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Semiweak Production of a Top Quark accompanied by a Bottom Quark and 2 Jets at Hadron Colliders

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Abstract

We present a computation for the process of semiweak production of a top and a bottom quark accompanied by two jets. The calculation is performed using the helicity formalism generalized to include massive fermions. The obtained matrix element squared allows for an accurate simulation of single top quark production in hadronic interactions with realistic transverse momentum and angular distributions for the jets in the event.

Differential cross-sections are compared with previously obtained distributions from $O(\alpha_s \alpha^2)$ calculations.

* Supported by the 'Netherlands Organization for Scientific Research (NWO)', The Netherlands

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1. Introduction

Huge charm and bottom quark cross-sections as observed in hadron collider experiments at the Tevatron at Fermilab [1] and Sp $\bar{p}$ S at Cern [2], have launched a series of dedicated studies on the heavy flavour discovery potential of existing and future colliders [3]. At present (recent) collider energies, only the Tevatron collider remains to have a potential for resolving the missing link in the third family, the top quark. The CDF collaboration at Fermilab has derived experimental limits which put a lower bound on the mass of this quark of roughly 91 GeV/c² [4], provided extensions to the Standard Model do not generate unexpected surprises. Combination of several electroweak observables, mainly derived from e⁺e⁻ precision experiments at LEP, indicates a 'preferred' mass range for $m_t \sim 140 \pm 45$ GeV/c² [5]. Part of this mass window may still be explored by the Fermilab experiments.

Charm and bottom quarks are predominantly produced in pairs through the usual QCD mechanisms [6], whereas contributions from charged weak boson decays are relatively small but become more important with increasing quark masses. Nevertheless, the above quoted limit on the top quark mass exceeds the available phase space in the decay of an on-shell $W^\pm$ produced in Drell-Yan type of processes and consequently, the cross-section of this production channel is significantly suppressed. For quark masses above 90 GeV/c² the QCD contribution clearly dominates heavy flavour cross-sections. However, with still increasing quark mass, phase space suppression due to a *pair* of heavy quarks becomes important and semiweak production of a single top quark accompanied by a bottom quark through the t-channel exchange of a charged weak boson becomes competitive [7]. Beside phase space considerations, for increasing quark masses, the first cross-section drops like $\sim 1/\hat{s}$ while the latter behaves as $\sim 1/M_W^2$. For a top quark mass larger than ~ 250 GeV/c² the weak boson-gluon fusion mechanism then even becomes dominant.

The paper by Yuan [7] initiated detailed studies of both theoretical and experimental aspects of the W-gluon mechanism by LHC working groups. These studies were presented at the Large Hadron Collider workshop in Aachen [8]. In a realistic simulation the question was addressed whether the so-called semiweak production mechanism could really help to improve precise determination of the top quark mass. Two different methods of calculating single top cross-sections and kinematical distributions were pursued. Either heavy flavour structure functions are convoluted with the lowest order (α^2) matrix element squared or in an explicit calculation of tree level boson-gluon fusion diagrams, the obtained $O(\alpha_s \alpha^2)$ matrix element squared is convoluted with 'light' quark and gluon structure functions. Both methods give reasonable agreement [8].

In a first attempt to derive an order of magnitude estimate for k-factors, we have presented a procedure in which we properly take into account mass singularities originating from terms in the expression for the $O(\alpha_s \alpha^2)$ matrix element which are proportional to m_t/m_b . Since $m_t \gg m_b$, these factors may become large and have to be resummed. Technical details of this procedure are described in ref. [9], while comparisons between QCD and semi-weak production of heavy flavours, for a wide range of heavy flavour masses at different collider energies, may be found in [6].

In order to arrive at a description of most common event topologies, the Aachen study group used

the Pythia 5.4 [10] event generator, which is based on the $O(\alpha^2)$ matrix element squared. Multi jet final states are then obtained by convolution with a parton shower algorithm. Event signatures turn out to be quite similar to an important background namely, $W + n$ jets, where the $W^\pm$ decays leptonically and n is an arbitrary number > 2 , which makes it very difficult to disentangle signal from background. Both mechanisms show a not too different jet energy and jet multiplicity distribution [11].

The method of applying parton shower algorithms is strictly valid in the collinear limit, a region in which important contributions to the cross-section are expected [12]. However, collinear emission only, will not drastically change event topologies. Hard jets, that is energetic parton emission at large angles with respect to the direction of the parent parton, indeed will. Parton shower models tend to underestimate the importance of large angle emission, while exact calculations of tree level matrix element generally give a better description of these regions in phase space. In order to be able to verify this statement, we have derived an expression for the $O(\alpha_s^2 \alpha^2)$ matrix element squared, which allows us to study kinematical distributions and four jet event topologies. The process has been introduced in the Eurojet package for simulation of hard hadronic interactions [13].

Here we limit ourselves to the qualitative description of our calculations. A comparison with parton shower calculations and its impact on mass determination and the separation of signal and background will be the subject of a future paper. In the next section we describe our method for computing the matrix element squared using the helicity amplitude formalism. Since the matrix element is calculated at tree level only (and contains divergencies), certain corners in phase space give non-reliable results and should not be accessed. A careful study of applicability as a function of phase space parameters is therefore unavoidable. Discussion of phase space cuts, kinematical distributions and a comparison with previously obtained $O(\alpha_s \alpha^2)$ differential cross-sections are presented in section 3.

2. Matrix element computation

Nowadays, with the help of algebraic manipulation programs, it is rather straightforward to perform trace calculations on complicated strings of gamma-matrices. However, with increasing number of Feynman diagrams, it is often difficult if not impossible to obtain a simple expression for the matrix element squared. For example, the lowest order expression for $|\mathcal{M}|^2$ in boson-gluon fusion is a rather simple expression when written as combinations of scalar products of parton 4-momenta [9]. The number of terms in the expression for the next-to-lowest order matrix element squared turns out to be 'orders of magnitude' larger, which may lead to numerical instabilities and exhaustive use of computer resources.

To overcome these problems, an alternative procedure has been developed in which the complete matrix element is evaluated numerically [14 - 18], thus before squaring the full expression. The idea behind this method is to abandon use of vector products of particle momenta and express the transition matrix element in terms of spinor products (p_1 and p_2 are defined as the 4-momenta of any massless particle '1' and '2'):

$$\bar{u}(p_1) u(p_2) \quad (2.1)$$

A spinor for a massless fermion can be written in terms of basic spinors $u_{\pm}(k_0)$ which satisfy the relation

$$u_{\pm}(k_0) \bar{u}_{\pm}(k_0) = \frac{1 \pm \gamma_5}{2} \not{k}_0 \quad (2.2)$$

where k_0 is again an arbitrary 4-vector. The spinor for a massless particle with momentum p^{μ} and helicity λ can then be expressed as

$$u_{\lambda}(p) = \frac{\not{p} u_{-\lambda}(k_0)}{\sqrt{2p \cdot k_0}} \quad (2.3)$$

Since k_0 can be chosen arbitrarily, one may as well chose a specific orientation such that formulae derived later reduce to a simple form. An efficient choice is the vector $k_0 = (1, 1, 0, 0)$, using the metric defined by $g^{\mu\nu} = (1, -1, -1, -1)$. The opposite helicity spinors for a particle may be related through the simple expression

$$u_{+}(k_0) = \not{k}_1 u_{-}(k_0) \quad (2.4)$$

implying the following relations among k_0 and k_1

$$k_0 \cdot k_0 = 0, \quad k_1 \cdot k_1 = -1 \quad \text{and} \quad k_0 \cdot k_1 = 0 \quad (2.5)$$

Thus the simplest choice for k_1 will be $k_1 = (0, 0, 1, 0)$. Next, we will derive the matrix element as a combination of spinor products, defined as

$$s_{\lambda}(p_1, p_2) \equiv \bar{u}_{\lambda}(p_1) u_{\lambda}(p_2) \quad (2.6)$$

which, by using equations (2.2) - (2.5), turn out to be a remarkably simple expression in terms of components of p_1 and p_2 . For $\lambda = +$ we obtain:

$$s_{+}(p_1, p_2) = (p_1^y + ip_1^z) \sqrt{\frac{\begin{smallmatrix} 0 & x \\ p_2 & -p_2 \end{smallmatrix}}{\begin{smallmatrix} 0 & x \\ p_1 & -p_1 \end{smallmatrix}}} - (p_1 \leftrightarrow p_2) \quad (2.7)$$

while for $\lambda = -$, we get

$$s_{-}(k_1, k_2) = -s_{+}(k_1, k_2)^* \quad (2.8)$$

Both expressions are as easy to handle as the conventional scalar products.

In the computation described next, we will closely follow the procedure presented in ref. [18]:

- Express the polarization vectors of a gluon with momentum p^{μ} in terms of massless fermion spinors

$$\epsilon_{\lambda}^{\mu}(p) = \frac{\lambda}{\sqrt{2}} \frac{\bar{u}_{\lambda}(p) \gamma^{\mu} u_{\lambda}(r)}{s_{\lambda}(p, r)} \quad (2.9)$$

where r^{μ} is an arbitrary auxiliary light-like vector different from p^{μ} .

- Write the solution of the Dirac equation for massive fermions (in our case the top and bottom quarks) as

$$u(p, s) = \frac{(\not{p} + m) u_{\lambda}(r)}{\sqrt{2p \cdot r}} \quad (2.10)$$

where r^{μ} is again any light-like vector and s^{μ} a unique spin vector [16].

Let us now turn to the process we are considering and introduce the definition for parton 4-momenta.

$$\begin{aligned} \bar{q}(k_1) g(k_2) &\rightarrow \bar{q}'(k_4) g(k_7) t(k_5) \bar{b}(k_8) \\ g(k_1) g(k_2) &\rightarrow q(k_4) \bar{q}(k_7) t(k_5) \bar{b}(k_8) \\ \bar{q}(k_1) \bar{q}(k_2) &\rightarrow \bar{q}'(k_4) \bar{q}(k_7) t(k_5) \bar{b}(k_8) \end{aligned} \quad (2.11)$$

The basic Feynman diagrams contributing to these processes are depicted in figure 1.

Fig. 1: Tree level Feynman diagrams contributing to the $O(\alpha_s^2 \alpha^2)$ process $p p \rightarrow t \bar{b} + 2 \text{ jets}$; $q g$ (a), $g g$ (b) and $q q$ (c) scattering.

The full set may then easily be obtained through symmetry transformations. Since we have two mas-

sive lines and four massless, we can use the light-like momenta as auxiliary vectors. For example for the top spinor we may write (see eq. 2.10))

$$u_{\lambda t}(k_5) = \frac{(\not{k}_5 + m_t) u_{\lambda t}(k_1)}{\sqrt{2k_5 \cdot k_1}} \quad (2.12)$$

Diagrams 1, 2, 3 and 6 in fig. 1^a are quite similar to those describing the process $g g \rightarrow Z^0 b \bar{b} \rightarrow l^+ l^- b \bar{b}$. Instead, we only have to replace the leptons by two massless quarks and change the Z^0 propagator by the $W^\pm$. The latter is a quite simple change of couplings, mass and width and simplifies the expression considerably due to the pure V - A nature of the W coupling to fermions. Since the basic building blocks in which the $Z^0 b \bar{b}$ computation is cast are already implemented in the Eurojet package and are described in detail in ref. [18], we refrain from further detailed discussion here. Let us recall the main observations. The matrix element for a generic graph can be written as:

$$M_i = C_i \frac{N_i}{D_i} \quad (2.13)$$

where C_i contains the coupling constants and SU(3) colour matrices. D_i is the product of denominators of propagators for the intermediate quarks and bosons and N_i a combination of spinors, gamma matrices, polarisations and fermion projectors. When we apply this formalism to for example the first graph in fig. 1 and define following notation for internal quark lines:

$$\begin{aligned} k_8 &= k_5 + k_7 \\ k_9 &= k_2 - k_6 \end{aligned} \quad (2.14)$$

we obtain expressions for N_1 and D_1 :

$$D_1 = (k_8^2 - m_t^2) (k_9^2 - m_b^2) \left[Q^2 + M_W^2 + iM_W \Gamma_W \right] \quad (2.15)$$

$$N_1 = \bar{u}(k_4) \gamma^\mu (1 - \gamma^5) u(k_1) \bar{u}_{\lambda t}(k_5) \gamma_\alpha (\not{k}_8 + m_t) \gamma_\mu (1 - \gamma^5) (\not{k}_9 + m_b) \gamma_\beta v_{\lambda b}(k_6) \varepsilon_{\lambda 1}^\alpha(k_7) \varepsilon_{\lambda 2}^\beta(k_2)$$

M_W , Γ_W , m_b and m_t denote mass and width of the W boson, bottom and top quark masses resp.. The invariant mass of the W is defined as $Q^2 = -(k_4 - k_1)^2$. Next, in order to obtain expressions which only contain massless spinors, we replace the polarisation vectors $\varepsilon_{\lambda 1}^\alpha(k_7)$, $\varepsilon_{\lambda 2}^\beta(k_2)$ by expressions defined in eq. (2.9) and substitute eq.(2.10) for the massive spinors. We can now write eqs (2.15) in terms of the functions:

$$F(\lambda_1, p_1, q, m, \lambda_2, p_2) \equiv \bar{u}_{\lambda 1}(p_1) (\not{q} + m) u_{\lambda 2}(p_2) \quad (2.16)$$

Defining a massless vector $\hat{q} = \sigma_q \tilde{q}$, where $\tilde{q}^\mu = q^\mu - \frac{q^2}{2q \cdot p_2} p_2^\mu$ and $\sigma_q = \text{sign}(\tilde{q}^0)$ we obtain

$$\begin{aligned}
F(\lambda, p_1, q, m, -\lambda, p_2) &= m s_\lambda(p_1, p_2) \\
F(\lambda, p_1, q, m, \lambda, p_2) &= \sigma_q s_\lambda(p_1, \hat{q}) s_{-\lambda}(\hat{q}, p_2)
\end{aligned}
\tag{2.17}$$

as combinations of spinor products $s_\lambda(k_1, k_2)$ (see eq. (2.6)). The matrix element for the graphs containing a 3-gluon vertex (6) can be derived from the symmetry between graphs 1, 2 and 3. The numerators for these graphs read:

$$\begin{aligned}
N_2 &= \dots * \not{q}_2 (-\not{k}_6 - \not{k}_7 + m_b) \not{q}_7 * \dots & \equiv \dots * B_2 * \dots \\
N_3 &= \dots * \not{q}_7 (-\not{k}_6 + \not{k}_2 + m_b) \not{q}_2 * \dots & \equiv \dots * B_3 * \dots \\
N_6 &= \dots * (-\not{k}_2 - \not{k}_7) \varepsilon_2 \varepsilon_7 + (2k_7 - k_2) \varepsilon_2 \not{q}_7 + (2k_2 - k_7) \varepsilon_7 \not{q}_2 * \dots & \equiv \dots * B_6 * \dots
\end{aligned}
\tag{2.18}$$

and we get:

$$B_2 - B_6 = B_3 + \frac{1}{2} \left[(k_6 + k_7 - k_2)^2 - m_b^2 \right] (\not{q}_2 \not{q}_7 - \not{q}_7 \not{q}_2) \tag{2.19}$$

So we have succeeded in describing the non-Abelian graphs as a combination of Abelian diagrams.

Graphs 7, 8 (fig. 1^b) and 9 (fig. 1^c) are computed using the same technique. For completeness, we write the denominator and numerator for the quark-quark scattering diagrams (9) below:

$$\begin{aligned}
D_9 &= (k_2 - k_7)^2 \left[(k_6 + k_7 - k_2)^2 - m_b^2 \right] \left[Q^2 + M_W^2 + iM_W \Gamma_W \right] \\
N_9 &\sim F(-\lambda_t, k_1, k_5, m_t, +, k_4) * \left[F(+, k_1, k_8, m_t, \lambda_q, k_2) * F(\lambda_q, k_7, k_6, m_b, -\lambda_b, k_4) \right. \\
&\quad \left. + F(+, k_1, k_8, m_b, -\lambda_q, k_7) * F(-\lambda_q, k_2, k_6, m_b, -\lambda_b, k_4) \right]
\end{aligned}
\tag{2.20}$$

The systematic notation of matrix elements allows for easy coding of formulae and debugging of computer code. Furthermore, the numerical efficiency of the expected cancellations between various contributions is better than that obtained by application of usual trace techniques. Cancellations among poles arising from light quark propagators now occur at the level of the matrix elements instead of the matrix elements squared [16, 17].

Gauge invariance of the complete set of Feynman diagrams can be easily verified. A closer look at formula (2.9) for the polarization vector of a gluon shows that if we choose for the arbitrary auxiliary vector the gluon momentum itself, namely $r = k_i$

$$e_\lambda^\mu(k_i) = \frac{\lambda}{\sqrt{2}} \frac{\bar{u}_\lambda(k_i) \gamma^\mu u_\lambda(k_i)}{s_{-\lambda}(k_i, k_i)} \tag{2.21}$$

the spinor product in the denominator will vanish. The matrix element with the $s_{-\lambda}(k, k)$ removed must then vanish as well in order to satisfy gauge invariance. Indeed, we observe these cancellations among different gauge invariant subsets of diagrams. For instance setting $r_7 = k_7$ leads to the following identities (G_i denotes the contribution of graph i , T is indicating the diagrams in which t and $\bar{b}$ lines are

interchanged):

$$\begin{aligned}
 G_1 + G_2 + G_3 &= G_1^T + G_2^T + G_3^T &= 0 \\
 G_1 - G_1^T + G_3 - G_3^T + G_2^T - G_2 + 2 G_6 + 2 G_6^T &= 0 \\
 G_4 + G_5 &= G_4^T + G_5^T &= 0
 \end{aligned} \tag{2.22}$$

while the choice $r_2 = k_2$ gives

$$\begin{aligned}
 G_1 + G_2 + G_3 + G_1^T + G_2^T + G_3^T &= 0 \\
 G_1 - G_1^T + G_3 - G_3^T + G_2^T - G_2 + 2 G_6 + 2 G_6^T &= 0 \\
 G_4 + G_4^T &= G_5 + G_5^T = 0
 \end{aligned} \tag{2.23}$$

A similar check can be made for the gluon - gluon scattering contributions. The choice $r_2 = k_2$ leads to:

$$G_7 + G_8 = G_7^T + G_8^T = 0 \tag{2.24}$$

and $r_1 = k_1$ gives:

$$G_7 + G_7^T = G_8 + G_8^T = 0 \tag{2.25}$$

In principle, for each numerical evaluation of the matrix elements, one should perform the sum over all possible helicity states. It turns out that, much alike the calculation presented in ref. [18], for fixed kinematics the matrix elements do not drastically change while varying helicity combinations. It is then appropriate to limit the calculation to a randomly chosen helicity configuration per matrix element evaluation.

3. Comparison of $O(\alpha_s \alpha^2)$ and $O(\alpha_s^2 \alpha^2)$ results

In ref. [6] we have presented a comparison of lowest order boson-gluon fusion cross-sections for heavy flavour production with QCD predictions. The $O(\alpha_s^3)$ heavy flavour pair production processes represent a non-negligible part of the total heavy flavour cross-section and (depending on assumptions with regards to scale definition, cuts etc.) compete with lowest order cross-section estimations. We expect that $O(\alpha_s^2 \alpha^2)$ semiweak processes will enhance boson-gluon fusion cross-sections as well.

Although the lowest order QCD and semiweak mechanisms, after integration over *full* phase space, produce perfectly *finite* results, the next-to-lowest order tree contributions are not unambiguously determined due to intrinsic divergencies. Since complete next-to-leading order QCD calculations are at hand [19], tree level calculations can be 'tuned' that after introduction of reasonable phase space cuts, both differential distributions and total cross-section match the exact calculation. In our case however, complete next-to-leading order calculations are not available [20]. Therefore, an 'educated' guess on phase space cuts has to be made. Similar to the QCD calculation, a cut on the transverse momentum (energy) of the 'light' jets seems to be an obvious choice. Gluon emission from heavy quark lines is

then perfectly regulated. In conjunction with a cut on the spatial separation between the two outgoing light partons, this mimics a cut-off on the invariant mass of the two jets and removes remaining collinear divergences. These particular choices have the advantage that they perfectly match the ones applied in the analysis of experimental data, where most jet-finding algorithms in hadronic interactions explicitly depend on definitions such as jet cone size and phase space separation between jets.

In figure 2 we present the effect on the total cross-section of varying both cuts on transverse momentum and on spatial separation between the two light jets, as calculated at the LHC beam-beam centre-of-mass energy ($\sqrt{s} = 16$ TeV). The spatial separation between the light partons is defined as

$$\Delta R = \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2} \quad (3.1)$$

where $\Delta\eta$ is the difference in pseudorapidity and $\Delta\phi$ the angle between the partons in the plane transverse to the beams.

Fig. 2: Semiweak $O(\alpha_s^2\alpha^2)$ top quark cross-sections as a function of transverse energy cut on the light jets (for fixed $\Delta R^{\text{cut}} = 1.0$) and as a function of spatial separation (for fixed $E_t^{\text{cut}} = 20$ GeV/c) at $\sqrt{s} = 16$ TeV, for different top quark masses ($m_t = 137$ GeV/c² - full curves, $m_t = 350$ GeV/c² - dashed curves).

The cross-section dependence is plotted for two different top quark masses namely, $m_t = 137 \text{ GeV}/c^2$ (full curves) and $m_t = 350 \text{ GeV}/c^2$ (dashed curves), while we fixed $m_b = 5 \text{ GeV}/c^2$. For each choice of top mass we present two different curves; one where E_t^{cut} , the transverse energy cut on the outgoing light partons is fixed at 20 GeV, varying ΔR^{cut} and one where the spatial separation between the light partons ΔR^{cut} is fixed at 1.0, varying E_t^{cut} . The larger effects are observed for energy cut-off changes. Requiring spatial separation appears to have only little influence on the total cross-section, which seems to indicate that after applying a reasonably large momentum cut-off, event topologies containing collinear phase space configurations are less preferred.

In the following computations, we have set cut-offs on the light partons $E_t > 20 \text{ GeV}/c$ and $\Delta R > 1.0$, which roughly resemble the expectation for resolving jets at LHC/SSC experiments and keep the matrix elements well behaved. In table 1 we list the lowest order and next-to-lowest order tree semiweak cross-sections (in nb) for two top quark masses, after demanding the light parton in the $O(\alpha_s \alpha^2)$ fusion process to have $E_t > 20 \text{ GeV}/c$ as well. The matrix elements are convoluted with the Eichten et al. [21] structure function parametrisations (set 1, $\Lambda = 0.2 \text{ MeV}$). The scale at which the structure functions and strong coupling constant are evaluated is defined as

$$Q^2 = \frac{\sum_{i=1}^n [(p_t^{(i)})^2 + m_i^2]}{n} \quad (3.2)$$

where i runs over the outgoing partons ($n = 3$ and $n = 4$ resp.).

	$O(\alpha_s \alpha^2)$	$O(\alpha_s^2 \alpha^2)$
$m_t = 137 \text{ GeV}/c^2$	0.36 ± 0.01	0.14 ± 0.01
$m_t = 350 \text{ GeV}/c^2$	0.090 ± 0.004	0.045 ± 0.002

Table 1: Contributions to single top quark cross-sections (in nb) for $m_t = 137 \text{ GeV}/c^2$ and $m_t = 350 \text{ GeV}/c^2$ at $\sqrt{s} = 16 \text{ TeV}$ ($E_t^{\text{cut}} = 20 \text{ GeV}$, $\Delta R^{\text{cut}} = 1.0$).

The next-to-leading order processes increase the total cross-section by roughly 50 %.

In fig. 3 we present top (a) and 'hardest' light jet (b) transverse momentum distributions for both production mechanisms, for different top quark masses. As one would naively expect, the additional jet in the $O(\alpha_s^2 \alpha^2)$ calculation is mainly located in the hemisphere opposite to the top and enhances the transverse momentum of both top and bottom quarks (distribution for the latter is not shown). The transverse momentum plot of the hardest light jet in the event shows no significant difference between the two production mechanisms. A similar remark holds for the pseudorapidity distribution of this jet (fig. 4), which peaks in both forward and backward directions, an important feature in the selection of semiweak production mechanisms experimentally [7].

The additional 'soft' jet has a rapidly falling transverse momentum distribution, which illustrates its Bremsstrahlung origin (fig. 5). If no b-tagging is at hand, this jet may easily be confused with the b-jet from semileptonic decay of the t-quark which, at least for m_t not too large, is expected to be rather soft.

Finally, in figure 6^{a-b} we have depicted parton multiplicity distributions. We demanded each parton (including the heavy flavour ones) to have a transverse energy larger than 20 GeV. Although some of the b-quarks do not satisfy this requirement, the bulk of the $O(\alpha_s^2\alpha^2)$ events will indeed contribute to 4-jet event topologies.

Conclusions

We have presented a calculation of QCD tree level corrections to semiweak production of heavy flavours through boson-gluon fusion in hadronic interactions. The calculation is based on the helicity formalism and provides an expression for the amplitude which is easy to evaluate numerically and has a rather simple structure. Cross-sections obtained after convoluting the matrix elements squared with parton densities and after phase space integration, significantly enhance lower order results.

Our calculation provides an alternative method to parton shower approaches in the aim to resolve heavy flavour events from background events with high jet multiplicity. In particular 3-jet and 4-jet final states are expected to be accurately described.

Fig. 3: Transverse momentum distributions for semiweak $O(\alpha_s \alpha^2)$ (full curves) and $O(\alpha_s^2 \alpha^2)$ (dashed curves) production of top quarks (a) and the accompanying 'hard' jet (b) for two different top quark masses, at the LHC ($\sqrt{s} = 16 \text{ TeV}$).

Fig. 4: Pseudorapidity distribution for the 'hardest' light jet in semiweak $O(\alpha_s \alpha^2)$ (full curves) and $O(\alpha_s^2 \alpha^2)$ (dashed curves) production of top quarks for two different top quark masses, at the LHC ($\sqrt{s} = 16 \text{ TeV}$).

Fig. 5: Transverse momentum distribution of 'soft' jets in semiweak $O(\alpha_s^2 \alpha^2)$ production of top quarks for $m_t = 137 \text{ GeV}/c^2$ (full curve) and $m_t = 350 \text{ GeV}/c^2$ (dashed curve), at the LHC ($\sqrt{s} = 16 \text{ TeV}$).

Fig. 6: Jet (parton) multiplicity distributions for the sum of semiweak $O(\alpha_s \alpha^2)$ and $O(\alpha_s^2 \alpha^2)$ production of top quarks (full curves) for two different top quark masses at $\sqrt{s} = 16 \text{ TeV}$. The $O(\alpha_s^2 \alpha^2)$ contributions are indicated separately (dashed curves).

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