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Effective field theory approach to the Higgs lineshape

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The phenomenology of an unstable particle, including searches and exclusion limits at the LHC, depends significantly on its lineshape. When the width of the resonance is large with respect to its mass, off-shell effects become relevant and the very same definition of width becomes nontrivial. Taking a heavy Higgs boson as an example, we propose a new formulation to describe the lineshape via an effective field theory approach. Our method leads to amplitudes that are gauge invariant, respect unitarity, and can appropriately describe the lineshape of broad resonances. The application of the method to the following relevant processes for the LHC phenomenology have been considered: gluon fusion, vector boson scattering, and $t\bar{t}$ production via weak boson fusion.

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I. INTRODUCTION

The CMS and ATLAS collaborations have announced the discovery of a resonance around 126 GeV compatible with the Higgs boson predicted by the Standard Model (SM) [1,2]. To eventually confirm the discovery of the SM Higgs boson, it will be necessary not only to measure the strength and the structure of its couplings to the fermions and vector bosons of the SM but also to exclude the existence of other heavier scalar states with similar properties.

The scalar sector of the SM is particularly simple, yet it does not provide any clue either on its origin, i.e., on the possible underlying dynamics, or on its stability. Alternative models that address these and other open questions often predict a richer structure for the scalar sector, such as in supersymmetry, technicolor theories, and models with extra dimensions. Two-Higgs doublet models provide the simplest and most often studied extensions of the scalar sector of the SM. In all these cases, heavier scalar and/or pseudoscalar partners of the “SM Higgs” boson can be expected below the TeV scale.

The width of such heavy states turns out to be sizable. A large width induces both a smearing and deformation of the signal lineshape as well as a sizable signal/background interference. For a heavy SM-like Higgs boson (e.g., $\Gamma \approx 100$ GeV for $m_H \approx 550$ GeV), the narrow width approximation (NWA) has been shown to be untenable, possibly invalidating the currently set exclusion limits on the heavy Higgs boson and beyond [3,4].

Going beyond the NWA, however, turns out not to be such an easy theoretical task. The challenges are manifold. First, the most accurate predictions for the signal cross sections, typically at the next-to-next-to-leading order (NNLO) in QCD and at NLO in electroweak (EW) interaction, assume a complete factorization between production and decay, i.e., they employ the NWA, and the $a$ posteriori inclusion of width effects is not uniquely defined [5]. Second, the very definition of a width, which amounts to a resummation of a specific subset of terms appearing at all orders in perturbation theory becomes problematic, leading to possible violations of gauge symmetry as well as of unitarity [6].

Currently, the most pragmatic and employed solution is the so-called complex mass scheme (CMS) [7]. In short, it amounts to analytically continue the parameters entering the SM Lagrangian that are related to the masses to complex values. Such a scheme maintains gauge invariance at all orders in perturbation theory and can be consistently employed in (N)NLO in EW computations [8,9]. However, since a fixed complex pole is assumed for any virtuality of the unstable particle, the resulting scattering amplitudes violate unitarity and do not correctly describe the lineshape of broad resonances.

An alternative to the CMS is the fermion-loop scheme [6,10], which offers a solution for the restoration of unitarity but yet does not give a complete description of unstable particles because the width of a heavy Higgs boson is dominated by its decay into gauge bosons. A consistent treatment of the bosonic contributions is possible in the framework of the background field gauge [11], which is equivalent to the pinch technique (PT) [12–15]. The resummation of the PT self-energy has been tackled in Refs. [16–20]. These approaches require a calculation of the complete radiative corrections at a fixed-loop order, greatly increasing the complexity of the calculation. Other suggested schemes for the treatment of unstable particles include the pole scheme [21], the Seymour scheme [22], the use of an effective Lagrangian including nonlocal interactions [23,24], and an approach based on collinear effective field theory [25].

Summarizing, in general and especially for a heavy Higgs boson, one would like to be able to include a running width in the propagator directly connected with the Higgs self-energy and at the same time to respect gauge invariance and unitarity.

In this work we tackle the problem from an effective field theory (EFT) point of view. We propose to systematically include width effects via a set of gauge invariant...
higher dimensional terms to the SM Lagrangian, along the lines of what was first proposed in Refs. [23,24]. Such new operators systematically encapsulate higher order terms coming from the self-energy and naturally allow a running and physical width for the Higgs boson in a gauge-invariant way. As we will show in the following, our scheme is consistent at higher orders, and it can be considered a generalization of the CMS as it reduces to it in the limit where the dependence on the virtuality of the Higgs self-energy is neglected.

II. SETTING UP THE STAGE

The two-point Green’s function for the Higgs boson is

$$\Delta_H(s) = s - m_{H,0}^2 + \Pi_{HH}(s),$$

(1)

where $m_{H,0}$ is the bare mass, and $\Pi_{HH}(s)$ is the Higgs self-energy. In the conventional on-shell definition, the mass and width are given by

$$m_{H,OS}^2 = m_{H,0}^2 - \text{Re}\Pi(m_{H,OS}),$$

(2)

$$m_{H,OS}(s) = \frac{\text{Im}\Pi(m_{H,OS})}{1 + \text{Re}\Pi(m_{H,OS})}.$$  

(3)

These definitions become gauge dependent at order $O(g^4)$.

In order to avoid the divergence of the tree-level propagator $D(s) = i/(s - m_{H,OS}^2)$, one performs the Dyson resummation to obtain

$$D(s) = \frac{i}{s - m_{H,OS}^2 + i m_{H,OS} \Gamma_{H,OS}}.$$  

(4)

To include the running effects of the width, one can further approximate the propagator by

$$D(s) = \frac{i}{s - m_{H,OS}^2 + i \text{Im}\Pi(s)},$$

(5)

where the imaginary part of $\Pi(s)$ is related to the Higgs boson width. The consistency of the above treatments of the Higgs propagator with the equivalence theorem and unitarity has been discussed by Valencia and Willenbrock [26].

Alternatively, as shown in a series of papers [21,27], a consistent, convenient, and resilient definition of mass $\mu$ and width $\gamma$ up to two loops is obtained by setting $s_H = \mu^2 - i \mu \gamma$ and then solving the implicit equation

$$s_H - m_{H,0}^2 + \Pi_{HH}(s_H) = 0$$

(6)
in terms of $s_H$. This gives a gauge-independent definition to all orders [28] [independent of the gauge choice present in the computation of $\Pi_{HH}(s_H)$] and in addition avoids unphysical threshold singularities [29].

The above definition is also consistent with the use of the CMS. In this scheme the propagator is $\Delta_H^{-1}(s) = s - s_H$. By definition this approach can give a good approximation of the full propagator

$$\Delta_H^{-1}(s) = \frac{1}{s - s_H + \Pi_{HH}^R(s)}$$

(7)

only close to the pole or, equivalently, for a small width, $\gamma/\mu \ll 1$. Here $\Pi_{HH}^R(s)$ is the renormalized self-energy, satisfying the following renormalization conditions:

$$\Pi_{HH}^R(s_H) = 0, \quad \Pi_{HH}^R(s_H) = 0.$$  

(8)

A natural improvement would consist of including the full resummed propagator in explicit calculations. This, however, leads to gauge violation already at the tree level. The reason being that in perturbation theory gauge invariance is guaranteed order by order while the presence of a width implies the resummation of a specific subset of higher order contributions, the self-energy corrections. This results in a mixing of different orders of perturbation theory. In particular, the following issues need to be addressed:

(1) In general $\Pi_{HH}(s)$ explicitly depends on the gauge-fixing parameter (GFP). To resum the self-energy correction to all orders, $\Pi_{HH}(s)$ must be extracted in a physically meaningful way.

(2) The resummed propagator spoils the gauge cancellation among different diagrams and eventually leads to the violation of the Goldstone boson equivalence theorem and unitarity bound.

Both issues can be tackled by the pinch technique [12–15]. In the PT framework, a modified one-loop self-energy for the Higgs boson can be constructed by appending to the conventional self-energy additional propagatorlike contributions concealed inside vertices and boxes. For the application of PT in resonant transition amplitude, and in particular, the extraction of a physical self-energy, we refer to the work of Refs. [16–20].

The modified self-energy correction for the Higgs boson is GFP independent and reflects properties generally associated with physical observables. At the one-loop level, we have the following expressions [19,30]:

$$\Pi_{HH}^{(WW)}(s) = \frac{\alpha_w}{16 \pi} \frac{m_H^4}{m_W^2} \left[ 1 + 4 \frac{m_W^2}{m_H^2} - 4 \frac{m_W^2}{m_H^4} (2s - 3m_W^2) \right] \times B_0(s, m_W^2, m_H^2),$$

(9)

$$\Pi_{HH}^{(ZZ)}(s) = \frac{\alpha_w}{32 \pi} \frac{m_H^4}{m_W^2} \left[ 1 + 4 \frac{m_Z^2}{m_H^2} - 4 \frac{m_Z^2}{m_H^4} (2s - 3m_Z^2) \right] \times B_0(s, m_Z^2, m_H^2),$$

(10)

$$\Pi_{HH}^{(ff)}(s) = \frac{3 \alpha_w}{8 \pi} \frac{m_f^2}{m_W^2} (s - 4m_f^2) B_0(s, m_f^2, m_f^2),$$

(11)

$$\Pi_{HH}^{(HH)}(s) = \frac{9 \alpha_w}{32 \pi} \frac{m_H^2}{m_W^2} B_0(s, m_H^2, m_H^2),$$

(12)

where the superscripts denote the contributions from the $W, Z$ fermions and Higgs loops, and
is the normal Passarino-Veltman function [31]. These results are independent of the GFP. Note that the expressions in Eqs. (9)–(12) coincide with the $\xi = 1$ result obtained in the background field gauge [32–34].

In addition, the gauge cancellation among different amplitudes can be restored by including certain vertex corrections obtained via the PT [30,35]. This is because in this framework the Green’s functions satisfy the tree-level-like Ward identities (WI), which are crucial for ensuring the gauge invariance of the resummed amplitude.

As an example, let us consider the Higgs-mediated part of same helicity fermion scattering into longitudinal W’s, $f_{\pm}f_{\mp} \rightarrow W^+_L W^-_L$. There are contributions from s-channel and t-channel diagrams, as is shown in Fig. 1. The contributions from t-channel and Higgs diagrams to the amplitudes coming from longitudinal components of the W’s and same helicity fermions (in the high-energy limit) read\(^1\)

\[
\mathcal{M}_h^f \equiv \mathcal{M}_h^{W^+W^-} = \frac{k_{1\mu}k_{2\nu}}{m_W^2} \bar{\psi}(p_2)\gamma_\mu\gamma_5 u(p_1)\Gamma_{\mu\nu}^{HWW}(q, k_1, k_2) k_{1\mu}k_{2\nu} \frac{1}{m_W^2},
\]

where $\Gamma_{\mu\nu}^{HWW}(q, k_1, k_2)$ is the $W^+W^-$ vertex. The ellipsis in $\mathcal{M}_h^f$ denotes terms that are not related to the Higgs exchange diagram. These terms come from the contribution of opposite helicity fermions and are supposed to cancel the bad high-energy behavior of the $\gamma/Z$ mediated diagrams.

Without the Higgs contribution $\mathcal{M}_L^f$ grows with energy and eventually violates unitarity. The cancellation of the bad high-energy behavior of each amplitude, and the equivalence theorem, are guaranteed by the following WI:

\[
\mathcal{M}_L^f \equiv \mathcal{M}_L^{W^+W^-} = -\frac{i g m_W^2}{4 m_W^2} \bar{\psi}(p_2)\gamma_\mu u(p_1) + \ldots,
\]

where $\phi_z$ are Nambu-Goldstone bosons. Only the leading terms at high energy are included. The relation above explicitly shows that the inclusion of higher order terms in the imaginary part of $\Delta_H(q^2)$ has to be related to the EW corrections of $\Gamma_{\mu\nu}^{HWW}$ and three scalar vertex. Only if both $\Delta_H(s)$ and $\Gamma_{\mu\nu}^{HWW}$ are computed in one loop via the PT, then the WI remains valid, and the gauge cancellation, as well as the equivalence theorem, are not spoiled. Besides, $\mathcal{M}_h^f$ is not affected by the Higgs width, and therefore the tree-level relations can be used. Thus the resummed propagator can be consistently included with the one-loop correction to $\Gamma_{\mu\nu}^{HWW}$ via the PT.

Even though correct, the solution outlined above for $f\bar{f} \rightarrow W^+_L W^-_L$ is not a general one. In $W^+_L W^-_L \rightarrow Z_L Z_L$, for example, it is not sufficient to include only the $HWW$ and $HZZ$ corrections. The triple and quartic vector-boson vertices at one loop are also required to cancel the bad high-energy behavior of the Higgs-mediated amplitude, and the overall procedure of analyzing the full set of WI’s becomes more and more involved. The goal of this work is to present a simple method to generate the needed corrections to the vertices and propagators so that the WI’s are automatically satisfied and unitarity is automatically ensured.

\section{III. The EFT Approach}

As explained above, we aim at finding a systematic approach to improve the Higgs propagator without breaking either gauge invariance or unitarity. In other words we are looking for a mechanism that guarantees the constraints imposed by the WI to be satisfied at any order in perturbation theory.

At one loop, the full calculation via the PT certainly provides an exact solution valid at NLO. The challenge is to achieve the same keeping the calculation at leading order, including only the necessary ingredients coming from NLO and resumming them into the propagator via a Dyson-Schwinger approach. The idea is to associate the corrections to an $ad$ $hoc$ constructed gauge-invariant operator and match the operator to the one-loop two-point function $\Delta_H(s)$ calculated via the PT. In so doing one aims at obtaining the exact resummed propagator already at the
leading order and, at the same time, the interactions modified to automatically satisfy the WT’s. The latter desired result ensures the gauge invariance of the amplitudes, and it can be considered as an approximation to a full one-loop calculation in PT.

To this aim, we consider the Taylor expansion of the function $\Pi(s) = \Pi_{HH}^R(s)$

$$\Pi(s) = \sum_{i=0}^{\infty} c_i s^i, \quad (16)$$

where $c_i$ are dimensionful constants and, as first attempt, we add the following infinite set of operators to the Lagrangian:

$$\hat{O}_{\Pi} = \sum_{i=0}^{\infty} c_i \phi^\dagger (-D^2)^i \phi \equiv \phi^\dagger \Pi (-D^2) \phi, \quad (17)$$

where $\phi$ is the Higgs doublet and $D^\mu$ is the covariant derivative. It is straightforward to check that $\hat{O}_{\Pi}$ modifies the Higgs propagator as desired: the two $\phi$’s contribute two Higgs fields, and each $-D^2$ contributes an $s$ leading to

$$\Pi(s) = \Pi_{HH}^R(s), \quad (18)$$

as desired. Note that in principle, $\hat{O}_{\Pi}$ is a nonlocal operator, yet by expanding it, we reexpress it in terms of an infinite series of local operators.$^2$

We remark that while very similar in spirit, our approach differs from that of Ref. [23]: the operator chosen there does not contain gauge fields, and it is therefore not sufficient to restore the gauge cancellation and fix the bad high-energy behavior in vector-vector scattering.

Equation (17) leads to the correct expression for the propagator. However, the first term $\Pi(0) \phi^\dagger \phi$ in the expansion corresponds to a tadpole contribution. This can be avoided if this term is replaced by

$$\Pi(0) \phi^\dagger \phi \to \Pi(0) \frac{1}{2 s^2} \left[ (\phi^\dagger \phi) - \frac{v^2}{2} \right]^2,$$

i.e., the Higgs self-interaction is suitably modified. As one can easily check, such a modification leaves the relation of Eq. (18) unchanged. The final form of the operator, which we dub $\hat{O}_{\Pi}$, is

$$\hat{O}_{\Pi} = \phi^\dagger [\Pi (-D^2) - \Pi(0)] \phi + \Pi(0) \frac{1}{2 s^2} \left[ (\phi^\dagger \phi) - \frac{v^2}{2} \right]^2. \quad (19)$$

The addition of this operator to the SM leads to several changes, which we now consider in detail. First of all, by construction, it gives rise to the propagator in Eq. (7), and a resummed propagator with the full one-loop self-energy via the PT at tree level is obtained. Second, it leads to modifications of the other interactions, in such a way that gauge invariance is maintained. For example, the $W$ and $Z$ two-point functions are modified by the addition of

$$i \Delta \Pi_{WW}^{\mu\nu}(q^2) = i \left( \frac{g}{2} \right)^2 \left[ \Pi(0) g_{\mu\nu} + \Pi'(q^2) q^\mu q^\nu \right],$$

$$i \Delta \Pi_{ZZ}^{\mu\nu}(q^2) = i \left( \frac{g}{2 c_w} \right)^2 \left[ \Pi(0) g_{\mu\nu} + \Pi'(q^2) q^\mu q^\nu \right],$$

where $\nu$ is the Higgs vacuum expectation value, and

$$\Pi'(x) = \Pi(x) - \Pi(0), \quad \Pi''(x) = \Pi'(x) - \Pi'(0). \quad (20)$$

The values for the $W$ and $Z$ masses are shifted

$$m_{\tilde{W}}^2 = \left( \frac{g}{2} \right)^2 (1 + \Pi'(0)), \quad m_{\tilde{Z}}^2 = \left( \frac{g}{2 c_w} \right)^2 (1 + \Pi'(0)), \quad (21)$$

as well as the propagators

$$\frac{i}{q^2 - m_{\tilde{W},\tilde{Z}}^2} \left[ - g_{\mu\nu} + \frac{\left( 1 + \Pi'(0) \right) q^\mu q^\nu}{m_{\tilde{W},\tilde{Z}}^2 + q^2 \Pi'(0)} \right]. \quad (22)$$

Let us first consider $f \bar{f} \to W^+_L W^-_L$ in the EFT approach. The operator modifies the $H W^+ W^-$ and the $H f \bar{f}$ interactions. The combined effect is a factor of $1 + \Pi'(s)$. Therefore in this process the EFT approach is equivalent to the following substitution of the Higgs propagator:

$$\Delta_H^{\mu}(s) = \frac{1 + \Pi'(s)}{s - s_H + \Pi(s)}, \quad (23)$$

which behaves like $1/s$ at large energy, and therefore exactly cancels the high-energy behavior from $\mathcal{M}_L^f$.

Note that in the limit of $s \to 0$, if we take $\Pi(s)$ to be the renormalized PT self-energy, then the above propagator is real and therefore does not violate the unitarity bound. This can be seen if we write the Higgs two-point function as

$$\Delta_H^{\mu}(s) = s - s_H + \Pi(s) \equiv Z \Delta_H^{\mu}(s) = Z(s - m_{\tilde{H},0}^2 + \Pi(s)), \quad (24)$$

where $\Delta_H^{\mu}(s)$ ($\Delta_H^0(s)$) is the renormalized (bare) two-point function, $Z$ is the wave function renormalization constant, and $\Pi(s)$ is the bare PT self-energy. The propagator Eq. (23) can then be written as

$$\Delta_H^{\mu}(s) = \frac{\Delta_H^R(s)}{\Delta_H^R(s)} = \frac{\Delta_H^0(s)}{\Delta_H^0(s)}, \quad (25)$$

which is real when $s \to 0$ because $\Pi_0(0)$ does not have an absorptive part.
It is also interesting to note that, if \( \Pi(s) \) has a linear dependence on \( s \), i.e.,

\[
\Pi(s) = i(s - \mu^2) \gamma \mu,
\]

(26)

Eq. (23) becomes

\[
\Delta_{\mu}^{(1)}(s) = \frac{1 + \frac{s}{s - \mu^2 + is \gamma \mu}}{s}
\]

(27)

and the EFT approach coincides with the scheme proposed by Seymour [22]. This makes sense because in the Seymour scheme the vector boson pair self-energy also has a linear dependence on \( s \). In our scheme we see that the numerator of Seymour’s propagator comes from the modified \( HW^+W^- \) vertex, as required by the WI.

We now turn to vector-vector scattering and, in particular, to \( W_L^+W_L^- \to Z_LZ_L \). This process features a pure gauge and a Higgs-mediated \( s \)-channel contribution; see Fig. 2. Both contributions do contain terms that increase quadratically with energy at high energy, whose cancellation is guaranteed by gauge invariance. To calculate \( W_L^+W_L^- \to Z_LZ_L \) amplitude in the EFT we need to extract the Feynman rules from \( \bar{O}_{H} \), i.e., the contributions that need to be added to the usual SM rules. This is straightforward and gives (all momenta incoming) the following:

\[
H(q)W^+(k_1)W^-(k_2), \quad ig \frac{m_W}{\sqrt{1 + \Pi'(0)}} \Pi'(q^2)g^{\mu\nu} + \cdots,
\]

\[
Z(k_1)W^+(k_2)W^-(k_3), \quad ig \frac{m_Z}{c_W^2} \frac{m_W^2}{1 + \Pi'(0)} \left[ \Pi''(k_3^2)g^{\mu\nu}k_3^\rho - \Pi''(k_3^2)g^{\mu\rho}k_3^\nu \right] + \cdots,
\]

\[
Z(k_1)Z(k_2)W^+(k_3)W^-(k_4), \quad ig^2 \frac{m_Z^2}{1 + \Pi'(0)} \left[ \Pi''(s)g^{\mu\nu}g^{\rho\sigma} + 4m_W^2 \left( \Pi''(t)g^{\mu\rho}g^{\nu\sigma} + \Pi''(u)g^{\mu\sigma}g^{\nu\rho} \right) \right] + \cdots,
\]

\[
H\phi^+\phi^-, \quad H\phi^0\phi^0, \quad -ig \frac{g[s_H - \Pi(0)]}{2m_W} \sqrt{1 + \Pi'(0)},
\]

\[
\phi^+\phi^- \phi^0\phi^0, \quad -ig^2 \frac{s_H - \Pi(0)}{4m_W^2} \left[ 1 + \Pi'(0) \right],
\]

where ellipses denote terms vanishing on shell and \( s = (k_1 + k_2)^2 \), \( t = (k_1 + k_3)^2 \), and \( u = (k_1 + k_4)^2 \). These Feynman rules are sufficient to calculate both \( W_L^+W_L^- \to Z_LZ_L \) and \( \phi^+\phi^- \to \phi^0\phi^0 \). At the leading order in \( m_W^2 \) and \( m_Z^2 \), we find for \( W_L^+W_L^- \to Z_LZ_L \),

\[
\mathcal{M}_{H}^{LLL} = -ig^2 \frac{s^2[1 + \Pi'(s)]^2}{4m_W^2 \left[ s - s_H + \Pi(s) \right] \left[ 1 + \Pi'(0) \right]},
\]

(29)

\[
\mathcal{M}_{gauge}^{LLL} = ig^2 \frac{s}{4m_W^2} \frac{1 + \Pi'(s)}{1 + \Pi'(0)},
\]

(30)

and for \( \phi^+\phi^- \to \phi^0\phi^0 \),

\[
\mathcal{M}_G = -ig^2 \frac{s + \Pi(s) - \Pi(0)}{4m_W^2 \left[ s - s_H + \Pi(s) \right]} \left[ s_H - \Pi(0) \right] \left[ 1 + \Pi'(0) \right],
\]

(31)

so that

\[
\mathcal{M}_H^{LLL} + \mathcal{M}_{gauge}^{LLL} = -ig^2 \frac{s + \Pi(s) - \Pi(0)}{4m_W^2 \left[ s - s_H + \Pi(s) \right]} \frac{s_H - \Pi(0)}{1 + \Pi'(0)}
\]

\[
= \left[ 1 + \Pi'(0) \right]^2.
\]

(32)

As expected, \( \mathcal{M}_H^{LLL} + \mathcal{M}_{gauge}^{LLL} \) does not grow with \( s \) and the equivalence theorem is recovered, up to a factor \( [1 + \Pi'(0)]^2 \), which exactly amounts to the wave function renormalization of the Goldstone fields.

An interesting feature of our approach is that in the limit where the dependence of \( \Pi(s) \) on \( s \) is neglected, \( \Pi(s) \equiv \Pi \) is a constant, then \( \Pi'(s) = \Pi''(s) = 0 \). The only effect of the operator is a shift in \( \lambda \), the coupling of the Higgs boson self-interaction. If \( m_H \) is the on-shell mass, this amounts to the replacement

\[
m_H^2 \to m_H^2 - \Pi.
\]

(33)

FIG. 2 (color online). Diagrams contributing to \( W_L^+W_L^- \to Z_LZ_L \).
i.e., given that $\Pi$ can be a complex number, it is equivalent to the CMS. In the complex mass renormalization, setting $\Pi = 0$ recovers the complex mass scheme.

The advantage of the EFT approach is the possibility of using an “arbitrary” functional form of the self-energy. We have shown that with special choices of $\Pi(s)$, the EFT approach can reduce to the Seymour scheme and the CMS scheme in certain cases. In particular, there is no need for spurious nonzero width for $t$-channel propagators as this can be easily imposed by always maintaining gauge invariance. Finally, we note that even though the restoration of gauge invariance and equivalence theorem is a general feature of our approach, one has to be careful in choosing the appropriate operator. For example, the following operator

$$O'_\Pi = \frac{1}{2v^2}(\phi^4 - \nu^2)\Pi(-\alpha^2)(\phi^4 - \nu^2),$$

introduced in Ref. [23], gives rise to the correct self-energy and the resummed propagator, but it does not modify the gauge contribution, so in $W^+_L W^-_L \rightarrow Z_L Z_L$, the gauge cancellation between the $s$-channel Higgs-mediated amplitude and the gauge amplitude is not restored. On the other hand, it modifies the Goldstone amplitude in a way so that the equivalence theorem is satisfied. As a result, both $W^+_L W^-_L \rightarrow Z_L Z_L$ and $\phi^+ \phi^- \rightarrow \phi^0 \phi^0$ have high-energy behavior and eventually break unitarity bounds. In general, adding higher dimensional operators to the Lagrangian leads to unitarity violation at some scale. We are going to show in the next sections that the operators we use do not have this problem.

Though the above operator $O'_\Pi$ solely does not treat the $HZZ$ and $HW^+ W^-$ correctly at high energy, when combined with $O_\Pi$, we can adjust them in a certain way to improve this method. We will discuss this in the following sections.

IV. UNITARITY

Adding operators of dimension $n > 4$ to the SM Lagrangian

$$\mathcal{L}_{\text{EFT}} = \mathcal{L}_{\text{SM}} + \sum_i c_i O_i[n]$$

is equivalent to recast the SM in terms of an effective field theory valid up to scales of order $\Lambda [37]$, beyond which the theory is not unitary. It is therefore mandatory to check whether this is the case for the operator $O_\Pi$. In fact, as we will see in the following section, a consistent perturbation theory implies that the same operator needs to also appear as a counterterm at higher orders. Overall we do not modify the theory and our procedure amounts to a reorganization of the perturbative expansion. However, we still need to make sure that neither unitarity is violated nor double counting happens at any given order in the perturbation theory. In this section we consider the first of these issues by showing that in sample calculations, $f \bar{f} \rightarrow VV$ and $VV \rightarrow VV$, at tree level the operator in Eq. (19) does not break unitarity at large energy.

In $f \bar{f} \rightarrow VV$ the change in $HVV$ vertex cancels the change in $H$ propagator at high $s$, independently of the helicities of $VV$, so the $s$-channel Higgs diagram does not lead to any bad high-energy behavior. The scattering of opposite helicity fermions does not entail the $s$-channel Higgs diagram and is the same as in the SM.

As we have already verified, in $W^+ W^- \rightarrow Z Z$ the longitudinal amplitude does not break unitarity, because the modification to the corresponding Goldstone interaction is finite $[s_H - \Pi(0)]$ and $1 + \Pi'(0)$. We now check the transverse amplitude $++ \rightarrow --, LL \rightarrow + +, ++ \rightarrow LL, ++ \rightarrow + +, in the limit$

$$s \sim |t| \sim |u| \gg m^2_H, \quad m^2_H \gg m^2_W.$$  

(Note that $+ -, + L, - L$ configurations do not feature a Higgs boson in the $s$-channel and therefore are left unchanged.) An explicit calculation for $W^+_L W^-_L \rightarrow Z_L Z_L$ gives

$$M^{++--} = M^{LLLL}_H \frac{4m^4_W}{s c_W} + O(m^4_W),$$

$$M^{++--} = M^{LLLL}_H \frac{4m^4_W}{s c_W} + O(m^4_W),$$

where $M^{LLLL}$ indicates the amplitude with four longitudinal vectors. For $W^+_L W^-_L \rightarrow Z_L Z_L$ we obtain

$$M^{H++L} = M^{LLLL}_H \frac{-2m^2_W}{s c_W} + O(m^2_W),$$

$$M^{L++L} = M^{LLLL}_H \frac{-2m^2_W}{s c_W} + O(m^2_W),$$

and for $W^+_L W^-_L \rightarrow Z_L Z_L$ we obtain

$$M^{H++L} = M^{LLLL}_H \frac{-2m^2_W}{s} + O(m^2_W),$$

$$M^{L++L} = M^{LLLL}_H \frac{-2m^2_W}{s} + O(m^2_W).$$

These results vanish faster than the longitudinal amplitude at large $s$. Finally for $W^+_L W^-_L \rightarrow Z_L Z_L$, we obtain

$$M^{H++} = -ig^2 m^2_Z \left[\frac{1 + \Pi'(s)}{s - s_H + \Pi(s)[1 + \Pi'(0)]} + O(m^4_W)\right],$$

$$= -ig^2 m^2_Z \Pi''(s) \sim s^{-1} \text{ at large } s,$$
\[ M_{\text{gauge}}^{+++} = i8g^2 \frac{e^2}{4\mu} s^2 + O(m_W^2). \] (44)

so at large energy the inclusion of \( \tilde{O}_H \) does not lead to any bad high-energy behavior.

**V. THE EFT APPROACH AT HIGHER ORDERS**

Starting at order \( \alpha_W \), the operator \( \tilde{O}_H \) is allowed in any leading order computation. At next-to-leading order in EW interactions, however, this is not necessarily consistent and possibly leads to double counting. In this section we argue that this is not a fundamental problem and can be dealt with by simply subtracting the same operator in a NLO as a counterterm, in full analogy to the procedure used in the CMS \[8\].

In the CMS, an imaginary part is added to the real mass and then subtracted as a counterterm at NLO. One can prove that this procedure does not spoil the WI’s, despite the fact that only a special class of higher order terms is resummed. As the EFT approach can be viewed as a generalization to the CMS, the same approach can be followed. The operator \( \tilde{O}_H \) corresponds to the imaginary part of the mass. It includes some of the higher order contribution and provides an improved solution to the WI’s. It enters the resummed propagator and other Feynman rules and needs to be subtracted at higher orders. The main difference is that in the CMS the propagator describes an unstable particle with a fixed width, while in the EFT approach one can resum an arbitrary part of the self-energy correction. This difference may be important when the width of the unstable particle is large, as in the case of a heavy Higgs boson, and the actual functional form of \( \Pi(s) \) becomes important.

In the pole-mass renormalization scheme, the two-point function of the Higgs can be written as

\[ \Delta_H(s) = s - s_H + \Pi^R_{HH}(s), \] (45)

where \( s_H \) is the pole and \( \Pi^R_{HH}(s) \) is the one-loop PT self-energy correction renormalized in the pole-mass scheme. We can now define the EFT approach by adding the operator in Eq. (19) and subtracting it as a counterterm:

\[ L_{\text{SM}} \rightarrow L_{\text{SM}} + \tilde{O}_H - \tilde{O}_H. \] (46)

In so doing the theory is exactly the same as before. Now Eq. (45) can be rewritten as

\[ \Delta_H(s) = s - s_H + \Pi(s) + \left[ \Pi^R_{HH}(s) - \Pi(s) \right], \] (47)

where the first three terms on the rhs start at leading order, while the last two terms with the bracket start at order \( \alpha_W \). The EFT approach then amounts to choose \( \Pi(s) \) in a way to capture the important part of (if not all of) \( \Pi^R_{HH}(s) \), so that this part of the self-energy correction is included at the leading order and will be resummed. In practice, one does not have to choose the exact PT self-energy, and gauge invariance is always guaranteed. In particular, choosing \( \Pi(s) = 0 \) corresponds to the CMS scheme.

In our scheme, EW NLO calculations are obviously more involved. The resummed propagator (7) and the modified Feynman rules do require extra work. However, one can also always employ a standard CMS at NLO and only include the full propagators and vertices in the LO result. In this way we can consistently have leading order calculated in the EFT approach and NLO in CMS but with counterterms from \( \tilde{O}_H \).

**VI. IMPROVED OPERATOR**

In the EFT approach, the exact functional form of the \( \Pi(s) \) is somewhat arbitrary. Ideally, one would like to use the imaginary part of the PT one-loop self-energy, as given in Eqs. (9)–(12). Unfortunately, in \( VV \rightarrow VV \) this choice produces an unphysical excess at low energy. As mentioned above, such an excess can be avoided by incorporating the operator \( O_{H}^{\Pi} \) in Eq. (34).

The reason for the excess is that the operator does not correctly describe the one-loop \( HVV \) vertices at low energy. This can be traced back to the WI in Eq. (15). The operator \( \tilde{O}_H \) in Eq. (19) does not modify \( \Gamma_{H}^{\phi^{+}\phi^{-}} \), hence, the equality relation is completely satisfied by the \( HWW \) vertex. At high energy this arrangement is satisfactory because we do not expect contributions from the Goldstone-Higgs vertex. However, at low energy, this vertex turns out to be relevant and thus, modifying only the \( HWW \) vertex gives rise to a bad behavior at low energy. On the other hand, the operator \( O_{H}^{\Pi} \) [Eq. (34)] satisfies the equation with the saturation of the Goldstone-Higgs vertex, yet fails at describing the high-energy behavior of weak boson scattering for example. The solution for this problem relies on the construction of a suitable combination of these two operators.

Let us rewrite the operator \( \tilde{O}_H \) in the following form:

\[ \tilde{O}_H = \phi^{+} \Pi_{1}(-D^{2}) \phi. \] (48)

Here we consider the operator of Eq. (19) in the simpler form, Eq. (17). This is justified by the fact that in practice we have to work only with the imaginary part of the self-energy, then \( \Pi(0) = 0 \). On the other hand, the operator proposed by Beenakker et al. [23] [Eq. (34)], can be written as

\[ O_{H}^{\Pi} = \frac{1}{2 \nu} (\phi^{+} \phi - v^{2}) \Pi_{2}(-\partial^{2}) (\phi^{+} \phi - v^{2}). \] (49)

To determine \( \Pi_{1} \) and \( \Pi_{2} \), we focus on the \( HZZ \) vertex and the \( H\phi^{0} \phi^{0} \) vertex. The inclusion of \( \tilde{O}_H \) and \( O_{H}^{\Pi} \) modifies these vertices, respectively:

\[^{3}\text{In fact the right-hand side has additional terms } igm_{H}^{2} [\Pi_{\phi^{+}\phi^{-}}(k_{1}^{2}) + \Pi_{\phi^{0}\phi^{0}}(k_{2}^{2})]/2 C_{W} \text{. They are zero for this process because we consider only the imaginary part.}\]

\[^{4}\text{Also because in the end we can always shift a constant part from } \tilde{O}_H \text{ to } O_{H}^{\Pi}.\]
This equation implies
\[ \Pi(q^2) = \Pi_1(q^2) + \Pi_2(q^2). \] (53)

We now need to find out the expressions for \( k_1 \mu k_2^\dagger \Gamma_{HZZ,\mu}(q, k_1, k_2) \) and \( \Gamma_{H\phi^0,\phi^0}(q, k_1, k_2) \), or equivalently, \( \Pi_1(q^2) \) and \( \Pi_2(q^2) \). Then the combined operator \( O_{\Pi_1} + O_{\Pi_2} \) should reproduce all three terms in Eq. (52) correctly. To this end, we calculate the absorptive part of both \( k_1 \mu k_2^\dagger \Gamma_{HZZ,\mu}(q, k_1, k_2) \) and \( \Gamma_{H\phi^0,\phi^0}(q, k_1, k_2) \) at one loop with PT. The results obtained are

\[ \Pi_{1,2}(s) = \sum_{(XX)} \Pi_{1,2}^{(XX)}(s), \] (54)

where XX are summed over WW, ZZ, tt, and HH, and

\[ \Pi_{1,2}^{(WW)}(s) = -\frac{\alpha_w}{4\pi} \left[ 2sB_0(s, m_W^2, m_W^2) + (s - 2m_W^2)(m_H^2 + 4m_W^2)C_0(m_Z^2, m_Z^2, s, m_W^2, m_W^2, m_W^2) \right], \] (55)

\[ \Pi_{1,2}^{(ZZ)}(s) = -\frac{\alpha_w}{16\pi m_W^2} \left[ [4m_Z^2 s + m_H^2 + m_H^2 - 2m_Z^2]B_0(s, m_Z^2, m_Z^2) + (4m_H^2 m_Z^2 - m_H^2 - 3m_H^2 m_Z^2 - 6m_Z^2)C_0(m_Z^2, m_Z^2, s, m_Z^2, m_H^2, m_Z^2) \right], \] (56)

\[ \Pi_{1,2}^{(tt)}(s) = \frac{3\alpha_w m_t^4}{8\pi m_W^2} \left[ sB_0(s, m_t^2, m_t^2) + 2m_t^2(s - 2m_t^2)C_0(m_Z^2, m_Z^2, s, m_t^2, m_t^2, m_t^2) \right], \] (57)

\[ \Pi_{1,2}^{(HH)}(s) = \frac{3\alpha_w m_H^4}{16\pi m_W^2} \left[ (m_H^2 - m_H^2)B_0(s, m_H^2, m_H^2) - (2m_H^4 s + m_H^4 - 2m_H^4 m_H^2 - 3m_H^4)C_0(m_Z^2, m_Z^2, s, m_H^2, m_H^2, m_Z^2) \right], \] (58)

\[ \Pi_{1,2}^{(WW)}(s) = \frac{\alpha_w}{16\pi m_W^2} \left[ [m_H^2 + 4m_W^2 m_H^2 + 12m_W^2]B_0(s, m_W^2, m_W^2) + 4m_W^2 (s - 2m_W^2)(m_H^2 + 4m_W^2)C_0(m_Z^2, m_Z^2, s, m_W^2, m_W^2, m_W^2) \right], \] (59)

\[ \Pi_{1,2}^{(ZZ)}(s) = \frac{\alpha_w}{32\pi m_W^2} \left[ [3m_H^2 + 6m_H^2 m_Z^2 + 8m_Z^2]B_0(s, m_Z^2, m_Z^2) + 2(4m_H^2 m_Z^2 s + m_H^4 - 3m_H^4 m_Z^2 - 6m_Z^4)C_0(m_Z^2, m_Z^2, s, m_Z^2, m_H^2, m_Z^2) \right], \] (60)

\[ \Pi_{1,2}^{(tt)}(s) = -\frac{3\alpha_w m_t^4}{4\pi m_W^2} \left[ 2B_0(s, m_t^2, m_t^2) + (s - 2m_t^2)C_0(s, m_Z^2, m_t^2, m_t^2) \right], \] (61)

\[ \Pi_{1,2}^{(HH)}(s) = \frac{3\alpha_w m_H^4}{32\pi m_W^2} \left[ (m_H^2 + 2m_Z^2)B_0(s, m_H^2, m_H^2) + 2(m_Z^2 s + m_H^4 - 2m_H^4 m_Z^2 - 3m_Z^4)C_0(m_Z^2, m_Z^2, s, m_H^2, m_Z^2, m_H^2) \right], \] (62)

where \( B_0 \) and \( C_0 \) are the Passarino-Veltman functions [31]. Note in the above equations, only the imaginary part of both sides will be used.

One can verify explicitly that the WI in Eq. (52) is satisfied, i.e., with the above definition we should have

\[ \Pi_{1,2}^{(XX)}(s) = \Pi_{1,2}^{(XX)}(s) + \Pi_{1,2}^{(XX)}(s), \] (63)

for (WW), (ZZ), (tt), and (HH), respectively.

Note that \( \Pi_1 \) corresponds to HZZ and \( \Pi_2 \) corresponds to \( H\phi^0\phi^0 \). In Fig. 3, we show a comparison of \( \Pi_{HH}, \Pi_1, \) and \( \Pi_2 \) at large energy, \( \Pi_{HH} \) and \( \Pi_1 \) display the same behavior, while \( \Pi_2 \) (i.e., \( H\phi^0\phi^0 \)) is negligible. This justifies the use of our operator \( O_{\Pi_1} \) in Eq. (17) at high energy, as it generates the right HVV vertex. At lower energy, however, \( \Pi_2 \) dominates and exhibits a “bump” above threshold. Because the operator given in Eq. (17) does
not give rise to the right $H\phi^0\phi^0$ vertex, it is clear that had we chosen this operator in the EFT scheme, we would have mistakenly considered this bump as part of $HVV$ vertex, resulting in an unphysical excess at low energy.

To reproduce the right behavior in both low- and high-energy regions, we use the combined operator

$$O_{\Pi} = O_{\Pi_1} + O_{\Pi_2}, \quad (64)$$

with $\Pi_1$ and $\Pi_2$ given by Eq. (54).

This operator gives a better description, as shown in Fig. 4 where we compare the longitudinal component of the $HZZ$ vertex derived from operator $O_{\Pi}$ [Eq. (17)] and the improved operator $\tilde{O}_{\Pi}$ [Eq. (64)] with the actual calculation at one loop with the PT. The improved operator describes the $HZZ$ vertex very well. Figure 5 compares the transverse component of the $HZZ$ vertex. The improved operator comes closer to the actual transverse component.

When applying our scheme with the operator $\tilde{O}_{\Pi}$ in Eq. (64), we prefer to work with the complex-pole renormalization. To do this we need the following counterterms for $\Pi_{HH}$:

$$\Pi_{HH}^R(s) = (\Pi_{HH}(s) - \delta s_H)Z + (s - s_H)\delta Z, \quad (65)$$

where $s_H$ is the pole. Here we take $Z = 1$ for the first term, which is already of order $O(\alpha_W)$. $\delta s_H$ and $\delta Z$ are given by

$$\delta s_H = \Pi_{HH}(s_H), \quad (66)$$

$$\delta Z = -\Pi_{HH}'(s_H). \quad (67)$$

For the Feynman rules, essentially we are going to replace all the $\Pi$'s in (28) by $\Pi_1$ and include a factor of $1 + \delta Z$ for the SM $HVV$ vertex to account for the wave function renormalization.
VII. APPLICATIONS

The treatment of the propagator of the Higgs boson is of immediate relevance for the LHC. As a simple testing ground of our proposal and comparisons to the conventional methods, we consider three processes of particular phenomenological importance at the LHC for a scalar boson (which for brevity, we identify with an hypothetical heavy Higgs boson): vector boson scattering, $t\bar{t}$ production via vector boson fusion, and Higgs production via gluon fusion. We have compared the effective approach described above in Eqs. (21), (22), and (28) with two other schemes:

1. A naive inclusion of the self-energy, i.e., using the following propagator:

$$i\Delta_H(s) = \frac{i}{s - s_H + \Pi_R(s)}.$$  \hspace{1cm} (68)

without changing anything else. Here $s_H = \mu^2 - i\mu\gamma$.

2. The CMS scheme,

$$i\Delta_H(s) = \frac{i}{s - s_H}.$$  \hspace{1cm} (69)

A modified version of MadGraph [38], with the implementation of the effective Lagrangian approach and the naive propagator with the PT self-energy, is used to generate events. As SM input parameters we take

$$m_Z = 91.188 \text{ GeV},$$  \hspace{1cm} (70)

$$G_F = 1.16639 \times 10^{-5} \text{ GeV}^{-2},$$  \hspace{1cm} (71)

$$\alpha^{-1} = 132.507,$$  \hspace{1cm} (72)

$$m_t = 173 \text{ GeV}.$$  \hspace{1cm} (73)

The pole mass is

$$\mu = 800 \text{ GeV},$$  \hspace{1cm} (74)

and $\Pi_R(s)$ is the imaginary part of the PT self-energy renormalized in the pole scheme. The factorization scale is set as the default dynamical scale of MadGraph and the PDF set is CTEQ6l1 [39].

A. Vector boson scattering

In $VV \rightarrow VV$ scattering processes, the effective description allows one to achieve a complete description of the Higgs lineshape at the resonance region and at the same time it corrects the bad high-energy behavior originated from the momenta-dependent part of the self-energy. As a bonus, we show that our definition avoids the need for including spurious $t$-channel widths that occur in the complex-mass scheme also affecting the high-energy behavior of the scattering amplitudes.

In Fig. 6 we show the energy behavior of the $ZZ \rightarrow ZZ$ scattering amplitude summed over helicities, $\sum_{\text{hel}}|M_{ZZ \rightarrow ZZ}|^2$, at scattering angle $\cos \theta = 0$. The fixed-width scheme, Eq. (69), naive propagator, Eq. (68), the effective description, and a case in which the width is set to zero are presented. The agreement between the effective scheme and the naive propagator at the resonance region is pretty good. The difference with respect to the fixed-width scheme is evident. At high energy, the naive propagator diverges, while the effective description behaves correctly. Similar comments can be made about $W^+W^- \rightarrow W^+W^-$ amplitude, shown in Fig. 7.

The fact that in both $ZZ \rightarrow ZZ$ and $W^+W^- \rightarrow W^+W^-$ the fixed-width scheme differs from the effective approach at the high-energy region indicates that the spurious $t$-channel width gives a non-negligible contribution.
The fact can be verified by comparing the different schemes with the no-width case. Moreover, in the case of \( W^+W^- \to W^+W^- \), shown in Fig. 8, the effective description and naive propagator are equivalent to the no-width case, and the excess observed in the amplitudes in the fixed-width scheme comes from the spurious width in the \( t \) and \( u \) channels.

At the LHC, the differences shown above may become important for a broad resonance. Despite the fact that a light Higgs boson has been observed, there is still room for new heavy and eventually broad resonances, e.g., in scaled-up QCD or in two-Higgs-doublet models. The VV scattering are embedded in more complex processes of the form \( qq \to qqVV \), where the two final state jets are emitted with high energy in the forward-backward region of the detectors and the vector bosons decay into two fermions with high \( p_T \) through the central region. We study the processes \( uc \to ucZZ \) and \( us \to dcW^+W^- \) assuming the nominal energy of LHC, \( E_{CM} = 14 \text{ TeV} \).

In Figs. 9 and 10, the distribution of the invariant mass of the ZZ system is shown. In Fig. 9, the resonant region is shown. A basic set of selection cuts to enhance vector boson scattering contribution, listed in the left column of Table I, has been applied. The effective description fits well with the running behavior of the Higgs propagator. In Fig. 10, the high-energy region is put in evidence. To better appreciate the differences between schemes at LHC energy, a further set of cuts has been added (right column).
As expected, the effective approach gives a well-behaved distribution at such energies contrary to the naive propagator and with a rate 10% lower than the fixed-width scheme. This difference amounts to the $t$-channel spurious contribution present in the fixed-width case. Similar conclusions can be drawn from Figs. 11 and 12, where the reconstructed $WW$ system invariant mass distribution for the $us \to dcW^+W^-$ process is shown.

FIG. 11 (color online). Mass distribution of $WW$ system in the process $us \to dcW^+W^-$ around the resonance peak. The cuts listed in the left column of Table I have been applied.

FIG. 12 (color online). Mass distribution of $WW$ system in the process $us \to dcW^+W^-$ at the high-energy region. All cuts listed in Table I have been applied.

t-channel spurious contribution present in the fixed-width case. Similar conclusions can be drawn from Figs. 11 and 12, where the reconstructed $WW$ system invariant mass distribution for the $us \to dcW^+W^-$ process is shown.

FIG. 13 (color online). Mass distribution of $t\bar{t}$ system in the process $us \to dctl$ around the resonance peak. The cuts listed in the left column of Table I have been applied.

FIG. 14 (color online). Mass distribution of $t\bar{t}$ system in the process $us \to dctl$ at the high-energy region. All cuts listed in Table I have been applied.
B. $W^+W^+ \to H \to t\bar{t}$ production

In $t\bar{t}$ production, we can observe a similar behavior with respect to $ZZ \to ZZ$ vector boson scattering. We have concentrated on the process $us \to dct\bar{t}$, in which the Higgs boson in produced by $W^+W^-$ fusion and decayed to a pair of top quarks. The energy in the center of mass is set to 14 TeV. In Fig. 13, the invariant mass distribution of $t\bar{t}$ at the resonant region is presented. The cuts shown in the left column of Table I have been applied in order to enhance the vector boson fusion contribution. Here again, the effective description describes the functional form of the propagator, which can go up to 5% of difference with respect to the fixed-width scheme. As seen in Fig. 14, in the high mass region, the effective description is dumped down by the effective $WWH$ vertex and does not grow with energy as is the case in which the naive propagator is adopted. The extra cuts shown in the right-hand column of Table I have been added in order to highlight the differences better.

C. Gluon-gluon fusion

For the study of a heavy Higgs boson produced via gluon-gluon fusion and decayed to a $W$-boson pair, $gg \to W^+W^- \to e^+\nu_e\mu^-\nu_\mu$, around the resonance peak. The cuts listed in Table II have been applied.

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|}
\hline
$p_T(\ell) > 2$ GeV & $\not{E}_T > 2$ GeV & $\eta(\ell) < 3$ \\
\hline
$\Delta R(\ell\ell) > 0.5$ & & \\
\hline
\end{tabular}
\caption{Cuts applied for $gg \to W^+W^-$.}
\end{table}

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig15}
\caption{(color online). Mass distribution of the reconstructed $WW$ system in the process $gg \to W^+W^-$ at the resonant region. The cuts listed in Table II have been applied.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig16}
\caption{(color online). Mass distribution of the reconstructed $WW$ system in the process $gg \to W^+W^-$ at the high-energy region. The cuts listed in Table II have been applied.}
\end{figure}

VIII. CONCLUSIONS

We have argued that it is possible to consistently and efficiently include running width effects for a heavy Higgs-like boson employing an EFT method. We can summarize the main points of our approach as follows:

1. Introducing a width for an unstable particle amounts to a rearrangement of the perturbative expansion where the corrections to the two-point function are resummed in the propagator. The addition of the operator $\hat{O}_{\Pi}$ defined in Eq. (64) allows one to effectively perform such resummation in a gauge-invariant...
and unitary way while keeping the full virtuality dependence of the self-energy. We have shown that in the limit where such dependence can be neglected our scheme is equivalent to the CMS.

(ii) At leading order, one has the freedom to choose the functional form of $\Pi(s)$. We propose to use the exact one-loop PT self-energy correction. The rationale is that such self-energies are gauge invariant and by exploiting the WI’s, we demand $\hat{O}_{\Pi}$ to mimic the most important one-loop corrections as much as possible. In practice, however, using any other form of $\Pi(s)$ does not break either gauge invariance or unitarity. In particular, one could avoid the need for a spurious nonzero width for $t$-channel propagators.

(iii) EW higher order corrections can still be performed in the CMS, without loss of accuracy or double counting issues. In practice, one can include the running width effects via the EFT at the leading order and neglect the virtuality dependence at NLO, i.e., employ the usual CMS for the NLO term.

In conclusion, in this work we have considered the case of how to consistently define a running width in the case of a heavy SM Higgs boson. The same approach can be used, for example, in the context of a two-Higgs-doublet model and applied to the current searches for new scalar states at the LHC. Extension to gauge vectors and heavy fermion states, on the other hand, are not straightforward and need further investigation.

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EFFECTIVE FIELD THEORY APPROACH TO THE HIGGS 


