

Chapter 2

Spin Determination of the LHC Higgs-Like Resonance

The Higgs discovery was initially made in the diboson modes $h \rightarrow ZZ^*, WW^*, \gamma\gamma$. The spin-parity assignment of the newly discovered boson h constitutes the first step to establish whether this particle is the SM Higgs responsible for electroweak symmetry breaking, or merely an impostor. We first discuss in Sect. 2.1 that the observation of the on-shell diphoton decay rules out the $J = 1$ assignment due to the well known Landau-Yang theorem. The most immediate cases alternate to the SM Higgs is that of a pseudoscalar ($J^P = 0^-$) and spin-2 ($J = 2$).

It was argued in [1, 2] that the alternate hypotheses are disfavoured by the analysis of the observed production and decay rates. In particular, the measured $h \rightarrow gg$ and $h \rightarrow \gamma\gamma$ rates are consistent with that of a SM Higgs:

$$\Gamma_{SM}(h \rightarrow gg) \approx 37 \Gamma_{SM}(h \rightarrow \gamma\gamma). \quad (2.1)$$

This is in tension with many spin-2 models with compactified extra-dimensions which predicts a universal coupling to photons and gluons $c_g = c_\gamma$ where a ratio of 8 should be observed between respective rates. Furthermore, it was noted in [1, 3] that universal couplings $h_{\mu\nu} D^\mu \Sigma^\dagger D^\nu \Sigma$ preserves gauge invariance and thus custodial symmetry. The approximate relation $c_W \approx c_Z$ was found to be consistent with the available data.

On the other hand, the methods adopted in the experimental analysis relies on studying the kinematic distributions (cf. [1–16] and references therein) which are expected to be highly anisotropic for higher spins. The reason why higher spins ($J \geq 3$) are not considered is because (i) of angular momentum conservation arguments, and (ii) that in these cases the number of independent helicity amplitudes for the ZZ^* decay does not exceed that of six as in $J = 2$ and can be distinguished based on the energy dependence [17, 18]. The framework can be made model independent by writing down the most general amplitude [17, 19, 20] or Lagrangian compatible with Lorentz and gauge symmetry. In particular, the use of effective field theory [8, 21, 22] has the following advantages: (i) it allows a hierarchy of new physics effects

to be ordered with respect to the energy scale, and (ii) provides a framework to also study the associated anomalous couplings. At the time of writing [23], ATLAS had placed 99.9% C.L. exclusions using the WW , ZZ and $\gamma\gamma$ channel regardless of the relative contribution of gg or $q\bar{q}$ production [24]. CMS reached a similar conclusion at 99.4% C.L but using only the WW and ZZ modes with the gg -initiated processes [25]. Whilst these constitute strong evidence against a spin-2 scenario under the assumption of minimal graviton couplings, they are not conclusive for a spin-2 impostor with generic couplings.

As explained in Sect. 1.2, a strong motivation for the existence of the Higgs boson is that it unitarises the high energy behaviour of vector boson scattering (see Appendix A.1 for a discussion of unitarity). The aim of our Letter (cf. Ref. [23]) is then to demonstrate that a $J = 2$ impostor mimicking the decay rates of the 125 GeV signal would jeopardise this feature. Rather than reproducing the entirety of the Letter, its motivation and calculations will be detailed in this chapter. Firstly the Landau-Yang theorem is presented in Sect. 2.1. Subsequently, the theoretical difficulty of constructing an interacting spin-2 theory is addressed in Sect. 2.2, which explains why the use of an EFT is necessary. In Sect. 2.3, we calculate the cutoff-scale for a theory with a spin-2 Higgs-impostor, h . This is taken to be the scale of unitarity violation for the longitudinal scattering of $hZ \rightarrow hZ$, which we found to be significantly lower than the LHC reach. We conclude with remarks on the Letter in Sect. 2.5.

2.1 Excluding $J = 1$

In this section we give a discussion on the Landau-Yang theorem, which is stated as follows:

Theorem 2.1.1 (Landau-Yang [26, 27]) *A spin-1 ($J = 1$) state cannot decay into two identical massless vectors.*

We inspect this briefly with an Abelian case where the vectors are photons. Interested readers are referred to [28] for the generalisation to non-Abelian cases. We first denote the amplitude $h(J = 1) \rightarrow \gamma\gamma$ as:

$$\mathcal{M}(1, 2) := \langle \gamma(k_1, \varepsilon_1) \gamma(k_2, \varepsilon_2) | h(p, \varepsilon) \rangle, \quad (2.2)$$

where $k_{1,2}$ and $\varepsilon_{1,2}$ are the four momenta and polarisation of the photons respectively. Transversality then requires that $\varepsilon \cdot p = 0$, whilst conservation of momenta requires that $p = k_1 + k_2$. We define $q := k_1 - k_2$ so that $p^2 = -q^2 = m_h^2$ and $p \cdot q = 0$. Gauge invariance is preserved when the vector polarisations are boosted by current-conserving objects:

$$\varepsilon'_i = \varepsilon_i - \frac{\varepsilon_i \cdot k_j}{k_i \cdot k_j} k_j, \quad (i, j) \in \{(1, 2), (2, 1)\}, \quad (2.3)$$

such that $p \cdot \varepsilon'_{1,2} = q \cdot \varepsilon'_{1,2} = 0$ is satisfied. The most general Lorentz invariant amplitude¹ must be linear in all the polarisation-vectors, taking the form:

$$\begin{aligned} \mathcal{M}(1, 2) := & F_1(p^2) (\varepsilon'_1 \cdot \varepsilon'_2) (\varepsilon \cdot q) + F_2(p^2) (\varepsilon \cdot q) \epsilon_{\mu\nu\rho\sigma} \varepsilon_1^{\prime\mu} \varepsilon_2^{\prime\nu} p^\rho q^\sigma \\ & + F_3(p^2) \epsilon_{\mu\nu\rho\sigma} \varepsilon_1^{\prime\mu} \varepsilon_2^{\prime\nu} \varepsilon^\rho p^\sigma. \end{aligned} \quad (2.4)$$

The Landau-Yang theorem then follows as the amplitude violates Bose symmetry, i.e. $\mathcal{M}(1, 2) = \mathcal{M}(2, 1)$, under the interchange of photons ($\varepsilon_1 \leftrightarrow \varepsilon_2$ and $q \rightarrow -q$). This is true whenever the functions $F_i(p^2)$ are not all zero [29]. Although it was pointed out that the theorem may be evaded² in the $J = 1$ hypothesis because of on-shell assumptions [32–34], this results in the lack of resonance structure which contradicts with the data [8].

2.2 Massive Spin-2

The construction of a consistent massive interacting spin-2 theory is met with various difficulties (see discussion in e.g. [35–40]), with the most prominent being the existence of ghost states [40]. Construction of a free Lagrangian for massive higher spins is done in [41, 42] and a unique Fierz-Pauli structure³ applies to the spin-2 case [44–46]:

$$\mathcal{L}_{J=2} \supset h^{\mu\nu} \mathcal{E}_{\rho\sigma}^{\mu\nu} h_{\rho\sigma} - \frac{1}{2} m_h^2 (h^{\mu\nu} h_{\mu\nu} - h^2), \quad (2.5)$$

where $h := h^\mu{}_\mu$ and the definition of the linearised Einstein operator $\mathcal{E}_{\rho\sigma}^{\mu\nu}$ is:

$$\mathcal{E}_{\rho\sigma}^{\mu\nu} h_{\rho\sigma} := \frac{1}{2} \square h_{\mu\nu} - \partial_\alpha \partial_{(\mu} h^\alpha{}_{\nu)} + \partial_\mu \partial_\nu h - \frac{1}{2} \eta_{\mu\nu} \square h. \quad (2.6)$$

Five constraint equations are then required to reduce the ten degrees of freedom in the symmetric $h_{\mu\nu}$ field, and can be obtained by acting with ∂_μ and $\frac{1}{2} \eta_{\mu\nu} + \frac{1}{m_h^2} \partial_\mu \partial_\nu$ on the equation of motion:

$$\mathcal{E}_{\rho\sigma}^{\mu\nu} - m_h^2 (h_{\mu\nu} - h \eta_{\mu\nu}) = 0. \quad (2.7)$$

One subsequently arrives at the gauge conditions:

¹Given the photon energy $\mathcal{E}$ in the rest frame of $h(J = 1)$, the momentum q assumes the form $(0, 0, 0, 2\mathcal{E})$. Furthermore, with $\varepsilon_{1,2}$ being the only possibilities for the transverse polarisations of the massless photon, the term $\epsilon_{\mu\nu\rho\sigma} \varepsilon_1^{\prime\mu} \varepsilon_2^{\prime\nu} \varepsilon^\rho p^\sigma$ must be proportional to $\varepsilon^0 \sim |\mathbf{p}|/m_h = 0$ in this rest frame and is hence neglected in (2.4).

²This is also possible if the functions $F_i(p^2)$ in (2.4) are anti-symmetric and hence recover Bose symmetry (cf. e.g. [30, 31] for a discussion on Bose symmetry violation).

³It was however shown in [43], that there can be a consistent theory that violates this structure.

$$\begin{aligned}\partial_\mu h^\mu{}_\nu - \partial_\nu h &= 0, \\ h &= 0.\end{aligned}\tag{2.8}$$

From this, it is seen that the coefficients (in particular the mass term) are tuned to avoid ghost terms. The gauge invariance condition $\delta h_{\mu\nu} = \partial_\mu \xi_\nu + \partial_\nu \xi_\mu$ can be restored in the $m = 0$ case by introducing extra degrees of freedom, A_μ and χ , following the same pattern (so $h_{\mu\nu} \rightarrow h_{\mu\nu} + \partial_{(\mu} A_{\nu)}$ and $A_\mu \rightarrow A_\mu + \partial_\mu \chi$):

$$h_{\mu\nu} \rightarrow h_{\mu\nu} + \frac{1}{m} \partial_{(\mu} A_{\nu)} + \frac{1}{3m^2} \left(\partial_\mu \partial_\nu \chi + \frac{1}{2} m^2 \eta_{\mu\nu} \chi \right).\tag{2.9}$$

The mass term in (2.5) then becomes:

$$\begin{aligned}\mathcal{L} \supset & -\frac{1}{8} F_{\mu\nu} F^{\mu\nu} + \frac{1}{12} \chi \square \chi - \frac{1}{2} m_h^2 (h_{\mu\nu} h^{\mu\nu} - h^2) \\ & + \frac{1}{6} m_h^2 \chi^2 + \frac{1}{2} \chi h + m_h (h \partial_\mu A^\mu - h^{\mu\nu} \partial_\mu A_\nu) + \frac{m_h}{2} \chi \partial_\mu A^\mu,\end{aligned}\tag{2.10}$$

where the helicity states— $h_{\mu\nu}$ for helicity-two, A_μ for helicity-one and χ for helicity-zero, of the massive spin two states are manifest. Technically, the gauge symmetry is augmented by the Stückelberg sector [40]:

$$\begin{aligned}h_{\mu\nu} &\rightarrow h_{\mu\nu} + \partial_{(\mu} \xi_{\nu)} + \frac{1}{2} \eta_{\mu\nu} m \Lambda, \\ A_\mu &\rightarrow \partial_\mu \Lambda - m \xi, \\ \chi &\rightarrow \chi - 3m \Lambda,\end{aligned}\tag{2.11}$$

where the shift in the metric is to decouple the kinematic mixing. The Stückelberg formalism [47] then emphasises that gauge invariance are not real in the sense that they are just redundancies in the description of a theory.

2.2.1 Couplings to Matter

There are several issues in constructing interaction terms for spin-2 states (for more pedagogical discussions, see [48, 49]). This is mainly attributed to the lack of currents invariant under gauge symmetry associated with such higher spin states (cf. e.g. [50]). We study the coupling of a linear theory to a source $T_{\mu\nu}$ of the form:

$$\mathcal{L} = \mathcal{L}_{J=2} + \frac{\kappa}{2} h_{\mu\nu} T^{\mu\nu}.\tag{2.12}$$

Taking the motivation from the Einstein-Hilbert action (augmented by the SM Lagrangian):

$$S = \int d^4x \sqrt{-g} (R + \mathcal{L}_{SM}), \quad (2.13)$$

one arrives at (2.5) by expanding $\sqrt{g}R$ with $g_{\mu\nu} = \eta_{\mu\nu} + \frac{\kappa}{2}h_{\mu\nu}$. Then (2.12) can be viewed as an expansion of (2.13):

$$S[g_{\mu\nu}] = S[\eta_{\mu\nu}] + \frac{\kappa}{2} \int d^4x h_{\mu\nu} T_{\mu\nu} + \mathcal{O}(\kappa^2). \quad (2.14)$$

with the source identified as follows:

$$T_{\mu\nu} := - \frac{2}{\sqrt{-g}} \frac{\delta S}{\delta g_{\mu\nu}} \Big|_{g_{\mu\nu}=\eta_{\mu\nu}}, \quad (2.15)$$

The divergencelessness of the current is crucial in keeping diffeomorphism⁴. To demonstrate this, take as an example the source term for a massless scalar in (2.5):

$$T_{\mu\nu}^{(0)} = \partial_\mu \phi \partial_\nu \phi - \frac{1}{2} \eta_{\mu\nu} \partial_\alpha \phi \partial^\alpha \phi, \quad (2.16)$$

where the conservation of the energy momentum tensor and gauge invariance is ensured by the on-shell condition $\square\phi = 0$. However, due to the additional term in (2.12), the Klein-Gordon equation is modified to:

$$\square\phi = \kappa \left(\partial^\mu (h_{\mu\nu} \partial^\nu \phi) - \frac{1}{2} \partial_\mu (h^\nu{}_\nu \partial^\mu \phi) \right). \quad (2.17)$$

The theory is then inconsistent, with the conservation of $T_{\mu\nu}^{(0)}$ broken at order $\mathcal{O}(\kappa)$:

$$\partial^\mu T_{\mu\nu}^{(0)} = \kappa \partial_\nu \phi \left(\partial^\alpha (h_{\alpha\beta} \partial^\beta \phi) - \frac{1}{2} \partial_\mu (h^\alpha{}_\alpha \partial^\mu \phi) \right). \quad (2.18)$$

One needs to introduce couplings of order $\mathcal{O}(h^2)$, so that the (non-linear) self-energy of the spin-2 will be included in $T_{\mu\nu}$, equivalent to a second-order expansion in (2.14). Iterative applications of these corrections should recover (2.13) but with the diffeomorphism promoted to the fully non-linearly realised gauge symmetry (coordinate covariance) [37, 51–56].

To make the dynamics of different helicity components of the massive spin-2 more obvious, one uses again the Stückelberg extension (2.9) to generate the following terms in the Lagrangian:

$$\mathcal{L} \supset \frac{\kappa}{2} h_{\mu\nu} T^{\mu\nu} + \frac{\kappa}{2} \chi T^\mu{}_\mu - \frac{\kappa}{m_h} A_\mu \partial_\nu T^{\mu\nu} + \frac{\kappa}{m_h^2} \chi \partial_\mu \partial_\nu T^{\mu\nu}. \quad (2.19)$$

⁴We note that in the case of a massless and sourceless Einstein-Hilbert action, this follows as a natural consequence of the linearised Bianchi identity.

We see that if $T_{\mu\nu}$ is not divergenceless, the scalar and vector modes will be strongly coupled to $T_{\mu\nu}$ in the $m_h \rightarrow 0$ limit⁵. In addition, the helicity-zero component never decouples due to the second term, and is the origin of what is known as the vDVZ discontinuity [57, 58]. This was suggested not to be a problem per se (see e.g. [59]), since the non-linearities of the theory should be important at scales $\Lambda_V \sim (m_h^4 \Lambda^2 \kappa)^{1/5}$ [60], with Λ being the intrinsic cutoff of the linear theory. The real issue is that when these non-linear effects are included by changing the Pauli-Fierz structure, Boulware-Deser ghost [61] propagate as a sixth degree-of-freedom. The construction of a ghost-free theory is contrived [62–64] but will not be further pursued in our work. Furthermore, non-universal couplings to $h_{\mu\nu}$ will be required in our case, in order to explain the already observed coupling in the VV ($V = W, Z, g, \gamma$) channels, further spoiling the conservation of $T_{\mu\nu}$ and thus gauge invariance [52, 65–67].

If one accepts the Pauli-Fierz Lagrangian coupled to the SM via (2.12) as a valid EFT in the linear regime, our paper then works to extrapolate the energy scale at which perturbative unitarity breaks down⁶. This is to say that we are ignorant of the pathologies which appear beyond this bound, by setting it as a UV cutoff [50].

2.3 $hZ \rightarrow hZ$

In Sect. 1.2, it was pointed out that the SM Higgs boson uniquely moderates all the high energy scattering amplitudes. In the scenario where a spin-2 impostor mimics the current LHC signal, the coupling κ_i of the impostor to the W and Z bosons may be determined by the condition:

$$\Gamma_{J=2}(h \rightarrow VV^*) = \Gamma_{J=0}(h_{SM} \rightarrow VV^*). \quad (2.20)$$

We simply quote the result from our letter [23]:

$$\begin{aligned} \kappa_W^2 &= 3.16 \times 10^{-5} \text{ GeV}^{-2}, \\ \kappa_Z^2 &= 4.42 \times 10^{-5} \text{ GeV}^{-2}. \end{aligned} \quad (2.21)$$

The Feynman diagrams for the tree level elastic scattering $h_{LL}Z_L \rightarrow h_{LL}Z_L$ are shown in Fig. 2.1. Since the final states are distinguishable, only the s - and t -channels

⁵We see in (2.19) that non-conserved currents can still be compatible with gauge invariance, so long as this vanishes in the $m_h \rightarrow 0$ limit, as noted in [43].

⁶Tree level unitarity constraints have been studied in massive spin-2 theories [68].

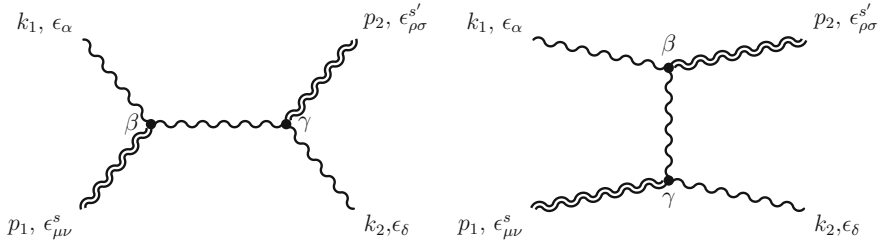


Fig. 2.1 Feynman diagrams depicting the scattering of the spin-2 impostor (*double wavy lines*) against the weak bosons (*single wavy lines*)

contribute. We now inspect the momentum dependences of various components of the Feynman diagram. Firstly, the appropriate hZZ Feynman rule may be obtained in [69]:

$$\begin{array}{c}
Z_\rho(k_1) \\
\diagdown \\
\text{---} \bullet \text{---} h_{\mu\nu}(p) \\
\diagup \\
Z_\sigma(k_2)
\end{array}
: -i \frac{\kappa}{2} \left((m_Z^2 + k_1 \cdot k_2) C_{\mu\nu,\rho\sigma} + D_{\mu\nu,\rho\sigma}(k_1, k_2) \right), \quad (2.22)$$

where:

$$\begin{aligned} C_{\mu\nu,\alpha\beta} &:= \eta_{\mu\alpha}\eta_{\nu\beta} + \eta_{\mu\beta}\eta_{\nu\alpha} - \eta_{\mu\nu}\eta_{\alpha\beta}, \\ D_{\mu\nu,\alpha\beta}(k_1, k_2) &:= \eta_{\mu\nu}k_{1\beta}k_{2\alpha} - \left[\eta_{\mu\alpha}k_{1\beta}k_{2\nu} + \eta_{\mu\beta}k_{1\nu}k_{2\alpha} - \eta_{\alpha\beta}k_{1\mu}k_{2\nu} + (\mu \leftrightarrow \nu) \right]. \end{aligned} \quad (2.23)$$

Then, by working in the rest frame of the propagator, the momenta of the incoming spin-1 and spin-2 state may be expressed as:

$$\begin{aligned} k_1^\mu &\approx p(1, 0, 0, 1), \\ p_1^\mu &\approx p(1, 0, 0, -1), \end{aligned} \quad (2.24)$$

respectively. The coordinate system can be chosen such that the final states lie in the $x - z$ plane, and that the outgoing spin-1 and spin-2 momenta are respectively given by:

$$\begin{aligned} k_2^\mu &\approx p(1, \sin\theta, 0, \cos\theta), \\ p_2^\mu &\approx p(1, -\sin\theta, 0, -\cos\theta), \end{aligned} \quad (2.25)$$

with θ measuring the scattering angle. Furthermore, the computation the scattering amplitude requires a specification of the polarisations vectors of the Z boson, which

have to satisfy:

$$\begin{aligned}\epsilon^\mu(p, \lambda) p_\mu &= 0, & \epsilon^\mu(p, \lambda) \epsilon_\mu(p, \lambda) &= 0, \\ \epsilon_\mu(p, -\lambda) &= \epsilon_\mu^*(p, \lambda), & \epsilon^\mu(p, \lambda) \epsilon_\mu(p, -\lambda) &= -1,\end{aligned}\tag{2.26}$$

where λ gives the corresponding helicity. The polarisation vectors for the incoming spin-1 states are then:

$$\begin{aligned}\epsilon_\mu^0 &\approx \begin{cases} \frac{p}{m_V}(1, 0, 0, -1), & \text{incoming,} \\ \frac{p}{m_V}(1, -\sin \theta, 0, -\cos \theta), & \text{outgoing,} \end{cases} \\ \epsilon_\mu^\pm &= \begin{cases} \frac{1}{\sqrt{2}}(0, -1, \mp i, 0), & \text{incoming,} \\ \frac{1}{\sqrt{2}}(0, -\cos \theta, \mp i, \sin \theta), & \text{outgoing.} \end{cases}\end{aligned}\tag{2.27}$$

On the other hand, the polarisation vectors for the spin-2 impostor read:

$$\epsilon_{\mu\nu}^s = \left\{ \epsilon_\mu^+ \epsilon_\nu^+, \frac{1}{\sqrt{2}} (\epsilon_\mu^+ \epsilon_\nu^0 + \epsilon_\mu^0 \epsilon_\nu^+), \frac{1}{\sqrt{6}} (\epsilon_\mu^+ \epsilon_\nu^- + \epsilon_\mu^- \epsilon_\nu^+ - 2\epsilon_\mu^0 \epsilon_\nu^0), \frac{1}{\sqrt{2}} (\epsilon_\mu^- \epsilon_\nu^0 + \epsilon_\mu^0 \epsilon_\nu^-), \epsilon_\mu^- \epsilon_\nu^- \right\},\tag{2.28}$$

which are in turns constructed from the corresponding spin-1 polarisation vectors:

$$\begin{aligned}\epsilon_\mu^0 &\approx \begin{cases} \frac{p}{m_G}(1, 0, 0, 1), & \text{incoming,} \\ \frac{p}{m_G}(1, \sin \theta, 0, \cos \theta), & \text{outgoing,} \end{cases} \\ \epsilon_\mu^\pm &= \begin{cases} \frac{1}{\sqrt{2}}(0, -1, \mp i, 0), & \text{incoming,} \\ \frac{1}{\sqrt{2}}(0, -\cos \theta, \mp i, \sin \theta), & \text{outgoing.} \end{cases}\end{aligned}\tag{2.29}$$

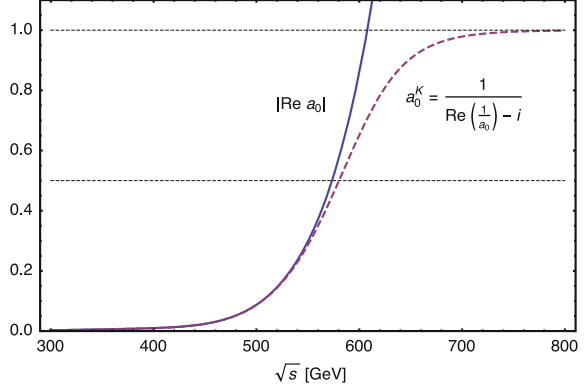
One can summarise the momentum dependences of the various components of the Feynman diagram as follows:

- Z longitudinal polarisation vector: $\mathcal{O}(p)$,
- h longitudinal polarisation tensor: $\mathcal{O}(p^2)$,
- hZZ vertex: $\mathcal{O}(p^2)$,
- Z propagator: $\mathcal{O}(1)$.

From these, one naively expects that $\mathcal{M}(Z_L h_{LL} \rightarrow Z_L h_{LL}) \sim \mathcal{O}(s^5)$ where $\sqrt{s}$ is the centre of momentum energy.

However, the calculation of the s - and t -channel processes yields a $\mathcal{O}(s^4)$ dependence (for more details, cf. Appendix A.2):

Fig. 2.2 The *solid line* shows the partial wave amplitude a_0 and the *dashed line* the unitarised amplitude using the K -matrix formalism



$$\begin{aligned}
 \mathcal{M}_s &= -\underbrace{\frac{32p^8\kappa^2}{3m_G^4(s-m_Z^2)}}_{\text{off diagonal}} + \frac{32p^6\kappa^2m_Z^2(1-\cos\theta)}{3m_G^4(s-m_Z^2)}, \\
 \mathcal{M}_t &= -\underbrace{\frac{2p^8\kappa^2(1+\cos\theta)^4}{3m_G^4(t-m_Z^2)}}_{\text{off diagonal}} - \frac{p^{10}\kappa^2\csc^2\frac{\theta}{2}\sin^6\theta}{3m_G^4m_Z^2(t-m_Z^2)},
 \end{aligned} \tag{2.30}$$

where terms from the off-diagonal part of the Z -propagator are indicated. The zeroth order partial wave was evaluated as per (1.48). It is clear from Fig. 2.2 that the unitarity bound (A.21) is violated at $\Lambda \sim 600$ GeV. We proceed to see what additional resonances may restore perturbative unitarity (cf. Sect. 1.1.2).

Whilst an addition resonance with $J = 1, 2$ or 3 are expected from naïve spin considerations, we restrict ourselves to the spin-1 case because the higher spin cases lead to amplitudes of higher momentum dependence and bring forward the onset of unitarity violation. Such a spin-1 Z' -boson is introduced by a kinetic-mixing term associated with the SM Z -boson:

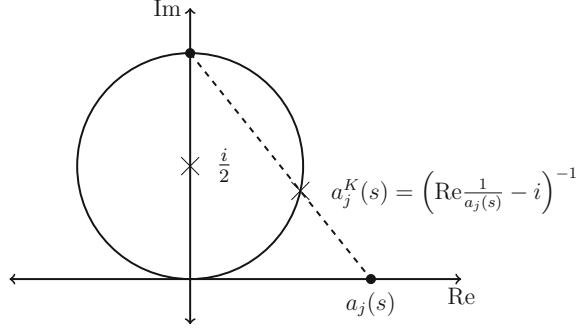
$$\mathcal{L} \supset -\frac{1}{2}Z_{\mu\nu}Z'^{\mu\nu}, \tag{2.31}$$

The Feynman rule of the hZZ' vertex can then be appropriately extracted as follows:

$$-i\kappa' \left[(k_Z \cdot k_{Z'}) C_{\mu\nu,\rho\sigma} + D_{\mu\nu,\rho\sigma}(k_Z, k_{Z'}) \right]. \tag{2.32}$$

The mixing term then contributes to $h_{LL}Z_L \rightarrow h_{LL}Z_L$ scattering through the s - and t - channels amplitudes which read:

Fig. 2.3 K -matrix formalism projects a real amplitude onto the Argand circle corresponding to (Appendix A.18)



$$\begin{aligned}\mathcal{M}_s^{Z'} &= \frac{8\kappa'^2 p^6 m_Z^2 (1 - \cos \theta)}{3m_G^4 (s - m_{Z'}^2)}, \\ \mathcal{M}_t^{Z'} &= -\frac{\kappa'^2 p^6 \csc \frac{\theta}{2} \sin^4 \theta (p^4 \sin^2 \theta + 2p^2 m_Z^2 \cos \theta + m_Z^4)}{m_G^4 m_Z^2 (t - m_{Z'}^2)}.\end{aligned}\tag{2.33}$$

The top two panels of Fig. 2.4 show the amplitudes with $\kappa = \kappa'$ and the bottom, with $\kappa = i\kappa'$. The latter relation between the couplings is chosen in attempt to cancel the leading order momentum dependence. It can be observed that there is no symmetry present in the $Z - Z'$ sector that results in the cancellation of the power law dependence. Subsequently, we investigated whether this problem can be cured with another unitarisation method.

The K -matrix ansatz [70–72] formally corresponds to adding infinitely heavy and wide resonances which unitarise the amplitude by projecting it onto the Argand circle as shown in Fig. 2.3. In relation to this, [73] showed that the amplitude from the higher order expansion of a chiral perturbation theory converges towards the unitary circle and that the inverse amplitude methods give a good parametric description. It was subsequently argued that when tree-unitarity violation is lower than the onset of new physics, a chiral theory will heal itself with a resonant structure with scalar quantum numbers. In the context of the electroweak sector, this corresponds to the introduction of the Higgs. The saturation observed in Fig. 2.4 beyond the Z' resonance is due to unitarisation and indicates that a UV completion of the theory is required to correctly extrapolate the physics in this region [74]. These results demonstrate that a theory describing the 125 GeV resonance as a spin-2 Higgs impostor is ill from a unitarity viewpoint. Furthermore, violation of perturbative unitarity cannot be amended by adding further resonances within the weakly coupled regimes of the theory.

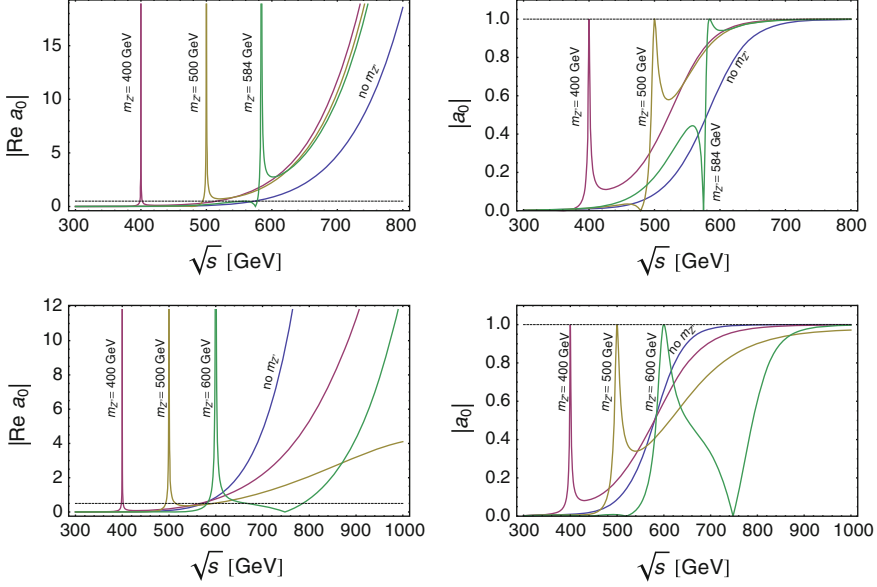


Fig. 2.4 The tree level $hZ \rightarrow hZ$ scattering amplitude with various masses of Z' added to the theory, assuming $\kappa = \kappa'$ (top) and $\kappa = i\kappa'$ (bottom). The corresponding unitarised amplitudes are shown on the right. Source [23]

2.4 STU Parameters

The contribution of the impostor to the STU parameters (cf. Sect. 1.2.2) is due to the rainbow diagram as shown in Fig. 2.5. We simply adapt the calculation done in [75] and remove the contributions from the massive KK tower [75]. This is detailed in our Letter and will not be reproduced here. Instead, we show some representative values of such modifications in Table. 2.1 and revised the associated figure from our later in Fig. 2.6. This makes clear that the measured oblique parameters cannot be made consistent with the measured value (cf. Ref. [76]) via removal of the SM Higgs contributions [77]:

Fig. 2.5 Rainbow diagram

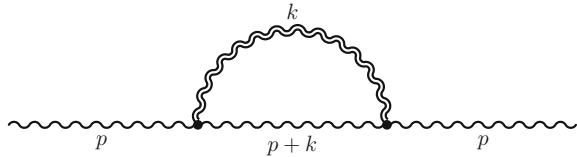


Fig. 2.6 Contribution of the spin-2 impostor and the SM Higgs to the STU oblique parameters. It is clear here that the impostor is not compatible with the experimental values

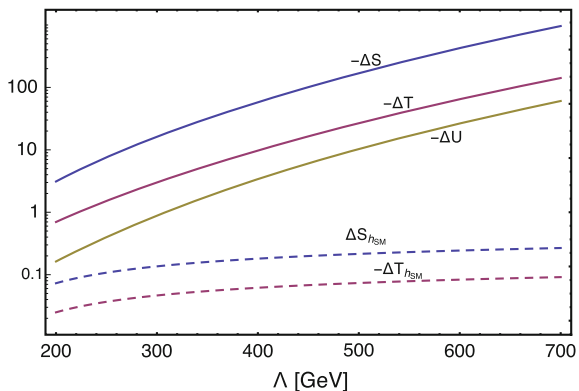


Table 2.1 Representative values of S , T and U parameters in the spin-2 Higgs impostor effective theory. Source [23]

Parameter/Energy cut-off (Λ)	200 GeV	400 GeV	600 GeV
ΔS	-3.11	-58.0	-426.0
ΔT	-0.698	-9.75	-64.6
ΔU	-0.163	-3.38	-26.5

$$\begin{aligned}
 \Delta S_{h_{SM}} &= \frac{1}{6\pi} \ln \left(\frac{\Lambda}{m_h^{ref}} \right), \\
 \Delta T_{h_{SM}} &= -\frac{3}{8\pi \cos^2 \theta_W} \ln \left(\frac{\Lambda}{m_h^{ref}} \right), \\
 \Delta U_{h_{SM}} &= 0,
 \end{aligned} \tag{2.34}$$

and replacement with those of a spin-2 impostor. It should be noted that here Λ denote the new physics scale expected at the unitarity violation scale ~ 600 GeV (cf. Sect. 2.3). The corrections to STU parameter due to the spin-2 impostor are sufficiently larger than those in [75] since the KK masses are sufficiently larger than that of the 125 GeV impostor ($m_{KK}^i \gg m_h$), and that these masses appear in the denominator of the self energy calculations.

2.5 Remarks

The argument against a minimally coupled spin-2 based on unitarity was generalised for more generic spin-2 couplings in our work [23]. Table.1 of [17] lists the most general tensor couplings for hZZ^* . The momentum dependences in the corresponding

amplitudes are the same, if not higher than the minimal case presented in this chapter, unless the coefficients are carefully tuned. We concluded that unitarity violation is expected below $\Lambda \lesssim 1$ TeV. This is applicable whenever the SM scalar is replaced with a spin-2 boson assuming generic couplings tuned to mimic the $h_{SM} \rightarrow WW$ and $h_{SM} \rightarrow ZZ$ decay rates. Furthermore, the resulting couplings between the spin-2 boson and the electroweak bosons cause severe disagreement with the measured STU precision parameters. The existence of a spin-2 Higgs impostor then cannot be accommodated in an effective theory and be made compatible with experiments.

References

1. J. Ellis, R. Fok, D.S. Hwang, V. Sanz, T. You, Distinguishing ‘Higgs’ spin hypotheses using $\gamma\gamma$ and WW^* decays. *Eur. Phys. J. C* **73**, 2488 (2013). [arXiv:1210.5229](#)
2. J. Ellis, V. Sanz, T. You, Associated production evidence against Higgs impostors and anomalous couplings. *Eur. Phys. J. C* **73**, 2507 (2013). [arXiv:1303.0208](#)
3. J. Ellis, V. Sanz, T. You, Prima facie evidence against spin-two Higgs impostors. *Phys. Lett. B* **726**, 244–250 (2013). [arXiv:1211.3068](#)
4. M.R. Buckley, M.J. Ramsey-Musolf, Diagnosing spin at the LHC via vector boson fusion. *JHEP* **09**, 094 (2011). [arXiv:1008.5151](#)
5. A. De Rujula, J. Lykken, M. Pierini, C. Rogan, M. Spiropulu, Higgs look-alikes at the LHC. *Phys. Rev. D* **82**, 013003 (2010). [arXiv:1001.5300](#)
6. U. De Sanctis, M. Fabbrichesi, A. Tonero, Telling the spin of the ‘Higgs boson’ at the LHC. *Phys. Rev. D* **84**, 015013 (2011). [arXiv:1103.1973](#)
7. C. Englert, D. Gonçalves, G. Nail, M. Spannowsky, The shape of spins. *Phys. Rev. D* **88**, 013016 (2013). [arXiv:1304.0033](#)
8. P. Artoisenet et al., A framework for Higgs characterisation. *JHEP* **11**, 043 (2013). [arXiv:1306.6464](#)
9. D. Boer, W.J. den Dunnen, C. Pisano, M. Schlegel, Determining the Higgs spin and parity in the diphoton decay channel. *Phys. Rev. Lett.* **111**, 032002 (2013). [arXiv:1304.2654](#)
10. J. Frank, M. Rauch, D. Zeppenfeld, Spin-2 resonances in vector-boson-fusion processes at next-to-leading order QCD. *Phys. Rev. D* **87**, 055020 (2013). [arXiv:1211.3658](#)
11. J. Frank, Higgs spin determination and unitarity of vector-boson scattering at the LHC. Ph.D. thesis, Karlsruhe Institute Technology, 2014
12. J. Frank, M. Rauch, D. Zeppenfeld, Higgs spin determination in the WW channel and beyond. *Eur. Phys. J. C* **74**, 2918 (2014). [arXiv:1305.1883](#)
13. A. Menon, T. Modak, D. Sahoo, R. Sinha, H.-Y. Cheng, Inferring the nature of the boson at 125–126 GeV. *Phys. Rev. D* **89**, 095021 (2014). [arXiv:1301.5404](#)
14. T. Modak, D. Sahoo, R. Sinha, H.-Y. Cheng, T.-C. Yuan, Disentangling the Spin-Parity of a resonance via the gold-plated decay mode. [arXiv:1408.5665](#)
15. J. Ellis, D.S. Hwang, V. Sanz, T. You, A fast track towards the ‘Higgs’ spin and parity. *JHEP* **11**, 134 (2012). [arXiv:1208.6002](#)
16. C. Englert, D. Gonçalves-Netto, K. Mawatari, T. Plehn, Higgs quantum numbers in weak boson fusion. *JHEP* **01**, 148 (2013). [arXiv:1212.0843](#)
17. S.Y. Choi, D.J. Miller, M.M. Muhlleitner, P.M. Zerwas, Identifying the Higgs spin and parity in decays to Z pairs. *Phys. Lett. B* **553**, 61–71 (2003). [arXiv:hep-ph/0210077](#)
18. G. Kramer, T.F. Walsh, Quasi two body e^+e^- annihilation. *Z. Phys.* **263**, 361–386 (1973)
19. Y. Gao, A.V. Gritsan, Z. Guo, K. Melnikov, M. Schulze, N.V. Tran, Spin determination of single-produced resonances at hadron colliders. *Phys. Rev. D* **81**, 075022 (2010). [arXiv:1001.3396](#)

20. S. Bolognesi, Y. Gao, A.V. Gritsan, K. Melnikov, M. Schulze, N.V. Tran et al., On the spin and parity of a single-produced resonance at the LHC. *Phys. Rev. D* **86**, 095031 (2012). [arXiv:1208.4018](#)
21. F. Maltoni, K. Mawatari, M. Zaro, Higgs characterisation via vector-boson fusion and associated production: NLO and parton-shower effects. *Eur. Phys. J. C* **74**, 2710 (2014). [arXiv:1311.1829](#)
22. F. Demartin, F. Maltoni, K. Mawatari, B. Page, M. Zaro, Higgs characterisation at NLO in QCD: CP properties of the top-quark Yukawa interaction. *Eur. Phys. J. C* **74**, 3065 (2014). [arXiv:1407.5089](#)
23. A. Kobakhidze, J. Yue, Excluding a generic spin-2 Higgs impostor. *Phys. Lett. B* **727**, 456–460 (2013). [arXiv:1310.0151](#)
24. ATLAS collaboration, Study of the spin of the new boson with up to 25 fb^{-1} of ATLAS data, ATLAS-CONF-2013-040
25. CMS collaboration, Combination of standard model Higgs boson searches and measurements of the properties of the new boson with a mass near 125 GeV, CMS-PAS-HIG-13-005
26. L.D. Landau, On the angular momentum of a system of two photons. *Dokl. Akad. Nauk Ser. Fiz.* **60**, 207–209 (1948)
27. C.-N. Yang, Selection rules for the dematerialization of a particle into two photons. *Phys. Rev.* **77**, 242–245 (1950)
28. M. Cacciari, L. Del Debbio, J. R. Espinosa, A. D. Polosa, M. Testa, A note on the fate of the Landau-Yang theorem in non-Abelian gauge theories. [arXiv:1509.07853](#)
29. W. Beenakker, R. Kleiss, G. Lustermans, No Landau-Yang in QCD. [arXiv:1508.07115](#)
30. AYu. Ignatiev, G.C. Joshi, M. Matsuda, The search for the decay of Z boson into two gammas as a test of Bose statistics. *Mod. Phys. Lett. A* **11**, 871–876 (1996). [arXiv:hep-ph/9406398](#)
31. S.N. Gninenko, AYu. Ignatiev, V.A. Matveev, Two photon decay of Z' as a probe of Bose symmetry violation at the CERN LHC. *Int. J. Mod. Phys. A* **26**, 4367–4385 (2011). [arXiv:1102.5702](#)
32. V. Pleitez, The angular momentum of two massless fields revisited. [arXiv:1508.01394](#)
33. J. P. Ralston, The need to fairly confront spin-1 for the new Higgs-like Particle. [arXiv:1211.2288](#)
34. S. Moretti, Variations on a Higgs theme. *Phys. Rev. D* **91**, 014012 (2015). [arXiv:1407.3511](#)
35. A. Koenigstein, F. Giacosa, D. H. Rischke, Classical and quantum theory of the massive spin-two field. [arXiv:1508.00110](#)
36. K. Hinterbichler, Theoretical aspects of massive gravity. *Rev. Mod. Phys.* **84**, 671–710 (2012). [arXiv:1105.3735](#)
37. C. de Rham, *Living Rev. Rel.* **17**, 7 (2014). [arXiv:1401.4173](#)
38. S.F. Hassan, R.A. Rosen, Resolving the ghost problem in non-linear massive gravity. *Phys. Rev. Lett.* **108**, 041101 (2012). [arXiv:1106.3344](#)
39. K. Hinterbichler, R.A. Rosen, Interacting spin-2 fields. *JHEP* **07**, 047 (2012). [arXiv:1203.5783](#)
40. S. Folkerts, A. Pritzel, N. Wintergerst, On ghosts in theories of self-interacting massive spin-2 particles. [arXiv:1107.3157](#)
41. L.P.S. Singh, C.R. Hagen, Lagrangian formulation for arbitrary spin. 1. The boson case. *Phys. Rev. D* **9**, 898–909 (1974)
42. L.P.S. Singh, C.R. Hagen, Lagrangian formulation for arbitrary spin. 2. The fermion case. *Phys. Rev. D* **9**, 910–920 (1974)
43. G. Dvali, O. Pujolas, M. Redi, Non Pauli-Fierz massive gravitons. *Phys. Rev. Lett.* **101**, 171303 (2008). [arXiv:0806.3762](#)
44. M. Fierz, W. Pauli, On relativistic wave equations for particles of arbitrary spin in an electromagnetic field. *Proc. Roy. Soc. Lond. A* **173**, 211–232 (1939)
45. W. Pauli, M. Fierz, On Relativistic Field Equations of Particles With Arbitrary Spin in an Electromagnetic Field. *Helv. Phys. Acta* **12**, 297–300 (1939)
46. P. Van Nieuwenhuizen, On ghost-free tensor lagrangians and linearized gravitation. *Nucl. Phys. B* **60**, 478–492 (1973)

47. E.C.G. Stueckelberg, Interaction forces in electrodynamics and in the field theory of nuclear forces. *Helv. Phys. Acta* **11**, 299–328 (1938)
48. R. Rahman, Higher Spin Theory - Part I. PoS **ModaveVIII**, 004 (2012). [arXiv:1307.3199](#)
49. A. Peach, Dissertation higher-spin gauge theories, Vasiliev theory and holography, Master's thesis, Durham University, 2013
50. M. Porrati, R. Rahman, Intrinsic cutoff and acausality for massive spin 2 fields coupled to electromagnetism. *Nucl. Phys. B* **801**, 174–186 (2008). [arXiv:0801.2581](#)
51. D.G. Boulware, S. Deser, Classical General Relativity Derived from Quantum Gravity. *Annals Phys.* **89**, 193 (1975)
52. M.D. Schwartz, *Quantum Field Theory and the Standard Model*, (Cambridge University Press, 2014)
53. N. Arkani-Hamed, H. Georgi, M.D. Schwartz, Effective field theory for massive gravitons and gravity in theory space. *Annals Phys.* **305**, 96–118 (2003). [arXiv:hep-th/0210184](#)
54. S. Deser, Selfinteraction and gauge invariance. *Gen. Rel. Grav.* **1**, 9–18 (1970). [arXiv:gr-qc/0411023](#)
55. R.M. Wald, Spin-2 fields and general covariance. *Phys. Rev. D* **33**, 3613 (1986)
56. T. Ortin, *Gravity and Strings*. Cambridge Monographs on Mathematical Physics (Cambridge University Press, 2015)
57. H. van Dam, M.J.G. Veltman, Massive and massless Yang-Mills and gravitational fields. *Nucl. Phys. B* **22**, 397–411 (1970)
58. V.I. Zakharov, Linearized gravitation theory and the graviton mass. *JETP Lett.* **12**, 312 (1970)
59. P. Creminelli, A. Nicolis, M. Papucci, E. Trincherini, Ghosts in massive gravity. *JHEP* **09**, 003 (2005). [arXiv:hep-th/0505147](#)
60. A.I. Vainshtein, To the problem of nonvanishing gravitation mass. *Phys. Lett. B* **39**, 393–394 (1972)
61. D.G. Boulware, S. Deser, Can gravitation have a finite range? *Phys. Rev. D* **6**, 3368–3382 (1972)
62. C. de Rham, G. Gabadadze, Generalization of the Fierz-Pauli action. *Phys. Rev. D* **82**, 044020 (2010). [arXiv:1007.0443](#)
63. C. de Rham, G. Gabadadze, A.J. Tolley, Resummation of massive gravity. *Phys. Rev. Lett.* **106**, 231101 (2011). [arXiv:1011.1232](#)
64. C. de Rham, G. Gabadadze, Selftuned massive spin-2. *Phys. Lett. B* **693**, 334–338 (2010). [arXiv:1006.4367](#)
65. S. Weinberg, E. Witten, Limits on massless particles. *Phys. Lett. B* **96**, 59 (1980)
66. S. Weinberg, Infrared photons and gravitons. *Phys. Rev.* **140**, B516–B524 (1965)
67. A. Jenkins, Constraints on emergent gravity. *Int. J. Mod. Phys. D* **18**, 2249–2255 (2009). [arXiv:0904.0453](#)
68. N. D. Christensen, Stefanus, On tree-level unitarity in theories of massive spin-2 bosons. [arXiv:1407.0438](#)
69. T. Han, J.D. Lykken, R.-J. Zhang, On Kaluza-Klein states from large extra dimensions. *Phys. Rev. D* **59**, 105006 (1999). [arXiv:hep-ph/9811350](#)
70. A.M. Badalian, L.P. Kok, M.I. Polikarpov, YuA Simonov, Resonances in coupled channels in nuclear and particle physics. *Phys. Rept.* **82**, 31 (1982)
71. W. Kilian, T. Ohl, J. Reuter, M. Sekulla, High-energy vector boson scattering after the Higgs discovery. *Phys. Rev. D* **91**, 096007 (2015). [arXiv:1408.6207](#)
72. S.U. Chung, J. Brose, R. Hackmann, E. Klempt, S. Spanier, C. Strassburger, Partial wave analysis in K matrix formalism. *Ann. Phys.* **4**, 404–430 (1995)
73. U. Aydemir, M.M. Anber, J.F. Donoghue, Self-healing of unitarity in effective field theories and the onset of new physics. *Phys. Rev. D* **86**, 014025 (2012). [arXiv:1203.5153](#)
74. A. Alboteanu, W. Kilian, J. Reuter, Resonances and unitarity in weak boson scattering at the LHC. *JHEP* **11**, 010 (2008). [arXiv:0806.4145](#)
75. T. Han, D. Marfatia, R.-J. Zhang, Oblique parameter constraints on large extra dimensions. *Phys. Rev. D* **62**, 125018 (2000). [arXiv:hep-ph/0001320](#)

76. M. Baak, M. Goebel, J. Haller, A. Hoecker, D. Kennedy, R. Kogler et al., The electroweak fit of the standard model after the discovery of a new boson at the LHC. *Eur. Phys. J. C* **72**, 2205 (2012). [arXiv:1209.2716](#)
77. M. Baak, M. Goebel, J. Haller, A. Hoecker, D. Ludwig, K. Moenig et al., Updated status of the global electroweak fit and constraints on new physics. *Eur. Phys. J. C* **72**, 2003 (2012). [arXiv:1107.0975](#)

Higgs Properties at the LHC

Implications for the Standard Model and for Cosmology

Yue, J.T.S.

2017, XVIII, 130 p. 35 illus., 21 illus. in color., Hardcover

ISBN: 978-3-319-63401-2