Prospects for Higgs physics at energies up to 100 TeV

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Abstract

We summarize the prospects for Higgs boson physics at future proton–proton colliders with centre of mass (c.m.) energies up to 100 TeV. We first provide the production cross sections for the Higgs boson of the Standard Model from 13 TeV to 100 TeV, in the main production mechanisms and in subleading but important ones such as double Higgs production, triple production and associated production with two gauge bosons or with a single top quark. We then discuss the production of Higgs particles in beyond the Standard Model scenarios, starting with the one in the continuum of a pair of scalar, fermionic and vector dark matter particles in Higgs-portal models in various channels with virtual Higgs exchange. The cross sections for the production of the heavier CP-even and CP-odd neutral Higgs states and the charged Higgs states in two-Higgs doublet models, with a specific study of the case of the Minimal Supersymmetric Standard Model, are then given. The sensitivity of a 100 TeV proton machine to probe the new Higgs states is discussed and compared to that of the LHC with a c.m. energy of 14 TeV and at high luminosity.

Keywords: Higgs, LHC, Hadron collider, standard model, supersymmetry, dark matter

(Some figures may appear in colour only in the online journal)

1. Introduction

It has been expected for a long time that the probing of the electroweak symmetry breaking (EWSB) mechanism would be at least a two-step process. The first step was the search and the (non) observation of a Higgs-like particle that would confirm (refute) the hypothesis of the standard model (SM)—that the electroweak symmetry is spontaneously broken by a scalar field that develops a non-zero vacuum expectation value—and many of its new physics extensions [1–4]; for a review see [5]. This test has been passed successfully with the historical discovery by the ATLAS and CMS collaborations [6, 7] of a Higgs boson with a mass of 125 GeV at the CERN LHC. A second step, that is as important as the first one, is to probe in all its facets the EWSB mechanism and to assess whether the Higgs sector of the theory is SM-like or involves new degrees of freedom. Compared to what has been achieved so far, this latter step needs first a much more precise determination of the basic properties of the observed Higgs particle, and second a more complete probe of the TeV scale in order to directly discover or definitely exclude the light new degrees of freedom that are expected to appear in almost all extensions of the SM. This is particularly true in models that allow for a natural Higgs boson with a mass that is protected against very high scales and, hence, solve the so-called hierarchy problem.
All the measurements in the Higgs sector performed so far [8] need to be significantly improved in order to better constrain the SM and to probe smaller effects of new physics. In particular, a more precise measurement of the Higgs couplings to the fermions and gauge bosons should be performed in the channels already probed. One thus needs to experimentally determine more precisely the cross sections in the dominant Higgs production modes [5, 9] such as the loop induced gluon fusion mechanism $gg \to H$ with the most easily accessible decay modes—e.g. $H \to \gamma\gamma$, $H \to ZZ^{*} \to 4\ell$, $H \to WW^{*} \to 2\ell 2\nu$ with $\ell = e, \mu$. Sub-dominant but very important channels such as vector boson fusion $qq \to Hqq$ with in particular the $H \to \tau\tau$ decay channel and Higgs-strahlung $q\bar{q} \to HW$, $HZ$ with $H \to bb$ decays need to be more thoroughly investigated.

The goal would be to reach an accuracy at the percent level, which is the size of the higher order electroweak corrections and presumably also that of the potential new physics effects (see for instance [10]).

Additional production channels—such as associated Higgs production with top quark pairs, $pp \to t\bar{t}H$, that would allow for a direct determination of the top-quark Yukawa coupling and might provide an unambiguous determination of the Higgs CP-properties, and rare decay modes such as $H \to Z\gamma$, which could give complementary information to the $H \to \gamma\gamma$ decay and $H \to \mu^{+}\mu^{-}$, which would test for the first time Higgs couplings to other fermions than those of the third generation—still need to be probed. Other higher order processes for single Higgs boson production, such associated production with two vector bosons or a single top quark might also provide interesting information.

Of prime importance would be the measurement of the Higgs self-coupling that allows for the complete determination of the Higgs potential—which is responsible for EWSB. The Higgs trilinear coupling can in principle be studied by measuring the rate for double Higgs production which, unfortunately, is very small at the LHC. The quartic Higgs couplings can only be accessed in triple Higgs production which is hopeless at the LHC.

Beyond the SM context, one needs to search for new Higgs bosons with mass scales that are larger than those probed so far and/or to detect them in modes that have not yet been considered. This would be the case, for instance, of the additional neutral and charged Higgs bosons predicted in two-Higgs doublet models (2HDM) [11, 12], as realized in the minimal supersymmetric extension of the SM (MSSM) [11, 13]. In these models, four additional physical Higgs bosons are present, besides the one already observed and that we will now denote by $h$: a heavier CP-even $H$ state, a CP-odd $A$ state and two charged $H^{\pm}$ bosons. The four Higgs masses $M_{h}, M_{H}, M_{A}$ and $M_{\mu}$, as well as the ratio of the vacuum expectation values of the two scalar fields $\tan \beta = v_{2}/v_{1}$ and the mixing angle $\alpha$ that diagonalizes the two CP-even Higgs states, are unrelated in a general 2HDM. In the MSSM however, supersymmetry imposes strong constraints on the parameters and, in fact, only two of them (e.g. $\tan \beta$ and $M_{A}$) are independent at tree level.

In the MSSM at high $A$ masses, one is in the decoupling regime in which the lighter CP-even $h$ state will have almost SM-like couplings while the four states $HA$ and $H^{\pm}$ become very heavy, degenerate in mass and decouple from the massive gauge bosons [14]. In the 2HDM, to cope naturally with the fact that the observed $h$ boson is SM-like, one invokes the so-called alignment limit [15–18] in which only one Higgs doublet gives mass to the $V = WZ$ bosons. In this case, the mixing angle $\alpha$ is such that the $h$ couplings are the same as the ones of the SM Higgs state, i.e. $\alpha = \beta = \frac{\pi}{2}$. In this case too, the $H$ state will no longer couple to massive gauge bosons as also does the pseudoscalar $A$ boson in CP-invariant theories. Hence, in both 2HDMs and the MSSM, the additional Higgs states need to be searched for at relatively high masses and in channels that do not involve the vector bosons and, hence, are rather complicated experimentally. Because there is a sum-rule requiring that the sum of the coupling squared of all the CP-even Higgs particles to $WZ$ bosons should add up to the SM-Higgs coupling squared $g_{HVV}^{2}$ [11], most other Higgs extensions in which the light $h$ state is forced to be SM-like will have a phenomenology that is similar to that of 2HDMs and the MSSM. This would be for instance the case of the next-to-MSSM (NMSSM) in which the additional singlet-like states almost decouple from fermions and gauge bosons if $h$ is SM-Higgs like [19, 20].

Another interesting example of a new physics extension is the so-called Higgs portal dark matter (DM) scenario [21–24]. One postulates the existence of a cosmologically stable, massive and weakly interacting neutral particle that would account for the DM in the universe and which couples only to the observed Higgs boson. This particle is thus undetectable and would appear only as missing energy when produced in association with SM particles. If the DM particle is light, $M_{DM} \lesssim \frac{1}{2}M_{h}$, it will appear in the decays of the observed Higgs boson, the latter being produced in the usual production channels, except for the dominant process $gg \to h$ in which an additional jet in the final state is required in order to make the process visible. These processes have been intensively discussed, either directly by searching for missing energy signals or indirectly through the measurement of the observed Higgs signal strengths; see e.g. [25] a recent review. In turn, if the DM particles are heavier than $\frac{1}{2}M_{h}$, the Higgs boson should be virtual, $h \to hDM$, implying that the production processes are at higher order in perturbation theory and have rather small production rates. Hence, as in the case of the MSSM or 2HDM, the production cross sections of the new states are too low at present energies and luminosities to allow for the probing of the entire parameter space allowed of these portal models.

To have more sensitivity to new physics phenomena, one needs a significantly larger sample of Higgs bosons to be produced, and two options are then at hand. The first is a large increase of the integrated luminosity. This is the high-luminosity LHC option (HL-LHC) on which there is presently a wide consensus in the community: collecting up to 3 ab$^{-1}$ of data at a centre of mass (c.m.) energy of $\sqrt{s} = 14$ TeV should be the next step for the LHC and the particle physics goal for the next decade. Prospects in the context of Higgs physics for the HL-LHC option have been discussed in detail in [26, 27].
However, in many cases and in particular for the probing of high mass scales (such as the production of new Higgs bosons with TeV masses) or very rare processes for the already observed Higgs state (as double Higgs production), this high-luminosity option will not be sufficient. A more radical option would be a significant increase in the c.m. energy. In this context, an upgrade of the LHC to an energy about 2–3 times higher has been discussed and, for instance, detailed studies of the physics of a $\sqrt{s} = 33$ TeV collider have been performed [28]. More recently, a Future Circular Collider (FCC-hh), a hadron collider with a c.m. energy of 100 TeV, has been proposed as a potential follow-up of the LHC at CERN\(^4\); such a very high energy machine is also under study in China [29].

The discovery potential of such machines\(^5\) is immense\(^6\) and, in the context of new Higgs states for instance, the search reach would be a factor three to six larger than the LHC depending on the considered (pair or singly produced) particle. Even in the case of the observed SM-like Higgs boson, the increase in the production cross sections allowed by the higher energy would be more than an order of magnitude compared to the 14 TeV LHC, and high luminosities would provide a sufficient data sample to probe very difficult channels such as Higgs pair production that cannot be seriously considered even at the HL-LHC. In the already probed Higgs channels, a very high precision can be achieved in the determination of the Higgs properties at such machines.

In fact, with the expected accuracy, major theoretical improvements in the determination of the Higgs cross sections and decay branching ratios (besides the recent improvements in the gluon fusion channel for single and double Higgs production that will be discussed in the present paper) need to be performed. In particular, more precise determinations of the gluon and quark structure functions and of the strong coupling $\alpha_s$ would be required to reduce the uncertainties in the cross sections [9, 32], along with a better determination of the bottom (and eventually charm) quark mass—which, together with $\alpha_s$, are the main source of uncertainties in the Higgs decay branching ratios [32–34]. Progresses in these directions are expected to occur in the coming decades. However, to discuss in detail the potential of the high-energy machine, it is still useful to provide an estimate of the theoretical precision with which all observables are currently known in order to quantify the possible improvement in the measurements at these colliders and to make comparisons with other options such as, for instance, an electron–positron collider with $\sqrt{s} \gtrsim 250$ GeV.

These are the issues that will be analyzed in this paper. While several analyses have addressed some of the processes [35–41], we will attempt to perform a comprehensive analysis of Higgs physics at a 100 TeV hadron collider, not only for Higgs production in the SM, but also in some well motivated extensions such as 2HDM/MSSM and Higgs-portal DM scenarios. The paper is organized as follows. In the next section, we discuss the production of the SM Higgs boson: single and double production in the main channels and single production in subleading channels such as associated production with two vector bosons or single top quarks; in the case of the main channels, an attempt to estimate the theoretical uncertainties is made. The production of DM particles through the already observed SM-like Higgs boson in Higgs-portal models will be analysed in section 3 and numerical results for the production cross sections in the various channels will be presented in the case of spin 0, 1 and 1 DM states. In section 4, we analyze the main processes for producing the heavier neutral and charged Higgs particles in 2HDMs like the MSSM and both single and pair production processes will be addressed; the parameter space that could be probed at 100 TeV will be delineated. A short conclusion will be given in the final section 5.

2. Production of the SM Higgs boson

2.1. Single Higgs production and higher order corrections

2.1.1. Gluon fusion. The main production process for a single Higgs boson at hadron colliders is gluon fusion, $gg \to H$. This process proceeds at the quantum level already at leading order (LO) [42] with triangular top and bottom quark loops from which the $H$ state is emitted. In the Standard Model, the top quark loop contribution is by far dominating; the bottom loop contribution is rather small and its interference with the dominant top loop does not exceed 10% of the total contribution. The next-to-leading order (NLO) QCD corrections were calculated two decades ago, first in the infinite top mass approximation $M_H \ll 2m_t$ [43, 44] by applying the low energy theorem which related the $Hgg$ amplitude to the QCD $\beta$ function [45, 46] and then using the exact quark mass dependence in the loop [47]. It was shown that if the LO cross section contains the exact top-quark mass dependence, the two results approximately agree with each other. This is particularly true in the Higgs mass range $M_H \lesssim 2m_t$ where the LO Hgg amplitude does not develop an imaginary part.

The NLO QCD corrections are found to be very large, with a $K$-factor\(^7\) of around 2 for $M_H = 125$ GeV at a c.m. energy $\sqrt{s} = 8$ TeV. The next-to-next-to-leading order (NNLO) QCD corrections were computed in the infinite top mass limit [48–50], leading to an increase of about 25% for the total cross section. Later the NNLO corrections have been evaluated in a top mass expansion [51–54] and it was found that the limit $M_H \ll 2m_t$ is again good for $M_H \lesssim 350$ GeV. The resummation of the soft gluons at next-to-next-to leading logarithm (NNLL) [55] leads only to a moderate increase of the cross section if the central value of the renormalisation and factorisation scales is chosen to be at $M_H = m_g = m_F = \frac{2}{3}M_H$ which, in passing improves the convergence behaviour of the perturbative series [56]. Note, however, that the calculation of the

\(^4\)See the web site https://fcc.web.cern.ch/Pages/default.aspx

\(^5\)These energies have been considered in the late 1980s: $\sqrt{s} = 33$ TeV is close to the energy that was foreseen for the late Superconducting Super Collider (SSC) ($\sqrt{s} = 40$ TeV) and $\sqrt{s} = 100$ TeV is a factor of two lower than the energy of the Eloisatron collider which was proposed in Europe at that time [30].

\(^6\)A review of the physics potential of such a machine in various scenarios, with very little overlap with our analysis here, has appeared very recently [31].

\(^7\)The $K$-factor is defined as the ratio of cross sections at the higher order (HO) and the lowest order (LO), $K = \sigma_{HO}/\sigma_{LO}$ when $\alpha_s$ and the PDFs are consistently evaluated at the respective perturbative orders.
subleading bottom quark loop contribution beyond the NLO approximation is still lacking as one cannot use the effective field theory (EFT) approach with an infinitely large loop mass as in the case of the top contribution.

The gluon fusion process has now reached an impressive accuracy with the very recent completion of the calculation of the next-to-next-to-next-to-leading order (N^3LO) QCD corrections [57] after a huge amount of theoretical efforts in the past few years [58–73]. For the $M_H = 125$ GeV Higgs boson, the N^3LO corrections amount to a $\approx 2\%$ increase of the cross section at both $\sqrt{s} = 8$ TeV and $\sqrt{s} = 14$ TeV but with a reduction of the scale uncertainty to about $3\%$ as will be discussed later. Improved predictions at N^3LO using renormalisation group equations have also been calculated very recently [74] and jet-vetoed cross section has also reached the N^3LO + NNLL accuracy [75]. Soft-gluon resummation has been recently extended to N^3LL order, leading to an increase of $\sim 5\%$ over the fixed-order results, while mass effects up to NLL are found to be negligible [76].

The electroweak (EW) radiative corrections have been computed at NLO, first in the infinite loop mass limit $m_t, M_V \gg 2M_H$ [77–79] and subsequently in an exact calculation [80, 81]. Approximate mixed QCD-EW corrections at NNLO are also available [56] in the EFT approach with the infinite loop mass limit, $M_H \ll M_V$ in this case. Both types of corrections amount to a few percent.

Currently no public release of the tools used to calculate the N^3LO corrections is available. Since the N^3LO QCD correction turns out to be quite small, we will thus stick to the NNLO QCD order supplemented with the mixed NNLO QCD + EW corrections and use the program iHixs [82] with a central scale $\mu_0 = \frac{1}{2} M_H$. This scale choice, which as mentioned earlier accounts for the next-to-next-to-leading logarithmic (NNLL) increase of the total cross section [55, 56, 83], is motivated by the better convergence of the perturbative expansion as it was shown in [56] and stressed again at N^3LO [57].

### 2.1.2. Vector boson fusion.

The vector boson fusion (VBF) channel, in which the Higgs boson is produced in association with two jets, $q q \rightarrow V V^* q q \rightarrow H q q$, is a pure electroweak process at LO [84–86]. The central scale in this process is usually chosen to be $\mu_0 = Q^0$, the momentum transfer of the fusing weak bosons. For the fully inclusive cross section, the NLO QCD corrections have been known for a while in both the structure function approach [87–89] and exactly [90]; they give rise to an $O(10\%)$ increase of the total cross section. The NNLO QCD corrections in the structure function approach were computed a few years ago and found to be rather small, below the percent level [91, 92] for the inclusive rate, while they can be much more sizeable when cuts are included and in the differential distributions as found recently [93]. NLO EW corrections are also available [94, 95] and yield a shift of the order of $5\%$ in the cross section. The radiative corrections to the fully inclusive cross section are thus moderate and well under control.

However, this process is interesting only when some specific cuts are employed to single out the VBF topology, namely forward jets with high transverse momenta, a large rapidity gap between the two forward jets, central Higgs decay products and no jet activity in this central area. This is necessary in order to suppress the QCD background and to minimize the contamination from the contribution at NNLO of the gluon-fusion process, $gg \rightarrow H j j$, which leads to the same final state. In this case the NLO corrections are also known exactly [90] and lead to the same enhancement of the cross section as for the inclusive case. However, a recent calculation of the QCD corrections at NNLO [93] shows that their contribution amounts to $O(10\%)$ and are thus non negligible, in contrast to the case for the inclusive cross section. More importantly, the variation of the VBF cross section with the renormalisation and factorisation scales is much larger than in the inclusive case.

Nevertheless, in our study, we will stick to the inclusive cross section and, given the smallness of the NNLO QCD corrections, we will give our predictions only at the NLO QCD + EW approximation using the Monte-Carlo program VBFNLO [96, 97].

#### 2.1.3. Higgs-strahlung channels.

The Higgs-strahlung processes $q \bar{q} \rightarrow V H$ with $V = W^\pm$ or $Z$ are also pure EW processes at LO. The NLO QCD corrections are actually pure Drell–Yan corrections [98] to the process $q \bar{q} \rightarrow V^* [88, 89], [99–101]$ and for the associated production with a $W$ boson this extends up to NNLO [102] by the use of the classic NNLO results for Drell–Yan [48, 103]. On top of these contributions there are NNLO QCD corrections where the Higgs is radiated off the top loops; they amount to $\sim 1\%$ for both processes [102]. In the case of the associated production with a $Z$ boson there is, in addition to the previously discussed contributions, the opening at NNLO of the channel $gg \rightarrow HZ$ not going through a virtual weak boson [104] that has been extended to include the real radiation effects that describe much better the kinematics of this subprocess [105]. The threshold N^3LO corrections are also available for the hadronic process [106].

The NLO + NNLO QCD Drell–Yan type corrections are moderate, of the order of at most $+35\%$ and the gluon fusion contributions in $ZH$ production are of the order of $10\%$ at $\sqrt{s} = 14$ TeV. In the past few years there have been improvements in the description of the $gg \rightarrow HZ$ subprocess with NLO QCD corrections and soft-gluon resummation that reduce the uncertainties in this subprocess [107, 108]. The NLO EW corrections are negative and reduce the cross section by $3–8\%$ at LHC energies [109]. They are computed in the so-called $G_F$ scheme where the electromagnetic coupling constant $\alpha$ is derived from the Fermi constant $G_F$ such that the corrections are not sensitive to the value of the light quark masses. The combination of the EW and NLO corrections follows [110]:

\[
\begin{align*}
\sigma^{QCD+EW}_{WH} &= \sigma^{QCD}_{WH} (1 + \delta^{EW}_{WH}), \\
\sigma^{QCD+EW}_{ZH} &= \sigma^{QCD,DY}_{ZH} (1 + \delta^{EW}_{ZH}) + \sigma_{gg\rightarrow ZH}.
\end{align*}
\]  

(2.1)
For the computation of the cross-section we will use the computer program \( vh@nnlo \) [111] which includes the full NNLO QCD+NLO EW corrections. We will then neglect the improvement in the gluon fusion subchannel for \( ZH \) production, which is formally a N\(^3\)LO contribution to the whole hadronic process.

2.1.4. Associated production with a heavy quark pair. Associated Higgs production with top or bottom quark pairs are the channels most affected by QCD backgrounds. The LO cross section for \( t\bar{t}H \) production was computed decades ago [112–115] and NLO QCD corrections are known to be modest provided that the central scale \( \mu_0 = \frac{1}{2}M_H + m_t \) [116–118] is used. In recent years the NLO EW corrections have been computed [119–121] and soft-gluon resummation is also available [122] and has extended to NNLL accuracy and has been used to obtain approximate NNLO results that are only slightly larger than NLO results but with an uncertainty band reduced by a factor of two [123].

The associated production with \( b\bar{b} \)–quark pairs [112, 113], in contrast, displays a different behaviour. The NLO QCD corrections can be calculated in the same way as for \( t\bar{t}H \) production but turn out to be rather large [124, 125]. This is due to the large logarithms generated by the integration of the transverse momenta of the final-state \( b\bar{b} \)-quarks.

These large logarithms can be resummed by considering the bottom quark as a massless constituent of the quark and using the Altarelli–Parisi evolution [126] of the bottom quark parton distribution function (PDF) to the scale of the process. Then one works in a five-flavour scheme (5FS) and the LO process that needs to be considered is \( bbH \) [127]. Requiring a high-\( p_T \) final-state \( b \)–quark requires the NLO QCD corrections [128–131] and the NNLO QCD corrections reintroduce the process \( gg \to bbH \) [132]. It turns out that choosing \( \mu_0 = \frac{1}{2}M_b \) as the factorisation scale and using the running bottom quark mass at the scale of the Higgs boson mass greatly improves the perturbative convergence of the series [131, 132]. The fully exclusive \( gg \to bbH \) process, calculated with four active parton flavours in a four-flavour scheme (4FS), and the 5FS process \( bb \to H \), should converge against the same value at higher perturbative orders. The NLO EW corrections were also calculated in the past decade and found to be modest [133]; we will then neglect them. We will also not include the very recent N\(^3\)LO corrections that are not yet implemented in a public code. They have been found to reduce further the scale uncertainty of the predictions [134, 135].

In our numerical analysis of the cross sections, we will use the program \( bbH@NNLO \) [132] for the associated production with bottom quarks at NNLO QCD and we will use MacGraph5 in the \( aMC@NLO \) framework [37] for the NLO QCD corrected cross sections of the associated production with top quarks.

2.2. The cross sections including the theoretical uncertainties

The production cross sections are affected by theoretical uncertainties which are in general divided into two categories, with an additional one in the \( gg \to H \) case.

(i) The scale uncertainty, which reflects the dependence of the truncated expansion of the cross section at a given perturbative order in \( \alpha_s \) on the renormalisation scale \( \mu_R \) that defines \( \alpha_s \) and on the factorisation scale \( \mu_F \) at which the matching of the perturbative calculation (the matrix element) and the (non perturbative) structure functions is made. This is generally estimated by varying the two scales in the interval

\[
\frac{1}{2} \mu_0 \leq \mu_R, \mu_F \leq 2 \mu_0 \tag{2.2}
\]

with \( \mu_0 \) being the central scale chosen for a given process and with some restrictions on the ratio \( \mu_R/\mu_F \) depending on the process\(^8\). For the various processes, we have adopted the following central scales in our analysis:

\[
\mu_0^{gg \to H} = \frac{1}{2}M_H, \quad \mu_0^{gg \to H|q\bar{q}} = Q_V, \quad \mu_0^{q\bar{q} \to VH} = M_{VH},
\]

\[
\mu_0^{q\bar{q}/gg \to H} = m_t + \frac{1}{2}M_H. \tag{2.3}
\]

In the case of \( bb \to H \) production we use the following set-up for the central scales:

\[
\mu_R^{bb \to H} = M_H, \quad \mu_F^{bb \to H} = M_H/4. \tag{2.4}
\]

For all processes, the scale uncertainty is derived by a variation of the renormalisation and factorisation scales within a factor of two from the central scale, as in equation (2.2). In the case of the \( bb \to H \) process however, following the LHC Higgs Cross Section Working Group framework [9] and given the asymmetry in the choice of the scales \( \mu_R \) and \( \mu_F \), we use the following domains of variation: \( \frac{1}{5}M_H \leq \mu_R \leq 5M_H, \quad \frac{1}{10}M_H \leq \mu_F \leq 0.7M_H \). The asymmetrical choice of the central renormalisation and factorisation scales for the \( bb \to H \) process is justified by the requirement of a good agreement between the 4FS and the 5FS [9]. This is also supported by the scale dependence studied in [132]. This requirement also drives the choice of a larger scale variation interval as exemplified in [138], and by the wish to have a reliable estimate of the scale uncertainties at NNLO given the very small dependence of the total cross section on the renormalisation scale.

(ii) The PDF and \( \alpha_s \) uncertainty stemming from the impact of the uncertainties in the modeling and the experimental data used in the fit that provides the PDF sets and the value of the strong coupling constant \( \alpha_s \). With the cross sections for most channels now known at NNLO and even beyond, thereby reducing the scale uncertainty, this is becoming the major source of theoretical errors. According to usual practice, we will consider the 90\% CL correlated PDF and 3\% \( \alpha_s \) uncertainties using the

\[\text{In some cases, like in the gluon fusion process which is subject to large QCD corrections, this domain of variation is sometimes considered to be too small. Other approaches to estimate the scale uncertainty have been proposed in this case; see for instance the approach of [136, 137] for which one obtains results that are similar to those obtained by extending the domain of scale variation to } \frac{1}{2} \mu_0 \leq \mu_R, \mu_F \leq 3 \mu_0.\]
MSTW2008 PDF sets [139, 140] and the following value for $\alpha_s$:

$$\alpha_s(M_Z^2) = 0.1171 \pm 0.0014 \text{ (68\% CL)} \text{ or } \pm 0.0032 \text{ (90\% CL)}$$

at NNLO

(2.5)

(iii) In the case of the gluon fusion channel, a third uncertainty related to the use of the effective field theory approach to account for the corrections beyond NLO should also be included as done in [32, 141] for the Tevatron and the LHC respectively. Indeed, the infinite quark mass limit is not adequate for the bottom loop, which generates a $\approx 10\%$ contribution when it interferes with the top loop contribution. The exact $b$-quark contribution is known at NLO but not at NNLO. In [32], the uncertainty generated by the use of the EFT approach in this case has been estimated to be of order 3–4% for $M_H = 125$ GeV (when the uncertainty due to the choice of a renormalisation scheme for the $b$-quark mass is also included). In addition, the mixed QCD-EW corrections at NNLO have been derived in the limit $M_W \gg M_H$, which is clearly not adequate. This leads to an additional error of a few percent [32], which generates a total EFT uncertainty of order of 5–7% for $M_H = 125$ GeV when all errors are summed.

Finally, the total uncertainty is obtained simply by adding linearly the scale and the PDF uncertainties and, in the case of $gg \rightarrow H$, also the EFT uncertainty on top of these. Nevertheless, for the dominant $gg \rightarrow H$ channel, as the results for the cross section at NLO are not yet publicly available at energies above $\sqrt{s} = 14$ TeV and the associated scale uncertainty is not given in [57] for such higher energies, we cannot provide an NLO description of $\sigma(gg \rightarrow H)$. We will therefore use the following approximation for the observed Higgs boson with a mass $M_H = 125$ GeV. At $\sqrt{s} = 14$ TeV, the scale uncertainty of $\sigma^{\text{NLO}}(gg \rightarrow H)$ is about $\pm 9\%$ and is thus of the same order as the sum of the scale uncertainty at NLO (which is about 3%) and the EFT uncertainty (which is about 5–7%). We will therefore consider that this trend is valid at all energies and assume that

$$\Delta\sigma^{\text{NLO}}_{\mu}(gg \rightarrow H) = \Delta\sigma^{\text{NLO}}_{\mu}(gg \rightarrow H) + \Delta\sigma^{\text{NLO}}_{\mu}(gg \rightarrow H).$$

(2.6)

The SM input parameters used in the calculations throughout this paper are as follows:

- $M_W = 80.385$ GeV, $M_Z = 91.1876$ GeV,
- $M_t = 173.2$ GeV,
- $M_H = 125$ GeV, $m_h^{\text{pole}} = 4.75$ GeV, $m_b(m_b) = 4.213$ GeV, $\alpha_s^{\text{LO}}(M_Z^2) = 0.13939$, $\alpha_s^{\text{NLO}}(M_Z^2) = 0.12018$, $\alpha_s^{\text{NNLO}}(M_Z^2) = 0.11707$.

According to the authors of [57] in which the tour de force of deriving the NLO corrections has been achieved, this generates an uncertainty of order 2% and the NNLO calculation with the exact $m_b$ dependence is now becoming the next big challenge in the field of higher order corrections.

The results for the cross section in all channels for the single production at a hadron machine of a Higgs boson with a mass $M_H = 100$ GeV are displayed in figure 1 as a function of the centre-of-mass energy, starting from $\sqrt{s} = 13$ TeV up to $\sqrt{s} = 100$ TeV. The total theoretical uncertainty bands are also displayed. The gluon fusion channel remains the dominant production mechanism up to the FCC-hh collider energies, with a cross section that ranges from $\sigma \approx 50$ pb at $\sqrt{s} = 13$–14 TeV to $\sigma \approx 800$ pb at $\sqrt{s} = 100$ TeV. It is followed by the VBF channel which has a cross section that is at least one order of magnitude smaller in the entire c.m. energy range. The Higgs-strahlung channels are the next important ones for energies below $\sqrt{s} \approx 30$ TeV with cross sections that are a factor 3–4 smaller than VBF. Above $\sqrt{s} \approx 30$ TeV, it is in fact the $pp \rightarrow t\bar{t}H$ process that becomes the third most important process. The $bb \rightarrow H$ process closes the list with a cross section that is slightly smaller than that of the $HZ$ process for $\sqrt{s} \gtrsim 14$ TeV. The numerical results are summarized in table 1 for all the considered channels; the values of the cross sections, together with the scale, PDF and total uncertainties are displayed.

In figure 2, we show the relative rise of the cross sections as a function of $\sqrt{s}$ when they are normalized to their values at $\sqrt{s} = 13$ TeV. One sees that compared to 13 TeV there is a gain of a factor $\approx 17$–18 at $\sqrt{s} = 100$ TeV in the case of gluon and vector boson fusion. Hence, with an integrated luminosity of a few ab$^{-1}$, the Higgs sample produced at a 100 TeV collider could allow for a measurement of some ratios of cross sections, such as $\sigma(pp \rightarrow H \rightarrow \gamma\gamma)/\sigma(pp \rightarrow H \rightarrow ZZ^*)$ that would be made free of theoretical uncertainties [142, 143], at the few per mille level. The gain in cross section when going from $\sqrt{s} = 13$ TeV to 100 TeV is smaller for the Higgs-strahlung...
process ($\approx 12$–$14$) but much larger ($\approx 72$) for $\bar{t}tH$ production as a result of the opening of the phase space. This clearly illustrates that the 100 TeV FCC-hh also has major potential to improve the understanding of the important bottom and top quark Yukawa couplings [144].

2.3. Single Higgs production in subdominant channels

We discuss briefly in this section the prospects for the associated production of one Higgs boson with a pair of massive weak bosons, namely $WW$, $WZ$ and $ZZ$ pairs as well as the associated production of a Higgs and a single top quark, $Wb \rightarrow Ht$ and considering the two channels $qb \rightarrow tHq'$ and $gq \rightarrow tHq'$.

2.3.1. Associated production with vector boson pairs. The channels $gq \rightarrow VH$ with $V = W, Z$ proceed through $s$-channel gauge boson and/or $t$-channel quark exchanges, in which the Higgs boson is radiated off the weak gauge bosons. The LO cross sections have been known for quite a while [145] and calculations with modern PDFs for the LHC were presented in [5]. A few years ago the NLO QCD corrections for $HWW$ and $HWZ$ channels were released [146, 147] showing that as in the case of $HV$ production, they increase the rates by approximately 50% (with a few percent contribution from the $gg$ fusion process). The cross sections are quite small but

Table 1. The total Higgs production cross sections (mostly at NNLO QCD + EW) in the main production processes for $M_H = 125$ GeV at a proton–proton collider (in pb) for given c.m. energies (in TeV) at the central scales given in the text.

| Channel   | $\sqrt{s}$ (TeV) | $\sigma$ (pb) | Scale (%) | PDF + $\alpha_s$ (%) | Total (%)   |
|-----------|------------------|---------------|-----------|-----------------------|-------------|
| $gg \rightarrow H$ | 13 | 46.68 | +8.6 | −9.3 | +7.5 | −8.0 | +16.1 | −17.3 |
|           | 14 | 52.43 | +9.4 | −9.2 | +7.5 | −7.5 | +16.9 | −16.6 |
|           | 33 | 189.5 | +8.6 | −7.7 | +7.5 | −7.3 | +16.1 | −15.0 |
|           | 100 | 788.6 | +7.1 | −6.1 | +8.3 | −8.0 | +15.4 | −14.1 |
| VBF       | 13 | 3.645 | +0.6 | −0.6 | +3.8 | −3.4 | +4.5 | −4.0 |
|           | 14 | 4.116 | +0.7 | −0.6 | +3.8 | −3.3 | +4.5 | −3.9 |
|           | 33 | 15.12 | +1.4 | −1.1 | +3.4 | −3.1 | +4.8 | −4.2 |
|           | 100 | 64.50 | +2.2 | −2.1 | +3.1 | −3.2 | +5.3 | −5.2 |
| $WH$      | 13 | 1.379 | +0.3 | −0.2 | +3.9 | −3.5 | +4.2 | −3.6 |
|           | 14 | 1.521 | +0.3 | −0.3 | +3.8 | −3.4 | +4.0 | −3.6 |
|           | 33 | 4.705 | +0.3 | −0.1 | +3.9 | −3.6 | +4.2 | −3.7 |
|           | 100 | 15.88 | +0.7 | −0.1 | +5.0 | −4.7 | +5.7 | −4.8 |
| $ZH$      | 13 | 0.8137 | +1.8 | −1.2 | +3.5 | −2.9 | +5.3 | −4.1 |
|           | 14 | 0.9037 | +1.8 | −1.3 | +3.3 | −2.9 | +5.1 | −4.3 |
|           | 33 | 2.969 | +2.0 | −1.6 | +3.7 | −3.6 | +5.7 | −5.2 |
|           | 100 | 11.28 | +1.8 | −1.7 | +4.5 | −4.3 | +6.3 | −6.0 |
| $t\bar{t}H$ | 13 | 0.514 | +6.7 | −9.7 | +6.9 | −7.1 | +13.6 | −16.8 |
|           | 14 | 0.623 | +6.8 | −9.7 | +7.0 | −6.9 | +13.8 | −16.6 |
|           | 33 | 4.51 | +8.5 | −9.1 | +6.4 | −5.8 | +14.9 | −14.9 |
|           | 100 | 37.0 | +9.5 | −10.2 | +6.3 | −5.4 | +15.8 | −15.6 |
| $b\bar{b} \rightarrow H$ | 13 | 0.529 | +12.6 | −35.3 | +4.9 | −6.0 | +17.6 | −41.3 |
|           | 14 | 0.598 | +12.7 | −35.8 | +5.1 | −5.9 | +17.7 | −41.6 |
|           | 33 | 2.20 | +13.7 | −41.2 | +4.3 | −5.7 | +18.0 | −46.9 |
|           | 100 | 8.91 | +15.9 | −47.4 | +5.2 | −5.9 | +21.1 | −53.4 |

Note: The corresponding shifts due to the theoretical uncertainties from the scale variation and the 90%CL MSTW PDF+$\alpha_s$ uncertainties are shown, as well as the total uncertainty when all errors are added linearly. Note that in the case of gluon fusion, the NNLO scale uncertainty is an estimation of the combined EFT+N3LO uncertainty.

Figure 2. The variation of the production cross sections for all channels with the c.m. energy relative to their values at $\sqrt{s} = 13$ TeV. The same graphical code as in figure 1 is used.

Rep. Prog. Phys. 79 (2016) 116201
In addition the associated production process, it is directly from the proton. The second one is the process in which a Higgs boson is produced in association with a weak boson pair at a proton–proton collider for given c.m. energies (in TeV) for $M_H = 125$ GeV. The MSTW2008 PDF set has been used.

### Table 2. The total cross section (in fb) at LO QCD for the Higgs production in association with a weak boson pair at a proton–proton collider for given c.m. energies (in TeV) for $M_H = 125$ GeV.

| $\sqrt{s}$ (TeV) | $\sigma_{pp \rightarrow HWW}^{LO}$ (fb) | $\sigma_{pp \rightarrow HWZ}^{LO}$ (fb) | $\sigma_{pp \rightarrow HZZ}^{LO}$ (fb) |
|------------------|-------------------------------------|-------------------------------------|-------------------------------------|
| 13               | 7.70                                | 3.28                                | 1.87                                |
| 14               | 8.72                                | 3.71                                | 2.12                                |
| 33               | 33.18                               | 13.96                               | 7.96                                |
| 100              | 164.1                               | 64.08                               | 36.98                               |

Note: $VW$ means $W^{+}W^{-}$, $W^{+}Z$ or $ZZ$. nevertheless they may provide additional tests and measurements at 100 TeV, for instance the $HWW$ coupling in the process $pp \rightarrow HWW \rightarrow 4W$. In addition the associate production with vector boson pairs is a background process for the production of a pair of Higgs bosons: For example the process $pp \rightarrow HWW \rightarrow b\bar{b}WW$ is a background of the search channel $pp \rightarrow HH \rightarrow b\bar{b}WW$.

We will not perform an error analysis and postpone it to a future publication [148] in which the NLO QCD corrections will be given for all three processes matched to parton shower in the POWHEG-BOX framework [149]. The numbers will then be given in this section at LO only and they have been obtained with a home-made computer program.

The results are displayed in figure 3. The production of a Higgs boson in association with a $W$ pair is the dominant, nearly three times the $HW^{+}Z$ production at 100 TeV. The detailed numbers are summarized in table 2. As a central scale we have used the invariant mass of the three-particle final state, $\mu_{Q} = M_{HWW}$.

### Table 3. The total cross section (in fb) at LO in QCD for the two Higgs subleading production channels $qb \rightarrow tHj$ and $qg \rightarrow tHb$ as well as for the subleading production channel $qg \rightarrow tHb$, at a few c.m. energies (in TeV) for $M_H = 125$ GeV at the central scale $\mu_{F} = \mu_{R} = \frac{1}{2}(M_H + M_j)$.

| $\sqrt{s}$ (TeV) | $\sigma_{qg \rightarrow tHb}^{LO}$ (fb) | $\sigma_{qb \rightarrow tHj}^{LO}$ (fb) | $\sigma_{qg \rightarrow tHb}^{LO}$ (fb) |
|------------------|-------------------------------------|-------------------------------------|-------------------------------------|
| 13               | 1.97                                | 60.54                               | 24.90                               |
| 14               | 2.17                                | 72.68                               | 30.90                               |
| 33               | 6.34                                | 475.8                               | 216.0                               |
| 100              | 19.61                               | 3453                                | 1676                                |

Note: The five-flavour scheme has been used and a sum over the final states $tH + X$ and $tH + \bar{X}$ is implicit.

2.3.2. **Associated production with a single top quark.** At the partonic level, the process in which a Higgs boson is produced in association with a single top quark is $Wb \rightarrow Ht$ and it proceeds through the s-channel exchange of a $t$-quark or the $t$-channel exchange of a $t$-quark or a $W$ boson. As in the case of the $tH$ associated production process, it is directly proportional to the top quark Yukawa coupling but it is in principle more favored by phase-space at low energies. It has however a much smaller cross section as it is formally of higher order in perturbation theory, since the $W$ boson should come for a splitting parton. Nevertheless, it has been shown in [150] that the process is extremely sensitive to the Yukawa coupling and, in particular, the $Wb \rightarrow Ht$ cross section can be enhanced by more than one order of magnitude compared to the SM case if the sign of the Yukawa coupling is reversed.

At the hadronic level, two processes can generate the $tH$ final state in a $t$-channel exchange. The first one is the $2 \rightarrow 3$ process $qb \rightarrow tHj$ with the $b$-quark treated as a parton taken directly from the proton. The second one is the $2 \rightarrow 4$ process $qg \rightarrow tHb$ in which the $b$ quark originates from gluon splitting. In principle both processes are equivalent (one is the NLO radiative correction of the other) and, as in the $bb \rightarrow H$ versus $gg \rightarrow bbH$ processes discussed in the previous section, one has to have a procedure to match the two processes
calculated with four or five active parton flavours in a four- or five-flavour scheme. However, as the process is subleading, we will not enter into such sophisticated details and calculate the two processes independently in a five-flavour scheme, following closely the analysis performed in [150]. As also stated in an earlier work [151] as well as in a very recent study [152] done at NLO in QCD, there is also a s-channel process $qq' \rightarrow tHb$. The main advantage of the five-flavour scheme is that the three production processes presented here are totally independent at LO. The separation between s-channel and t-channel productions still remains at NLO in the five-flavour scheme [152].

The cross sections for the three processes are calculated using the program MadGraph 5 with the factorisation and renormalisation scales set to $\mu_0 = \frac{1}{2}(M_H + M_b)$, a scale that minimizes the higher-order corrections and reduces the scale uncertainty [152]; the MSTW2008 PDFs have been adopted. While the inclusive rate is calculated in the $qb \rightarrow tHj$ and $qq' \rightarrow tHb$ cases, we use the following cuts for the detection of the $b$-quark (which allows an additional means to suppress the QCD background) in the case of the $qg \rightarrow tHj b$; $p_T^{b} > 25$ GeV and $|p_T^{b}| < 2.5$. This has allowed us to make a cross check by comparing our results with those of [150] given at $\sqrt{s} = 8$ and 14 TeV using their setup: a very good agreement has been found. Our results for the two other processes have been compared to [152], again using their set-up, and a perfect agreement has been found. The results for the cross sections in the three processes are shown in figure 4 as a function of $\sqrt{s}$ and some numerical values are displayed in table 3 for some specific c.m. energies. At $\sqrt{s} = 14$ TeV, one has $\sigma(qg \rightarrow tHj b) \approx 73$ fb, i.e. it is almost two orders of magnitude smaller than $\sigma(pp \rightarrow tHj)$, while $\sigma(qg \rightarrow tHj b)$ is approximately three times smaller. The s-channel process $qq' \rightarrow tHb$ is totally subdominant, of the order of 2 fb. At $\sqrt{s} = 100$ TeV, the cross section for $qb \rightarrow tHj$ has increased by a factor or $\approx 57$.

2.4. SM multi-Higgs production

We will briefly review in this section the results for the pair production of the SM Higgs boson, which allows for the determination of the important trilinear couplings among the Higgs states. These were already obtained up to 100 TeV in a previous publication [153] (which itself was an update of the older analyses of [154, 155] made for the 14 TeV LHC and for an electron–positron collider) and we update in particular the gluon fusion channel in which new results have appeared quite recently and the $tHH$ channels. We will present both the cross sections at the various c.m. energies and the theoretical uncertainties, the calculation of which follows the main lines already presented in the previous section for single Higgs production in the dominant channels. The case of triple Higgs production has already been studied in the literature with hadronic energies up to 200 TeV and it has been found that the cross section is so small for such c.m. energies that it will be extremely difficult to observe it [156–158]. Nevertheless, a very recent study [39], which relies on very optimistic assumptions about the experimental setup, considers this process to be interesting at 100 TeV despite the aforementioned limitations. This calls for further investigations in the context of the FCC-hh (see also [40, 41]) and, with theoretical progress as well as experimental improvements, one might hope for a potential observation of this process at 100 TeV. Therefore, in this section, although we will not include a complete and detailed analysis, we will give the rates for the largely dominant triple production process $gg \rightarrow HHH$ for energies up to 100 TeV.

2.4.1. Status of the higher order corrections in Higgs pair production. The gluon fusion channel, $gg \rightarrow HH$, is the dominant Higgs pair production process, in much the same way as it is the dominant process for single Higgs production. It is also mediated by loops of heavy quarks which are of two types in this case: triangular in the case of the diagram in which a virtual $H$ boson is produced and splits to real Higgs bosons (and which involves the trilinear Higgs coupling that needs to be determined) and box-type in the case of the diagrams in which both Higgs bosons are emitted from the internal quark lines (and hence, do not involve the triple Higgs couplings and can be considered as an irreducible background in the measurement of the latter). Contrary to single Higgs production, bottom quark loops have a negligible contribution to the total cross section in this case: less than 1% at LO [159–162].

The process was known for a long time at NLO QCD in the infinite top mass approximation [163], and progress was made only quite recently—progress that pushed this process to NNLO QCD in again the infinite top mass approximation [164, 165]. The study of the NLO QCD structure has also made progress and top mass expansion was performed in 2013 [166] while the exact real corrections were computed in the aMC@NLO framework [158, 167]. A full NLO calculation including the exact quark mass dependence is still missing though. Soft-gluon resummation was also performed up

![Figure 5. The total cross sections for Higgs pair production at a proton–proton collider, including higher-order corrections discussed in the text, in the main production channels as a function of the c.m. energy with $M_H = 125$ GeV. The MSTW2008 PDF set has been used and theoretical uncertainties are included as corresponding bands around the central values.](image-url)
The production of a Higgs pair in association with a weak boson, \( q\bar{q} \to HHV \), was calculated for the first time a while ago [173] and shares common aspects with the single Higgs-strahlung process. In particular, the NLO and NNLO QCD corrections can be implemented in complete analogy to the single Higgs process, i.e., by adapting the corrections to the Drell–Yan mechanism [98, 103] to this case. In addition, one needs to consider again a new subchannel at NNLO for \( ZHH \) production, namely \( gg \to ZHH \), was calculated for the first time a while ago [168] and then matched with the NNLO result [169], where it has been found that the resummation increases the cross section over the NNLO result by \( \sim+7\%\) at 14 TeV while reducing the scale uncertainty down to \( \sim\pm5\%\). The VBF process, \( qq \to VVq \), has also been known for quite a while only at LO [159, 161, 170, 171]. It is very similar to the single Higgs production case, and the NLO QCD corrections follow essentially the same trend and can be derived simply by turning the tensor structure of the fusion of the weak bosons from \( VV \to H \) into \( VV \to HH \) while using the exact same QCD corrections to the quark lines as for single Higgs production. These NLO QCD corrections were calculated quite recently [153, 167] and are now available in two different codes not only for the total rates but also for the differential distributions [97, 167]. The NLO QCD corrections amount to a \( \sim+7\%\) correction. The approximate NNLO QCD corrections have also been computed in the structure function approach [172] and similarly to the case of single Higgs production the increase of the rate is very modest, half a percent at most, while it decreases the scale uncertainty down to the percent level for the inclusive cross section.

The production of a Higgs pair in association with a weak boson, \( q\bar{q} \to HHV \), was calculated for the first time a while ago [173] and shares common aspects with the single Higgs-strahlung process. In particular, the NLO and NNLO QCD corrections can be implemented in complete analogy to the single Higgs process, i.e., by adapting the corrections to the Drell–Yan mechanism [98, 103] to this case. In addition, one needs to consider again a new subchannel at NNLO for \( ZHH \) production, namely \( gg \to ZHH \), which proceeds through triangle, box and pentagon loops of heavy quarks. In sharp contrast to single Higgs production, this subchannel gives a significantly more important contribution, which can amount to 30% of the NNLO total rate [153]. The NLO QCD corrections were calculated in [153, 167] and are available also for the differential distributions, while the NNLO QCD corrections are only available for the total rates [153].

The last double Higgs production channel that we discuss is the associated production with a top-quark pair. This process has been known only at LO and it is only recently that the NLO QCD corrections have been made available [167]. This is due to the complexity of the calculation, which needs modern tools as one has to calculate QCD corrections to a \( 2 \to 4 \) process involving massive quarks in the final states. Depending on the chosen central scale, they can reduce or enhance the total rate. In our case we have chosen the central scale \( \mu_k = \mu_F = M_{HH} \) and we see a very modest increase of the rates over the LO results at 14 TeV, while the increase is larger at 100 TeV, of the order of +10%.

### Table 4. The total Higgs pair production cross section (in fb) for given c.m. energies (in TeV) for \( M_H = 125 \text{ GeV} \).

| Channel | \( \sqrt{s} \) (TeV) | \( \sigma \) (fb) | scale (%) | PDF (%) | PDF+\( \alpha_s \) (%) | total (%) |
|---------|-------------------|----------------|-----------|---------|-----------------|----------|
| \( gg \to HH \) | 13 | 34.56 | +8.5 | -8.5 | +4.0 | -4.0 | +7.1 | -6.2 | +25.6 | -24.7 |
| 14 | 41.11 | +8.3 | -8.3 | +3.9 | -4.0 | +7.0 | -6.2 | +25.3 | -24.5 |
| 33 | 247.93 | +7.0 | -7.0 | +2.5 | -2.7 | +6.2 | -5.4 | +23.2 | -22.4 |
| 100 | 1670.83 | +5.8 | -5.8 | +2.0 | -2.7 | +6.2 | -5.7 | +22.0 | -21.5 |
| \( VBF \to HH \) | 13 | 1.73 | +1.7 | -1.1 | +4.6 | -4.2 | +5.9 | -4.2 | +7.6 | -5.2 |
| 14 | 2.01 | +1.7 | -1.1 | +4.6 | -4.1 | +5.9 | -4.1 | +7.6 | -5.1 |
| 33 | 12.05 | +0.9 | -0.5 | +4.0 | -3.7 | +5.2 | -3.7 | +6.1 | -4.2 |
| 100 | 79.55 | +1.0 | -0.9 | +3.5 | -3.2 | +5.2 | -3.2 | +6.2 | -4.1 |
| \( WHH \) | 13 | 0.30 | +0.1 | -0.3 | +3.6 | -2.9 | +3.6 | -3.0 | +3.7 | -3.3 |
| 14 | 0.57 | +0.1 | -0.3 | +3.6 | -2.9 | +3.6 | -3.0 | +3.7 | -3.3 |
| 33 | 1.99 | +0.1 | -0.1 | +2.9 | -2.5 | +3.4 | -3.0 | +3.5 | -3.1 |
| 100 | 8.00 | +0.3 | -0.3 | +2.7 | -2.7 | +3.8 | -3.4 | +4.2 | -3.7 |
| \( ZHH \) | 13 | 0.36 | +4.0 | -2.9 | +2.8 | -2.3 | +3.0 | -2.6 | +7.0 | -5.5 |
| 14 | 0.42 | +4.0 | -2.9 | +2.8 | -2.3 | +3.0 | -2.6 | +7.0 | -5.5 |
| 33 | 1.68 | +5.1 | -4.1 | +1.9 | -1.5 | +2.7 | -2.6 | +7.9 | -6.7 |
| 100 | 8.27 | +5.2 | -4.7 | +1.9 | -2.1 | +3.2 | -3.2 | +8.4 | -8.0 |
| \( \tilde{t}HH \) | 13 | 0.77 | +2.4 | -7.3 | +3.8 | -4.3 | +6.7 | -6.6 | +9.2 | -13.9 |
| 14 | 0.95 | +3.1 | -7.5 | +3.6 | -4.3 | +6.5 | -6.6 | +9.6 | -14.1 |
| 33 | 8.13 | +6.9 | -8.2 | +2.5 | -3.0 | +6.2 | -6.0 | +13.1 | -14.2 |
| 100 | 79.86 | +7.4 | -7.5 | +1.6 | -2.0 | +5.1 | -4.2 | +12.5 | -11.6 |

Note: The central scale is \( \mu_k = \mu_F = M_{HH} \).

### Table 5. The total triple Higgs production cross section \( gg \to HHH \) (in fb) for given c.m. energies (in TeV) for \( M_H = 125 \text{ GeV} \).

| \( \sqrt{s} \) (TeV) | \( \alpha_{LO} \to HHH \) (fb) |
|-------------------|----------------|
| 13 TeV | 0.033 |
| 14 TeV | 0.040 |
| 33 TeV | 0.33 |
| 100 TeV | 3 |

Note: The central scale is \( \mu_k = \mu_F = M_{HH} \).
2.4.2. Numerical results for pair production. We present our numerical results, an update of [153], using the following central scales:

\[
\mu_{\text{gg fusion}} = \frac{M_{\text{Hgg}}}{Q_V}, \quad \mu_{\text{VBF}} = \frac{M_{\text{VBF}}}{Q_V},
\]

where \(Q_V\) is the c.m. energy given in TeV. We also account for the scale uncertainty using the formula (21) in [165]. The uncertainty related to the use of the effective approach for the top loop is still kept this time contrary to the single Higgs case as it has significant impact. For the Higgs-strahlung and VBF processes, we simply collect our previous results obtained in [153] and replace the 8 TeV LHC predictions by the case of the LHC at \(\sqrt{s} = 14\) TeV. The VBF results are obtained using VBFNLO [97] while the Higgs-strahlung predictions are obtained with a computer program of our own that was developed especially for this purpose. The results for \(\bar{t}tH\) have been obtained with MadGraph5/mc@NLO [37]. The reduction of the scale uncertainty is spectacular, from \(\sim 30\%\) down to \(\sim 2\%\) at 14 TeV.

The results are displayed in figure 5 as a function of the c.m. energy starting from LHC 13 TeV up to the FCC-hh energy of 100 TeV. As stated above this clearly shows that the gluon fusion channel is the dominant production on the full c.m. energy range, as in single Higgs production. This is followed by the VBF channel up to 100 TeV, but in this

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\(\mu_{\text{gg fusion}} = M_{\text{Hgg}}, \quad \mu_{\text{VBF}} = M_{\text{VBF}}\)
region it is actually the associated production with a top quark pair that dominates over VBF. Higgs-strahlung channels are one order of magnitude smaller than VBF over the whole c.m. energy range. The numerical results are summarized in table 4 for all channels. Hence, clearly, the gluon fusion channel will be the most relevant channel at 100 TeV in much the same way as at 14 TeV. Nevertheless the VBF contribution can be important in the analyses including two jets in the final state and is already included in various 14 TeV analyses [175, 176], and given the higher number of events at 100 TeV, this VBF contribution could be even more interesting.

2.4.3. Expectations for triple Higgs production. As stated already in the literature [156–158] the cross section for triple Higgs production at the LHC is very small (note that the results in [158] include approximate NLO QCD corrections). We display in table 5 the production rates in the dominant gg → HHH channel at a central scale \( \mu \) for the parameter set of equation (2.7). We have used our own implementation and checked our code against the numbers given in [157] and found full agreement. The rates are indeed negligible at the LHC, being three orders of magnitude smaller than pair production. Nevertheless, at 100 TeV, the cross section reaches the level of 3 fb (maybe a factor 2 higher if the \( K \)-factors are the same as in single and pair production) and could thus lead to a few thousand events with a luminosity of 1 ab\(^{-1}\). The extraction of the signal, in particular in the 6\( b \) [39], 4\( b \gamma \) [40] and 4\( b \tau \) [41] detection channels, requires very large integrated luminosities and is nevertheless extremely challenging in view of the formidable backgrounds.

3. Dark matter and the Higgs-portal

3.1. Higgs-portal dark matter models

A very interesting scenario would be that the particles that form partly or entirely the dark matter in the universe interact only with the Higgs sector [21–23] (see also [24] for a review and more references). The DM particles can be made stable by a \( Z_2 \) symmetry and annihilate to SM states through the exchange of Higgs bosons. These Higgs portal scenarios for DM can be of several kinds, depending on whether the models contain additional Higgs and/or matter particles or not, but the simplest one would be the scenario in which the SM is extended to contain only the DM particle while its minimal Higgs sector with one scalar doublet is kept unchanged. The DM particles will then interact only with the observed \( h \) state and their annihilation into SM fermions and bosons, for instance, can occur only through the \( h \) exchange in the s-channel.

To describe the DM properties in this minimal Higgs-portal scenario, it is convenient to work in a quite model-independent framework (although the origin of the \( Z_2 \) parity that ensures the stability of the DM particle is still model-dependent) in which the particle consists of a real scalar \( S \), a vector \( V \) or a Majorana fermion \( f \) that interacts with the SM fields only through the \( h \) exchange in the s-channel.
addition to those of the SM. These are, besides the three spin assignment possibilities, the mass of the DM state and its coupling to the $h$ boson. The relevant terms in the Lagrangians describing the spin-0, spin-1/2 and spin-1 cases have a general form that can be simply written as

$$
\Delta L_{S} = -\frac{1}{2} m_{S}^{2} S^{2} - \frac{1}{4} \lambda_{S} S^{4} - \frac{1}{4} \lambda_{\text{SS}} H^{+} H S^{2},
$$

$$
\Delta L_{V} = \frac{1}{2} m_{V} V^{*} V + \frac{1}{4} \lambda_{V} (V_{\mu} V^{\mu})^{2} + \frac{1}{4} \lambda_{VV} (H^{+} H V^{*} V^{\mu},
$$

$$
\Delta L_{f} = -\frac{1}{2} m_{f}^{2} f^{2} - \frac{1}{4} \lambda_{\text{ff}} H^{+} H f^{2}. \tag{3.1}
$$

where the self-interaction terms $S^{2}$ in the scalar and the $(V_{\mu} V^{\mu})^{2}$ term in the vector cases are not essential for our discussion and can be ignored. In the fermionic case, the form that we adopt here for the Higgs-DM coupling is not renormalisable, but as it is a rather convenient parametrisation, we keep it in our discussion. After electroweak symmetry breaking when the neutral component of the doublet field $H$ is shifted to $(\nu + h)^{\sqrt{2}}$ with $\nu = 246$ GeV, the physical masses of the DM particles will be given by

$$
M_{S}^{2} = m_{S}^{2} + \frac{1}{4} \lambda_{\text{SS}} v^{2},
$$

$$
M_{V}^{2} = m_{V}^{2} + \frac{1}{4} \lambda_{VV} v^{2},
$$

$$
M_{f}^{2} = m_{f}^{2} + \frac{1}{4} \lambda_{\text{ff}} v^{2}. \tag{3.2}
$$

Since the cosmological relic density of the DM particles is obtained by means of the annihilation to SM particles through the exchange of the $h$ boson, there is in principle a relation between the coupling $\lambda_{\text{SS}}$, and the mass $M_{S}$ of the DM particle if the Planck satellite constraint, $\Omega_{DM} h^{2} = 0.1186 \pm 0.020$ [184] with $h$ being the reduced Hubble constant, is imposed. However, this is only true if the $\chi$ particle is absolutely stable and has to account for all the DM in the universe. For a more general discussion in the context of collider physics that we are discussing here, we will ignore this constraint. Nevertheless, there are also constraints from the rates for the direct and indirect detection of the $\chi$ particles in astrophysical experiments. In particular, the elastic DM interaction with nuclei occurs through the $t,s$-channel exchange of the $h$ boson and the resulting nuclear recoil or the spin-independent DM-nucleon cross section can be interpreted in terms of the DM mass and coupling; see for instance [178] in which the present constraints from the two most sensitive experiments XENON100 [185] and LUX [186] have been summarized.

The issue with which we will be concerned here is how to observe directly the Higgs-portal DM particles at high energy proton colliders. There are essentially two ways, depending on the $h$ versus $\chi$ particle masses. If these particles are light enough, $M_{S} \lesssim \frac{1}{2} M_{h}$, the invisible $h \rightarrow \chi \chi$ decay can occur and its partial width will contribute to the total decay width of the observed Higgs boson. It will then alter the branching ratios for the visible decays and hence the signal rates in the various channels in which the $h$ boson has been detected at the LHC. In the previous LHC runs with $\sqrt{s} = 7 + 8$ TeV, the Higgs cross sections times branching ratios in some channels like $pp \rightarrow h \rightarrow \gamma\gamma$ and $h \rightarrow ZZ \rightarrow 4\ell^{\pm}$ have been measured at the $\approx 20\%$ level and, as they agree with the SM expectations, they set an upper bound $\text{BR}(h \rightarrow \text{inv.}) \lesssim 0.3$ for SM-like $h$ couplings; see e.g. [187]. The sensitivity will certainly improve in the coming LHC run at $\sqrt{s} = 13$ and 14 TeV, but the observation of the invisible Higgs branching ratio would be extremely difficult if $\text{BR}(h \rightarrow \text{inv.}) \lesssim 0.1$, in view of the large QCD uncertainties that affect the Higgs production cross sections, in particular in the main $gg \rightarrow h$ production channel, as discussed earlier.

On the other hand, DM particles could be directly detected by studying the vector boson fusion and the Higgs-strahlung processes in which the Higgs boson decays invisibly [191, 192]. The ATLAS and CMS searches at the $7 + 8$ TeV run in these missing energy channels give the constraint $\text{BR}(h \rightarrow \text{inv.}) \lesssim 30\%$ [193, 194] if the $h$ production and visible decay rates are SM-like. From parton level analyses, one does not expect that this limit on the invisible branching ratio will be significantly improved at the 14 TeV LHC upgrade even with a sufficiently large amount of data [195].

In turn, if the mass of the DM particle is larger than half the $h$ mass, $M_{S} \gtrsim \frac{1}{2} M_{h}$, there is no invisible Higgs decay and the detection of the $\chi$ particles in collider experiments becomes much more difficult. In fact, the only possible way to observe the $\chi$ states would be through their pair production in the continuum via the exchange of the Higgs boson [179, 196, 197]. The latter needs to be produced in association with visible particles and at hadron colliders, three main processes are at hand: (i) double production in the Higgs-strahlung process $gg \rightarrow V^{\ast} \rightarrow V\chi\chi$ with $V$ being either a $W$ or a $Z$ boson, (ii) the vector boson fusion processes which lead to two jets and missing energy $qq \rightarrow V^{\ast}V^{\ast}qq \rightarrow \chi\chi qq$, in the final state and (iii) the gluon fusion mechanism, which is mainly mediated by loops of the heavy top quark that couples strongly to the Higgs boson, $gg \rightarrow h^{\ast} \rightarrow \chi\chi$, but in which additional jets are emitted in the final state in order to make the process visible.

11 One could have a renormalisable Higgs coupling to spin-1/2 Majorana DM particles similar to the one above and a good example would be the case of the MSSM where the $h\chi\chi$ coupling (as well as the physical mass $M_{\chi}$) can be defined in terms of the elements of the matrix that diagonalizes the $4 \times 4$ mass matrix of the four neutralino states (the bino, wino, and the two higgsinos) with the lightest one being identified with the DM particle. In the case where all the additional Higgs bosons except for $h$ and all the SUSY particles except for the stable lightest neutralino are very heavy, one would be in a situation [180–183] that is quite similar to the $h$-portal fermionic DM scenario that we are discussing here.

12 The indirect observation of invisible Higgs decays would be much easier at a future $e^{+}e^{-}$ collider and it has been shown that at $\sqrt{s} \approx 500$ GeV with a sample $100$ fb$^{-1}$ data, the Higgs cross section times the visible branching fractions can be determined with a percent level accuracy [188–190].

13 Here again, the situation is more favorable at an $e^{+}e^{-}$ collider with a c.m. energy $\sqrt{s} \gtrsim 250$ GeV, as invisible decays at the level of a few percent can be observed in the process $e^{+}e^{-} \rightarrow hZ$ by simply analyzing the recoils of the leptonically decaying $Z$ boson [188–190].
3.2. The cross sections for DM production through the Higgs portal

We now present numerical results for the DM pair production cross sections through Higgs splitting in the three possible processes discussed above. Results are shown for a proton collider at $\sqrt{s} = 100$ TeV and compared with those obtained at $\sqrt{s} = 14$ TeV.

For DM double production in the Higgs-strahlung process with either a W or a Z boson, we have used the package Feynrules [198] in which we implemented the Lagrangian equation (3.1) describing the DM interactions in the three spin cases and exported the results into the MadGraph5/aMC@NLO framework [37]. The cross sections include thus the NLO QCD corrections and we adopted the MSTW set for the PDFs and the MadGraph5 default cards (including kinematical cuts). The DM pair production rates are shown in figure 6. The left column corresponds to energy $\sqrt{s} = 14$ and that on the right to 100 TeV. The cross sections are plotted as a function of the mass of the generic DM particle $\chi$ and we have set the DM couplings to the $h$ portal to $\lambda_{h\chi\chi} = 1$ (in the fermionic case, we also set $\Lambda = 1$ TeV). For other couplings one simply has to multiply the rates by $\lambda_{h\chi\chi}^2$.

One see that for the three types of particles the cross sections are extremely small at the LHC even for $M_\chi = 100$ GeV. For such a $\chi$ mass, they do not exceed the fb level in the spin-1 case and they are one and two orders of magnitude smaller in, respectively, the spin-1/2 and spin-0 cases (with the rate for $W\chi\chi$ being twice as large as the one for $Z\chi\chi$ as is usually the case for such processes). At $\sqrt{s} = 100$ TeV, there is an increase of the cross section by almost two orders of magnitude at low masses. The rates remain thus too low for the processes to be very interesting, in particular at $M_\chi$ significantly above 100 GeV and for the spin-0 and $1/2$ cases.

For the vector boson fusion case in which the pair of escaping DM particles is produced in association with two jets, we also used the same strategy as above for the calculation of the cross section and we use Feynrules [198] to export our simple $h$-portal extension of the SM into the MadGraph5 framework [37]. The cross sections are calculated at LO only and we again adopt the MSTW PDFs and the MadGraph5 default cards (but removing the cuts on the jet transverse momentum). The production rates at $\sqrt{s} = 14$ and 100 TeV are presented in figure 7 in the same configuration as Higgs-strahlung that is, as functions of the mass $M_\chi$, setting the coupling to $\lambda_{h\chi\chi} = 1$ and considering the three spin configurations.

The cross sections are one order of magnitude larger than in the Higgs-strahlung case for the three spin-configurations and the hierarchy is the same: one order of magnitude larger for spin-1 than for spin-0, and than for spin $1/2$ if $\Lambda = 1$ TeV. Again, the rates are 100 times larger at $\sqrt{s} = 100$ TeV than at 14 TeV and, for instance, in the vector case they almost reach the picobarn level for $M_\chi = 100$ GeV and hence, provide a chance to observe the process. The cross sections fall steeply with the $\chi$ mass and even in the optimistic spin-1 case, they drop to below the femtobarn level for $M_\chi = 300$ GeV.

Finally, concerning DM double production in association with one jet either from gluon fusion $gg \rightarrow \chi\chi g$ or $qg$ annihilation $gg \rightarrow q\bar{q}\chi\chi$, we have used a modified version of the program HIGLU [199] that describes Higgs production and where the couplings of equation (3.1) are implemented. The cross sections include the NLO QCD corrections (which are equivalent to the NNLO ones in the SM Higgs case and are thus rather large) and we adopt the MSTW set for the PDFs. The renormalisation and factorisation scales have been set to $M_\chi$. The production rates are shown in figure 8 again at $\sqrt{s} = 14$ TeV (left) and 100 TeV (right) in the same configurations as in the two previous processes.

Here, the cross sections are an additional order of magnitude larger than in the vector boson fusion case (except at very low $M_\chi$ for the spin-0 case) and approximately follow the same trend: about a factor of 50 larger at 100 TeV than at 14 TeV and larger by a factor of 10 and 100 in the vector case than in, respectively, the scalar and the fermionic cases. At $\sqrt{s} = 100$ TeV, one has large production rates at reasonably low $\chi$ masses. This is particularly the case for a vector DM particle where one obtains a few picobarns in the low mass range, $M_\chi = 100$ GeV. Here again, the cross sections fall steeply with the $\chi$ mass.

4. The heavier Higgs bosons of 2HDMs and the MSSM

4.1. The physical set-up

A very good benchmark for studying extended Higgs sectors are two Higgs doublet models or 2HDMs for short [11, 12]. Compared to the SM with its unique Higgs particle, the Higgs sector of 2HDMs involves five physical states after electroweak symmetry breaking and has a phenomenology that is much richer. Indeed, the model has two CP-even neutral states $h$ and $H$ that mix and share the properties of the SM Higgs boson, a CP-odd or pseudoscalar $A$ boson with properties that are completely different from that of the SM Higgs and, above all, it has two charged Higgs states $H^\pm$ that provide a unique signature for new physics. While the pattern of the couplings of the two doublet Higgs fields to the gauge sector is somewhat fixed and relatively simple, there are several possibilities for the structure of the couplings to standard fermions, leading to several types of 2HDMs [11, 12].

The Minimal Supersymmetric standard model or MSSM is a specific type of 2HDM, the so-called type II in which one Higgs field doublet gives mass to up-type fermions and the other gives mass to down-type fermions [11]. However, there are strong constraints on the Higgs sector that reduce the number of free parameters to only two inputs at tree-level. Nevertheless, if the SUSY spectrum is light, a number of complications and new features are introduced as the superparticles can substantially affect the phenomenology of the Higgs bosons, either indirectly through loop corrections or new processes or directly by allowing, for instance, new production and decay mechanisms. However, as the LHC data indicate that this scale is rather high, $M_3 \gtrsim 1$ TeV, it is very likely that the SUSY particles will not affect the MSSM Higgs sector.
making the phenomenology in this model quite similar to that of 2HDMs.

In both the general 2HDM and the MSSM, the most important production mechanisms of the neutral CP-even Higgs bosons are simply those of the SM Higgs particle that have been discussed in detail in the previous sections of this paper. However, major quantitative differences compared to the SM case can occur since the cross sections will depend on the specific Higgs mass and coupling patterns, which can be widely different; this is particularly the case for the pseudoscalar Higgs boson. In the case of the charged Higgs particles, new production mechanisms not discussed before occur. Another major difference between the SM and these models is the Higgs decay pattern, which can be much more involved.

These are the issues that we will discuss in this section, in which the potential of a 100 TeV proton collider to probe this extended Higgs sector is analysed in detail. But before that, let us summarize the physical spectrum in these two models.

### 4.1.1. The case of 2HDMs

In our analysis, we will consider the CP-conserving two-Higgs-doublet-model with a softly broken $Z_2$ symmetry and here, we briefly highlight its main features; for more details we refer the reader to [200, 201], whose approach and notations we adopt. The scalar potential of this model, in terms of the two Higgs doublet fields $\Phi_1$ and $\Phi_2$, is described by three mass parameters and five quartic couplings and is given by

$$V = m_{11}^2 |\Phi_1|^2 + m_{22}^2 |\Phi_2|^2 - m_{12}^2 (\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1) + \frac{1}{2} \lambda_1 (\Phi_1^\dagger \Phi_1)^2 + \frac{1}{2} \lambda_2 (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 (\Phi_1^\dagger \Phi_1)(\Phi_2^\dagger \Phi_2) + \frac{1}{2} \lambda_4 (\Phi_2^\dagger \Phi_2)^2 + (\Phi_1^\dagger \Phi_1)^2].$$

The masses of the two CP-even neutral states $h$ and $H$, as well as that of the CP-odd neutral $A$ and two charged $H^\pm$ states that are present in the model, $M_h$, $M_H$, $M_A$, and $M_{H^\pm}$, are free parameters and we will assume here that the lighter CP-even $h$ boson is the observed Higgs resonance with a mass of $M_h = 125$ GeV. Three other parameters describe the model and among them there are two mixing angles $\beta$ and $\alpha$: $\tan \beta = v_2/v_1$ where $v_1/\sqrt{2}$ and $v_2/\sqrt{2}$ are the vacuum expectation values of the neutral components of the fields $\Phi_1$ and $\Phi_2$ (with $v_1^2 + v_2^2 = v^2 = (246$ GeV$)^2$) and the angle $\alpha$ that diagonalizes the mass matrix of the two CP-even and $H$ bosons. Finally, there is another mass parameter $m_{12}$ which enters only in the quartic couplings among the Higgs bosons,

$$\lambda_{\phi_1 \phi_2} = g_{HVV}^{\text{SM}} / g_{HH}^{\text{SM}} = f(\alpha, \beta, m_{12}).$$

In this parametrisation, the neutral CP-even $h$ and $H$ bosons share the coupling of the SM Higgs particle to the massive gauge bosons $V = W,Z$ and one has at tree level,

$$g_{hVV} = g_{hVV}^{\text{SM}}, g_{hVV}^{\text{2HDM}}, g_{HHV} = \sin(\beta - \alpha),$$

$$g_{HHVV} = g_{HHVV}^{\text{SM}}, g_{HHVV}^{\text{2HDM}} = \cos(\beta - \alpha).$$

while, as a consequence of CP invariance, there is no coupling of the CP-odd $A$ to vector bosons, $g_{AVV} = 0$. There are also couplings between two Higgs and a vector boson which, up to a normalization factor$^{14}$, are complementary to the ones above. For instance one has, using a normalisation that gives similar couplings as in the previous case,

$$g_{HAA} = g_{hAA} = \cos(\beta - \alpha), g_{HAA} = g_{HAA}^{\text{SM}} = \sin(\beta - \alpha).$$

For completeness, additional couplings of the charged Higgs boson which will be needed in our discussion. They do not depend on any SUSY parameter and are simply given by (here, we include the pre-factors in the couplings except for the momenta)

$$g_{HAA} = \frac{\alpha}{2} \sin \theta_W, g_{HAA} = -e,$$

$$g_{HAA} = -e \cos 2\theta_W/\sin 2\theta_W.$$ 

In a general 2HDM, the interaction of the Higgs states with fermions are model-dependent and there are generally two options which are discussed in the literature [12]. In Type II models, the field $\Phi_1$ generates the masses of isospin down-type fermions and $\Phi_2$ the masses of up-type quarks. In turn, in Type I models, the field $\Phi_2$ couples to both isospin up- and down-type fermions. The couplings of the neutral Higgs bosons to gauge bosons and fermions are given in table 6 in the two models. The couplings of the charged Higgs boson to fermions follow that of the CP-odd Higgs boson.

The Higgs couplings to fermions and gauge bosons depend only on the ratio $\tan \beta$ and the difference $\beta - \alpha$ between the two mixing angles. If one enforces the fact that the couplings of the observed $h$ boson should be SM-like, as the LHC Higgs data strongly indicate [8, 25], one simply needs to set $\beta - \alpha = \pi/2$. This is called the alignment limit [15–18, 202]. In this limit, the $h$ couplings are SM-like $g_{HVV} = g_{HAA} = g_{Hdd} \to 1$, while the couplings of the CP-even $H$ state reduce to those of the pseudoscalar $A$ boson. In particular, there is no $H$ coupling to vector bosons, $g_{HVV} \to g_{HVV}^{\text{SM}} = 0$ and the couplings to up-type fermions are $g_{Huu} = \cot \beta$ while those to down-type fermions are $g_{Hdd} = \cot \beta$ and $g_{Hdd} = \tan \beta$, respectively, type I and II models.

In addition to $\tan \beta$, at least the Higgs masses $M_H, M_A$ and $M_{H^\pm}$ will enter the 2HDM phenomenology. These are in principle free parameters and can have arbitrary values, except for the $H$ state that was assumed to be heavier than $h$. This makes any phenomenological analysis rather complicated and to simplify our discussion here, we will make the additional assumption that they are of the same order$^{15}$,

$$M_H \approx M_A \approx M_{H^\pm}.$$ 

$^{14}$ The complete set of Feynman rules can be found in appendix A of [11]

$^{15}$ More precisely, the discussion that we will have in the next sections will hold if the differences of the masses of the $\phi_1, \phi_2 = H, A, H^\pm$ states, $|M_h - M_{h_0}| \lesssim M_W$, is satisfied. This pattern will be more or less equivalent to the one in the MSSM close to the decoupling limit as will be seen later.
Finally, in the alignment limit \( \beta - \alpha = \frac{\pi}{2} \) the expressions of two important triple couplings among the CP-even Higgs bosons simplify to

\[
\lambda_{hhh} = 1, \quad \lambda_{Hhh} = 0. \tag{4.7}
\]

This means that the triple \( h \) coupling is SM like, while there is no \( Hhh \) coupling at the tree-level. The other couplings depend on the additional parameter \( m_{12} \) and they will not affect the discussion that we will have in this paper, so we will ignore them.

### 4.1.2. The SUSY case and the hMSSM approach.

The MSSM is essentially a two Higgs doublet model of type II, that is, the field \( \Phi_2 \) generates the masses of isospin down-type fermions and \( \Phi_1 \) the masses of up-type quarks. However, supersymmetry imposes strong constraints on the Higgs sector and among the four Higgs boson masses \( M_h, M_H, M_A \) and \( M_{H^\pm} \) (as well as the parameter \( m_{12} \)) and the two mixing angles \( \alpha \) and \( \beta \), only two of them are in fact independent at the tree level. These are in general taken to be \( M_A \) and \( \tan \beta \). Nevertheless, when the radiative corrections are included in the Higgs sector, in particular the dominant loop contributions from the top and stop quarks that have strong couplings to the Higgs bosons, many supersymmetric parameters will enter the game \([203–208]\). This is for instance the case of the SUSY scale, taken to be the geometric average of the two stop masses \( M_{\tilde{t}_1, \tilde{t}_2} = \sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}} \), the stop/sbottom trilinear couplings \( A_{\tilde{t} \tilde{b}} \) or the higgsino mass parameter \( \mu \) (other corrections, that involve the gaugino mass parameters \( M_{1,2,3} \) for instance, are rather small).

In particular, the radiative corrections in the CP-even neutral Higgs sector are extremely important and shift the value of the lightest \( h \) boson mass from the tree-level value predicted to be \( M_h \leq M_H | \cos 2\beta | \leq M_T \) to the value \( M_h = 125 \) GeV that has been measured experimentally. In the basis \((\Phi_2, \Phi_1)\), the CP-even Higgs mass matrix including the radiative corrections can be written as:

\[
M_{h}^2 = M_{h}^2 \begin{pmatrix}
\frac{c_{\beta}^2}{s_{\beta}^2} & -s_{\beta}c_{\beta} \\
-s_{\beta}c_{\beta} & s_{\beta}^2
\end{pmatrix} + M_{A}^2 \begin{pmatrix}
\frac{2 s_{\beta}^2}{c_{\beta}^2} & -s_{\beta}c_{\beta} \\
-s_{\beta}c_{\beta} & s_{\beta}^2
\end{pmatrix} \\
+ \begin{pmatrix}
\Delta M_{11} & \Delta M_{12} \\
\Delta M_{12} & \Delta M_{22}
\end{pmatrix},
\tag{4.8}
\]
where we have used the short-hand notation $c_{ij} \equiv \cos \beta_{ij}$ etc. and introduced the radiative corrections by a general $2 \times 2$ matrix $\Delta M_{ij}$. The neutral CP-even Higgs boson masses $H$, $h$, $A$ that diagonalizes the $h$, $H$ states, $H = \cos \alpha \Phi_0^0 + \sin \alpha \Phi_1^0$ and $h = -\sin \alpha \Phi_0^0 + \cos \alpha \Phi_1^0$ can then be easily derived. It is well known that in the $2 \times 2$ matrix for the radiative corrections, only the $\Delta M_{22}$ entry, which involves the by far dominant stop-top sector correction [203–205],

$$\Delta M_{22}^2 \approx \Delta M_{h}^{2 \, \text{loop}} = \frac{3 m_{\tilde{t}}^4}{8 \pi^2 v^2} \left[ \log \frac{M_S^2}{m_{\tilde{t}}^2} + \frac{X_t^2}{M_S^2} - \frac{X_t^4}{12 M_S^2} \right]$$

(4.9)

where $M_S$ is the SUSY scale and $X_t = A_t - \mu \tan \beta$ the stop mixing parameter, is relevant that is, $\Delta M_{22}^2 \gg \Delta M_{11}^2, \Delta M_{12}^2$. It has been recently advocated [209–214] that in this case, one can simply trade $\Delta M_{22}^2$ for the by now known $M_h$ value using

$$\Delta M_{22}^2 = \frac{M_A^2 (M_A^2 - M_h^2) - M_h^2 M_{\tilde{c}_{2\beta}}^2}{M_{\tilde{c}_{2\beta}}^2 + M_{\tilde{t}_{2\beta}}^2 - M_h^2}$$

(4.10)

and write the parameters $M_H$ and $\alpha$ in terms of $M_A$, $\tan \beta$ and $M_h$ in the simple form

$$M_H = \frac{(M_A^2 + M_Z^2 - M_h^2)(M_{\tilde{c}_{2\beta}}^2 + M_{\tilde{t}_{2\beta}}^2) - M_h^2 M_{\tilde{c}_{2\beta}}^2}{M_{\tilde{c}_{2\beta}}^2 + M_{\tilde{t}_{2\beta}}^2 - M_h^2}$$

$hMSSM :$

$$\alpha = -\arctan \left( \frac{M_{\tilde{t}_{2\beta}}^2 + M_{\tilde{c}_{2\beta}}^2}{M_{\tilde{c}_{2\beta}}^2 + M_{\tilde{g}_{2\beta}}^2 - M_h^2} \right)$$

(4.11)

This is the so-called $hMSSM$ approach [212, 214] which has been shown to provide a very good approximation of the MSSM Higgs sector [215].

In the case of the $H^\pm$ masses, the radiative corrections are very small at high enough $M_A$ and one has to a good approximation [216]

$$M_{H^\pm} \approx \sqrt{M_A^2 + M_W^2}.$$  

(4.12)
In this hMSSM approach, the MSSM Higgs sector with only the largely dominant $\Delta M_{12}^2$ correction, can again be described with only the two parameters $\tan \beta$ and $M_{h}$, as the loop corrections are fixed by the value of $M_{h}$. Another advantage of this approach is that it allows description of the low $\tan \beta$ region of the MSSM (see also [215]), which was overlooked as for SUSY scales of order 1 TeV, values $\tan \beta \lesssim 3$ were excluded because they lead to $M_{h} < 125$ GeV. The price to pay is that for such low $\tan \beta$ values, one has to assume $M_{h} \gg 1$ TeV and, hence, the model is excessively fine-tuned.

In fact, the possibility that the SUSY scale is rather high was an implicit assumption in this hMSSM approach. Indeed, one needs that $M_{h}$ is much larger than the other SUSY parameters, and in particular $M_{h} \gg |\mu|$, in such a way that $\Delta M_{12}^2 \gg \Delta M_{11,12}^2$ is indeed a good approximation. In this case, the subleading corrections, e.g. $\propto |\mu|/M_{h}$, that enter $\Delta M_{11,12}^2$ are too small and can be ignored to write equation (4.10). Another implicit assumption is that the CP-even Higgs sector can indeed be described by equation (4.8) also at very high $M_{h}$, which is a non-trivial statement but which has been verified in most cases [217].

In the MSSM, the couplings of the CP-even $h$ and $H$ and the CP-odd $A$ states to fermions and vector bosons are given by the type II entries of the 2HDM couplings shown in table 6 (again, the strength of the $H^\pm$ couplings to fermions will be similar to that of $A$). The only difference is that now the angle $\alpha$ included the radiative correction is fixed by the hMSSM relation equation (4.11). We note here that additional direct corrections should in principle enter the Higgs couplings but because $M_{h}$ is taken to be very large, they are assumed to have a small impact in the hMSSM and will be ignored.

Another important set of couplings are the Higgs self-couplings and in the hMSSM, they are again given in terms of $\beta$ and $\alpha$, with the latter fixed by $\tan \beta$ and $M_{A}$. For instance, the $hhh$ and $Hhh$ self-couplings, up to a normalization factor, read

$$\lambda_{hhh} = 3 \cos 2\alpha \sin (\beta + \alpha) + 3 \frac{\Delta M_{12}^2 \cos \alpha}{M_{Z}^2} \cos^2 \alpha$$
$$\lambda_{Hhh} = 2 \sin 2\alpha \sin (\beta + \alpha) - \cos 2\alpha \cos (\beta + \alpha) + 3 \frac{\Delta M_{12}^2 \sin \alpha}{M_{Z}^2} \cos^2 \alpha.$$  

A last remark is that when $M_{h} \gg M_{A}$, one is in the so-called decoupling regime [14] in which the $h$ state is light and as $\alpha = \beta - \pi/2$, has almost exactly the SM-Higgs couplings, $g_{hVV} = g_{hff} = 1$. The other CP-even $H$ and the charged $H^\pm$ bosons become heavy and degenerate in mass with the $A$ state, $M_{H} \approx M_{H^\pm} \approx M_{h}$, and decouple from the massive gauge bosons. The intensity of the couplings of the $H$ and $A$ states are the same. In this regime, the MSSM Higgs sector thus looks almost exactly like that of the 2HDM of type II in the alignment limit, especially if the additional assumption on the Higgs masses that we made to further simplify the model, equation (4.6), is used. The only exception will be the trilinear couplings, which are different because in the MSSM, there are additional loop corrections that make, for instance, $\lambda_{Hhh}$ non-zero in the decoupling limit.

4.2. Production of the Higgs bosons at hadron colliders

4.2.1. The gluon fusion process. In 2HDMs, the neutral Higgs bosons can be produced in the gluon fusion mechanism, $gg \rightarrow H, A$ as well as $h$, via loops involving mainly the heavy bottom and top quarks. At the one-loop level, i.e. the Born approximation in this case, the cross sections will depend on the magnitude of the Higgs couplings to quarks but there is also a difference between the CP-even and CP-odd cases. The $gg$ amplitudes with $\Phi = H, A$ follow the same trend except for the form factors, which are different for the two CP Higgs cases. In the alignment limit of 2HDMs, the lightest $h$ boson has SM-like couplings and its cross section follows exactly
that of the SM Higgs particle, which was also described in section 2.1.1. In the case of the heavier $\Phi = H, A$ states, the rates depend on $\Phi M$ and strongly on $\beta \tan \beta$. For small values, $\beta \approx \tan 1$, the dominant contribution to the cross section comes from top quark loops as the $\bar{\Phi} t t$ couplings, $/\beta \propto \Phi g_{1t \tan \beta}$, are strong. For low $\Phi$ masses, $\Phi M \lesssim 2m_t$, one could use the EFT approach in which the heavy top quark is integrated out and include not only the NLO QCD corrections [43, 44, 47] but also the NNLO corrections, which are also known in this case [48–50]. In fact, it was shown that it is a good approximation to incorporate the NLO corrections in this limit even for $M_b \gtrsim 2m_t$, provided that the Born term contains the full quark mass dependence [47]. Note that in the heavy top limit, while the LO loop amplitudes are different for the $Agg$ and $Hgg$ cases, as one has form factors for spin-$1/2$ particles $A_{jz}^{1/2} = \mp 2$ and $A_{jz}^{3/2} = \pm 2$ respectively for the CP-even and CP-odd cases, the QCD corrections at NLO and NNLO are about the same [218–220].

At high values of $\tan \beta$, $\tan \beta \gtrsim 10$, the $\Phi$ couplings to top quarks are strongly suppressed while those to the bottom quarks, $\Phi \bar{b} b \propto \tan \beta$, are enhanced. This implies that the contribution of the $b$-quark loop to the $gg \rightarrow \Phi$ processes (which was less than $10\%$ in the SM case) will become the dominant one. In fact, for extremely large $\tan \beta$ values, the cross section, which grows as $\tan^2 \beta$ and is enhanced by large logarithms $\log (m_t^2/M_\Phi^2)$, can be much larger than for an SM-like Higgs of the same mass. In this case, as $M_b \gg 2m_b$, one is in the chiral limit, in which the rates are approximately the same in the CP-even and CP-odd Higgs cases. In this limit, one can no longer use the EFT approach and integrate out the bottom quark to simplify the calculation of the higher order terms. The QCD corrections can thus be included only in NLO where they have been calculated keeping the exact quark mass dependence [47]. At LHC energies, the $K$-factors are much smaller in this case, $K_{NLO} \approx 1.2$, than in the case of the top quark loop, $K_{NLO} \approx 2$ [9].

For small to intermediate $\tan \beta$ values, $\tan \beta \approx 3–10$, the suppression of the $\Phi tt$ coupling is already effective while the $\Phi bb$ coupling is not yet strongly enhanced, resulting in production cross sections that are smaller than in the SM case. In fact, the minimum of the cross section is obtained for the value $\tan \beta \approx \sqrt{m_t/m_b} \approx 7$ as one has $m_t \approx 173$ GeV and $m_b \approx 3$ GeV for, respectively, the top-quark mass measured at colliders and the $b$-quark mass in the $\overline{\text{MS}}$ scheme evaluated at the scale of the Higgs mass. Here again, because the top and bottom loop contributions have a comparable weight, one can include only

Figure 12. The cross sections for associated neutral Higgs pair production in the $q\bar{q}$ annihilation channels, $\sigma(q\bar{q} \rightarrow hA)$ (upper curves) and $\sigma(q\bar{q} \rightarrow HA)$ (lower curves). The rates are as functions of $M_A$ in the hMSSM approach with $M_A = 125$ GeV and $\tan \beta = 1, 7$ and 30. The c.m. energies of $\sqrt{s} = 14$ TeV (left) and 100 TeV (right) have been assumed and the MSTW PDFs have been used.
those NLO QCD corrections which are known exactly (and of course, not the EW corrections, which are known only for an SM-like Higgs and not applicable in this case).

In the case of the MSSM, the same discussion that we had above on a 2HDM of type II approximately holds, but with two differences. The first one is that for a light H boson, say $M_H \lesssim 300$ GeV for low values of $\tan \beta$ and $M_H \lesssim 150$ GeV for $\tan \beta \gg 1$, we are not yet in the decoupling regime with $\alpha = \beta - \pi/2$ and the couplings of the $H$ state to fermions are not exactly the same as those of the pseudoscalar $A$ state. However, as discussed earlier, the difference between these two couplings should be small as the current Higgs data from the LHC indicate that we are close to this decoupling regime [25]. Nevertheless, present direct limits from SUSY searches at LHC and indirect limits from the mass of the observed $h$ state indicate that these particles should be rather heavy and, hence, their impact on the $gg \rightarrow H$ production cross section should be limited. In any case, in the $h$MSSM approach that we adopt here, these SUSY effects are ignored.

We perform the numerical analysis of this process taking the example of the MSSM and for this purpose, we use the code $\text{SusHi}$ [221] in which the $h$MSSM approach is implemented. The code implements the full top and bottom-loop contributions at NLO in QCD from [47, 222], NNLO-QCD top contributions in the heavy-top limit from [48, 220], and electroweak contributions by light quarks from [79, 223]. The cross sections for the production of the heavier CP-even $H$ (upper plots) and the CP-odd $A$ bosons are shown as functions of the respective Higgs mass in figure 9 for the LHC at $\sqrt{s} = 14$ TeV (left) and 100 TeV (right), with the MSTW PDFs.

Figure 13. The cross sections for neutral MSSM Higgs pair production in gg fusion, $gg \rightarrow hA$ (top) and $H A$ (bottom) as a function of $M_A$ for $\tan \beta = 1, 7$ and 30. Shown are the rates for $\sqrt{s} = 14$ TeV (left) and 100 TeV (right), with the MSTW PDFs.
and 30 are chosen to illustrate the low, intermediate and high tan β regimes. The MSTW PDF set has been adopted and the factorisation and renormalisation scales have been set to $\mu = \mu_F = \mu_R = \frac{1}{2} M_t$. 

As can be seen from the figure, and when comparing with the SM Higgs case discussed in section 2.1, the production rates for the MSSM CP-even H state are smaller than for a SM-like state at low tan β when the suppressed top quark loop

**Figure 14.** The production rates of the charged Higgs bosons as functions of $M_{H^\pm}$ for $\sqrt{s} = 14$ TeV (left) and $\sqrt{s} = 100$ TeV (right) for tan β = 1, 7 and 30. Three processes are considered: associated production with a top quark $qb \rightarrow H^\pm t$ (upper plots), pair production in $\bar{q}q$ annihilation and loop induced $gg$ fusion $pp \rightarrow q\bar{q} + gg \rightarrow H^+H^-$ (middle plots) and associated production with $h, H, A$ states $\bar{q}q \rightarrow H^\pm + h, H, A$ for tan β = 1 only (lower plots).
contribution is still dominant, and much larger at high $\beta$ values, when the $b$-quark loop contribution is strongly enhanced. They are small for $\beta \sim \tan 7$ when one has a maximal $\Phi_{g tt}$ suppression and a minimal $\Phi_{g bb}$ enhancement. For the value $\beta = \tan 30$ used for illustration, the $g g H$ cross sections are one to two orders of magnitude smaller than in the SM with a dominating top loop contribution. At $\sqrt{s} = 100$ TeV, the rates are enhanced with respect to LHC by only one order of magnitude for $M_H = 200$ GeV but by three orders of magnitude for $M_H = 2$ TeV due to the more favorable phase-space.

The cross sections for $g g \to A$ are about the same as for $H$ production for $M_A \gtrsim 200$ GeV, an approximation which improves at higher $\tan \beta$, for which the decoupling regime is more quickly reached and the $b$-loop contributions more important, resulting in almost equal $A_{g g}$ and $H_{g g}$ amplitudes in the chiral limit. A noticeable difference also occurs near the $2m_t$ threshold where the CP-odd amplitude $A_{H/2}^{g g}$ develops a singularity (that is unphysical and due to the QCD corrections in the absence of a regulating finite width [47]) while the CP-even one $A_{H/2}^{g g}$ simply reaches a maximum. For low $\tan \beta$ values, however, the amplitudes are slightly different for $H$ and $A$: first, because the couplings $g_{g \eta H}$ are not the same at low $M_A$ values and second, because of the different one-loop $A_{12}^{\Phi}$ form factors.

In 2HDMs of type II in the alignment limit, the gluon fusion cross sections are the same as in figure 9 for the CP-odd $A$ and also the CP-even $H$ boson in the decoupling limit, i.e. for $M_A \gtrsim 200$ GeV at sufficiently high $\tan \beta$ when $g_{H g H} \approx g_{A g H}$. In the case of type I 2HDMs in the alignment limit, since one has $g_{H g H} = g_{A g H} = 1/\tan \beta$ to both top and bottom quarks, the rates are simply given by the ones for $\beta = 1$ (when the top contribution dominates the $g g H$ loop) in figure 9 divided by $\tan^2 \beta$. The rates are thus approximately the same as in type II models for $\tan \beta \lesssim 3$ and much smaller at high $\tan \beta$.

4.2.2. Associated production with heavy quarks. Many of the features discussed above for the gluon fusion process appear in associated production of the neutral Higgs particles $\Phi = H, A$ (and even $h$, which is SM-like and thus has already been discussed before) with top and bottom quark pairs, $pp \to q\bar{q}, gg \to t\bar{t} \Phi$ and $bb \Phi$. The two processes have been analyzed in section 2 and the discussion holds in the
case of both the CP-even and CP-odd Higgs bosons of the 2HDMs. The only differences are that in 2HDMs one has to consider heavier particles and multiply the SM Higgs cross sections by the squares of the reduced Higgs Yukawa couplings

\[ \sqrt{\sigma_\Phi} = \sqrt{\sigma_{\Phi}}. \]  

(4.14)

Since \( g_{\Phi q q} \propto \tan^2 \beta \), with \( I_0 = \frac{1}{2} \), for isospin up (down)-type quarks, the cross sections for the \( tt\Phi \) and \( bb\Phi \) processes are the same for type II 2HDMs in the alignment limit and for the MSSM in the decoupling limit\( ^{16} \).

Figure 16. Projections for the HL-LHC with \( \sqrt{s} = 14 \) TeV (upper plot) and at \( \sqrt{s} = 100 \) TeV (lower plot) with 3000 fb\(^{-1}\) data for the 2\( \sigma \) sensitivity in the type II 2HDM \( [\tan \beta, M_0] \) plane (with \( M_\Phi = M_A = M_{h^\pm} = M_{\gamma^\pm} \)) when the ATLAS and CMS searches for the \( A/H/H^\pm \) states in their fermionic decays are combined in the alignment limit.

In the case of \( tt\Phi \) however, there is an additional difference at low Higgs masses and moderate centre of mass energies where different but small mass effects, \( O(m_t^2M_\Phi^2) \), in the matrix element squared appear between the CP-even and CP-odd cases as one is not close enough to the chiral limit, \( M_\Phi \gg M_t \). Since one has \( g_{\Phi q q} \propto 1/\tan \beta \), the cross section for the \( pp \rightarrow tt\Phi \) process is significant only for \( \tan \beta \) close to unity or lower, \( \tan \beta \lesssim 3 \).

Note that because at the high collider energies that we are considering here, the gluon luminosities are much larger compared to the quark luminosities, the cross sections for these channels are dominantly generated by the gluon fusion and not the \( q\bar{q} \) annihilation subprocesses. The QCD corrections have been known to NLO in the case of a CP-even Higgs state for a decade, and lead to a \( K \)-factor that is of order unity, \( K_{\text{NLO}} \approx 1.1 \) for \( M_{\text{H}} \approx 125 \) GeV at \( \sqrt{s} = 14 \) TeV \( [117, 118] \). In the case of a CP-odd Higgs boson, the QCD corrections have been derived more recently \( [226] \) and they lead to a \( K \)-factor that is only

\[ 16 \] In fact, at high values of \( \tan \beta \), there are additional corrections that one should take into account in the MSSM: the so-called \( \Delta \) corrections (see e.g. \[224, 225\]) that affect the \( \Phi bb \) couplings and which grow as \( \mu \tan \beta \); these corrections are small for heavy sbottoms and gluinos and are ignored in the hMSSM approach that we adopt here; for a discussion on this issue, see \[211, 214\].
slightly higher, \( K_{\eta A} \approx 1.18 \) for the same Higgs mass and c.m. energy. Considering the process at LO only, as will be done here, is therefore a reasonable approximation in this case.

The cross sections \( \sigma(pp \to t\bar{t}\Phi) \) with \( \Phi = H, A \) are shown in the upper part of figure 10 as a function of \( M_\Phi \) at both \( \sqrt{s} = 14 \, \text{TeV} \) (left) and \( \sqrt{s} = 100 \, \text{TeV} \) (right) in the hMSSM with \( \tan\beta = 1, 7 \) and 30. They have been obtained with the help of a modified version of the LO program H\( \Phi \)Q [88, 174] with the renormalisation and factorisation scales fixed to \( \mu_y = m_y + \frac{1}{2} M_\Phi \) and again using the MSTW set of structure functions. As can be seen, the cross sections are sizeable only for \( \tan\beta = 1 \) and not too high values of \( M_\Phi \). For \( M_\Phi = 1 \, \text{TeV} \) for instance, they are at the fb level at LHC but reach the pb level at \( \sqrt{s} = 100 \, \text{TeV} \). Note the difference between the rates of \( A \) and \( H \) at low \( M_A \) and \( \tan\beta \) as a result of the different masses and couplings in this area, which is outside the decoupling regime. One has approximately the same rates in 2HDMs in the alignment limit.

In contrast, the production rates for the \( pp \to bb\Phi \) process with \( \Phi = H, A \) are strongly enhanced at high \( \tan\beta \) values in type II 2HDMs like the MSSM since one has \( g_{b\bar{b}\Phi} \propto \tan\beta \). The cross sections are almost identical for the case of the CP-even \( H \) and CP-odd \( A \) bosons as chiral symmetry holds since \( M_A \gg m_b \) and as one has \( g_{\bar{b}b\Phi} = g_{b\bar{b}\Phi} \) in the alignment limit of 2HDMs and the MSSM at \( \tan\beta \gtrsim 5 \) where the decoupling limit is quickly reached. In type I models, \( g_{b\bar{b}\Phi} \propto \cot\beta \) and the cross section follows that of \( t\bar{t}\Phi \); it is thus negligible for \( \tan\beta > 1 \) as the \( bb\Phi \) Yukawa coupling, \( \propto m_b \), is small.

The NLO QCD corrections to the \( pp \to bb\Phi \) processes [124] are the same as those discussed in the SM case. Since \( m_b \) is very small compared to \( M_\Phi \) and chiral symmetry holds, the corrections are now the same for the CP-even and CP-odd states. This small \( m_b \) value leads to another major difference between the \( b\bar{b}\Phi \) and \( t\bar{t}\Phi \) cases: the cross sections \( \sigma(gg \to bb\Phi) \) develop (from the splitting of gluons to \( bb \) pairs) large logarithms, \( \log(Q^2/m_b^2) \), with the scale \( Q \) being typically of the order of the factorisation scale and the \( b\bar{b} \) quark transverse momentum, \( Q \approx \mu_F \approx p_T^b \). Hence, while one has reliable results for \( \sigma(gg \to bb\Phi) \) at high \( p_T^b \), the convergence of the perturbative series is poor in the opposite case unless the large logarithms are resummed. As discussed in section 2, this resummation is performed by treating the \( b\bar{b} \) as a massless parton in the proton, using the bottom PDF in a five-flavour scheme (5FS) and considering instead the fusion process \( b\bar{b} \to \Phi \) [127]. The requirement of one or two additional high-\( p_T \) final state quarks is fulfilled by considering the NLO [128–131] or NNLO [132] QCD corrections to this process.

The cross sections for the fusion process \( bb \to \Phi \) with \( \Phi = H, A \) are displayed as a function of \( M_\Phi \) in the lower part of figure 10 at the usual c.m. energies, \( \sqrt{s} = 14 \, \text{TeV} \) (left) and \( \sqrt{s} = 100 \, \text{TeV} \) (right) and adopting again the MSTW parton densities. We have used a modified version of the public code SusHi [221] in which the hMSSM approach was implemented and chosen again \( \tan\beta = 1, 7 \) and 30 for illustration. The NNLO QCD corrections, which are known in this case [132], are included with renormalisation and factorisation scales set at \( \mu_R = \mu_F = \frac{1}{2} M_H \), and we use \( m_b(M_\Phi) \) in the \( b\bar{b} \) Yukawa coupling, which significantly improves the convergence of the perturbative series. As expected, the production cross sections are extremely large at the value \( \tan\beta = 30 \): at the LHC, the production rates are approximately the same as in the \( gg \to \Phi \) fusion process at low \( M_\Phi \) but decrease less steeply with increasing Higgs mass. At \( \sqrt{s} = 100 \, \text{TeV} \), the cross sections increase by one or two orders of magnitude, respectively, depending whether we are at low or high Higgs masses. In type II 2HDMs like the MSSM and at high \( \tan\beta \), the \( pp \to bb\Phi \) processes are the dominant ones for neutral Higgs production at hadron colliders.

### 4.2.3. Other processes for single production

In the SM case, there were two additional processes for single Higgs production: Higgs-strahlung and vector boson fusion. Since these two processes directly involve the Higgs couplings to the massive gauge bosons \( V = W \) or \( Z \), they do not occur in the case of the pseudoscalar Higgs particle as the AVV coupling is forbidden by CP-invariance, \( g_{AVV} = 0 \). In fact, even in CP-violating MSSMs and more generally 2HDMs, these couplings are absent at tree-level and can be generated only at higher orders; they are thus also very small and lead to negligible cross sections for these processes at the LHC and beyond [17].

The Higgs-strahlung and vector boson fusion processes are in principle allowed only for the CP-even \( H \) state, except of course for the SM-like \( h \) boson which was discussed in section 2.1. Nevertheless, because in the alignment limit of 2HDMs and in the decoupling limit of the MSSM one has \( g_{HVV} = \cos(\beta - \alpha) \), the HVV coupling is also very small or (nearly) vanishing so that the cross sections \( \sigma(qq \to HV) \) and \( \sigma(qq \to VV'qq \to Hqq) \) are also tiny. The observation of such processes thus signals a departure from decoupling or alignment in two-Higgs doublet extensions and would allow for a direct measurement of the HVV coupling as both cross sections scale as \( g_{HVV}^2 \) (in addition to the indirect measurement that can be made by studying the rates of the lighter \( h \) boson, whose coupling is \( g_{HVV}^2 = 1 - g_{HVV}^2 = \sin^2(\beta - \alpha) \), in SM-like processes).

In figure 11 we assume a full strength HVV coupling \( g_{HVV} = 1 \) for illustration and plot the cross sections for the production of the CP-even \( H \) boson in the Higgs-strahlung and vector boson fusion processes as a function of \( M_H \) at \( \sqrt{s} = 14 \, \text{TeV} \) and 100 TeV. The rates follow the same trend as for the SM-Higgs boson and are much larger in the VBF than in the Higgs-strahlung case and drop less steeply with \( M_H \). In the VBF case and at \( \sqrt{s} = 100 \, \text{TeV} \), the cross section stays at the picobarn level for masses up to \( M_H \approx 2 \, \text{TeV} \) and, even for a departure from the alignment or decoupling limits as small
as $s_{\text{INV}}^2 = 10^{-3}$, one still obtains reasonable production cross sections if the luminosity is high enough (a few fb).

4.2.4. Neutral Higgs boson pair production. In a general two Higgs doublet model, the production of pairs of neutral Higgs particles in the continuum can be achieved in two main processes: $q \bar{q}$ annihilation, leading to $hA$ and $HA$ final states through the exchange of a virtual $Z$ boson,

$$q \bar{q} \rightarrow Z^* \rightarrow hA, HA,$$

or $gg$ fusion induced by heavy box and triangle diagrams, the latter being sensitive to the triple Higgs couplings, leading to various Higgs final states,

$$gg \rightarrow hh, HH, hH, AA \text{ and } hA, HA.$$

The partonic cross sections for neutral Higgs pair production in $q \bar{q}$ annihilation, $q \bar{q} \rightarrow \phi A$ with $\phi = h$ or $H$ are, up to coupling factors, those of the associated $\phi$ production with a $Z$ boson with another change in the phase-space factor to account for the production of two spin-zero bosons instead. Hence, as in Higgs-strahlung, the QCD corrections at NLO are simply those of the Drell–Yan process with off-shell $Z$ boson to be evaluated at the optimal scales $\mu_R = \mu_F = M_\phi$. The cross sections are proportional to the square of the reduced $\phi AZ$ coupling which, in 2HDMs, are simply given by

$$s_{\phi AZ}^2 = 1 - s_{\phi VW}^2.$$ 

Thus, in the alignment, or the decoupling limit of 2HDMs, one would have $g_{\phi AZ} = \cos(\beta - \alpha) \rightarrow 0$ and $g_{\phi VW} = \sin(\beta - \alpha) \rightarrow 1$. The crossection for $pp \rightarrow hA$ is thus expected to be very small (and would be another way to measure the departure from the decoupling limit) while that of $pp \rightarrow HA$ would be suppressed only by the phase-space. An analysis of gluon-fusion Higgs pair production in 2HDM has been presented in [228].

Note that $A + h/H$ production, as well as the production of all possible combinations of pairs of Higgs bosons, are also accessible in bottom quark fusion, $b \bar{b} \rightarrow \Phi \phi \Phi_2$ with $\Phi = h, H, A$ (which are equivalent to $gg \rightarrow b \bar{b} \phi \phi_2$ since in the former processes, $b$-quarks also come from gluon splitting). The lower $b$-quark luminosities compared to those of light quarks may be compensated for by large values of $\tan \beta$, which in principle strongly enhance the cross sections. Nevertheless, the rates stay at a rather modest level even for very high energies and very large $\tan \beta$ values.

The cross sections for the $q \bar{q} \rightarrow hA$ and $HA$ processes are shown in figure 12 again as a function of $M_A$ for $\tan \beta = 1, 7, 30$ assuming the $h$MSSM; the c.m. energies are also fixed at 14 TeV (left) and 100 TeV (right). In the case of $Ah$ production, the cross sections are sizeable only far from the decoupling limit, i.e. low $\tan \beta \sim 1$ and not too high $M_A$. The rates for $Ah$ do not depend on $\tan \beta$ in general and are significant for $M_A \lesssim 0.5$ TeV, especially at $\sqrt{s} = 100$ TeV where they still can be at the few fb level. Note that in 2HDMs, one can relax the mass equality $M_H \approx M_A$ that holds in the MSSM so that phase-space effects can be more (or less) favorable than in the $h$MSSM.

In the $gg$ fusion mechanism, a large number of processes for Higgs pair production is accessible. The corresponding Feynman diagrams involve top and bottom quark loops (and in SUSY theories, possibly their scalar partners when these particles are relatively light) that appear in box and triangular loops. In the latter channels, the pair production proceeds with the virtual exchange of the neutral CP-even $h$ and $H$ states (for $AA, Hh, HH$ and $hh$ production) or the CP-odd state $A$ (for $Ah$ and $AH$ production). The continuum production can be supplemented by resonant production with the Higgs boson decaying into pairs of lighter ones when phase-space allows (in addition to channels in which the resonant Higgs state decays into a lighter Higgs and a $Z$ boson as mentioned previously). In the MSSM, the only possibility would be a the resonant production of the $H$ state $gg \rightarrow H$, which then decays into two lighter CP-even Higgs bosons, $H \rightarrow hh$. However, in the decoupling limit corresponding to $M_\chi \gtrsim 350$ GeV for $\tan \beta \approx 1$, both the $Hhh$ coupling and the branching fraction for the $hh$ decay ($H$ decays are dominated by final states that are either $bb$ at high $\tan \beta \geq 1$ or $t\bar{t}$ at low $\tan \beta$) are very small. In 2HDMs in the alignment limit, the coupling $g_{Hhh}$ vanishes and there are no $H \rightarrow hh$ decays. Instead, because in this case the hierarchy of Higgs masses are different from the MSSM, one could have additional resonant channels such as $gg \rightarrow H \rightarrow AA$ (or $H \rightarrow H^*H^*$) for instance.

For the numerical analysis, we will however stick to the $hA$ and $HA$ processes discussed above for $q \bar{q}$ annihilation, since the results for $hH$ and $HH$ production are similar to the former and latter cases respectively and the SM-like $hh$ case has been discussed in section 2. In both cases, there is no resonant process as the particle that is exchanged in the s-channel is the $A$ boson. The rates are evaluated at LO using the program HPAIR [174] (the NLO corrections are not known for the dominant $b$-loop contributions at high $\tan \beta$ with the scales fixed to the invariant mass of the two final Higgses.

The cross sections $\sigma(gg \rightarrow hA)$ and $\sigma(gg \rightarrow HA)$ in the $h$MSSM are shown in figure 13 again as a function of $M_A$ for $\tan \beta = 1, 7$ and 30. The rates for $HA$ are most significant either at low or at high $\tan \beta$ when one of the $tt$ or $\Phi b\bar{b}$ couplings is sufficiently large, but one is limited by the phase-space (and by the fact that the process is of high order) and for $M_b = 500$ GeV, the cross section is at the few ten fb level even at $\sqrt{s} = 100$ TeV for $\tan \beta \approx 1$ or 30. Again, in 2HDMs with $M_A \approx M_H$, the rates can be larger if $H$ is much lighter than $A$. For the $gg \rightarrow Ah$ process, the phase space is more favorable but the cross sections are significant only for $\tan \beta \approx 1$ since $h$ has SM-like couplings and they are mainly generated by the box diagram with top quark exchange (there is no enhancement of the $hb\bar{b}$ coupling at high $\tan \beta$).

4.2.5. The production of the charged Higgs bosons. By far the dominant process for charged Higgs production at hadron colliders is certainly the top and antitop quark decays $t \rightarrow H^+b$ and $\bar{t} \rightarrow H^-\bar{b}$ for sufficiently light $H^\pm$ states, $M_{H^\pm} \lesssim m_t - m_b \sim 170$ GeV. The branching fractions of the decay $t \rightarrow H^+b$ (which has to compete only with the dominant
$t \rightarrow bW$ decay) are large at low or high $\tan \beta$ values when either the top or bottom quark component of the $H^{\pm} b \bar{b}$ coupling, $g_{H^{\pm}b} \approx m_t \tan \beta + m_b \tan \beta$, is significant, but even at intermediate values $\tan \beta \approx 7$, it is above the percent level. As the production cross sections for top quark pairs $q \bar{q}, gg \rightarrow t \bar{t}$ (with gluon fusion largely dominating at LHC energies and beyond) are extremely large at hadron colliders they allow probing all $\tan \beta$ values provided that the decay is not phase-space suppressed. Through the dominant $H^{\pm}$ decay for $\tan \beta \gtrsim 1$, $H \rightarrow \tau \nu$, searches have been conducted at the previous LHC run and masses $M_{H^{\pm}} \lesssim 140$ GeV have been excluded for any $\tan \beta$ value by the ATLAS and CMS collaborations. Masses up to the kinematical limit of $M_{H^{\pm}} \sim 170$ GeV could be probed at the LHC with $\sqrt{s} = 14$ TeV and high luminosity as will be discussed in the next section, leaving little space for a 100 TeV collider.

If the charged Higgs boson is heavier than the top quark, one has to resort to direct production processes, the most relevant one at high energies being associated production with top and bottom quarks in $q \bar{q}$ annihilation and mainly $gg$ fusion, $pp \rightarrow gg, q \bar{q} \rightarrow H^{\pm} b + H^{\mp} b$. Here again, the production sections are most significant at low or high $\tan \beta$ when the $H^{\pm}$ fusion coupling is large, and are minimal at intermediate $\tan \beta \approx m_t \tan \beta$ when the top component of the $H^{\pm}$ coupling is suppressed and the bottom component not sufficiently enhanced. The NLO QCD corrections to this process are significant [229–231] and exhibit logarithms that involve the ratio of the $b$-quark mass and the factorisation scale, $\log(m_b/\mu_F)$. As in the case of associated $HH$ production with $b \bar{b}$ pairs, these large logarithms can be resumed by treating the bottom quark as a parton and considering the process $gb \rightarrow H^{\pm} t$ in a five-flavour scheme. The NLO QCD corrections in this case are also known and parts of the corrections correspond in fact to the original process in the four-flavour scheme, $gg \rightarrow tbH^{\pm}$. The optimal choices for the renormalisation and factorisation scales have been shown to be $\mu_F = \frac{1}{2} \mu_R = \frac{1}{2}(M_{H^{\pm}} + m_t)$.

The numerical results for the cross sections in this process are displayed in the upper part of figure 14 as a function of $M_{H^{\pm}}$ again at $\sqrt{s} = 14$ TeV (left) and 100 TeV (right) for the usual three values $\tan \beta = 1, 7$ and 30. They have been evaluated using the program Prospino of [232, 233] and include the NLO QCD corrections in the five-flavour scheme with the scales set to the values given above. Here we used the CTEQ PDF set [234] instead. For the low and high $\tan \beta$ values, as they scale as $m^2 \cot^2 \beta$ or $m^2 \tan^2 \beta$ respectively, the cross sections exceed the 10 pb level only for $H^{\pm}$ masses up to $M_{H^{\pm}} \sim 200$ GeV at $\sqrt{s} = 14$ TeV but up to $M_{H^{\pm}} \sim 700$ GeV for $\sqrt{s} = 100$ TeV and they drop quickly with increasing charged Higgs mass.

Note that since the $H^{\pm}$ properties depend only on $\tan \beta$ and $M_{H^{\pm}}$, the discussion above holds in both the MSSM and type II 2HDMs. However, in the MSSM there are SUSY-QCD corrections that are proportional to $m_t \tan \beta$ and can be large for not too heavy SUSY particles; the notorious $\Delta_R$ corrections [224, 225] that appear also in $gg \rightarrow b \bar{b} + H/A$, which were discussed previously. We also ignore these corrections here as in the hMSSM approach we assume a high SUSY scale. In addition, in 2HDMs of type I, the cross sections simply scale as $(m^2 + m^2) \cot^2 \beta$ and hence are significant only at small $\tan \beta$ values, with QCD corrections that do not involve large logarithms and are hence moderate.

Charged Higgs bosons can also be produced in pairs. At LO, the process proceeds via $q \bar{q}$ annihilation with the exchange of a virtual photon and Z boson and the cross section depends only on $M_{H^{\pm}}$ and no other MSSM parameter. The QCD corrections are again those of Drell–Yan [98] which at NLO lead to an increase of the rate by $\approx 30\%$ at scales $\mu_F = \mu_R = 2M_{H^{\pm}}$. Another important pair production process at high energies is the $gg$ fusion mechanism, $gg \rightarrow H^{\pm} H^{\mp}$, which proceeds through loops involving top and bottom quarks, including contribution from the channels $gg \rightarrow h, H$ with $h, H \rightarrow H^{\pm} H^{\mp}$. In the MSSM both $s$-channel particles are off shell but in more general 2HDM, one could have $M_{H^{\pm}} \gtrsim 2M_{H^{\pm}}$ so that the $H^{\pm} H^{\mp}$ decay is resonant, a situation similar to $gg \rightarrow AA$. The process is known at LO and we take the scales to be $\mu_F = \frac{1}{2} \mu_R = \frac{1}{2}M_{H^{\pm}}$.

The combined cross sections at NLO for $q \bar{q}$ and LO for $gg$, evaluated with the program Prospino of [232, 233], are shown at $\sqrt{s} = 14$ TeV and 100 TeV in the middle frames of figure 14 as a function of $M_{H^{\pm}}$ for the three values $\tan \beta = 1, 7$ and 30. Here again, we use the CTEQ PDF set [234] and take the renormalisation and factorisation scales equal to the values mentioned above. The rates are dominated by $gg$ fusion and are about two orders of magnitude higher at 100 TeV than at 14 TeV. For either very low or very high $\tan \beta$, the rates can be at the level of 10 fb for $M_{H^{\pm}} = 1$ TeV at the highest energy.

Finally, associated $H^{\pm}$ production with a neutral Higgs boson, $q^c q^c \rightarrow \Phi H^{\pm}$ with $\Phi = h, H$ and $A$, is mediated by virtual $W$ exchange and the cross section is again simply the one in SM Higgs-strahlung $q \bar{q} \rightarrow H_{SM} W$, with the proper change of the coupling and phase-space factors. For associated production with the CP-even $h, H$ states $q^c q^c \rightarrow h, H \rightarrow H^{\pm} \bar{H}^{\mp}$, the cross sections follow exactly the same trend as the ones for $hA, HA$ production except for the overall normalization factor: once the two $H^{\pm}$ charges are summed, the rates are larger than for $A$ by about a factor of two for the same mass $M_{H^{\pm}} \sim M_h$. The rates for $AH^{\pm}$ production follow those of $HH^{\pm}$ in the decoupling limit of the MSSM or the alignment limit of 2HDMs. The rates are shown at NLO in the lower part of figure 14 at $\sqrt{s} = 14$ and 100 TeV for $\tan \beta = 1$ only. For $\tan \beta = 7$ and 30, they are almost the same for $AH^{\pm}$ and $HH^{\pm}$ production while they become negligible for $hH^{\pm}$ production.

4.3. The sensitivity on the extended Higgs sectors

In this section, we will attempt to quantify the increase in sensitivity on the heavier Higgs states that one can obtain by moving from a c.m. energy of 14 TeV to 100 TeV. We will assume that $3000 \text{ fb}^{-1}$ of data will be collected both at FCC-hh and the high luminosity (HL) LHC option. We will illustrate this gain in sensitivity in the hhMSSM, following exactly the results obtained in [214] where the HL-LHC case was discussed in detail (see also [235] for an another recent
analysis at 14 and 100 TeV). We will also discuss the case of 2HDMs of type II in the alignment limit and in the simple case where the masses of the three heavy Higgs states $H, H, H^+$ are comparable. We will concentrate on the direct searches of the heavy Higgs particles since in both the alignment and decoupling limits of the two models, the lighter Higgs particle will behave as the SM Higgs boson discussed in section 2 and we ignore the (complementary) indirect constraints that can be obtained from measuring its couplings to fermions and gauge bosons in the various production and decay channels.

To analyze the discovery reach of the heavy states at the LHC and beyond, it is necessary to know their various decay modes. In the case of the MSSM, the decay modes have been discussed in great detail and we refer the reader to the recent account [214] that was made in the context of the hMSSM and which we will follow closely here. In the case of 2HDMs, the situation will be more difficult to describe than in the hMSSM since the masses of the four Higgs states, as well as the two mixing angles $\alpha$ and $\beta$, are free parameters, making any attempt to perform a full analysis a daunting task. Even when including the results on the observed $h$ state, i.e. by fixing its mass to $M_h = 125$ GeV and forcing its couplings to be SM-like by adopting the alignment limit that leads to $\alpha = \beta = -\frac{1}{2} \pi$, one still has four input parameters to deal with, instead of two in the hMSSM for instance.

In the present analysis, we will further simplify our discussion in type II 2HDMs by assuming that, similarly to the hMSSM case, the heavier $H, A$, and $H^+$ states have a comparable mass, $M_H^+ \approx M_B \approx M_h$ in such a way that decay of one Higgs particle to another one and a gauge boson is kinematically not allowed. This ensures that complicated decays such as $H \rightarrow AZ, H^+W^\mp, A \rightarrow HZ, H^+W^\pm$ and $H^+ \rightarrow AW^\mp, HW^\pm$ are kinematically not allowed at the two-body level. Our justification is that in fact, even if these modes are allowed, they cannot compete with the fermionic decays of the $H^+, H, A$ states involving top or bottom quarks (and tau-leptons) at either low or high $\tan \beta$. Adding the fact that, in the alignment limit, decays such as $H \rightarrow WW, ZZ$ or $A \rightarrow hZ$, which involve the reduced couplings $g_{HVV} = g_{AVV} = \cos(\beta - \alpha) = 0$ as well as the decay $H \rightarrow hh$, are absent or have small branching ratios, one will have a simple decay pattern.

Indeed, in this alignment limit with $M_{H^+} \approx M_B \approx M_h$, the only important decays of the neutral $HA$ states will be into $t\bar{t}$, $b\bar{b}$ and $\tau^+\tau^-$ final states, while those of the charged Higgs boson will be into $tb$ and $\tau
u$. In fact, at low $\tan \beta$ only the decays $A, H \rightarrow t\bar{t}$ for masses $M_A \approx M_h \gtrsim 2m_t$, and $H^+ \rightarrow tb$ for $M_{H^+} \gtrsim 180$ GeV are relevant with branching ratios of order unity, while at high $\tan \beta$, one would have $\text{BR}(A(H \rightarrow b\bar{b}) \approx \text{BR}(H^+ \rightarrow t\bar{b}) \approx 90\%$ and $\text{BR}(A(H \rightarrow \tau^+\tau^-) \approx \text{BR}(H^+ \rightarrow \tau\nu) \approx 10\%$, a simple reflection of $3m_t^2/m_h^2 \approx 10$ (with 3 being the colour factor). At intermediate values of $\tan \beta$ and above the top threshold the two sets involving top and bottom + tau decays would have comparable branching rates.

In the hMSSM outside the decoupling regime, i.e. for low $\tan \beta$ values with $M_h \lesssim 350$ GeV, the decay pattern can be rather involved as discussed in [214] for instance and a summary is as follows: (i) above the $2M_V$ threshold, the $H \rightarrow WW$ and $ZZ$ decay rates are still significant as $g_{HVV}$ is not completely suppressed; (ii) for $2M_h \lesssim M_t \lesssim 2m_t$, $H \rightarrow hh$ is the dominant $H$ decay mode at low $\tan \beta$ as the $Hhh$ self-coupling is large in this case; (iii) for $M_h + M_Z \lesssim M_s \lesssim 2m_t$, the $A \rightarrow hZ$ decays occur with large rates at low $\tan \beta$; (iv) this is also the case for the channel $H^+ \rightarrow Wh$ which is important for $M_{H^+} \lesssim 250$ GeV, but at intermediate $\tan \beta$ values this time.

In the context of the hMSSM, the impact of the various searches that have been performed by the ATLAS and CMS collaborations at $\sqrt{s} = 7 + 8$ TeV with up to $\approx 25$ fb$^{-1}$ data has been used to constrain the $[M_h, \tan \beta]$ parameter space of the model [214]. The searches that have been considered are essentially the fermionic Higgs decays $H/A \rightarrow \tau\tau$ and $H^+ \rightarrow \tau\nu$ and the bosonic ones, $H \rightarrow WW, ZZ, hh$ and $A \rightarrow Zh$. The channels $H^+ \rightarrow tb$ as well as the $H, A \rightarrow t\bar{t}$ in some approximation have also been included. As experimental input, the following Higgs searches and measurements, published by the ATLAS and CMS collaborations, have been used.

| Search channel | ATLAS | CMS |
|----------------|-------|-----|
| $A/H \rightarrow \tau\tau$ | [236] | [237] |
| $A \rightarrow Zh$ | — | [238] |
| $H \rightarrow hh$ | [239] | [240] |
| $H \rightarrow WW$ | — | [241] |
| $H \rightarrow ZZ$ | — | [242, 243] |
| $H^+ \rightarrow \tau\nu/ltb$ | [244, 245] | [246, 247] |

Concerning the $H/A \rightarrow t\bar{t}$ analysis at low $\tan \beta$ with the $A/H$ states dominantly produced in the $gg$ fusion mechanism, there was no dedicated analysis from the ATLAS and CMS collaborations. As a first approximation, we thus estimated the sensitivity in this channel by considering searches that have been performed for high mass spin-1 electroweak resonances (in particular new Z’ and Kaluza–Klein electroweak gauge bosons) that decay into top quarks and adapting them to our case. One should note however that this adaptation is quite naive as for spin-1 particles the main production channel is $qq \rightarrow t\bar{t}$ background; the resonances show up as peaks in the $t\bar{t}$ invariant mass distribution. In the higgs case, the main process is $gg \rightarrow H/A$ and would thus interfere with the $gg \rightarrow t\bar{t}$ background [248–251]. This interference depends on the mass and total width of the Higgs and on their CP-nature, making it either constructive or destructive. This leads to a rather complex signature with a peak-dip structure of the $t\bar{t}$ mass distribution that has not yet been experimentally addressed.

The constraints obtained in [214] from the ATLAS and CMS searches with the 25 fb$^{-1}$ data collected at 7 + 8 TeV in the channels above are quite impressive and a large part of the parameter space (in particular at high $\tan \beta$ for $M_h \lesssim 300$ GeV) has been already excluded. If no new signal is again observed, they can be still vastly improved at the next LHC phase with a c.m. energy of $\sqrt{s} = 14$ TeV and with one or two orders of magnitude accumulated data. The projections
for this case, and the procedure to obtain them, have been discussed in [214] to which we refer for the details. Here, we simply summarise how the projections are obtained at a given c.m. energy and luminosity.

For a specific search channel, one starts with the expected median 95%CL exclusion limits that have been given by the ATLAS and CMS collaborations in the searches performed at $7 + 8$ TeV with $\sim 5 + 20$ fb$^{-1}$ data. One then assumes that the sensitivity will approximately scale with the square root of the number of expected events and does not include any additional systematical effect. In addition, having only the information on the signal cross sections for a given Higgs mass, and not the corresponding background rates for the same mass bin, one needs to make the additional naive assumption that the background also simply scales as the signal cross sections (which is true for many channels).

The output of the projections following this procedure is presented in the $[\tan \beta, M_A]$ hMSSM plane in the lower part of figure 15 in the fermionic and bosonic Higgs search channels, including our naive treatment of the $gg \rightarrow H/A \rightarrow t\bar{t}$ mode. The c.m. energy of $\sqrt{s} = 100$ TeV and a luminosity of 3000 fb$^{-1}$ have been assumed as well as the same signal and background efficiencies for the high energy hadron collider as at the 8 TeV LHC. For the sake of comparison, we also show in the upper plot of the figure the sensitivity at the HL-LHC with $\sqrt{s} = 14$ TeV and 3000 fb$^{-1}$ data that was obtained in [214].

Figure 16 shows the same sensitivity in the $[\tan \beta, M_A]$ plane in our simplified type II 2HDM in which we assume the alignment limit $\alpha = \beta - \frac{1}{2}$ and approximately mass degenerate heavy Higgs bosons $M_A = M_H = M_{H^\pm} = M_{\Phi}$. The same collider parameters as for the hMSSM case have been adopted and we also show for the sake of comparison the sensitivity in 2HDMs at the HL-LHC.

As can be seen, a vast improvement of the sensitivity to the MSSM and 2HDM parameter spaces is expected at 100 TeV compared to the HL-LHC with the same luminosity. In the very low and very high tan $\beta$ regions, masses close to 3 TeV can be probed at a 100 TeV collider in, respectively, the $H/A \rightarrow t\bar{t}$ and $H/A \rightarrow \tau^+\tau^-$ modes, compared to only 1.5 TeV at the HL-LHC. The $H/A \rightarrow t\bar{t}$ and $\tau^+\tau^-$ channels intersect at $M_A = 1.5$ TeV for FCC-hh (instead of $M_A = 1.5$ TeV for HL-LHC), a mass value below which the entire hMSSM parameter space is fully covered by the searches. In the case of the charged Higgs state, the FCC-hh mass reach is lower: $M_{H^\pm} = 2$ TeV (1.2 TeV) at $\tan \beta = 60(1)$, but is significantly higher than at HL-LHC. In all cases, the sensitivity of a 100 TeV collider in probing the parameter space is twice as large as HL-LHC with the same luminosity.

Note that for the fermionic channels, the discussion is qualitatively the same in our adopted 2HDM as for the hMSSM, with the difference that the sensitivity is slightly higher at intermediate tan $\beta$ values in the former case. Indeed, in the hMSSM, the branching fractions in these fermionic channels are slightly suppressed by some bosonic modes such as $H \rightarrow hh, WW, ZZ$ and $A \rightarrow hZ$ that still survive.

In turn, for these bosonic modes, there is some sensitivity only in the hMSSM. The modes $H \rightarrow hh$ and $A \rightarrow hZ$, besides the decays $H \rightarrow WW$ and $ZZ$, could be observed up to $M_A = 1$ TeV (0.5 TeV) for $\tan \beta \approx 7$, a significant improvement over the LHC at least in the low tan $\beta$ range.

5. Conclusions

We have analyzed the prospects of future high-energy proton–proton colliders to probe the Higgs sectors of the standard model and of some of its new physics extensions. In the SM context, we have studied the production of the observed Higgs particle in the dominant channels and have shown that the one to two orders of magnitude increase of the cross sections (depending on the channels considered) at energies close to $\sqrt{s} = 100$ TeV would allow for a much more accurate determination of the Higgs couplings to gauge bosons and fermions. Some observables, like the ratio of partial widths for Higgs decays into two photons and into four leptons, could then be measured at the per-mille level with a few ab$^{-1}$ data. We then analyzed the various processes for Higgs pair production that allow for the measurement of the triple Higgs coupling, probably the only parameter that would remain undetermined after the high luminosity option of the LHC. Again, the rates at 100 TeV are so large that a relatively precise measurement would be possible. The last SM parameter to be probed would then be the quartic Higgs coupling, which can be accessed only in triple Higgs production. At $\sqrt{s} = 100$ TeV, the rates for the dominant gluon-fusion process are not completely negligible and a luminosity of a few ten ab$^{-1}$ could make the formidable challenge of observing three Higgs particles not entirely hopeless.

In a second step, we have considered the production of the invisible particles that form the cosmological dark matter. We have worked in a model-independent effective framework in which the DM particle is either a spin-zero, a spin-one or a spin-half Majorana fermion that interacts only through the Higgs portal. If the DM particles are heavier than $\frac{1}{2} M_h$, the only way to observe them would be through Higgs exchange in the continuum. We have thus evaluated the cross sections in three processes with missing energy: gluon fusion with an extra jet, vector boson fusion and Higgs-strahlung, and shown that at $\sqrt{s} = 100$ TeV one could probe DM particles with a few 100 GeV mass for favorable couplings.

Finally, we have evaluated the potential of a 100 TeV proton collider in probing the heavy Higgs bosons that are present in extended Higgs scenarios, taking the example of two Higgs doublet models and their minimal supersymmetric SM incarnation. We have discussed in a comprehensive manner the production of the heavier CP-even, the CP-odd and the charged Higgs states in all possible channels: single production, associated production with massive fermions or gauge bosons and pair production. Taking the examples of a 2HDM in the alignment limit and the so-called hMSSM, we have shown that a collider with $\sqrt{s} = 100$ TeV and 3 ab$^{-1}$ data could cover the entire parameter space of the models for Higgs masses up to 1 TeV and that some channels could be observed for masses of the additional Higgs bosons up to 3 TeV if their couplings to fermions are significant.
The sensitivity of such a high energy collider in probing the electroweak symmetry breaking mechanism is thus far superior to that of the LHC with high luminosity.

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