_{q_2 q_1}^L$ |
| $\tilde{V}_{LR}^{q_1\bar{q}_1 q_2 \bar{q}_2}$ | 0 | $N_c (A_{q_2} D_{q_1 q_1}^R + A_{q_1} D_{q_2 q_2}^R)$ |
| $V_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | $\frac{m_g}{m_q} D_{q_1 q_1}^R$ | $A_{q_1} D_{q_1 q_1}^R$ |
| $\tilde{V}_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | 0 | $N_c A_{q_1} D_{q_1 q_1}^R$ |
| $V_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | 2$\frac{m_g}{m_q} D_{q_1 q_1}^R$ | 2$A_{q_1} D_{q_1 q_1}^R$ |
| $\tilde{V}_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | 0 | 2$N_c A_{q_1} D_{q_1 q_1}^R$ |

TABLE II: One-loop QED shifts for VLR semi-leptonic and four-lepton operators. The shifts are given in units of $(\frac{e^2}{m^2})$ for QED, whereas the QCD shifts all vanish.

| Operator         | QCD shift | QED shift |
|------------------|-----------|-----------|
| $V_{LR}^{e_1\bar{e}_1 e_2 \bar{e}_2}$ | 0 | $N_c A_{q_1} D_{e_1 e_2}^R$ |
| $V_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | $A_{q_1} D_{q_1 q_1}^R$ | $A_{q_1} D_{q_1 q_1}^R$ |
| $V_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | 0 | $A_{q_2} D_{q_1 q_2}^R$ |
| $V_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | 0 | $A_{q_2} D_{q_1 q_2}^R + A_{q_1} D_{q_2 q_1}^L$ |
| $V_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | 0 | $A_{q_2} D_{q_1 q_1}^R + A_{q_1} D_{q_2 q_1}^L$ |
| $V_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | 0 | $A_{q_2} D_{q_1 q_1}^R + A_{q_1} D_{q_2 q_1}^L$ |
| $V_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | 0 | 2$A_{q_1} D_{q_1 q_1}^R$ |
| $V_{LR}^{q_1\bar{q}_1 q_1 \bar{q}_2}$ | 0 | 2$A_{q_1} D_{q_1 q_1}^R$ |

At the one-loop level, physical operators mix into the evanescent operators in Eqs. (29) and (31) with $1/\epsilon$ poles. Since only the tensor operator insertions into closed-penguin topologies feature such divergent mixings, these poles are not cancelled when subtracting Fierz-conjugated insertions.

They are, however, treated by enforcing tree-level relations $\mathcal{F}E_B^{q_1\bar{q}_1 G} = \mathcal{F}E_B^{q_2\bar{q}_2 G} = 0$, consistent with the requirement that evanescent operators vanish in the limit $D \to 4$. In general, these relations will also obtain loop-level shifts which become relevant at higher orders.

### III. RESULTS

In this section we report the obtained shifts from dipole operators to all four-fermion operators given in Eqs (10)-(24). We find no additional shifts to any four-fermion operators with $V^{AA}$ or $S^{AA}$ Dirac structures as well as to operators with all different fermion flavors. Therefore, we only report QCD and QED shifts for the $V^{LR}$ and $T^{LL}$ operators with at least two equal flavors. The results for the chirality-flipped operators $V^{LR}$ and $T^{RR}$ are obtained by the replacement $P_L \leftrightarrow P_R$. Results are presented in Tabs. III-V for vector operator insertions and in Tabs. III-V for tensor operator insertions. To shorten the notation we introduce the abbreviation:

$$A_f = \frac{m_f}{m_q} Q_f.$$  

Results for $S^{AA}$ and $T^{AA}$ operator insertions can be obtained from Tabs. III-V in conjunction with corresponding tree-level Fierz relations. Additionally, the scheme-dependencies arising from the definition of the evanescent operators in Eqs. (30) and (31) are set to zero, i.e. $a = b = 0$. Only tensor and scalar operators obtain

---

2 Here, one may also compute such diagrams in the Larin scheme, introducing the evanescent operator

$$E_{f_1 f_2 f_3}^{L/R} = \frac{m_f}{2e} \epsilon_{\mu\nu\rho\sigma} (f_1 \sigma^{\mu\nu} P_{L/R} f_2) F_{\rho\sigma} \pm (1 \pm e) D_{f_1 f_2 f_3}^{L/R},$$

and analogous for the gluon operator. The resulting shifts are identical at one-loop, though differences can arise at higher orders from one-loop evanescent operator insertions.
scheme-dependent shifts, e.g.\(^3\)

\[
\langle T_{q_1 q_2 q_3 q_4}^{AB} \rangle_{1}^{(1)} = \frac{-\alpha_s}{2\pi} \frac{m_{q_3}}{m_{q_4}} \langle T_{q_1 q_2 q_3 q_4}^{AA} \rangle_{1}^{(0)} + 6 \langle T_{q_1 q_2 q_3 q_4}^{AA} \rangle_{1}^{(1)} = \frac{-\alpha_s}{4\pi} \left(1 + a\right) \frac{m_{q_3}}{m_{q_4}} \langle T_{q_1 q_2 q_3 q_4}^{AA} \rangle_{1}^{(0)} .
\]

Note that operators with Dirac structure \(T^{AB}\) with \(A \neq B\) obtain non-vanishing scheme-dependent shifts from projections onto evanescent dipole operators and therefore should be included in the basis of operators in calculations where \(a, b \neq 0\).

\[\text{IV. APPLICATIONS}\]

As an application of the results presented in Sec.\(^4\) we discuss the simple example of computing one-loop corrections to the muonic electric dipole operator \(D_{\mu e}^{\gamma}\) in the Weak Effective theory (WET) below the electroweak scale. Matching conditions from the SM Effective Theory (SMEFT) onto the WET, known at tree-level\(^16\)\(^17\) and since recently also at one-loop\(^18\), as well as the one-loop WET renormalization group equations\(^19\)\(^20\) are traditionally given in the JMS basis introduced in\(^17\). This contains the semi-leptonic tensor operators of the form\(^4\)

\[O_{ijkl}^{T,RR} = (\bar{\tau}^i \sigma_{\mu\nu} P_R e^j \bar{\tau}^k \sigma_{\nu\mu} P_R e^l) .
\]

\[\text{V. CONCLUSIONS}\]

In this letter, we have expanded upon the results in Ref.\(^11\) by including the one-loop shifts in the four-
TABLE V: One-loop QED shifts for leptonic TLL operators. The shifts are given in units of \( \frac{\alpha}{\pi} \) for QED, whereas the QCD shifts all vanish.

| Operator | QCD shift | QED shift |
|----------|-----------|-----------|
| \( T_{12}^{LL} \) | 0 | \(-A_{1} D_{12}^{L} \) |
| \( T_{21}^{LL} \) | 0 | \( \frac{1}{2} A_{1} D_{12}^{L} + A_{1} D_{12}^{L} \) |
| \( T_{12}^{LL} \) | 0 | \(-A_{1} D_{12}^{L} \) |
| \( T_{21}^{LL} \) | 0 | \(-A_{1} D_{12}^{L} \) |

fermion Fierz relations including QCD and QED dipole operators.

We illustrated the usefulness of our results by employing the one-loop shifts in the computation of a one-loop matrix element. This allowed us to use a simpler operator basis that avoided complications involving traces including \( \gamma_{5} \), while still using the known matching results for the JMS basis. The one-loop Fierz transformations can also be used in matching calculations, where the model matches onto a basis which needs to be Fierz-conjugated. This is the case, for instance, in LQ models, where the matching is performed in the LQ basis and the results are then transformed into the SM basis using Fierz relations.

Another interesting application is to include Fierz-conjugation of four-fermion operators using shifts in codes such as abc_eft [24]. One could additionally extend this work to include two-loop shifts, which are relevant for two-loop dipole calculations, where both tensor and scalar operator insertions into closed-penguin topologies lead to problematic traces involving \( \gamma_{5} \). Operator insertions with at least one different fermion flavor can be exchanged for their open-penguin counterparts using Fierz relations along with corresponding shifts.

VI. NOTE ADDED

The present work resulted after common discussions with the authors of [13]. The results presented in this article are complementary to the ones in Ref. [13] in the sense that they are applicable in the WET below the EW scale, whereas in [13] the SMEFT together with the full SM gauge group was considered.

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Appendix A: Tree-Level Fierz Relations

In this appendix, we collect the tree-level Fierz relations necessary for the calculation of the one-loop shifts arising from dipole operators. Using the notation of Eqs. (1024) with \( A, B \in \{ R, L \} \), we use the four-quark operator relations

\[
V_{AA}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow V_{AA}^{q_{1}q_{2}q_{3}q_{4}} ,
\]

\[
V_{AB}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -2S_{AB}^{q_{1}q_{2}q_{3}q_{4}} \quad (A \neq B) ,
\]

\[
S_{AA}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -\frac{1}{2} S_{AB}^{q_{1}q_{2}q_{3}q_{4}} - \frac{1}{8} T_{AA}^{q_{1}q_{2}q_{3}q_{4}} ,
\]

\[
S_{AB}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -\frac{1}{2} V_{AA}^{q_{1}q_{2}q_{3}q_{4}} \quad (A \neq B) ,
\]

\[
T_{AA}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -6S_{AA}^{q_{1}q_{2}q_{3}q_{4}} + \frac{1}{2} T_{AA}^{q_{1}q_{2}q_{3}q_{4}} ,
\]

semi-leptonic operator relations

\[
V_{AA}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow V_{AA}^{q_{1}q_{2}q_{3}q_{4}} ,
\]

\[
V_{AB}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -2S_{AB}^{q_{1}q_{2}q_{3}q_{4}} \quad (A \neq B) ,
\]

\[
S_{AA}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -\frac{1}{2} S_{AA}^{q_{1}q_{2}q_{3}q_{4}} - \frac{1}{8} T_{AA}^{q_{1}q_{2}q_{3}q_{4}} ,
\]

\[
S_{AB}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -\frac{1}{2} V_{AA}^{q_{1}q_{2}q_{3}q_{4}} \quad (A \neq B) ,
\]

\[
T_{AA}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -6S_{AA}^{q_{1}q_{2}q_{3}q_{4}} + \frac{1}{2} T_{AA}^{q_{1}q_{2}q_{3}q_{4}} ,
\]

four-lepton operator relations

\[
V_{AA}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow V_{AA}^{q_{1}q_{2}q_{3}q_{4}} ,
\]

\[
V_{AB}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -2S_{AB}^{q_{1}q_{2}q_{3}q_{4}} \quad (A \neq B) ,
\]

\[
S_{AA}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -\frac{1}{2} S_{AA}^{q_{1}q_{2}q_{3}q_{4}} - \frac{1}{8} T_{AA}^{q_{1}q_{2}q_{3}q_{4}} ,
\]

\[
S_{AB}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -\frac{1}{2} V_{AA}^{q_{1}q_{2}q_{3}q_{4}} \quad (A \neq B) ,
\]

\[
T_{AA}^{q_{1}q_{2}q_{3}q_{4}} \rightarrow -6S_{AA}^{q_{1}q_{2}q_{3}q_{4}} + \frac{1}{2} T_{AA}^{q_{1}q_{2}q_{3}q_{4}} ,
\]
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