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NLO corrections to differential cross sections for pseudo-scalar Higgs boson production

B. Field [a], J. Smith [a], M.E. Tejeda-Yeomans [a,1], W.L. van Neerven [b]

[a] *C.N. Yang Institute for Theoretical Physics, State University of New York at Stony Brook, New York 11794-3840, USA*
[b] *Instituut-Lorentz, University of Leiden, P.O. Box 9506, 2300 RA Leiden, The Netherlands*

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Abstract

We have computed the full next-to-leading (NLO) QCD corrections to the differential distributions $d^2\sigma/(dp_T\,dy)$ for pseudo-scalar Higgs (A) production at large hadron colliders. This calculation has been carried out using the effective Lagrangian approach which is valid as long as the mass of the pseudo-scalar Higgs boson m_{A} and its transverse momentum p_T do not exceed the top-quark mass m_t. The shape of the distributions hardly differ from those obtained for scalar Higgs (H) production because, apart from the overall coupling constant and mass, there are only small differences between the partonic differential distributions for scalar and pseudo-scalar production. Therefore, there are only differences in the magnitudes of the hadronic differential distributions which can be mainly attributed to the unknown mixing angle β describing the pseudo-scalar Higgs coupling to the top quarks.

The scalar Higgs boson H, which is the corner stone of the standard model, is the only particle which has not yet been observed. Its discovery or its absence will shed light on the mechanism how particles acquire mass as well as answer questions about super-symmetric extensions of the standard model or about the compositeness of the existing particles and the Higgs boson. Among these two alternatives super-symmetry is the most appealing one, in particular, the minimal super-symmetric extension of the standard model. The latter version contains two complex Higgs doublets instead of one and it is therefore called the Two-Higgs-Doublet Model (2HDM). Here the scalar particle spectrum contains both the Higgs boson H and another neutral scalar boson h. Furthermore, it contains two charged scalar bosons $\mathrm{H}^{\pm}$ and a neutral pseudo-scalar Higgs boson A. The tree-level masses are expressed in two independent parameters, namely, the mass m_{A} and the ratio of the vacuum expectation values of the two Higgs doublets defined by $\tan\beta = v_2/v_1$ (see, e.g., [1]). According to the experiments at LEP their parameter ranges are restricted so that $m_{\mathrm{A}} < 91.9\ \mathrm{GeV}/c^2$ and $0.5 < \tan\beta < 2.4$ [2] are excluded. In this Letter we study A-production which in lowest order proceeds via gluon–gluon fusion where the gluons are coupled to the A via a heavy flavour triangular loop. This is similar to H-production except that now the coupling constant describing the interaction

E-mail address: neerven@lorentz.leidenuniv.nl (W.L. van Neerven).
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of the A with the quarks depends on both the masses of the quarks and on the angle β. This follows from the 2HDM where the coupling constants of the up and down quarks behave like $g_{\rm up} \sim m_u \cot\beta$ and $g_{\rm down} \sim m_d \tan\beta$, respectively [1]. Since the effective Lagrangian approach below is only valid in the case the mass of the quark appearing in the triangular loop satisfies the condition $m_q \gg m_{\rm A}$, the bottom quark is excluded. However, then we have to require that in the 2HDM the coupling of the A to the top-quark is stronger than to the bottom-quark which implies the condition

$$\frac{m_t}{m_b} \gg \tan^2\beta. \tag{1}$$

If we choose $m_b = 4.5\ \text{GeV}/c^2$ and $m_t = 173.4\ \text{GeV}/c^2$ one obtains the inequality $\tan\beta \ll 6.21$. In view of the experimental boundaries above one can conclude that the results of the calculation below can be only applied for the regions $\tan\beta < 0.5$ and $2.4 < \tan\beta < 6.21$.

In the effective Lagrangian approach scalar H-production is described by the Lagrangian density [3,4]

$$\mathcal{L}_{\rm eff}^{\rm H} = G_{\rm H}\,\Phi^{\rm H}(x)\,O(x), \quad \text{with} \quad O(x) = -\frac{1}{4}\,G_{\mu\nu}^a(x)\,G^{a,\mu\nu}(x), \tag{2}$$

whereas pseudo-scalar A-production is obtained from [5–7]

$$\mathcal{L}_{\rm eff}^{\rm A} = \Phi^{\rm A}(x)\big[G_{\rm A}\,O_1(x) + \widetilde{G}_{\rm A}\,O_2(x)\big],$$
$$\text{with} \quad O_1(x) = -\frac{1}{8}\,\epsilon_{\mu\nu\lambda\sigma}\,G_a^{\mu\nu}(x)\,G_a^{\lambda\sigma}(x), \qquad O_2(x) = -\frac{1}{2}\partial^\mu \sum_{i=1}^{n_f} \bar{q}_i(x)\gamma_\mu\gamma_5 q_i(x), \tag{3}$$

where $\Phi^{\rm H}(x)$ and $\Phi^{\rm A}(x)$ are the scalar and pseudo-scalar fields, respectively, and n_f denotes the number of light flavours. Up to NLO the second operator $O_2(x)$ contributes and in the case of massless quarks it cannot be neglected in higher orders. The effective couplings $G_{\rm B}$ ($\rm B = H, A$) are determined by the top-quark triangular graph describing the decay process $B \to g + g$ in the limit $m_t \to \infty$

$$G_{\rm B}^2 = 4\sqrt{2}\left(\frac{\alpha_s(\mu_r^2)}{4\pi}\right)^2 G_F\,\tau_{\rm B}^2\,F_{\rm B}^2(\tau_{\rm B})\,\mathcal{C}_{\rm B}^2\left(\alpha_s(\mu_r^2), \frac{\mu_r^2}{m_t^2}\right), \qquad \tau_{\rm B} = \frac{4\,m_t^2}{m_{\rm B}^2}, \quad \rm B = H, A, \tag{4}$$

and the functions $F_{\rm B}$ are defined by

$$F_{\rm H}(\tau) = 1 + (1-\tau)\,f(\tau), \qquad F_{\rm A}(\tau) = f(\tau)\cot\beta,$$
$$f(\tau) = \arcsin^2\frac{1}{\sqrt{\tau}}, \quad \text{for } \tau \geqslant 1, \qquad f(\tau) = -\frac{1}{4}\left(\ln\frac{1-\sqrt{1-\tau}}{1+\sqrt{1-\tau}} + \pi i\right)^2, \quad \text{for } \tau < 1. \tag{5}$$

In the large m_t-limit $F(\tau)$ behaves as

$$\lim_{\tau\to\infty} F_{\rm H}(\tau) = \frac{2}{3\,\tau}, \qquad \lim_{\tau\to\infty} F_{\rm A}(\tau) = \frac{1}{\tau}\cot\beta. \tag{6}$$

Here m and m_t denote the masses of the (pseudo-)scalar Higgs boson and the top quark respectively. The running coupling constant is given by $\alpha_s(\mu_r^2)$ where μ_r denotes the renormalization scale and G_F is the Fermi constant. The coefficient functions $\mathcal{C}_{\rm B}$ originate from the corrections to the top-quark triangular graph provided one takes the limit $m_t \to \infty$. We have presented the couplings $G_{\rm B}$ in Eq. (4) for general m_t on the Born level only in order to keep some part of the top-quark mass dependence. This is an approximation because the gluons which couple to the (pseudo-)scalar Higgs boson via the top-quark loop in the partonic subprocesses are very often virtual. The virtual-gluon momentum dependence is neither described by $F_{\rm B}(\tau)$ nor by $\mathcal{C}_{\rm B}$. For on-mass-shell gluons the latter quantity has been computed in the large m_t limit up to order α_s in [3,5,6] and up to order α_s^2 in [4,7]. For our NLO calculations we only need these coefficient functions corrected up to order α_s and they read

$$\mathcal{C}_{\mathrm{H}}\left(\alpha_s(\mu_r^2), \frac{\mu_r^2}{m_t^2}\right) = 1 + \frac{\alpha_s^{(5)}(\mu_r^2)}{4\pi}(11) + \cdots, \tag{7}$$

$$\mathcal{C}_{\mathrm{A}}\left(\alpha_s(\mu_r^2), \frac{\mu_r^2}{m_t^2}\right) = 1, \tag{8}$$

where $\alpha_s^{(5)}$ is presented in a five-flavour number scheme. Notice that Eq. (8) holds in all orders because of the Adler–Bardeen theorem [8]. The effective Lagrangian approach has been successfully applied to compute the total cross section of scalar Higgs production in hadron–hadron collisions in NLO [3] and NNLO [9–13]. In the case of pseudo-scalar Higgs production this cross section was computed in NLO in [5,6] and in NNLO in [14,15].

In this Letter we study the semi-inclusive reaction with one pseudo-scalar Higgs boson A in the final state which is given by

$$\mathrm{H}_1(p_1) + \mathrm{H}_2(p_2) \to \mathrm{A}(-p_5) +' X', \tag{9}$$

where H_1 and H_2 denote the incoming hadrons and X represents an inclusive hadronic final state. Further we define the following kinematical invariants

$$S = (p_1 + p_2)^2, \qquad T = (p_1 + p_5)^2, \qquad U = (p_2 + p_5)^2. \tag{10}$$

The latter two invariants can be expressed in terms of the transverse momentum p_T and rapidity y variables as

$$\begin{aligned} T &= m^2 - \sqrt{S}\sqrt{p_T^2 + m^2}\cosh y + \sqrt{S}\sqrt{p_T^2 + m^2}\sinh y, \\ U &= m^2 - \sqrt{S}\sqrt{p_T^2 + m^2}\cosh y - \sqrt{S}\sqrt{p_T^2 + m^2}\sinh y, \end{aligned} \tag{11}$$

where m is the mass of the pseudo-scalar Higgs boson. The hadronic cross section is given by

$$\begin{aligned} S^2\frac{d^2\sigma^{\mathrm{H}_1\mathrm{H}_2}}{dT\,dU}(S,T,U,m^2) &= \sum_{a,b=q,g}\int_{x_{1,\min}}^{1}\frac{dx_1}{x_1}\int_{x_{2,\min}}^{1}\frac{dx_2}{x_2}\,f_a^{\mathrm{H}_1}(x_1,\mu^2)f_b^{\mathrm{H}_2}(x_2,\mu^2)s^2\frac{d^2\sigma_{ab}}{dt\,du}(s,t,u,m^2,\mu^2), \\ \text{with} \quad x_{1,\min} &= \frac{-U}{S+T-m^2}, \qquad x_{2,\min} = \frac{-x_1(T-m^2)-m^2}{x_1S+U-m^2}, \end{aligned} \tag{12}$$

where s, t and u are the partonic analogues of S, T and U in Eq. (10) where p_1 and p_2 now represent the incoming parton momenta. Further $f_a^{\mathrm{H}_i}$ denotes the parton density corresponding to hadron H_i and μ stands for the factorization scale which for convenience is set equal to the renormalization scale μ_r appearing in Eq. (4). The NLO corrections to the partonic cross section $d^2\sigma/(dt\,du)$ in the case of H-production based on the effective Lagrangian in Eq. (2) are presented in [16] and [17]. Here we will give the corresponding results for the A described by the Lagrangian in Eq. (3). The calculation proceeds in the same way as presented in [17]. We use n-dimensional regularization in order to compute the loop and phase space integrals which contain ultraviolet, infrared and collinear singularities. However, there is one extra complication in the pseudo-scalar case. This concerns the Levi-Civita tensor in Eq. (3) which is essentially a four dimensional object. Here we follow the same prescription as in Eq. (4) in [15] where the product of two Levi-Civita tensors is contracted in n dimensions if one sums over dummy Lorentz indices. The LO subprocesses contributing to the partonic cross section are given by

$$g + g \to g + \mathrm{A}, \qquad q + \bar{q} \to g + \mathrm{A}, \qquad q(\bar{q}) + g \to q(\bar{q}) + \mathrm{A}. \tag{13}$$

The matrix elements squared do not differ from those derived for the scalar H provided $n = 4$, see [5,6], which implies that the LO double differential partonic cross sections are the same for both bosons except for an overall constant given by $F_{\mathrm{B}}(\tau)$ in Eq. (5). In NLO one has to compute the one-loop virtual corrections to the processes in Eq. (13) above and to add the contributions from the following two-to-three-body reactions

$$g+g\to g+g+\mathrm{A},\qquad g+g\to q_i+\bar{q}_i+\mathrm{A},\tag{14}$$

$$q+\bar{q}\to g+g+\mathrm{A},\qquad q_1+\bar{q}_2\to q_1+\bar{q}_2+\mathrm{A},\quad q_1\neq q_2,$$

$$q+\bar{q}\to q_i+\bar{q}_i+\mathrm{A},\quad q_i\neq q,\qquad q+\bar{q}\to q+\bar{q}+\mathrm{A},\tag{15}$$

$$q_1+q_2\to q_1+q_2+\mathrm{A},\quad q_1\neq q_2,\qquad q+q\to q+q+\mathrm{A},\tag{16}$$

$$q(\bar{q})+g\to q(\bar{q})+g+\mathrm{A}.\tag{17}$$

After renormalization of the strong coupling constant α_s and mass factorization which are carried out in the $\overline{\mathrm{MS}}$-scheme we obtain the NLO corrected coefficient functions according to the procedure in [17]. The coefficient functions are as long as in the case of H-production so that they cannot be explicitly presented. However, the differences between the results for the H and the A are so small that we can show them below. If we put for simplicity $G_{\mathrm{H}}=G_{\mathrm{A}}=G$ and $m_{\mathrm{H}}=m_{\mathrm{A}}=m$ the differences between the soft-plus-virtual differential cross sections are given by

$$s^2\frac{d^2\sigma^{\mathrm{S+V}}_{gg\to g\mathrm{A}}}{dt\,du}-s^2\frac{d^2\sigma^{\mathrm{S+V}}_{gg\to g\mathrm{H}}}{dt\,du}=\pi\delta(s+t+u-m^2)G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2\frac{N}{(N^2-1)^2}\Big[2\big|M^{(1)}_{gg\to g\mathrm{B}}\big|^2\Big],\tag{18}$$

$$s^2\frac{d^2\sigma^{\mathrm{S+V}}_{q\bar{q}\to g\mathrm{A}}}{dt\,du}-s^2\frac{d^2\sigma^{\mathrm{S+V}}_{q\bar{q}\to g\mathrm{H}}}{dt\,du}$$
$$=\pi\,\delta(s+t+u-m^2)G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2\frac{1}{N^2}\Big[2\,C_A\big|M^{(1)}_{q\bar{q}\to g\mathrm{B}}\big|^2+(C_F-C_A)\big|MB^{(1)}_{q\bar{q}\to g\mathrm{B}}\big|^2\Big],\tag{19}$$

$$s^2\frac{d^2\sigma^{\mathrm{S+V}}_{qg\to q\mathrm{A}}}{dt\,du}-s^2\frac{d^2\sigma^{\mathrm{S+V}}_{qg\to q\mathrm{H}}}{dt\,du}$$
$$=\pi\,\delta(s+t+u-m^2)G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2\frac{1}{N(N^2-1)}\Big[2C_A\big|M^{(1)}_{qg\to q\mathrm{B}}\big|^2+(C_F-C_A)\big|MB^{(1)}_{qg\to q\mathrm{B}}\big|^2\Big].\tag{20}$$

The colour factors of the group $SU(N)$ are given by $C_A=N$ and $C_F=(N^2-1)/(2N)$ and the Born matrix elements squared belonging to the processes in Eq. (13) are equal to

$$\big|M^{(1)}_{gg\to g\mathrm{B}}\big|^2=N(N^2-1)\frac{1}{stu}\big[s^4+t^4+u^4+m^8\big],\tag{21}$$

$$\big|M^{(1)}_{q\bar{q}\to g\mathrm{B}}\big|^2=C_AC_F\,\frac{1}{s}\big[t^2+u^2\big],\tag{22}$$

$$\big|M^{(1)}_{qg\to q\mathrm{B}}\big|^2=C_AC_F\frac{1}{u}\big[-s^2-t^2\big].\tag{23}$$

The differences above can be wholly attributed to the virtual corrections and not to the soft gluon contributions which are the same for both H- and A-production. These virtual corrections also entail some extra terms denoted by

$$\big|MB^{(1)}_{gg\to g\mathrm{B}}\big|^2=\frac{2}{3}N(N^2-1)\frac{m^2}{stu}\big[stu+m^2(st+su+tu)\big],\tag{24}$$

$$\big|MB^{(1)}_{q\bar{q}\to g\mathrm{B}}\big|^2=C_AC_F(-t-u),\tag{25}$$

$$\big|MB^{(1)}_{qg\to q\mathrm{B}}\big|^2=C_AC_F(s+t).\tag{26}$$

Denoting the two-to-three-body reactions by

$$a(p_1) + b(p_2) \to c(-p_3) + d(-p_4) + \mathrm{A}(-p5), \qquad s_4 = (p_3 + p_4)^2, \tag{27}$$

then the differences between the partonic cross sections due to the subprocesses in Eqs. (14)–(17) are equal to

$$s^2\frac{d^2\sigma^{\mathrm{HARD}}_{gg\to gg\mathrm{A}}}{dt\,du} - s^2\frac{d^2\sigma^{\mathrm{HARD}}_{gg\to gg\mathrm{H}}}{dt\,du} = \pi G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2 \frac{N^2}{N^2-1}\left[-4\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)} - \frac{17}{3}\right], \tag{28}$$

$$s^2\frac{d^2\sigma^{\mathrm{HARD}}_{gg\to q\bar{q}\mathrm{A}}}{dt\,du} - s^2\frac{d^2\sigma^{\mathrm{HARD}}_{gg\to q\bar{q}\mathrm{H}}}{dt\,du} = \pi G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2 \frac{n_f}{N^2-1}\left[C_A\left\{\frac{2}{3}\right\} + C_F\left\{2\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)} + 2\right\}\right], \tag{29}$$

$$s^2\frac{d^2\sigma^{\mathrm{HARD}}_{q\bar{q}\to gg\mathrm{A}}}{dt\,du} - s^2\frac{d^2\sigma^{\mathrm{HARD}}_{q\bar{q}\to gg\mathrm{H}}}{dt\,du} = \pi G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2 \frac{C_AC_F}{N^2}\left[C_A\left\{\frac{2}{3}\right\} + C_F\left\{2\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)} + 2\right\}\right], \tag{30}$$

$$s^2\frac{d^2\sigma^{\mathrm{HARD}}_{q_1\bar{q}_2\to q_1\bar{q}_2\mathrm{A}}}{dt\,du} - s^2\frac{d^2\sigma^{\mathrm{HARD}}_{q_1\bar{q}_2\to q_1\bar{q}_2\mathrm{H}}}{dt\,du} = \pi G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2 \frac{C_AC_F}{N^2}\left[-2\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)} - 1\right], \tag{31}$$

$$s^2\frac{d^2\sigma^{\mathrm{HARD}}_{q_1\bar{q}_1\to q_i\bar{q}_i\mathrm{A}}}{dt\,du} - s^2\frac{d^2\sigma^{\mathrm{HARD}}_{q_1\bar{q}_1\to q_i\bar{q}_i\mathrm{H}}}{dt\,du} = \pi G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2 \frac{(n_f-1)C_AC_F}{N^2}\left[-\frac{2}{3}\right], \tag{32}$$

$$\begin{aligned} s^2\frac{d^2\sigma^{\mathrm{HARD}}_{q\bar{q}\to q\bar{q}\mathrm{A}}}{dt\,du} - s^2\frac{d^2\sigma^{\mathrm{HARD}}_{q\bar{q}\to q\bar{q}\mathrm{H}}}{dt\,du} &= \pi G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2 \frac{C_F}{N^2}\left[C_A\left\{-2\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)} - \frac{5}{3}\right\} + \frac{ss_4((s-m^2)^2+s_4^2-2tu)}{8(s_4-t)^2(s_4-u)^2}\right. \\ &\quad \left. + \frac{(s-m^2)^2+s_4^2-2tu}{4(s_4-t)(s_4-u)} + \frac{(s-m^2)^2+s_4^2-2tu+6ss_4}{4ss_4}\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)} + \frac{9}{4}\right], \end{aligned} \tag{33}$$

$$s^2\frac{d^2\sigma^{\mathrm{HARD}}_{q_1q_2\to q_1q_2\mathrm{A}}}{dt\,du} - s^2\frac{d^2\sigma^{\mathrm{HARD}}_{q_1q_2\to q_1q_2\mathrm{H}}}{dt\,du} = \pi G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2 \frac{C_AC_F}{N^2}\left[-2\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)} - 1\right], \tag{34}$$

$$\begin{aligned} s^2\frac{d^2\sigma^{\mathrm{HARD}}_{qq\to qq\mathrm{A}}}{dt\,du} - s^2\frac{d^2\sigma^{\mathrm{HARD}}_{qq\to qq\mathrm{H}}}{dt\,du} &= \pi G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2 \frac{C_F}{N^2}\left[C_A\left\{-2\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)} - 1\right\}\right. \\ &\quad \left. + \frac{s^2+s_4^2}{4(s_4-t)(s_4-u)}\ln\frac{ss_4}{tu-m^2s_4} - \frac{3}{2}\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)}\right], \end{aligned} \tag{35}$$

$$s^2\frac{d^2\,\sigma^{\mathrm{HARD}}_{qg\to qg\mathrm{A}}}{dt\,du} - s^2\frac{d^2\sigma^{\mathrm{HARD}}_{qg\to qg\mathrm{H}}}{dt\,du} = \pi G^2\left(\frac{\alpha_s(\mu^2)}{4\pi}\right)^2 \frac{1}{N}\left[C_A\left\{-2\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)} - 1\right\} + C_F\left\{\ln\frac{tu-m^2s_4}{(s_4-t)(s_4-u)} - \frac{1}{2}\right\}\right], \tag{36}$$

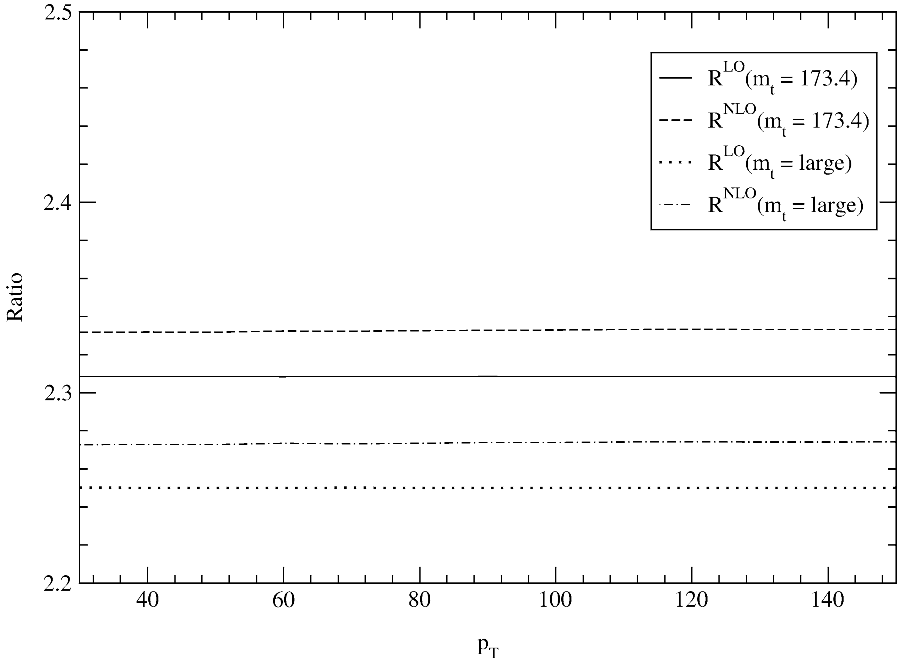

Fig. 1. The ratio R in Eq. (37) plotted as a function of p_T for $\sqrt{S} = 14$ TeV and $\mu^2 = p_T^2 + m_B^2$ with $m_H = m_A = 120\ \text{GeV}/c^2$; $R^{LO}(m_t = \infty)$ (dotted line), $R^{LO}(m_t = 173.4)$ (solid line), $R^{NLO}(m_t = \infty)$, (dot-dashed line) $R^{NLO}(m_t = 173.4)$ (dashed line).

where the meaning of the superscript HARD is explained in [17]. From these expressions we infer that the partonic cross sections (coefficient functions) for H and A are equal in LO and almost equal in NLO. This means that apart from the overall normalization due to the constant G_B there will not be any difference in the shapes of the double differential cross sections. We show this in Fig. 1 where we plot the ratio

$$R = \frac{d\sigma_A}{d\sigma_H}, \tag{37}$$

for $d\sigma_B = d\sigma_B/dp_T$ and proton–proton collisions at the LHC with $\sqrt{S} = 14$ TeV. For these and the next plots we have adopted the parton density set MRST98 (LO, lo05a.dat) [18] for the LO calculations with $\Lambda_5^{NLO} = 130.5$ MeV as input for the leading log running coupling constant. For the NLO cross sections we have chosen the set MTST99 (NLO, cor01.dat) [19] with $\Lambda_5^{NLO} = 220$ MeV as input for the next-to-leading log running coupling constant. Furthermore, the factorization/renormalization scale is chosen to be $\mu^2 = \mu_r^2 = p_T^2 + M_B^2$. For the masses of the Higgs bosons we take $m_H = m_A = 120\ \text{GeV}/c^2$ and the top quark mass is set equal to $m_t = 173.4\ \text{GeV}/c^2$. Further, we have put $\tan\beta = 1$. In the case of an infinite top quark mass (here we choose $m_t = 173.4 \times 10^3\ \text{GeV}/c^2$), we get $R^{LO} = 9/4$ irrespective of the values of m_H and m_A. This follows from Eq. (6) and the fact that the LO partonic cross sections are the same for H-production and A-production. A finite m_t as given above introduces a small effect and one gets $R^{LO} = 2.31$ which amounts to a shift upwards of 0.06 (see Fig. 1). In NLO the partonic cross sections differ a little bit and $\mathcal{C}_H^2 = [1 + 22\alpha_s/(4\pi)]\mathcal{C}_A^2$ (see Eqs. (7) and (8)). Therefore, we expect a deviation from the R^{LO} result when m_t is taken infinite in both the LO and NLO reactions. However, it turns out that both differences compensate each other. The NLO corrected partonic cross section for A is larger than the one for H and one obtains an upward shift $\Delta R^{NLO} = 0.26$. The shift due to the coefficient function in Eq. (7) is negative and amounts to $\Delta R^{NLO} = -0.24$. Hence the actual value becomes $R^{NLO} = 2.27$ (see Fig. 1) which is very close to $R^{NLO} = 9/4$. If m_t is finite one gets again an upward shift of 0.06 like in LO and $R^{NLO} = 2.33$ (see Fig. 1). One can make

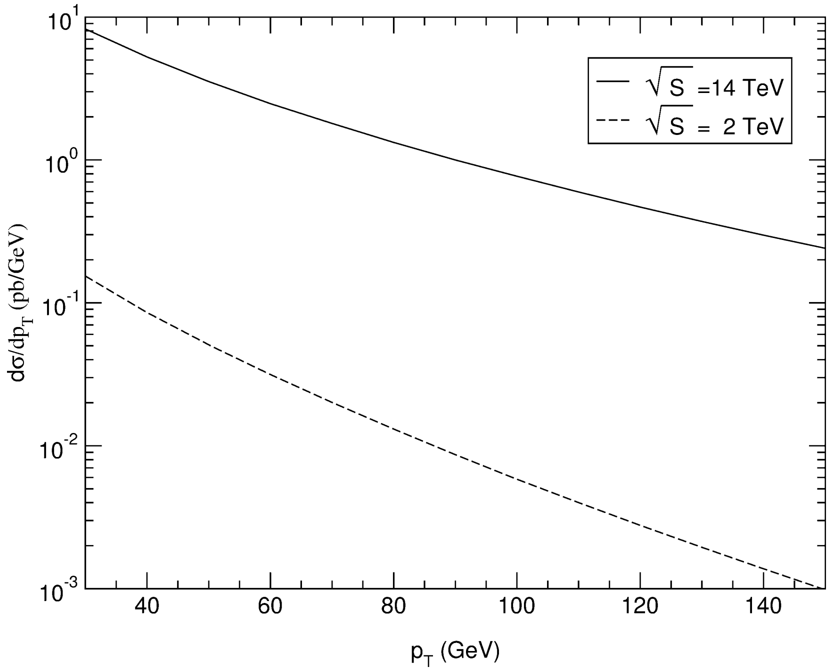

Fig. 2. The transverse momentum distribution $d\sigma_A/dp_T$ with $\mu^2 = p_T^2 + m_A^2$, $m_A = 91.9$ GeV/c^2, $\tan\beta = 0.5$; $\sqrt{S} = 14$ TeV (solid line), $\sqrt{S} = 2$ TeV (dashed line).

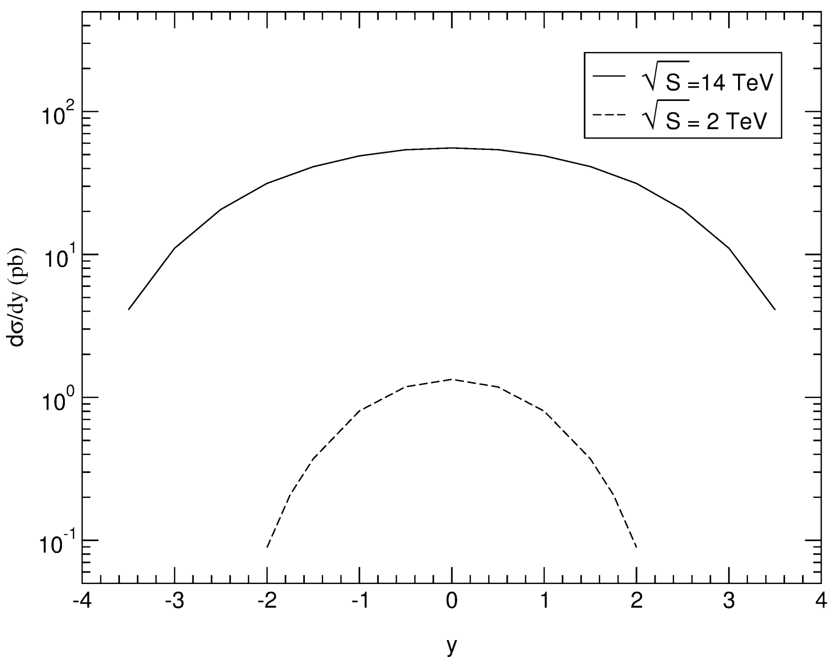

Fig. 3. The rapidity distribution $d\sigma_A/dy$ calculated from the integral of $d^2\sigma_A/(dp_T\,dy)$ between 8 $p_{T,\min} > p_T > p_{T,\min}$ and $p_{T,\min} = 30$ GeV/c. Input parameters are $\mu^2 = p_{T,\min}^2 + m_A^2$, $m_A = 91.9$ GeV/c^2, $\tan\beta = 0.5$; $\sqrt{S} = 14$ TeV (solid line), $\sqrt{S} = 2$ TeV (dashed line).

similar plots for the rapidity y distributions which yield the same ratios as shown in Fig. 1 for the p_T distributions. The most important feature is that the ratios are independent of p_T and y showing the shape independence of the distributions on the parity of the Higgs boson (scalar versus pseudo-scalar). This behaviour was discovered for both the (pseudo-)scalar p_T distributions and for the opening angle distribution between the (pseudo-)scalar bosons and the highest p_T-jet in the reaction $p + p \to (\mathrm{H\ or\ A}) + \mathrm{jet} + \mathrm{jet} +' X'$ in [20]. From Fig. 1 and the observations made above it is clear that the ratios between the NLO and LO corrected cross sections (K-factors) are the same for H-production and A-production. This also holds for the variation of the NLO cross sections with respect to the mass factorization/renormalization scales. They are given for H-production in [16,17] and we do not have to show them again for A-production. In Fig. 2 we present the p_T distributions in NLO for A-production in proton–antiproton collisions at $\sqrt{S} = 2$ TeV (Fermilab Tevatron, Run II) and in proton–proton collisions at $\sqrt{S} = 14$ TeV (LHC). Further we have chosen $m_{\mathrm{A}} = 91.9\ \mathrm{GeV}/c^2$ and $\tan\beta = 0.5$. The parton density set and the factorization scale are given above. From Fig. 2 we infer that the p_T-distributions decrease rather slowly as p_T increases and that the differential cross section for the Tevatron is two orders of magnitude smaller than the one predicted for the LHC. The latter observation also holds for the corresponding rapidity distributions shown in Fig. 3. They are obtained by integrating $d^2\sigma_{\mathrm{A}}/(dp_T\,dy)$ over the range $p_{T,\min} < p_T < 8\ p_{T,\min}$ with $p_{T,\min} = 30\ \mathrm{GeV}/c$. The cross section for $p_T > 8\ p_{T,\min}$ is negligible. Notice that the range of the rapidity for A-production at the Tevatron is rather small. Finally, we want to comment on the relative importance of the partonic subprocesses contributing to the hadronic differential cross section in Eq. (12). For the LHC ($\sqrt{S} = 14$ TeV) the gg-channel dominates and the qg-subprocess contributes about one third of the cross section. This is because at these high energies the x-values of the gluon density $f_g^{\mathrm{P}}(x)$ is so small that it becomes much larger than the quark densities. At lower energies like $\sqrt{S} = 2$ TeV (Tevatron) the x-values are larger so that the valence quark densities also play a role. This explains why the contribution of the qg-subprocess is of the same magnitude as the one from the gg-channel for A-production at the Tevatron.

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