The rapidity gap Higgs signal at LHC

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Abstract

We compare the $WW$ and pomeron-pomeron fusion mechanisms for the double-diffractive production of a Higgs boson. We determine the suppression of the ‘rapidity gap’ pomeron-pomeron fusion events due to QCD radiative effects. In particular we use leading log techniques to estimate the cross sections for both exclusive and inclusive double-diffractive Higgs production at LHC energies. The same approach can be applied to the double-diffractive central production of large $E_T$ dijets. These two processes provide one of the most justified applications of various aspects of leading logarithm QCD techniques.
The biggest challenge facing the experiments at the forthcoming very high energy proton-proton collider (LHC) is the search for possible Higgs bosons. The present estimates based on the Standard Model and its minimal supersymmetric extension favour the existence of an intermediate mass Higgs boson \(M_H \lesssim 2M_W\) \(^1\). In this case the best signals, which are based on the decay modes \(H \to b\bar{b}\) or \(\gamma\gamma\), will be extremely difficult to isolate from the background.

One novel possibility to reduce the background is to study the central production of the Higgs in events with a large rapidity gap on either side. Such rapidity gaps appear automatically if the Higgs is produced via WW boson fusion, \(pp \to WW \to H\) \(^2\); recent developments are given in \(^3\), \(^4\), \(^5\). There have also been several discussions about the possible advantage of using a similar rapidity gap signal in which the \(W\) boson is replaced by the pomeron, \(\mathcal{I}P\), see for example \(^6\), \(^7\), \(^8\), \(^9\), \(^10\), \(^11\).

The motivation of these \(\mathcal{I}P \mathcal{I}P \to H\) studies starts with the observation that at LHC energies \(gg\) fusion is the dominant mechanism with a Higgs production cross section up to a factor of 10 larger than that of WW fusion. However, although WW and \(gg\) fusion mechanisms appear to have a similar structure, in \(gg\) fusion the colour flow induces many secondaries which completely fill the original (partonic level) rapidity gaps. For this reason \(\mathcal{I}P \mathcal{I}P \to H\) mechanism has been proposed \(^7\), \(^8\) instead of \(gg \to H\). The idea is based on the hope that on the one hand \(\sigma(\mathcal{I}P \mathcal{I}P \to H)\) will be at least of the order of a few percent of \(\sigma(gg \to H)\), while on the other hand the colour flow is screened in \(\mathcal{I}P \mathcal{I}P\) fusion leading to rapidity gaps.

In order to make estimates of \(\sigma(\mathcal{I}P \mathcal{I}P \to H)\) it is necessary to invoke a model for the pomeron. One possibility is the non-perturbative approach of refs. \(^7\), \(^8\), which we call examples of the “soft” pomeron. Another possibility is to consider the perturbative QCD so-called “hard” pomeron, see for example \(^10\), \(^11\). The literature shows a wide range of predictions, which may be expressed in terms of two extreme estimates. The “soft” pomeron-like models give the upper extreme with

\[
\sigma_{\text{max}}(\mathcal{I}P \mathcal{I}P \to H) \sim \sigma(gg \to H) \frac{\sigma_{\text{el}}}{\sigma_{\text{tot}}}^2
\]

where the “suppression” factor containing the elastic and total \(pp\) cross sections is the probability of having two rapidity gaps, one either side of the Higgs. The low extreme, based on the “hard” pomeron \(^10\), \(^11\), is

\[
\sigma_{\text{min}}(\mathcal{I}P \mathcal{I}P \to H) \sim \sigma(gg \to H) (M_H^2 \sigma_{\text{tot}})^{-2}
\]

where now the “suppression” factor is the probability to have a point-like two-gluon configuration (with \(\lambda \sim 1/M_H\)) in each pomeron so that they have sufficient chance to fuse into the Higgs. These simple estimates of the suppression factor range from \(10^{-1}\) to \(10^{-12}\). Although naive, these results are in fact quite representative of the range of values that may be found in the literature.

\(^1\)For each reaction the short-distance process is \(gg \to H\), where we regard \(g\) in the case of \(\mathcal{I}P \mathcal{I}P\) fusion as a gluonic constituent of the pomeron.
Let us start from the ordinary $gg \rightarrow H$ fusion process. A relevant Feynman diagram for ‘rapidity gap’ production is shown in Fig. 1, where the additional $t$ channel gluon is needed to screen the colour. (The reason for the dashed and dotted gluon lines will be explained below.) Within this two-gluon exchange picture of the pomeron it is clear that the most optimistic scenario is first to assume that the gluon, which screens the colour, does not couple to the Higgs, and second, to assume that it has small virtuality $Q_T^2$ to enhance the probability of screening via a large value of $\alpha_s$. This idea was invoked in an attempt to describe the diffractive events in small $x$ deep inelastic scattering \[12\]. The simplest and most extreme prediction is given in ref. \[13\]. It was assumed that the ‘screening’ gluon is so soft that there is no suppression, apart from a factor of $1/N_c^2$ which is the probability of forming a colour singlet $gg$ $t$-channel state. The perturbative realisation of the soft screening approach has been studied for Higgs production \[9\] and for dijet production \[14\].

An important question, which has not yet been addressed in the literature, concerns the probability of relatively hard gluon emission coming from distance scales $\lambda > 1/M_H$ shorter than the characteristic transverse size ($\sim 1/Q_T$) of the pomeron at which the colour flow is screened. Such gluons could fill up the rapidity gaps. The goal of the present paper is to estimate the suppression of the rapidity gap events due to these effects. We will show that the typical values of $Q_T$ of the ‘screening’ gluon are indeed much smaller than $M_H$, but nevertheless are sufficiently large for perturbative QCD to be applicable.

Of course, there is also a suppression of rapidity gap events due to parton-parton rescattering and to the possibility of multiple (or ‘pile-up’) interactions at high luminosities \[3, 4, 6, 15\]. For example, a rough estimate of the former suppression is \[16\]

\[
[1 - 2(\sigma_{el} + \sigma_{SD})/\sigma_{tot}]^2 \sim 1 - 10\%
\]

depending on the value of the cross section, $\sigma_{SD}$, for single diffraction. These suppressions are common effects for any Higgs production model, including $ppP$ and $WW$ fusion, as well as for the background processes. Such effects will not be discussed further.

We calculate the rate of both exclusive and inclusive Higgs production. In the exclusive process, $pp \rightarrow ppH$, only the Higgs and the recoil protons occur in the final state. Due to the presence of the proton form factors, the Higgs is produced with small transverse momentum $q_T$. We find that the production cross section is negligibly small. On the other hand in the inclusive process, $pp \rightarrow X + \text{gap} + H + \text{gap} + X'$, the initial protons are destroyed. The phase space available for Higgs production, and hence $q_T$, are large. The cross section is found to be comparable to that for $WW$ fusion.

We will work in the double logarithmic approximation (DLA) and even for the Born amplitude we will use the leading power of all logarithms to simplify the calculations. Due to the large value of the Higgs mass $M_H$, this approach is rather well justified. Indeed double-diffractive Higgs production provides probably one of the most justified applications of various aspects of leading log techniques to date.
Exclusive production

We start with the calculation of the (double-diffractive) exclusive process which, at the quark level, is shown in Fig. 1. The Born amplitude for the process (shown by the solid lines in the figure) is of the form

\[ M(qq \rightarrow qHq) = \frac{2}{9} 2A \int \frac{d^2Q_T}{Q^2k_1^2k_2^2} 4\alpha_s^2(Q^2) (k_1,k_2), \]  

(3)

where \( \frac{2}{9} \) is the colour factor for this colour-singlet exchange process, and the factor of 2 takes into account that both of the \( t \) channel gluons can emit the Higgs boson. In the Standard Model the \( gg \rightarrow H \) vertex factor is, after convolution with the gluon polarisations, given by

\[ A^2 = \sqrt{2} G_F \alpha_s^2(M_H^2) N^2/9\pi^2 \]  

(4)

where \( G_F \) is the Fermi coupling, and where \( N \approx 1 \) provided that we are away from the threshold \( M_H = 2m_t \). Note that in the forward direction [where \( t_i = (q_i - q_i')^2 \rightarrow 0 \) for \( i = 1,2 \) and \( k_{1T} = k_{2T} = Q_T \)] the integral over the gluon loop reduces to \( \int d^2Q_T/Q^4 \). Hence, as mentioned above, small values of \( Q_T \) of the screening gluon are favoured.

In order to make the (Born) calculation more realistic we first have to include the ladder 'evolution' gluons (shown by the dashed lines in Fig. 1) and to consider the process \( pp \rightarrow pHp \) at the proton, rather than the quark, level. This is achieved by the replacements \[ f(x,Q^2) = \frac{\partial(xg(x,Q^2))}{\partial \ln Q^2} \]  

(5)

where \( x = x_1 \) or \( x_2 \) for the upper or lower ladders in Fig. 1 respectively, and where \( f(x,Q^2) \) is the unintegrated gluon density\(^2\) of the proton.

The second correction to the Born formula, \[ F_S \], is the inclusion of the Sudakov form factor, which is shown schematically by the dotted curved line in Fig. 1. \( F_S \) is the probability not to emit bremsstrahlung gluons (one of which is shown by \( p_T \)) in the interval \( Q_T < p_T < M_H/2 \). Clearly the upper bound of the interval is \( p_T \approx M_H/2 \). The lower bound, \( Q_T < p_T \), occurs because there is destructive interference of the amplitude in which the bremsstrahlung gluon is emitted from a 'hard' gluon \( k_i \) with that in which it is emitted from the soft 'screening' gluon \( Q \). That is there is no emission when \( \lambda \approx 1/p_T \) is larger than the separation, \( \Delta \rho \sim 1/Q_T \), of the two \( t \)-channel gluons in the transverse plane since then they act as a single coherent colour-singlet system. The Sudakov form factor for the above interval of \( p_T \) is

\[ F_S = \exp\left(-S(Q_T^2,M_H^2)\right) \]  

(6)

\(^2\)Strictly speaking even at zero transverse momentum, \( q_{1T} - q_{2T} = 0 \), we do not obtain the exact gluon structure function, as a non-zero component of longitudinal momentum is transferred through the two-gluon ladder. However, in the region of interest, \( x \sim 0.01 \), the value of \( |t_{\min}| = m_p^2x^2 \) is so small that we may safely put \( t = 0 \) and identify the ladder coupling to the proton with the unintegrated gluon distribution \( f(x,Q^2) \) \[ \text{[18].} \]
where \( S \) is the mean multiplicity of bremsstrahlung gluons

\[
S(Q_T^2, M_H^2) = \int_{Q_T^2}^{M_H^2/4} \frac{C_A \alpha_S(p_T^2)}{\pi} \frac{dE}{E} \frac{dp_T^2}{p_T^2} = \frac{3\alpha_S}{4\pi} \ln^2 \left( \frac{M_H^2}{4Q_T^2} \right). \tag{7}
\]

Here \( E \) and \( p_T \) are the energy and transverse momentum of an emitted gluon in the Higgs rest frame. The last equality assumes a fixed coupling \( \alpha_S \), and is shown only for illustration.

Inserting corrections (5) and (6) into the Born amplitude (3) gives

\[
M(pp \rightarrow pHp) = A \pi^3 \int \frac{dQ^2}{Q^4} e^{-S(Q_T^2, M_H^2)} f(x_1, Q_T^2) f(x_2, Q_T^2) \tag{8}
\]

in the leading log approximation. The integral has a saddle point given by

\[
\ln(M_H^2/4Q_T^2) = (2\pi/N_c \alpha_S(Q^2)) (1 - 2\gamma) \tag{9}
\]

where \( \gamma \) is the anomalous dimension of the gluon, \( g(x, Q^2) \propto (Q^2)^\gamma \). Suppose that we were to assume a constant \( \gamma = 0.15 \). Then for \( M_H = 100 \) (200) GeV the saddle point would occur at \( Q^2 = 7.3 \) (14) GeV\(^2\), well into the perturbative region, and the Sudakov suppression of the cross section would be \( (F_S)^2 = 0.04 \) (0.025). However, a more realistic evaluation using, say, the MRS(R2) set of partons [19] shows that the integrand reaches its maximum at \( Q^2 \sim 2 \) GeV\(^2\), where the suppression is \( (F_S)^2 = 0.003 \) (0.0004). If \( \gamma \) were frozen in the region \( Q^2 \leq 4 \) GeV\(^2\) then the cross section would be decreased by a further factor of 2 — a factor which is typical of the uncertainty.

Table 1 shows the values of the exclusive cross section,

\[
\frac{d\sigma}{dy}(pp \rightarrow p + H + p) = \frac{|M|^2}{16\pi^3 b^2}, \tag{10}
\]

calculated from (8). We have integrated over the \( dt_i \) assuming form factors \( \exp(-bt_i/2) \) at the proton-pomeron vertices, with \( b = 5.5 \) GeV\(^{-2}\). We find that the cross section is more than a factor of \( 10^5 \) smaller than the inclusive \( pp \rightarrow gg \rightarrow H \) cross section, without rapidity gaps; and even a factor 10 less than the \( \gamma \gamma \rightarrow H \) cross section [20]. Exclusive double-diffractive Higgs production is thus only of academic interest.

**Inclusive production**

We find that the cross section for inclusive double-diffractive Higgs production is much larger. Here the initial protons may be destroyed and the transverse momentum of the Higgs is no longer limited by the proton form factor, and so the Sudakov suppression is weaker. The process is shown in Fig. 2 in the form of the amplitude multiplied by its complex conjugate. The partonic quasielastic subprocess is \( ab \rightarrow a' + \text{gap} + H + \text{gap} + b' \). If the partons \( a, b \) are quarks then the Born amplitude for the subprocess is given by (3). However, the form factor suppressions are more complicated than for the exclusive process. As the momenta transferred,
\[ t_i = (Q - k_i)^2, \] are large we can no longer express the upper and lower ‘blocks’ in terms of the gluon structure function, but instead they are given by BFKL non-forward amplitudes.

We begin with the expression for the Born cross section for the subprocess \( gg \rightarrow g + H + g \)

\[
\frac{d\sigma}{dy} = A^2 \alpha_s^4 \frac{81}{2^8 \pi} I
\]  
(11)

with

\[
I = \frac{1}{\pi^2} \int \frac{dQ^2}{Q^2} \frac{dQ'^2}{Q'^2} \frac{d^2k_{1T}}{k_{1T}^2} \frac{d^2k_{2T}}{k_{2T}^2} (k_{1T}.k_{2T})(k'_{1T}.k'_{2T}),
\]  
(12)

where the six propagators of Fig. 2 are evident. As before the leading log contribution comes from the region where the screening gluons are comparatively soft. That is \( Q_T \ll k_{iT} \) and \( Q'_T \ll k'_{iT} \), and so

\[
t_i = (Q - k_i)^2 \simeq -k_{iT}^2 \simeq -k'_{iT}^2
\]  
(13)

for \( i = 1, 2 \). After performing the azimuthal integrations, the \( gg \rightarrow H \) vertex factors become

\[
(k_{1T}.k_{2T})(k'_{1T}.k'_{2T}) \rightarrow \frac{1}{2} k_{1T}^2 k_{2T}^2 \simeq \frac{1}{2} t_1 t_2
\]  
(14)

and we see that (12) indeed yields the maximum number (four) of logarithms.

Again we must estimate the suppression due to gluon bremsstrahlung filling up the rapidity gaps. Now the mean number of gluons emitted, with transverse momenta \( Q_T < p_T < k_{iT} \), in the rapidity interval \( \Delta \eta_i \) is

\[
n_i = \alpha_s N_c \frac{\pi}{\pi} \Delta \eta_i \ln \left( \frac{k_{iT}^2}{Q_T^2} \right).
\]  
(15)

The amplitude for no emission in the gap \( \Delta \eta_i \) is therefore \( \exp(-n_i/2) \). In this way we see that the Born integral (12) is modified to

\[
I = \frac{1}{2} \int \frac{dQ^2}{Q^2} \frac{dQ'^2}{Q'^2} \frac{dt_1}{t_1} \frac{dt_2}{t_2} \exp \left( -\frac{n_1}{2} - \frac{n'_1}{2} - \frac{n_2}{2} - \frac{n'_2}{2} - S - S' \right)
\]  
(16)

where the exponential factor represents the total form factor suppression in order to maintain the rapidity gaps \( \Delta \eta_1 \) and \( \Delta \eta_2 \) in Fig. 2. The Sudakov form factor, \( \exp(-S(k_{1T}^2, M_H^2)) \), arises from the insistence that there is no gluon emission in the interval \( k_T < p_T < \frac{1}{2}M_H \), see (3) and (4).

The justification of the non-Sudakov form factors, \( \exp(-n_i/2) \) is a little subtle. First we notice from (13) that due to the asymmetric configuration of the \( t \)-channel gluons, \( Q_T \ll k_{iT} \), we have, besides \( \Delta \eta_i \), a second logarithm, \( \ln(k_{iT}^2/Q_T^2) \), in the BFKL evolution. These double logs are resummed\(^3\) to give the BFKL non-forward amplitude \( \exp(-n_i/2)\Phi(Y_i) \), where the remaining factor \( \Phi(Y_i) \) accounts for the usual longitudinal BFKL logarithms\(^4\).

\[
Y_i \equiv (\alpha_s N_c/2\pi) \Delta \eta_i.
\]  
(17)

\(^3\)The resummation corresponds to the Reggeization of the \( t \)-channel gluons.

\(^4\)Here \( \Delta \eta_i \) (or \( Y_i \)) plays the role of \( \ln(1/x) \) in the BFKL evolution.
In the region of interest at LHC energies, $Y_i < 0.3$, it is sufficient to include only the $O(Y_i)$ term, which gives $\Phi \simeq 1 + Y_i Q^2 / k_{iT}^2 \simeq 1.1 \pm 0.1$ [21]. At our level of accuracy we may neglect the enhancement due to $\Phi$, and hence we obtain (16), which is valid in the double log approximation.

To evaluate $I$ of (14) we first perform the $Q^2$ and $Q'^2$ integrations and obtain $(Y_1 + Y_2)^{-2}$. Then we integrate over $\ln(t_1/t_2)$ which gives $\frac{1}{2}(1/Y_1 + 1/Y_2)$ where, at large $\Delta \eta_i$, we neglect the $t_i$ dependence of $S_i$. Thus (14) becomes

$$I = \frac{1}{4Y_1 Y_2(Y_1 + Y_2)} \int \frac{dt}{t} \exp(-2S(t, M_H^2)).$$

(18)

For fixed $\alpha_s$ the final ($dt$) integration gives $\pi (2N_c \alpha_s) - \frac{1}{2}$ in the DLA. However, to predict the cross section for inclusive production at the LHC we must convolute the parton-parton cross sections with the parton densities $a(x_i, t)$ of the proton, with $a = g$ or $q$, and evaluate the $dt$ integral numerically. There is a subtlety when we come to include these parton luminosity factors

$$\int_{x_{min}}^1 dx_a \ a(x_a, k_{iT}^2) \ldots,$$

with $a = g, q$. At first sight we might expect $x_{min} = x_H \equiv M_H / \sqrt{s}$ for central Higgs production. However, at large $k_{iT}$ the rapidities of the $a', b'$ jets are small in the Higgs rest frame; $\eta_{a'} = \ln(x_{a'} / \sqrt{s} / k_{iT})$. Thus in order to maintain the rapidity gaps ($\eta_{a'} > \Delta \eta_i$), we must take

$$x_{min} = x_H (1 + k_{iT} \exp(\Delta \eta_1) / M_H).$$

(19)

The results for $\sigma_{in}(\mathcal{P}\mathcal{P})$ shown in table 1 are the sum over all types of initial partons, and correspond to $\Delta \eta_1 = \Delta \eta_2 = \Delta \eta$ where $\Delta \eta$, the parton level rapidity gap, is taken to be either $\Delta \eta = 2$ or $\Delta \eta = 3$. From (18) we see that the rapidity gap cross section decreases as $1/Y^3$, that is as $(1/\Delta \eta)^3$, if $\Delta \eta_1 = \Delta \eta_2$. As expected, the suppression decreases with increasing $\alpha_s$ (like $\alpha_s^{-3.5}$ in the DLA). For comparison we give the estimates for the $WW \rightarrow H$ cross section for the same rapidity gap configuration. From table 1 we see that the $\mathcal{P}\mathcal{P}$ fusion Higgs signal is comparable to that of $WW$ fusion for $M_H \simeq 100$ GeV, but is of decreasing importance as the value of $M_H$ (or $\Delta \eta$) increases. However NLO corrections (which are not included in table 1) may increase the value of $\sigma_{in}(\mathcal{P}\mathcal{P} \rightarrow H)$ by a factor 2–4.

It is interesting to note that, due to the strong ordering of $k_T$ in the leading log approximation, almost all the momentum transfer $k_{iT}$ is balanced by the $k_{jT}$ of the parton which borders the rapidity gap. Thus, in principle, the $\mathcal{P}\mathcal{P}$ and $WW$ signals could be distinguished by the transverse momentum $k_{jT}$ of the jets which border the rapidity gap. Also note that we have

$^5$The $K$ factor enhancement (analogous to that in Drell-Yan production) is expected to be 1.6–2 [22], and there could be a factor of up to 2 from the single log BFKL enhancement term $\Phi^4$.

$^6$For $WW$ fusion one half of the cross section comes from events with $k_{iT} < M_W$, while for $\mathcal{P}\mathcal{P}$ fusion one half comes from $k_T < 13$ (25) GeV for $M_H = 100$ (300) GeV if $\Delta \eta = 2$. Indeed we could reduce $WW$ fusion in comparison with $\mathcal{P}\mathcal{P}$ fusion by about a factor $(k_{iT}^2 / M_W^2)^2$ by selecting events with both ‘border’ jets satisfying the cut-off $k_{jT}^2 < k_0^2$, with $k_0^2$ taken to be much less than $M_W^2$. For example, for $M_H = 100$ GeV, $\Delta \eta = 2$ if we take $k_0 = 20$ GeV then we find $d\sigma_{in}(\mathcal{P}\mathcal{P} \rightarrow H)/dy = 200$ fb as compared to $d\sigma_{in}(WW \rightarrow H)/dy = 4$ fb.
implicitly assumed that there is no interference between the $IPIP \rightarrow H$ and $WWW \rightarrow H$ amplitudes. This is a good approximation since the two amplitudes (i) are essentially out of phase, (ii) produce Higgs with different $q_T$ distributions, and (iii) involve different partons (namely $a = g$ for $IPIP$ fusion and $a = q$ for $WWW$ fusion.)

**Discussion**

Recall that our DLA approach to $IPIP \rightarrow H$ is only justified for the asymmetric configuration of the $t$ channel gluons, $Q_T^2 \ll k_{iT}^2$. We must check that this is in fact the case. We have seen above that typically $k_{iT} \sim 20$ GeV at LHC energies. Now, taking $\alpha_s = 0.2$, we have $Y \simeq 0.1\Delta\eta \simeq 0.25$. Thus, using (15), we find $\ln(k_{iT}^2/Q_T^2) \simeq 1/2Y \simeq 2$. So indeed $Q_T^2 < 0.15 k_{iT}^2$. Since $k_{iT}$ is rather large, the suppression due to the Sudakov form factor is not so strong for inclusive production, $(F_S)^2 \simeq 0.5$. We conclude that for relatively small $Y$, say $Y < 0.3$, the approach is self-consistent and we may use the DLA expressions, $\exp(-Y_i \ln(k_{iT}^2/Q_T^2))$, for the BFKL non-forward amplitudes, see (15) and (16). Moreover we have seen that the suppression has a clear physical interpretation.

At large $Y$, say $\Delta\eta > 5$, the situation is different. As $\Delta\eta$ increases we enter the symmetric BFKL gluon configuration, $Q_T^2 \sim k_{iT}^2$. We no longer have double logs (and moreover we lose three logs from the $Q^2, Q'^2$ and $d(t_1-t_2)$ integrations in $I$ of (14)). Instead, at large $Y$ and $t \neq 0$, we have the familiar exponential growth of the BFKL amplitude arising from the resummation of the (single) longitudinal logs. We obtain

$$\Phi(Y) \sim \exp(\lambda \Delta\eta)/(\Delta\eta)^{3/2}$$

(20)

where $\lambda$ is the BFKL intercept. Due to the $\Delta\eta$ term in the denominator, the growth only starts at $\Delta\eta \sim \frac{3}{2}\lambda^{-1} > 5$ (if we take $\lambda \sim 0.3$ from the rise of $F_2$ observed at HERA with decreasing $x$). This rapidity gap configuration is beyond the LHC energy range and is not discussed further here, although it could become important at very high energies.

Our conclusion is that the interesting proposal “that the Higgs signal could be improved by studying production in the double-diffractive configuration” does not look so optimistic as it first seemed. Exclusive production is negligibly small and even the inclusive cross section is of the same order as the cross section for the more familiar $WWW \rightarrow H$ process. The problem is that QCD radiation has a large probability to fill the parton-level rapidity gaps.

Finally we note that our approach may be used to estimate the cross section for the central production of a pair of high $E_T$ jets with a rapidity gap either side of the pair. We simply need to replace the $gg \rightarrow H$ cross section by that for $gg \rightarrow$ dijet. Since the latter cross section is much larger, and since we have an extra parameter $E_T$, such dijet production (even at Fermilab

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7 For example, in a recent study [21] of $J/\psi$ electroproduction at large $t$ (but in the symmetric gluon configuration $Q_T^2 \sim k_{iT}^2$) it was found for $Y = 0.35$ that the factor $\Phi(Y)$ enhanced the cross section by about a factor of 5; see also [9].
energies) offers an excellent opportunity to study QCD (double and single) leading log techniques. Moreover, estimates of this dijet production will be important to determine the level of the background to the $H \rightarrow b\bar{b}$ signal.

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Table 1: The cross sections $d\sigma/dy|_0$ (in fb) for the double-diffractive central production of a Higgs boson in $pp$ collisions at $\sqrt{s} = 14$ TeV via $IPIP$ or $WW$ fusion. Here $IP$ denotes the ‘hard’ QCD pomeron, and $\sigma_{ex, in}$ refer to exclusive, inclusive production respectively. The MRS(R2) set of partons [19] is used. We take running $\alpha_S$ in the evaluation of the Sudakov form factors, with $\alpha_S(M^2_Z) = 0.118$, but fixed $\alpha_S = 0.2$ in the BFKL amplitude.

| $M_H$ (GeV) | $\sigma_{ex}(IPIP)$ | $\sigma_{in}(IPIP)$ $\Delta\eta = 2$ (3) | $\sigma_{in}(WW)$ $\Delta\eta = 2$ (3) |
|-------------|---------------------|--------------------------------|--------------------------------|
| 100         | $1.8 \times 10^{-2}$ | 300 (33)                      | 220 (60)                      |
| 200         | $5 \times 10^{-3}$  | 85 (9)                        | 180 (50)                      |
| 300         | $4.4 \times 10^{-4}$| 36 (4)                        | 140 (40)                      |
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Figure Captions

Fig. 1 The Born amplitude for the exclusive double-diffractive production of a Higgs boson of transverse momentum $q_T$, shown together with the QCD radiative corrections arising from ‘evolution’ gluons (dashed lines) and the Sudakov form factor (curved dotted lines). The ‘soft screening’ gluon has four-momentum $Q$.

Fig. 2 The amplitude multiplied by its complex conjugate for the inclusive central production of a Higgs boson with rapidity gaps $\Delta \eta_1$ and $\Delta \eta_2$ on either side. The suppression due to QCD radiative effects comes from the double log resummations $\exp(-n_i/2)$ in the BFKL non-forward amplitudes and from the Sudakov form factors $\exp(-S)$ shown by the dotted curves; see eq. (16).
Fig. 1
Fig. 2