Vol. 50 (2019) *Acta Physica Polonica B* No 11

EW EXTENSIONS OF THE STANDARD MODEL AND THE SUPPRESSION OF FERMIONIC EFT INTERACTIONS*

Fernando Alvarado

Instituto de Física Corpuscular (Centro Mixto CSIC-UV)
Institutos de Investigación de Paterna
Apartado 22085, 46071 Valencia, Spain

Adolfo Guevara

Departamento de Física Atómica, Molecular y Nuclear
and Instituto Carlos I de Física Teórica y Computacional
Universidad de Granada, 18071 Granada, Spain

Juan José Sanz-Cillero

Departamento de Física Teórica and IPARCOS
Universidad Complutense de Madrid
Plaza de las Ciencias 1, 28040 Madrid, Spain

(Received November 26, 2019)

The low-energy effective field theory for electroweak (EW) interactions is studied here. It embeds the Standard Model (SM) as a particular limit and parametrizes new physics deviations. We discuss some experimental resonant diboson searches and four-fermion operator analyses that seem to push the new physics scale well over the TeV. On the other hand, the more precise oblique parameter determinations allow new physics resonances in the few TeV range. This apparent contradiction is easily solved by postulating a Standard Model extension that only couples directly to the bosonic degrees of freedom of the Standard Model but not to its fermions.

DOI:10.5506/APhysPolB.50.1937

1. Introduction

These proceedings are mainly based on the results of Refs. [1–7]. We will be discussing a series of experimental results that seems to point out that new physics states must be well over the TeV. However, when studied in

* Presented by J.J. Sanz-Cillero at the XLIII International Conference of Theoretical Physics "Matter to the Deepest", Chorzów, Poland, September 1–6, 2019.

deeper detail, one finds in the bibliography analyses on strongly interacting theories beyond the Standard Model (BSM) where resonances in the range of $M_R \sim 1$–3 TeV would not at present produce any experimentally measurable deviation with respect to the SM [4, 5].

LHC resonant diboson searches at 8 TeV and 13 TeV have pointed out lower-bounds of the order of a few TeV on the mass of the lightest spin-1 triplets (under the EW group). Figure 1 shows one of the strongest bounds, $M_V \gtrsim 4$ TeV [8] (see Refs. [3, 9] and references therein for further details). However, these spin-1 resonance analyses, in general, rely in two very specific benchmark points of the heavy vector triplet model [10]: HVT-$\mathrm{A}_{g_V=1}$ and HVT-$\mathrm{B}_{g_V=3}$. In all cases, the latter provides the more stringent mass bound. However, in spite of accounting for the production from all possible initial mechanisms (including EW gauge boson fusion) and having an almost 100% branching ratio into a diboson BB' pair (with B and B' being EW gauge bosons or a Higgs), this benchmark is actually dominated by the Drell–Yan resonance production $\bar{q}q' \to R$. Hence, this fermionic vertex is what this type of analyses are actually testing [3].

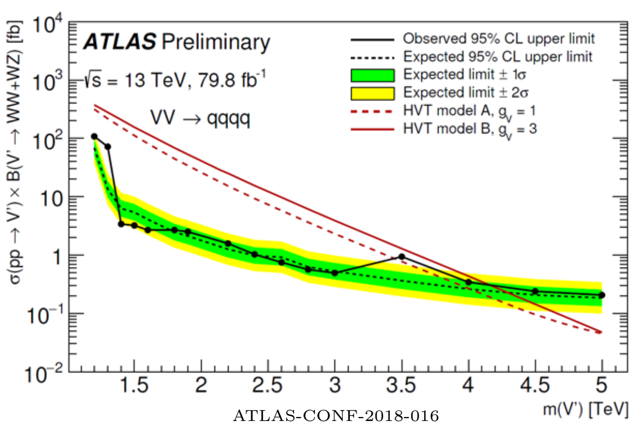

Fig. 1. This plot shows one of the most stringent WW, WZ and ZZ diboson resonant searches based on fat-jet reconstruction techniques [8]. The experimental results are compared to the benchmark models HVT-$\mathrm{A}_{g_V=1}$ and HVT-$\mathrm{B}_{g_V=3}$ [10], which assume a spin-1 EW triplet. This leads to the lower bound $M_V > 4.15$ TeV at the 95% C.L.

Another type of searches looks instead for a non-resonant increasing in the cross section due to possible four-fermion operators in the low-energy effective field theory (EFT). However, as one can see in Fig. 2, both collider analyses of dijet and dilepton events at the LHC and low-energy hadronic experiments have led to really strong bounds on the suppression of these fermionic EFT operators. The four-fermion terms in the EFT Lagrangian are found to be suppressed by $\Lambda \gtrsim 10$–20 TeV in the most stringent cases. Similar bounds were obtained at LEP and Tevatron. See Refs. [3, 11] and references therein.

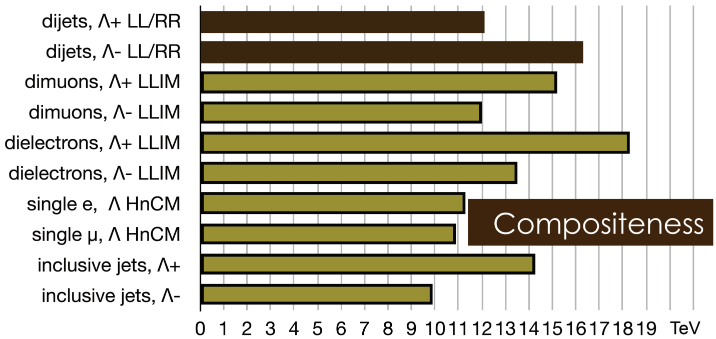

Fig. 2. Summary on four-fermion operators presented at ICHEP 2016 by the CMS Exotica Physics Group Summary. Bounds on four-fermion effective Lagrangians are expressed in terms of the sometimes-called "compositeness" scale Λ [12, 13].

These results are sometimes argued to imply that new-physics strongly coupled theories and composite particles must be well over the TeV. However, EW precision tests still allow resonances with masses in the few TeV range: in Fig. 3, one can observe that the experimental determinations [14, 15] of

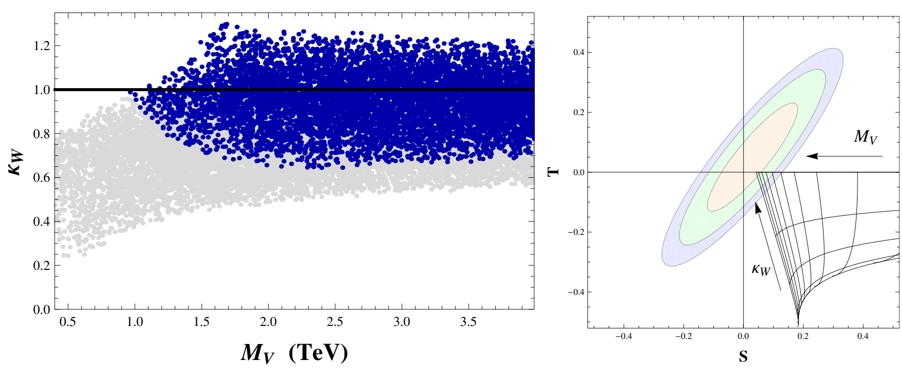

Fig. 3. (Color online) Left-hand side: Scatter plot for the 68% C.L. region, in the case when only the first VV–AA WSR is assumed [16, 17]. The black (dark blue) and light gray regions correspond, respectively, to $0.2 < M_V/M_A < 1$ and $0.02 < M_V/M_A < 0.2$. We consider $M_A > M_V > 0.4$ TeV in the plot. Right-hand side: NLO determinations of S and T, imposing the two WSRs [16, 17]. The grid lines correspond to M_V values from 1.5 to 6.0 TeV, at intervals of 0.5 TeV, and the hW^+W^- coupling $\kappa_W = 0, 0.25, 0.50, 0.75, 1$. The arrows indicate the directions of growing M_V and κ_W. The ellipses give the experimentally allowed regions at 68% (inner/orange), 95% (middle/green) and 99% (outer/blue) C.L.

the S and T oblique parameters [18, 19] lead to an $h \to W^+W^-$ coupling κ_W close to the SM one and spin-1 resonance masses with $M_V \gtrsim 1$ TeV (in less constraining theories, where only one VV–AA Weinberg sum rule, WSR, applies) or $M_V \gtrsim 4$ TeV (in more constraining scenarios with two Weinberg sum rules) [16, 17].

In Sec. 2, we remind the basic principles for the construction of the low-energy EFT based on the non-linear $\mathrm{SU(2)_L} \times \mathrm{SU(2)_R}$ representation of the EW Goldstone bosons, including only the SM particles. Section 3 incorporates heavier BSM states. The lightest multiplets of heavy resonances are included in the present discussion. We will see that the HEFT low-energy couplings (LECs) are typically $\mathcal{O}(10^{-3})$ in the case of resonances in the few TeV range. However, if so, one may have expected to detect these resonances in LHC resonant searches, which currently seem to exclude masses below 4 TeV. Nonetheless, these experimental results heavily rely on a large coupling of the SM fermions and the BSM resonances. In Refs. [4, 5, 20–22], it is shown that if these new states are only produced via intermediate EW gauge bosons, the cross section is small enough to remain undetected at present. We mention here two examples: production via WZ vector boson scattering (VBS) [4, 20–22], and the Drell–Yan (DY) resonant production thanks to a mixing term between the heavy resonances and the EW gauge bosons [5, 23]. Thus, in Sec. 4, we show that in a general model that couples the new physics only to the SM bosonic sector (EW gauge bosons and the scalar sector), one observes a strong suppression of new physics in SM fermion interactions. We note that this does not imply the absence of the latter interactions. Thus, one obtains, *e.g.*, four-fermion effective operators in the low-energy HEFT but with much smaller LECs. The same happens in this type of models with operators that violate custodial symmetry, such as those contributing to the oblique T parameter, which are very much suppressed. Some final conclusions are given in Sec. 5.

2. (Non-linear) HEFT (aka EWχL, aka EWET)

The low-energy effective theory under discussion is indistinctly denoted as EW effective theory (EWET), EW Chiral Lagrangian (EWχL) and Higgs Effective Field Theory (HEFT). This HEFT describing the interactions between the known SM particles at $E \ll 1$ TeV is based on three aspects:

1. **Symmetries:**

   — **SM gauge symmetries (exact in the SM):** the EFT action is invariant under the EW group $G_{\mathrm{SM}} = \mathrm{SU(2)_L} \times \mathrm{U(1)}_Y$, which gets spontaneously broken down to $H_{\mathrm{SM}} = \mathrm{U(1)_{EM}} \subset G_{\mathrm{SM}}$.

— **SM scalar sector global symmetry (approximate in the SM):** the EFT action will be based on the global chiral symmetry $\mathcal{G} = \mathrm{SU}(2)_{\mathrm{L}} \times \mathrm{SU}(2)_{\mathrm{R}} \times \mathrm{U}(1)_{B-L} \supset G_{\mathrm{SM}}$ of the SM scalar sector, which gets spontaneously broken down to its custodial subgroup $\mathcal{H} = \mathrm{SU}(2)_{\mathrm{L+R}} \times \mathrm{U}(1)_{B-L} \supset H_{\mathrm{SM}}$. However, this symmetry is explicitly broken in the SM by the left↔right asymmetry between the two SU(2) sectors as only one generator of the $\mathrm{SU}(2)_{\mathrm{R}}$ component of $\mathcal{G}$ is gauged (B^μ), on the contrary to $\mathrm{SU}(2)_{\mathrm{L}}$, which has its three generators gauged ($W^\mu_{1,2,3}$). Notice that this symmetry becomes again exact in the gauge and scalar sectors of the SM in the limit $g' \to 0$. The other source of custodial symmetry breaking in the SM is the mass (and Yukawa coupling) asymmetry between the top-type and bottom-type fermions in the SU(2) doublet. The EFT will implement this soft explicit breaking in exactly the same way it occurs in the SM. For more details, see Ref. [2, 3, 24].

2. **SM content:** the EFT Lagrangian contains the G_{SM} gauge bosons. The Higgs particle h is incorporated as a chiral singlet under $\mathcal{G}$, while the EW Goldstones ω^a from the $\mathcal{G}/\mathcal{H}$ spontaneous symmetry breaking are non-linearly realized via the unitary matrix $U(\omega) = 1 + i\omega_a\sigma_a/v + \mathcal{O}(\omega^2)$, where ω_a stands for the triplet of EW Goldstones — which transforms non-linearly under $\mathcal{G}$ — σ_a stands for the Pauli matrices and $v \simeq 246$ GeV stands for the Higgs vacuum expectation value (v.e.v.). There is an additional freedom in the choice of the unitary matrix, which is usually taken in the exponential form $U(\omega) = \exp\{i\omega_a\sigma_a/v\}$ or the spherical representation of the $\mathcal{G}/\mathcal{H}$ coset, $U(\omega) = \sqrt{1-\omega_a\omega_a/v^2} + i\omega_a\sigma_a/v$ (see, *e.g.*, Refs. [25, 26]). The HEFT also includes all the SM fermion fields (quarks and leptons) as light degrees of freedom (d.o.f.).

3. **Chiral power counting:** the different EFT building blocks carry an infrared "chiral" scaling [2, 24, 27, 28] where boson fields scale as $\mathcal{O}(p^0)$, and covariant derivatives, gauge and Yukawa couplings and SM particle masses scale as $\mathcal{O}(p)$. Regarding SM fermion bilinears, although a naive dimensional analysis [2, 28] would assign them a scaling $\mathcal{O}(p)$, the EFT stemming from this power counting does not agree with the phenomenology, as four-fermion operators should be then part of the leading order (LO) EFT Lagrangian, they would be suppressed by a scale $\Lambda \sim v$ and they should already have been experimentally observed. For this reason, fermion bilinears $\bar\psi\Gamma\psi$ in the interaction Lagrangian need to be assigned a higher scaling, at least the $\mathcal{O}(p^2)$ usually assumed [2, 3], which accounts for an implicit ad-

ditional weak coupling suppression that must always go with the SM fermion interactions[1]. Additionally, as custodial symmetry is explicitly — but softly — broken in the SM, we assign an additional suppression $\mathcal{O}(p)$ to the building blocks and effective operators that break custodial symmetry. This is necessary in order to avoid large violations of this symmetry which would lead to large contributions to the oblique T parameter [18, 19] which would contradict the experimental data [14, 15].

This allows one to sort out the HEFT Lagrangian in terms of effective operators with increasing order p^d [29–34]:

$$\mathcal{L}_{\rm HEFT} \;=\; \sum_{d\geq 2} \mathcal{L}^{(d)}_{\rm HEFT}\,. \tag{2.1}$$

The LO Lagrangian occurs at $\mathcal{O}(p^2)$ and contains indeed the SM one, this is, $\mathcal{L}_{\rm SM}\subset \mathcal{L}^{(2)}_{\rm HEFT}$. It has the general form [3]:

$$\begin{aligned}\mathcal{L}^{(2)}_{\rm HEFT} \;=\; & \sum_\xi \left[i\,\bar\xi\gamma^\mu d_\mu \xi - v\left(\,\bar\xi_{\rm L}\,\mathcal{Y}\,\xi_{\rm R} + {\rm h.c.}\right)\right] \\ & -\frac{1}{2g^2}\left\langle \hat W_{\mu\nu}\hat W^{\mu\nu}\right\rangle_2 - \frac{1}{2g'^2}\left\langle \hat B_{\mu\nu}\hat B^{\mu\nu}\right\rangle_2 - \frac{1}{2g_s^2}\left\langle \hat G_{\mu\nu}\hat G^{\mu\nu}\right\rangle_3 \\ & +\frac{1}{2}\partial_\mu h\,\partial^\mu h - V(h/v) + \frac{v^2}{4}\,\mathcal{F}_u\,\langle u_\mu u^\mu\rangle_2\,,\end{aligned} \tag{2.2}$$

where ξ describes the SM fermion fields, d_μ is the covariant derivative with the corresponding $G_{\rm SM}$ gauge connection and $\langle\dots\rangle_2$ and $\langle\dots\rangle_3$ refer to traces under SU(2) and SU(3). In the second line, we have the Yang–Mills Lagrangian for $G_{\rm SM}$. In the third line, one has the Higgs potential and the interaction with the EW Goldstone bosons parametrized by the tensors $u_\mu = iu(D_\mu U)^\dagger u = -\partial_\mu\omega_a\sigma_a/v+\dots$, with $U=u^2$. Since the Higgs h is a chiral singlet, the Yukawa $\mathcal{Y}$ and the factor $\mathcal{F}_u$ can be arbitrary functions of h with an analytical expansion around $h=0$, with the HEFT action

[1] In the case of the mass term, the coupling in front of the $-\bar\psi\psi$ operator, the fermion mass, is experimentally known and of the order of the external momenta, $m_\psi=\mathcal{O}(p)$, so there is no need to assume a further implicit suppression of the $\bar\psi\Gamma\psi$ bilinears in this term. On the other hand, in the $h\bar\Psi\psi$ Yukawa interaction, the naive dimensional analysis tells us that this interaction operator would be $\mathcal{O}(p)$. This would lead to predictions in contradiction with the experiment. On the other hand, counting this $h\bar\Psi\psi$ operator as $\mathcal{O}(p^2)$, like the other terms in the LO Lagrangian, leads to a consistent agreement. Note that in the SM, this operator is always accompanied by an explicit weak factor, the Yukawa coupling, which increases the naive scaling of the Yukawa interaction and makes it $\mathcal{O}(p^2)$.

still remaining invariant under $\mathcal{G}$ transformations. For instance, symmetry invariance allows the general structure $\mathcal{F}_u = 1 + 2ah/v + bh^2/v^2 + \mathcal{O}(h^3)$, with the SM corresponding to the precise values $a_{\rm SM} = b_{\rm SM} = 1$. Further details on the notation can be found in Refs. [2, 3].

The next-to-leading order (NLO) effective Lagrangian is $\mathcal{O}(p^4)$ and has the general form of [3]

$$\mathcal{L}^{(4)}_{\rm HEFT} = \sum_{i=1}^{12} \mathcal{F}_i(h/v)\, \mathcal{O}_i + \sum_{i=1}^{3} \widetilde{\mathcal{F}}_i(h/v)\, \widetilde{\mathcal{O}}_i + \sum_{i=1}^{8} \mathcal{F}_i^{\psi^2}(h/v)\, \mathcal{O}_i^{\psi^2}$$
$$+ \sum_{i=1}^{3} \widetilde{\mathcal{F}}_i^{\psi^2}(h/v)\, \widetilde{\mathcal{O}}_i^{\psi^2} + \sum_{i=1}^{10} \mathcal{F}_i^{\psi^4}(h/v)\, \mathcal{O}_i^{\psi^4} + \sum_{i=1}^{2} \widetilde{\mathcal{F}}_i^{\psi^4}(h/v)\, \widetilde{\mathcal{O}}_i^{\psi^4}\,. \qquad (2.3)$$

This is the CP even effective Lagrangian. Operator with (without) tilde are P odd (even) and contain purely bosonic operators $(\mathcal{O}_i, \widetilde{\mathcal{O}}_i)$, operators with one fermion bilinear $(\mathcal{O}_i^{\psi^2}, \widetilde{\mathcal{O}}_i^{\psi^2})$ and four-fermion operators $(\mathcal{O}_i^{\psi^4}, \widetilde{\mathcal{O}}_i^{\psi^4})$. The full list of the 30 (8) P-even (P-odd) effective operators can be found in Ref [3].

In principle, one may construct more and more complex operators of higher and higher chiral dimension. However, the importance of this classification of the effective operators in terms of its chiral dimension p^d is that amplitudes have then the general form (*e.g.*, for a $2 \to 2$ scattering) of

$$\mathcal{M} \approx \underbrace{\frac{p^2}{v^2}}_{\text{LO (tree)}} + \left(\underbrace{\frac{\mathcal{F}_k^r(\mu)\, p^4}{v^2}}_{\text{NLO (tree)}} - \underbrace{\frac{\Gamma_k p^4}{16\pi^2 v^2} \ln\frac{p^2}{\mu^2} + \ldots}_{\text{NLO (1-loop)}} \right) + \mathcal{O}\left(p^6\right), \qquad (2.4)$$

where p stands for external momenta in the process or any equivalent soft scale of the EFT like, *e.g.*, the masses of the SM particles. The LO contribution, $\mathcal{O}(p^2)$, is provided by all the tree-level diagrams with $\mathcal{L}^{(2)}_{\rm HEFT}$ vertices. NLO corrections arise at $\mathcal{O}(p^4)$. The one-loop diagrams with $\mathcal{L}^{(2)}_{\rm HEFT}$ contribute at $\mathcal{O}(p^4)$ with non-analytic terms like, *e.g.*, the generic logarithm in the expression above, a finite $\ln\mu^2$ dependence and possible ultraviolet (UV) divergences [35–40]. These UV divergences are renormalized through the other type of $\mathcal{O}(p^4)$ contributions arising in the computation: tree-level diagrams with $\mathcal{L}^{(2)}_{\rm HEFT}$ vertices and one $\mathcal{L}^{(4)}_{\rm HEFT}$ vertex, where the couplings $\mathcal{F}_k$ of these $\mathcal{L}^{(4)}_{\rm HEFT}$ operators renormalize the $\mathcal{O}(p^4)$ one-loop UV divergences. The dots stand for any other one-loop contribution at NLO that is UV finite and renormalization-scale-independent.

We note that in the case of the SM, this chiral expansion is simply the usual loop expansion: the suppression factors $\frac{p^2}{16\pi^2 v^2}$ do not have the naively expected power-like growing with the energy but correspond to the standard expansion parameters $\frac{p^2}{16\pi^2 v^2} \sim \frac{g^{(\prime)\,2}}{16\pi^2}, \frac{\lambda^2_{\rm Fer}}{16\pi^2}, \frac{\lambda}{16\pi^2}$, which remain much smaller than one all the way up to the Planck scale (barring energies below the GeV for the QCD coupling constant).

3. HEFT + resonances: what might we expect?

The HEFT can be extended up to higher energies by incorporating the lightest multiplets of heavy resonances to the Lagrangian. Reference [3] incorporated the lightest multiplets of vector, axial vector, scalar, pseudoscalar and spin-1/2 fermion resonances. EW singlets, doublets and triplets and color singlets and octets were considered. These are all the resonances that may contribute at tree-level to the low-energy HEFT Lagrangian at NLO [3].

The high-energy Lagrangian can be sorted out in the (symbolic) form of

$$\begin{aligned} \mathcal{L}^{\rm HE}[R,\ell{\rm ight}] &= \mathcal{L}^{(2)}[\ell{\rm ight}] + \mathcal{L}_R[R,\ell{\rm ight}]\,, \\ \text{with } \mathcal{L}_R &= \mathcal{L}_R^{\rm Kin} + R\,\chi_R[\ell{\rm ight}] + \mathcal{O}\left(R^2\right)\,. \end{aligned} \tag{3.1}$$

The extended Lagrangian [2, 3] is based on the same symmetry breaking pattern $\mathcal{G}/\mathcal{H}$ and approximate symmetries of the HEFT. Only the interaction Lagrangian with one resonance field is needed for the contributions to the $\mathcal{O}(p^4)$ HEFT. The precise form of the $\chi_R[\ell{\rm ight}]$ tensors for the various resonances R are provided in Ref. [3], being all of them $\mathcal{O}(p^2)$.

In order to integrate out the heavy resonances and extract the low-energy effective Lagrangian, one needs first to solve the heavy resonance equations of motion (EoM) in terms of the light d.o.f. The low-energy limit of this classical solution has the form of

$$R_{c\ell}[\ell{\rm ight}] \;\sim\; \frac{1}{M_R^2}\chi_R[\ell{\rm ight}] \;+\; \mathcal{O}\left(\frac{p^4}{M_R^4}\right)\,. \tag{3.2}$$

We make then the substitution $R \to R_{c\ell}$ in the full high-energy action $\mathcal{L}^{\rm HE}$ [41]. This provides a non-local effective Lagrangian which, nonetheless, accepts at low energies an expansion in terms of an infinite series of local operators, where the first contribution arises at $\mathcal{O}(p^4)$ in the form of [2, 3, 41]

$$\Delta\mathcal{L}^{(4)}_{\rm HEFT} \;\sim\; \sum_R \frac{1}{M_R^2}\left(\chi_R[\ell{\rm ight}]\right)^2\,. \tag{3.3}$$

The various $\mathcal{O}(p^4)$ couplings of the low-energy EFT receive contributions proportional to the resonant couplings $(F_V, G_V, F_A\ldots)$ and suppressed by $1/M_R^2$. To illustrate this, let us show, for instance the contribution from a vector multiplet V and an axial vector A, both being EW triplets

$$\mathcal{L}_{\rm HEFT}^{(4)} = \frac{\mathcal{F}_1}{4}\langle f_+^{\mu\nu} f_{+\,\mu\nu} - f_-^{\mu\nu} f_{-\,\mu\nu}\rangle + \mathcal{F}_4\langle u_\mu u_\nu\rangle\langle u^\mu u^\nu\rangle + \dots \tag{3.4}$$

$$\text{with } \mathcal{F}_1 = -\frac{F_V^2}{4M_V^2} + \frac{F_A^2}{4M_A^2} = -\frac{v^2}{4}\left(\frac{1}{M_V^2} + \frac{1}{M_A^2}\right), \tag{3.5}$$

$$\mathcal{F}_4 = \frac{G_V^2}{4M_V^2} = \frac{v^2}{4}\left(\frac{1}{M_V^2} - \frac{1}{M_A^2}\right), \qquad \dots \tag{3.6}$$

where the $f_\pm^{\mu\nu}$ building blocks are P-even or P-odd covariant combinations of the W and B field-strength tensors [2]. In these equations, we are assuming that the underlying resonance theory only has P-even interactions. Otherwise, there are additional contributions [2, 3]. Finally, some UV completion hypotheses are assumed (like, *e.g.*, VV–SAA WSRs [42, 43]) which lead to phenomenological predictions like those on the right-hand side of Eqs. (3.5) and (3.6) [1, 41, 44]. Here, we are just showing a pair of selected coupling, $\mathcal{F}_1$ and $\mathcal{F}_4$ with $h = 0$ (also denoted as a_1 and a_4), related to the oblique S parameter [25, 45] and $WW \to WW$ scattering [4, 20–22], respectively. Figure 4 shows the values for these two LECs that derived from the experimental measurements of the S and T parameters for asymptotically-free BSM theories [18, 19], *i.e.*, accepting two WSRs [1] (black areas).

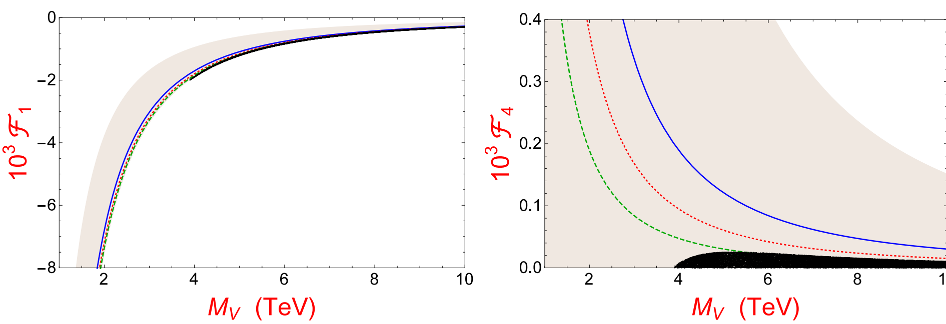

Fig. 4. (Color online) Predictions for the $\mathcal{O}(p^4)$ LECs $a_1 = \mathcal{F}_1(0)$ and $a_4 = \mathcal{F}_4(0)$ for asymptotically-free BSM theories (accepting two WSRs), as a function of M_V [1]. The light-shaded regions cover all possible values for $M_A > M_V$, while the blue (top solid), red (middle dotted) and green (bottom dahed) lines correspond to $M_V^2/M_A^2 = 0.8$, 0.9 and 0.95, respectively. The oblique S and T constraints restrict the allowed values to the dark areas.

What we would like to emphasize from these and similar other results in the bibliography [4, 5, 20–22, 25] is that vector triplets with a mass in the range of a few TeV are associated with LECs $|\mathcal{F}_j| \sim 10^{-3}$–$10^{-4}$ [1]. Recent LHC experimental analyses on WW and WZ VBS have been able to reduce the uncertainty in a_4 and a_5 ($\mathcal{F}_4$ and $\mathcal{F}_5$ for $h = 0$) down to

$|a_{4,5}| \lesssim 10^{-1}$, making use of large-radius jet substructure techniques [46] (see Fig. 5). These bounds are still very far away from the $\mathcal{O}(10^{-3})$ order of magnitude one would expect for resonance masses in the 1–3 TeV range[2]. LEP EW precision data on the S-parameter [14, 15] provides the best known LEC, $a_1 = \mathcal{F}_1(0)$, which reaches an uncertainty of $|a_1| \lesssim 10^{-3}$ [22]. Future e^+e^- colliders are expected to cut this uncertainty down by at least a factor five [15, 48].

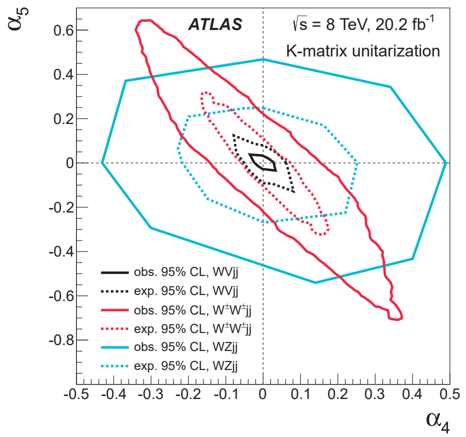

Fig. 5. Constraints on the $\alpha_4 = a_4 = \mathcal{F}_4(0)$ and $\alpha_5 = a_5 = \mathcal{F}_5(0)$ LECs from WZ and WW VBS analyses. Figure borrowed from Ref. [46].

4. Suppressing HEFT fermion operators: a simple model

At this point, one may wonder whether it is possible to reconcile the seemingly contradictory results from LEC analyses (with precisions in general far from the requested $|\mathcal{F}_j| \sim 10^{-3}$ for resonances in the TeV) and the diboson resonant and four-fermion operator searches (with lower bounds $M_R \gtrsim 4$ TeV and $\Lambda \gtrsim 10$–20 TeV, respectively). The simple solution proposed in Refs. [6, 7] was partly motivated by the previous analysis [23] for $q\bar{q}' \to R$ Drell–Yan production of spin-1 resonances R at the LHC. Therein, a quark and an antiquark merged into an EW gauge boson which induced a resonant signal through a kinetic mixing with a spin-1 EW triplet R. This mixing introduced, however, a very strong suppression of the fermionic branching ratios with respect to the dominant $R \to WW, WZ$ diboson

[2] The latest CMS study on VBS [47] leads to bounds on these couplings two orders of magnitude more precise, $|a_{4,5}| \lesssim 10^{-3}$ [20]. However, it must be taken with a grain of salt since this analysis does not include unitarization: the WW, WZ and ZZ scattering amplitudes eventually violate the unitarity bound and become unphysical in the TeV region under study. This unitarity violation, very likely, leads to an overestimate of the precision of these $a_{4,5}$ measurements [20].

channel. As it happened with more recent VBS and DY production studies [4, 5, 20–23], in order to generate the spin-1 resonance, one needs to first produce intermediate EW gauge bosons, suppressing the cross sections by additional factors of the EW gauge couplings. In Refs. [6, 7], we go one step further in this direction and extract the general impact on the HEFT Lagrangian, at low energies, stemming from not having a direct interaction between the SM fermions and the spin-1 resonances. Thus, we consider the vector and axial-vector resonance Lagrangians discussed in Sec. 3, but restricted to interactions with just the SM bosons, *i.e.*, disconnecting the SM fermion vertices [2, 3, 6, 7]

$$\begin{aligned} \mathcal{L}_V &= \mathcal{L}_V^{\rm Kin} + \langle V_{\mu\nu}\chi_V^{\mu\nu}\rangle_2 \,, \qquad (4.1)\\ \chi_V^{\mu\nu} &= \left(\frac{F_V f_+^{\mu\nu}}{2\sqrt{2}} + \frac{\widetilde{F}_V f_-^{\mu\nu}}{2\sqrt{2}}\right) + \frac{iG_V}{2\sqrt{2}}[u^\mu,u^\nu] + \frac{\widetilde{\lambda}_1^{hV}}{\sqrt{2}}\left((\partial^\mu h)u^\nu - (\partial^\nu h)u^\mu\right)\,, \end{aligned}$$

with an analogous Lagrangian for the axial vector EW triplet A (with appropriate replacements like, *e.g.*, $F_V V_{\mu\nu} f_+^{\mu\nu} \to F_A A_{\mu\nu} f_-^{\mu\nu}$, *etc.*) [6].

At low energies, this induces and effective Lagrangian of the form of

$$\Delta\mathcal{L}_{\rm HEFT}\Big|_R = \mathcal{L}_R\Big|_{R\to R_{c\ell}} = \mathcal{L}_{\rm HEFT}^{(4)} + \mathcal{L}_{\rm HEFT}^{(6)} + \ldots \qquad (4.2)$$

The LO, $\mathcal{O}(p^2)$, HEFT Lagrangian is not corrected by these resonances. At NLO, $\mathcal{O}(p^4)$, the V and A contribution has the form of [2, 3]

$$\mathcal{L}_{\rm HEFT}^{(4)} = -\sum_{R=V,A}\frac{1}{M_R^2}\langle\chi_R^{\mu\nu}\chi_{R\,\mu\nu}\rangle_2\,. \qquad (4.3)$$

It is important to note that, in addition to the absence of bilinear and four-fermion operators, the effective Lagrangian does not receive custodial-breaking terms at NLO.

The $\mathcal{O}(p^6)$ contribution to the HEFT Lagrangian is given by [6, 7]

$$\mathcal{L}_{\rm HEFT}^{(6)} = -\sum_{R=V,A} 2\left\langle\nabla^\rho\left(\frac{\chi_{R\,\rho\nu}}{M_R^2}\right)\nabla_\mu\left(\frac{\chi_R^{\mu\nu}}{M_R^2}\right)\right\rangle_2\,, \qquad (4.4)$$

where the $F_{V,A}$ and $\widetilde{F}_{V,A}$ couplings induce a $g^{(\prime)\,2}$ suppressed low-energy operator with four derivatives that contributes to the EW gauge boson self-energy

$$\mathcal{L}_{\rm HEFT}^{(6)} \supset \mathcal{L}_{\rm HEFT}^{(6)}\Big|_{(\nabla f)^2} = -\frac{2F_V^2}{M_V^4}\langle\nabla^\rho f_{+\,\rho\nu}\nabla_\mu f_+^{\mu\nu}\rangle_2 + \ldots \qquad (4.5)$$

However, these type of terms induce contributions in a whole set of observables that seem to have little to do with the original operators, with the structure shown on the right-hand side of Eq. (4.5). The key-point is that they may be simplified by means of the gauge field EoM and redefinitions of their fields into interaction operators with at least three or more fields already present in $\mathcal{L}^{(4)}_{\mathrm{HEFT}}$, where two covariant derivatives turn into two powers of $g^{(\prime)}$, *i.e.*, into two powers of $m_{W,Z}$. Thus, although some of the EoM-simplified operators have the structure of the $\mathcal{O}(p^4)$ HEFT Lagrangian, they are still $\mathcal{O}(p^6)$ suppressed; the additional suppression is now hidden within the LEC instead of being in the operator structure. In particular, since the gauge field EoM has the form $\nabla_\mu f_\pm^{\mu\nu} \sim g^{(\prime)\,2} J^\nu + \ldots$, the $\mathcal{O}(p^6)$ Lagrangian hides indeed four-fermion operators with a very suppressed coefficient. This is illustrated diagrammatically in Fig. 6.

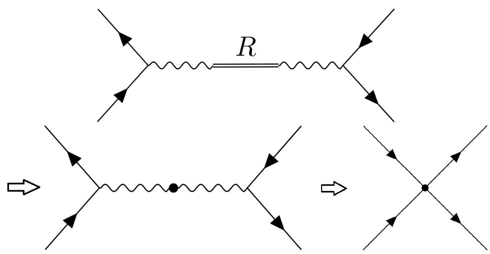

Fig. 6. The $u\bar{d} \to W^+ \to R^+ \to W^+ \to u\bar{d}$ scattering (1st diagram) turns into an amplitude with an $\mathcal{O}(g^2 p^4)$ correction to the W self-energy at $E \ll M_R$ (2nd diagram). The EoM-simplifications of these $\mathcal{L}^{(6)}_{\mathrm{HEFT}}$ terms just tell us that the latter 2nd diagram with two W propagators ($\Delta_W(q) \sim -i/q^2$) is equivalent to the local four-fermion scattering in the 3rd diagram [7].

Observing the experimental four-fermion analyses (see Refs. [3, 11] and references therein) in the usual parametrization $\mathcal{L}_{\mathrm{eff}} = \frac{2\pi}{\Lambda^2}\sum_{i,j=\ell,r}\eta_{ij} J_i^\mu J_{j\,\mu}$, the most stringent bounds are, in general, obtained for $\eta_{rr} = \eta_{\ell\ell} = \eta_{r\ell} = \eta_{\ell r} = -1$: $\Lambda \gtrsim 20$ TeV. The prediction from a vector and an axial–vector triplet with interaction given by Eq. (4.1) yields the prediction

$$\frac{2\pi}{\Lambda^2} = \frac{1}{M_V^2} \times \underbrace{\frac{4m_Z^4 - 8m_Z^2 m_W^2 + 7m_W^4}{24v^2 M_V^2}}_{= 9.6\times 10^{-5}\left(\frac{1\,\mathrm{TeV}^2}{M_V^2}\right)} \times \underbrace{\frac{r^3+1}{r^2(r-1)}}_{>1}\,, \tag{4.6}$$

where we have assumed here a P-even interaction. We have used the LO relations $m_W^2 = g^2v^2/4$ and $m_Z^2 = (g^2 + g^{\prime\,2})v^2/4$. Finally, in this expression, we assume two VV–AA WSRs [1], which lead to $r = M_A^2/M_V^2 > 1$, $F_V^2 =$

$v^2 r/(r-1)$ and $F_A^2 = v^2/(r-1)$. Hence, Λ is not simply the mass of the lightest BSM states but it is highly increased by a large additional factor. Indeed, when the strongest four-fermion bounds $\Lambda \gtrsim 20$ TeV are translated into a vector mass lower-bound, one obtains values that go from $M_V \gtrsim 1.9$ TeV for V and A extremely degenerate case (with $r = 1 + 10^{-3}$), $M_V \gtrsim 0.6$ TeV for the slightly degenerate scenario (with $r = 1.1$), and the really loose bound $M_V \gtrsim 0.3$ TeV for an V–A splittings with $r = 2$ or larger [6, 7]. Therefore, in spite of the large scales Λ found by these four-fermion analyses, it is very easy to devise theoretical frameworks where they actually imply very loose constraints on the masses of new physics states.

On the other hand, the EW precision observables and their expected improvement in future colliders [15, 48] might be a more efficient approach if we happen to be in this type of non-direct SM-fermion coupling scenarios. The first custodial breaking operators appear at $\mathcal{O}(p^6)$ in these model in Eq. (4.1): after employing the gauge field EoM and field redefinitions, we find that there is a contribution to Longhitano's a_0 term in $\mathcal{L}_{\mathrm{HEFT}}^{(4)}$ but carrying the original $\mathcal{O}(p^6)$ suppression, this is, with a very suppressed contribution to the T-parameter. On the other hand, the S-parameter is ruled by Longhitano's a_1 coupling [25, 45] which arises already at $\mathcal{O}(p^4)$ without any additional implicit suppression. At tree-level, we have the predictions

$$\begin{aligned} S &= -16\pi^2 a_1 = \frac{4\pi v^2}{M_V^2}\,\frac{(r+1)}{r}\,, \\ T &= \frac{2}{\alpha} a_0 = -\frac{4\pi v^2}{M_V^2} \times \underbrace{\frac{(m_Z^2 - m_W^2)}{4M_V^2}\,\frac{m_Z^2}{m_W^2}} \times \underbrace{\frac{r^3+1}{r^2(r-1)}}_{>1}\,, \\ &= 6.0 \times 10^{-4} \left(\frac{1\,\mathrm{TeV}^2}{M_V^2}\right) \end{aligned} \qquad (4.7)$$

where again we have assumed that the BSM extension obeys two VV–AA WSRs (which imply $r = M_A^2/M_V^2 > 1$). We have compared these results with the experiment [14, 15] in Fig. 7 and the corresponding M_V bounds can be seen in Table I [6, 7]. One can see that these experimental determinations are much more constraining and lead to masses $M_V \gtrsim 3$ TeV in this most constraining scenario with two VV–AA WSRs. These bounds become softer if one assumes less stringent BSM theories where only the first WSR applies [16, 17]. From this point of view, a factor five improvement in the precision of the oblique parameters measurements at future colliders [15, 48] seems to be a very promising way to either discover new physics in the few TeV range or definitely exclude it up to masses $M_V \sim 10$ TeV. Note that we are just showing tree-level predictions and one-loop corrections are not discussed here.

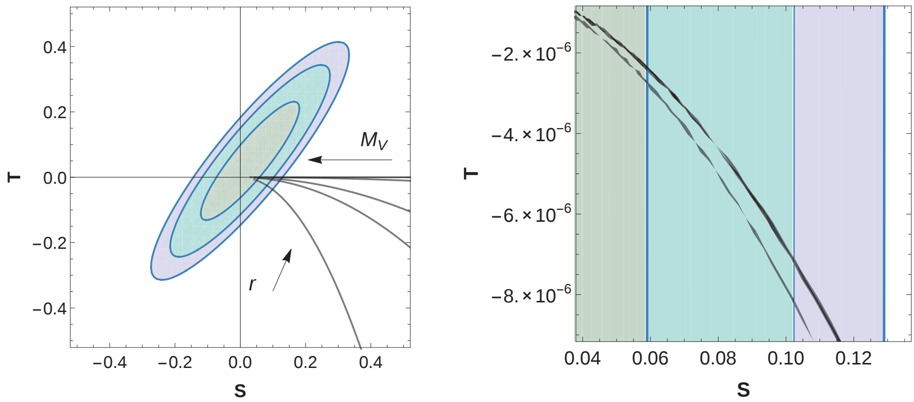

Fig. 7. (Color online) Allowed values for S and T for LEP data at 68% (light gray), 95% (medium gray/pale green) and 99% C.L. (dark gray/violet). Left-hand side: S *vs.* T for $r-1 = 10^{-4}, 10^{-3}, 5\times10^{-3}, 10^{-2}$ and 0.5. Right-hand side: S *vs.* T for $r = 2,3,4,5$ (zoom of the left-hand side figure). Notice that the plots converge rapidly for $r \geq 3$ to the thick line.

TABLE I

Bounds on M_V for different values of r using the allowed region for S and T at 68% and 95% C.L. from LEP [14, 15].

$r = M_A^2/M_V^2$	Lower bound 68% C.L.	For M_V 95% C.L.
$1+10^{-3}$	5.2 TeV	4.0 TeV
1.1	5.1 TeV	3.9 TeV
2	4.5 TeV	3.4 TeV
∞	3.7 TeV	2.8 TeV

5. Conclusions

We would like to end these proceedings with an optimistic message: new physics may be just around the corner, at a few TeV, crouching. The "bosonic" EW precision measurements do not disfavour this statement. All one needs is a proper (and strong) suppression of the $R \to f\bar{f}'$ interaction. Theories where the BSM sector only couples to the SM bosons, and has no direct coupling to SM fermions, naturally reproduce these features without any necessity of further tuning, while they still allow for measurable deviations from the SM in appropriate "bosonic" observables.

Resonances with $M_R \sim 2$ TeV are hence still perfectly allowed. At LHC searches, resonance production via vector boson scattering or Drell–Yan are naturally small in these models. They are difficult to pin down and they require large integrated luminosities. Four-fermion operator studies with dijet and dilepton signatures (sometimes called "compositeness" studies) may not lead to any successful result even in the long term for this type of scenarios. On the other hand, low-energy searches and an important improvement in the bosonic EW precision observables may be our best option to discover new physics if it has the structure of this type of models.

We thank the organizers for the nice scientific environment and their patience. This work was supported in part by the Spanish MINECO project FPA2016-75654-C2-1-P, FIS2017-58053-C2-1-P and FIS2017-84038-C2-1-P, and by CONACYT project No. 250628 ('Ciencia Básica') and the support 'Estancia Posdoctoral en el Extranjero'.

REFERENCES

[1] A. Pich, I. Rosell, J. Santos, J.J. Sanz-Cillero, *Phys. Rev. D* **93**, 055041 (2016) [arXiv:1510.03114 [hep-ph]].

[2] A. Pich, I. Rosell, J. Santos, J.J. Sanz-Cillero, *J. High Energy Phys.* **1704**, 012 (2017) [arXiv:1609.06659 [hep-ph]].

[3] C. Krause *et al.*, *J. High Energy Phys.* **1905**, 092 (2019) [arXiv:1810.10544 [hep-ph]].

[4] R.L. Delgado *et al.*, *J. High Energy Phys.* **1711**, 098 (2017) [arXiv:1707.04580 [hep-ph]].

[5] A. Dobado, F.J. Llanes-Estrada, J.J. Sanz-Cillero, *J. High Energy Phys.* **1803**, 159 (2018) [arXiv:1711.10310 [hep-ph]].

[6] A. Guevara, F. Alvarado, J.J. Sanz-Cillero, arXiv:1909.00875 [hep-ph].

[7] F. Alvarado, A. Guevara, J.J. Sanz-Cillero, in preparation.

[8] ATLAS Collaboration, ATLAS-CONF-2018-016.

[9] T. Dorigo, *Prog. Part. Nucl. Phys.* **100**, 211 (2018) [arXiv:1802.00354 [hep-ex]].

[10] D. Pappadopulo, A. Thamm, R. Torre, A. Wulzer, *J. High Energy Phys.* **1409**, 060 (2014) [arXiv:1402.4431 [hep-ph]].

[11] J.A. Aguilar-Saavedra *et al.*, arXiv:1802.07237 [hep-ph].

[12] CMS Exotica Physics Group Summary–ICHEP 2016, Technical Report, CERN, Geneva, http://cms-results.web.cern.ch/cms-results/public-results/publications/

[13] S. Rappoccio, *Rev. Phys.* **4**, 100027 (2019) [arXiv:1810.10579 [hep-ex]].

[14] M. Baak *et al.*, *Eur. Phys. J. C* **72**, 2205 (2012) [arXiv:1209.2716 [hep-ph]]; http://gfitter.desy.de/

[15] M. Baak *et al.*, arXiv:1310.6708 [hep-ph].

[16] A. Pich, I. Rosell, J.J. Sanz-Cillero, *Phys. Rev. Lett.* **110**, 181801 (2013) [arXiv:1212.6769 [hep-ph]].

[17] A. Pich, I. Rosell, J.J. Sanz-Cillero, *J. High Energy Phys.* **1401**, 157 (2014) [arXiv:1310.3121 [hep-ph]].

[18] M.E. Peskin, T. Takeuchi, *Phys. Rev. Lett.* **65**, 964 (1990).

[19] M.E. Peskin, T. Takeuchi, *Phys. Rev. D* **46**, 381 (1992).

[20] C. Garcia-Garcia, M. Herrero, R.A. Morales, *Phys. Rev. D* **100**, 096003 (2019) [arXiv:1907.06668 [hep-ph]].

[21] R.L. Delgado, C. Garcia-Garcia, M.J. Herrero, arXiv:1907.11957 [hep-ph].

[22] M. Fabbrichesi, M. Pinamonti, A. Tonero, A. Urbano, *Phys. Rev. D* **93**, 015004 (2016) [arXiv:1509.06378 [hep-ph]].

[23] O. Cata, G. Isidori, J.F. Kamenik, *Nucl. Phys. B* **822**, 230 (2009) [arXiv:0905.0490 [hep-ph]].

[24] D. de Florian *et al.* [LHC Higgs Cross Section Working Group], DOI:10.2172/1345634, 10.23731/CYRM-2017-002 [arXiv:1610.07922 [hep-ph]].

[25] R.L. Delgado, A. Dobado, M.J. Herrero, J.J. Sanz-Cillero, *J. High Energy Phys.* **1407**, 149 (2014) [arXiv:1404.2866 [hep-ph]].

[26] A. Castillo, R.L. Delgado, A. Dobado, F.J. Llanes-Estrada, *Eur. Phys. J. C* **77**, 436 (2017) [arXiv:1607.01158 [hep-ph]].

[27] S. Weinberg, *Physica A* **96**, 327 (1979).

[28] A. Manohar, H. Georgi, *Nucl. Phys. B* **234**, 189 (1984).

[29] T. Appelquist, C.W. Bernard, *Phys. Rev. D* **22**, 200 (1980).

[30] A.C. Longhitano, *Phys. Rev. D* **22**, 1166 (1980).

[31] A.C. Longhitano, *Nucl. Phys. B* **188**, 118 (1981).

[32] G. Buchalla, O. Catà, C. Krause, *Nucl. Phys. B* **880**, 552 (2014) [*Corrigendum ibid.* **913**, 475 (2016)] [arXiv:1307.5017 [hep-ph]].

[33] G. Buchalla, O. Catà, *J. High Energy Phys.* **1207**, 101 (2012) [arXiv:1203.6510 [hep-ph]].

[34] R. Alonso *et al.*, *Phys. Lett. B* **722**, 330 (2013) [*Corrigendum ibid.* **726**, 926 (2013)] [arXiv:1212.3305 [hep-ph]].

[35] F.K. Guo, P. Ruiz-Femenía, J.J. Sanz-Cillero, *Phys. Rev. D* **92**, 074005 (2015) [arXiv:1506.04204 [hep-ph]].

[36] R. Alonso, E.E. Jenkins, A.V. Manohar, *Phys. Lett. B* **754**, 335 (2016) [arXiv:1511.00724 [hep-ph]].

[37] R. Alonso, E.E. Jenkins, A.V. Manohar, *J. High Energy Phys.* **1608**, 101 (2016) [arXiv:1605.03602 [hep-ph]].

[38] G. Buchalla *et al.*, *Nucl. Phys. B* **928**, 93 (2018) [arXiv:1710.06412 [hep-ph]].

[39] R. Alonso, K. Kanshin, S. Saa, *Phys. Rev. D* **97**, 035010 (2018) [arXiv:1710.06848 [hep-ph]].

[40] O. Catà, C. Müller, arXiv:1906.01879 [hep-ph].

[41] G. Ecker, J. Gasser, A. Pich, E. de Rafael, *Nucl. Phys. B* **321**, 311 (1989).

[42] S. Weinberg, *Phys. Rev. Lett.* **18**, 507 (1967).

[43] C.W. Bernard, A. Duncan, J. LoSecco, S. Weinberg, *Phys. Rev. D* **12**, 792 (1975).

[44] G. Ecker *et al.*, *Phys. Lett. B* **223**, 425 (1989).

[45] D. Espriu, M.J. Herrero, *Nucl. Phys. B* **373**, 117 (1992).

[46] M. Aaboud *et al.*, [ATLAS Collaboration], *Phys. Rev. D* **95**, 032001 (2017) [arXiv:1609.05122 [hep-ex]].

[47] A. M. Sirunyan *et al.*, [CMS Collaboration], *Phys. Lett. B* **798**, 134985 (2019) [arXiv:1905.07445 [hep-ex]].

[48] J. B. Guimarães da Costa *et al.*, [CEPC Study Group], arXiv:1811.10545 [hep-ex].