Higgs production in pp collisions by double-pomeron exchange

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It is estimated that, for Higgs production in high-energy pp collisions, in some 1% of the events the initial protons will emerge with only a small change of momentum and well separated in rapidity from the other final-state particles. The calculation is based on the exchange of a pair of nonperturbative gluons, which simulate double-pomeron exchange.

1. Introduction

Production of Higgs particles by conventional hard-scattering mechanisms has been well studied [1]. Recently, there has been interest in a different mechanism, double-pomeron exchange [2,3]. The event structure associated with this mechanism is distinctive, in that the initial hadrons emerge from the reaction with only a small loss of momentum and well separated in rapidity from the rest of the final-state particles [4]. So far, the calculations of the double-pomeron mechanism have been based on a guess for the structure function of the pomeron; in this note we replace this guess with a more theoretically-based input. We confirm that, at high enough energy, some 1% of Higgs-production events should originate from double-pomeron exchange.

Although pomeron exchange is known to give an excellent phenomenological description [5] of hadronic total cross sections and of elastic scattering, the theory of the pomeron is far from fully worked out. Nevertheless, we do have the beginnings of a theory, in which in some approximation pomeron exchange corresponds to the exchange of a pair of nonperturbative gluons, which takes place between a pair of quarks [6]: see fig. 1. These diagrams give an amplitude which behaves as a fixed power of the energy:

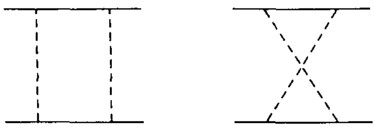

Fig. 1. Simplest contribution to pomeron exchange between quarks, from the exchange of a pair of nonperturbative gluons.

$$i\beta_0^2(t)\gamma_\lambda\cdot\gamma^\lambda \tag{1.1}$$

and so does not display Regge behaviour. The phenomenology requires rather the behaviour

$$-\exp[-\tfrac{1}{2}i\pi\alpha(t)]s^{\alpha(t)-1}\beta^2\gamma_\lambda\cdot\gamma^\lambda \tag{1.2}$$

for the elastic quark–quark scattering amplitude, with

$$\alpha(t)=1+\epsilon+\alpha' t\,, \tag{1.3}$$

where

$$\epsilon\approx 0.08,\quad \alpha'=0.25\ \mathrm{GeV}^{-2},$$
$$\beta^2\approx 4\ \mathrm{GeV}^{-2}\ (\text{constant})\,. \tag{1.4}$$

It is rather sure that this behaviour is obtained by making gluon and quark loop insertions in fig. 1, though the calculation of this is fraught with difficulties [7]. We have argued [5] that, at least at $t=0$, it seems plausible to assume that

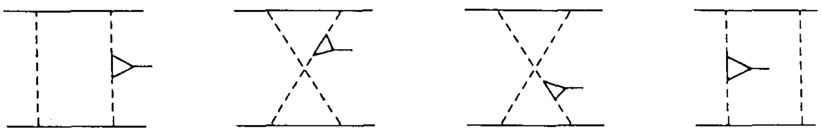

Fig. 2. Simplest contributions to Higgs production. The Higgs is coupled to one of the nonperturbative gluons through a t-quark loop.

$$\beta^2=\beta_0^2+O(\epsilon)\ , \tag{1.5}$$

so that calculating fig. 1 at $t=0$ gives a good estimate of β^2.

We shall apply the same approach to the Higgs-production amplitude. There the simplest gluon-exchange diagrams are those shown in fig. 2. As usual, we have coupled the Higgs H to two gluons through a t-quark loop. The complete two-pomeron exchange amplitude is drawn in fig. 3 and has the behavior

$$\left(\frac{s}{s_2}\right)^{\alpha_1-1}\left(\frac{s}{s_1}\right)^{\alpha_2-1} F_{\alpha_1\alpha_2}(z)\beta^2\gamma_\lambda\cdot\gamma^\lambda\ , \tag{1.6}$$

where

$$\alpha_1=\alpha(t_1),\quad \alpha_2=\alpha(t_2),\quad z=\frac{s_1 s_2}{s m_{\rm H}^2}\ . \tag{1.7}$$

The general structure of the pomeron–pomeron–Higgs coupling $F_{\alpha_1\alpha_2}$ is known [8]; we assume that it is valid to normalise it at $t_1=t_2=0$ by calculating the sum of the four diagrams in fig. 2. The form (1.6) is an asymptotic one: it neglect lower powers of s/s_1 and s/s_2. Thus one requires

$$\frac{s_1}{s},\frac{s_2}{s}<\delta\ . \tag{1.8}$$

It is usually assumed that δ must be no more than 0.1

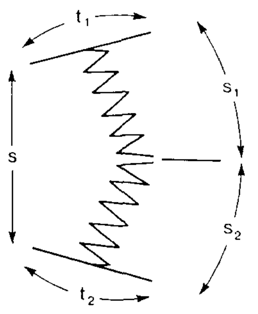

Fig. 3. The complete double-diffractive diagram for Higgs production.

for the formula (1.6) to be a good approximation.

Fig. 3 appears to describe *exclusive* Higgs production, pp→ppH, with the Higgs being produced with a transverse momentum of only a few hundred MeV/c. However, additional nonperturbative initial and final state interactions, while not changing the magnitude of the cross section, will generate extra particles [9]. Further, the coupling of the Higgs (to two gluons) which we use is calculated to lowest order in perturbative QCD. Higher-order effects will result in the radiation of gluons, which will also have the consequence that the transverse momentum of the Higgs will not be as small as fig. 3 appears to give. Compare the corresponding situation in the case of the Drell–Yan process [10]. Thus our calculation really is an inclusive one; indeed, in ref. [3] it was argued that the exclusive cross section will be negligibly small. Nevertheless, we expect that in the production we calculate, the Higgs will be accompanied by rather less radiation of other particles than in normal inclusive events, and its average transverse momentum will be smaller.

2. Calculation of gluon exchange

The sum of the two diagrams of fig. 1 is equal to half of te s-channel discontinuity of the first one. This is much simpler to calculate than the separate diagrams [6]. Likewise, as we show in Appendix 1, the sum of the four diagrams of fig. 2 is equal to the s-channel discontinuity of the first one. Without this, it is very difficult to calculate the diagrams. (In their original preprint, the authors of ref. [2] suggested calculating the diagrams of fig. 2, but concluded it was too difficult because they had failed to appreciate the discontinuity technique.)

As we have explained, we need the discontinuity for $t_1=t_2=0$. To calculate this discontinuity, we must put the lines q_1 and q_2 "on shell" (see the momentum labelling in fig. 4). The calculation is along the lines

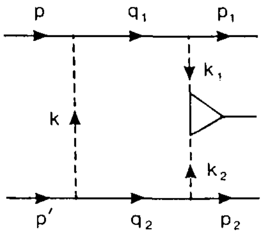

Fig. 4. The first diagram of fig. 2. Putting the quark lines q_1 and q_2 "on-shell" is equivalent to calculating the sum of al four diagrams of fig. 2.

previously used for the elastic amplitude [6]. Take the masses of the light quarks to be approximately zero and write

$$k=\frac{\bar{x}p}{s}+\frac{\bar{y}p'}{s}+v\,,$$

$$p_1=x_1 p+\frac{\bar{y}_1 p'}{s}+v_1\,,$$

$$p_2=\frac{\bar{x}_2 p}{s}+y_2 p'+v_2\,, \tag{2.1}$$

where v, v_1, v_2 are transverse to p and p' and so effectively are two-dimensional. Then

$$t_1=-v_1^2/x_1\,,\quad t_2=-v_2^2/y_2\,,$$

$$s_1\sim s(1-y_2),\quad s_2\sim s(1-x_1)\,. \tag{2.2}$$

As we explained, we are interested in calculating the sum of the diagrams of fig. 2 for $t_1=t_2=0$, so we set $v_1=v_2=0$. We then have

$$\int d^4k\,\delta(q_1^2)\delta(q_2^2)$$

$$\sim\frac{1}{2s}\int d\bar{x}\,d\bar{y}\,d^2v\,\delta(\bar{y}-v^2)\delta(-\bar{x}-v^2)\,,$$

$$\int d^4p_1\,\delta(p_1^2)\sim\tfrac{1}{2}\int dx_1\,d\bar{y}_1\,\delta(x_1\bar{y}_1)\,,$$

$$\int d^4p_2\,\delta(p_2^2)\sim\tfrac{1}{2}\int d\bar{x}_2 dy_2)\,\delta(\bar{x}_2 y_2)\,. \tag{2.3}$$

Along the upper quark line we have the Dirac matrices $\gamma^\mu\gamma\cdot q_1\gamma^\lambda$. This is approximately equivalent to

$$2q_1^\mu\gamma^\lambda\,, \tag{2.4}$$

because to calculate the differential cross section we also need $\gamma\cdot p_1$ on the left of this expression and $\gamma\cdot p$ on the right, and when these are included the difference between (2.4) and the original expression gives contributions to the cross section that are of order δ (see (1.8)) [#1]. Similarly, from the lower quark line we obtain $2q_2^\mu\gamma_\lambda$.

We write the coupling of the Higgs to two gluons as $\delta^{ab}V_{\mu\nu}$, where a and b are colour indices. For singlet exchange, the over-all colour factor in the diagram is $\frac{2}{9}$. Altogether, the s-channel discontinuity of the diagram is thus

$$M\gamma^\lambda\cdot\gamma_\lambda\,, \tag{2.5a}$$

with

$$M=\frac{ig^4}{9\pi^2 s}\int d^2v\,WD(-v^2)D(-x_1v^2)D(-y_2v^2)\,, \tag{2.5b}$$

where D is the (Feynman-gauge) nonperturbative gluon propagator and $W=q_1^\mu V_{\mu\nu}q_2^\nu$.

The Higgs coupling is taken to be through a t-quark loop. This has the property, to lowest order in perturbative QCD, that $k_1^\mu V_{\mu\nu}=0=V_{\mu\nu}k_2^\nu$. Hence the general structure of $V^{\mu\nu}$ is

$$V^{\mu\nu}=(k_1\cdot k_2 g^{\mu\nu}-k_2^\mu k_1^\nu)\frac{V_1}{m_H^2}+(k_1^2k_2^2g^{\mu\nu}$$

$$-k_2^2k_1^\mu k_1^\nu-k_1^2k_2^\mu k_2^\nu+k_1\cdot k_2k_1^\mu k_2^\nu)\frac{V_2}{m_H^4}\,, \tag{2.6}$$

V_1 and V_2 are functions of k_1^2 and k_2^2, but are effectively constant when these variables are $\ll m_H^2$, m_t^2, as is the case because in the loop integration they are kept to less than 1 GeV^2 or so because of the rapid fall-off of the propagators of the attached nonperturbative gluons [6]. V_1 is related to the decay width Γ of the Higgs to two gluons [1]:

$$\Gamma=\frac{|V_1|^2}{4\pi m_H}\,. \tag{2.7}$$

Hence [2]

$$|V_1|^2=\frac{\sqrt{2}\,G_F\alpha_s^2m_H^4}{18\pi^2}|N|^2\,, \tag{2.8}$$

where $|N|^2$ is a function of m_t/m_H which is close to 1 unless the Higgs is much heavier than the t-quark.

[#1] This was checked using the symbolic manipulation program FORM by J. Vermaseren.

V_2 does not seem to have been evaluated, but fortunately we do not need it. It is not too hard to show that it is of the same order of magnitude as V_1, but because it is dived in (2.6) by an extra power of m_{H}^2 its contribution to $W=q_1^\mu V_{\mu\nu} q_2^\nu$ is smaller by a factor of order $(1\ \mathrm{GeV}^2)/m_{\mathrm{H}}^2$.

When we calculate W, we find that there is a delicate cancellation between the contributions from the two terms in the parentheses that multiplies V_1 in (2.6). We obtain

$$W=\frac{V_1 s v^2}{2m_{\mathrm{H}}^2}. \tag{2.9}$$

Hence the amplitude at $t_1=t_2=0$ is

$$M=\frac{\mathrm{i}g^4\sqrt{G_{\mathrm{F}}}\,\alpha_{\mathrm{s}} N}{2^{1/4}54\pi^2}\int \mathrm{d}v^2 v^2 D(-v^2) \times D(-x_1 v^2) D(-y_2 v^2). \tag{2.10}$$

Because we are interested only in the leading powers of s/s_1 and s/s_2, we may not set $x_1=y_2=1$. As usual [6], we choose a form for D which makes the integration easy:

$$g^2 D(-v^2)=A\exp(-v^2/\mu^2). \tag{2.11}$$

(There is no reason to believe that the true form of D is as simple as this.) We have found previously [6,9] that $\mu\approx 1$ GeV and, from the magnitude of the total cross section (see (1.2)),

$$A^2\mu^2=72\pi\beta^2. \tag{2.12}$$

With the choice (2.11), (2.10) becomes

$$M=\frac{2^{13/4}\mathrm{i}\mu\beta^3\sqrt{G_{\mathrm{F}}}\,\alpha_{\mathrm{s}} N}{9\sqrt{\pi}\,g^2}. \tag{2.13}$$

In this expression, α_{s} is the perturbative coupling evaluated at a scale m_{H}^2 or m_{t}^2, and so is about 0.1, while g^2 is the coupling at a small scale and therefore nonperturbative – we take it to be [11] about 4π. The result (2.13) is for quark–quark scattering. For nucleon–nucleon scattering [12] at $t_1=t_2=0$ we must multiply by 3^2. So,

$$M=\mathrm{i}CN,\quad C\approx 1\times 10^{-3}\ \mathrm{GeV}^{-3}. \tag{2.14}$$

3. Reggeisation

We now turn our attention to the double-Regge formula (1.6). As we explained in section 1, we assume that it is a good approximation to set (2.14) equal to $\beta^2 F_{\alpha_1\alpha_2}$ at $t_1=t_2=0$. In order to extend to nonzero values of t_1 and t_2 we need the analysis in ref. [8] of the general structure of $F_{\alpha_1\alpha_2}$.

We modify that analysis to take account of the quark spins: we have one power less of s/s_1 and s/s_2, which is compensated by the Dirac matrices. (To calculate differential cross sections, one needs to take traces, with extra factors of the type $\gamma\cdot p$, so that the resulting power of s comes out the same as in the spinless case.) So, with some manipulation of the form given in ref. [8],

$$\begin{aligned}F_{\alpha_1\alpha_2}(z)=&f_{\alpha_1\alpha_2}(z+\mathrm{i}\epsilon)\exp(-\mathrm{i}\pi\alpha_1)\exp(-\mathrm{i}\pi\alpha_2)\\&+f_{\alpha_1\alpha_2}(z+\mathrm{i}\epsilon)\exp(-\mathrm{i}\pi\alpha_1)\\&+f_{\alpha_1\alpha_2}(z+\mathrm{i}\epsilon)\exp(-\mathrm{i}\pi\alpha_2)+f_{\alpha_1\alpha_2}(z-\mathrm{i}\epsilon),\end{aligned} \tag{3.1}$$

where $f_{\alpha_1\alpha_2}$ has the general structure

$$\begin{aligned}f_{\alpha_1\alpha_2}(z)=&\int \mathrm{d}u\,\phi(u)(-\mathrm{i}zu)^{\alpha_2-1}\\&\times\Gamma(1-\alpha_1)\Gamma(\alpha_1-\alpha_2)\\&\times{}_1F_1(-\alpha_1,\alpha_2-\alpha_1+1;-\mathrm{i}zu)\\&+(\alpha_1\leftrightarrow\alpha_2).\end{aligned} \tag{3.2}$$

The function $\phi(u)$ is unknown; it can also depend on t_1 and t_2.

Because the transverse momentum of the Higgs is much less than its mass, we have $z\approx 1$. Because of the requirement (1.8) that s/s_1 and s/s_2 are large, the presence of the Regge powers in the amplitude (see (1.6)) ensures that the main contribution comes from small t_1 and t_2, where α_1 and α_2 are close to 1. In spite of appearances, $F_{\alpha_1\alpha_2}(z)$ is finite when $\alpha_1=\alpha_2=1$: there is a cancellation between the divergences in the two terms in (3.2) and the various terms in (3.1). So it seems reasonable to approximate it by a constant – anyway, this is the best we can do. To fix this constant, we use the result (2.14) of the calculation of the simple diagrams of figs. 1, 2. To obtain the amplitude for nucleon–nucleon collisions at nonzero momentum transfer, we also [12] have to include an elastic form factor F_1 for each nucleon. So the amplitude is taken to be

$$T=\mathrm{i}CN\left(\frac{s}{s_2}\right)^{\alpha_1-1}\left(\frac{s}{s_1}\right)^{\alpha_2-1}F_1(t_1)F_1(t_2)\gamma_\lambda\cdot\gamma^\lambda . \quad (3.3)$$

4. Cross section

With (3.3), the cross section is

$$\sigma=\frac{2C^2|N|^2 s}{(2\pi)^5}\int \mathrm{d}^4p_1\,\mathrm{d}^4p_2\;\delta(p_1^2)\delta(p_2^2)$$
$$\times\delta((p+p'-p_1-p_2)^2-m_\mathrm{H}^2)$$
$$\times\left(\frac{s}{s_2}\right)^{2\alpha_1-2}\left(\frac{s}{s_1}\right)^{2\alpha_2-2}[F_1(t_1)F_1(t_2)]^2 . \quad (4.1)$$

We use the first three equations of section 2 and the expression (1.3) for the pomeron trajectory to write this as

$$\sigma=\frac{C^2|N|^2}{2(2\pi)^5}\left(\frac{s}{m_\mathrm{H}^2}\right)^{2\epsilon}$$
$$\times\int\frac{\mathrm{d}x_1}{x_1}\frac{\mathrm{d}y_2}{y_2}\,\delta\left((1-x_1)(1-y_2)-\frac{m_\mathrm{H}^2}{s}\right)$$
$$\times\int \mathrm{d}^2v_1\,\mathrm{d}^2v_2(1-x_1)^{2\alpha' v_1^2}(1-y_2)^{2\alpha' v_2^2}$$
$$\times\exp[-2\lambda(v_1^2+v_2^2)] , \quad (4.2)$$

where we have introduced an exponential approximation for the nucleon form factor at small t:

$$F_1(t)\approx\exp(\lambda t) \quad (4.3)$$

with

$$\lambda\approx 2\ \mathrm{GeV}^{-2} . \quad (4.4)$$

We recall that $1-x_1$ and $1-y_2$ are the fractional longitudinal momentum losses of the two incident nucleons. Their difference is the fractional longitudinal momentum of the Higgs. They must be small for the analysis to be valid. The constraint (1.9) is

$$x_1, y_2>1-\delta . \quad (4.5)$$

If we integrate (4.2) over this range we find approximately

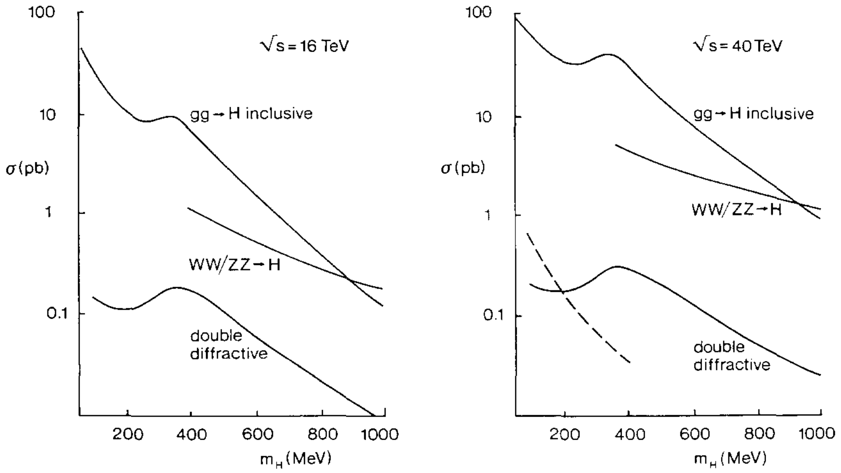

Fig. 5. Calculations of double-diffractive Higgs production at the LHC and the SSC. The two upper curves in each case correspond to the conventional inclusive mechanisms and are from ref. [14]. The broken curve is the calculation of double-diffractive Higgs production from ref. [2].

$$\sigma=\frac{C^2|N|^2}{128\alpha'\pi^3}\left(\frac{s}{m_{\rm H}^2}\right)^{2\epsilon}\frac{1}{\alpha'\log(s/m_{\rm H}^2)+2\lambda}$$

$$\times\log\left(\frac{\alpha'\log(\delta s/m_{\rm H}^2)+\lambda}{\alpha'\log(1/\delta)+\lambda}\right), \tag{4.6}$$

provided $s>m_{\rm H}^2/\delta^2$. Using the explicit form for N in ref. [13] we find at LHS and SSC energies cross sections shown in fig. 5. We have taken $\delta=0.1$ and assumed that the t-quark mass is not too far from 150 GeV. Also shown in the figure are the cross sections of ref. [14] calculated from the inclusive gg→H, WW→H and ZZ→H mechanism. At Tevatron energy, $\sqrt{s}=1.8$ TeV,

$$m_{\rm H}=50\ \text{GeV},\quad \sigma=5\times10^{-38}\ \text{cm}^2,$$

$$m_{\rm H}=100\ \text{GeV},\quad \sigma=2\times10^{-38}\ \text{cm}^2,$$

$$m_{\rm H}=150\ \text{GeV},\quad \sigma=7\times10^{-39}\ \text{cm}^2,$$

$$m_{\rm H}=175\ \text{GeV},\quad \sigma=1\times10^{-40}\ \text{cm}^2. \tag{4.7}$$

At this energy, Higgs particles heavier than this cannot be produced by double pomeron exchange, because of the constraint (1.9).

The broken line in fig. 5 is the calculation of double-diffractive Higgs production of ref. [2]. These authors use a pomeron structure function which is "soft", that is they assume that there is only very small probability that any gluon carries a substantial fraction of the pomeron's momentum. Our calculation corresponds to a much harder pomeron structure function, which is why our cross section falls off much more slowly with increasing Higgs mass. This issue will be pursued further in a subsequent paper.

Our cross sections are calculated for the case where both initial hadrons are scattered quasi-elastically. In practice, it will be difficult to establish that they have not broken up, with their fragments going down the beam pipes. To a first approximation, one can calculate this possibility [15] by omitting the elastic form factors in (3.3), that is set $\lambda=0$ in the result (4.6). This gives cross sections an order of magnitude greater.

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Appendix

We show that the sum of the four diagrams in fig. 2 is equal to the discontinuity of the first one is the s-channel at fixed s_1, s_2, t_1 and t_2.

Comparing the general form (3.3) with the result of the calculation in section 2, we see that we need to take the limit $\alpha_1=\alpha_2=1$ in (3.1). We first write (3.1) as the sum of two terms:

$$F_{\alpha_1\alpha_2}(z)=D_{\alpha_1\alpha_2}(z)+R_{\alpha_1\alpha_2}(z), \tag{A1}$$

where

$$D_{\alpha_1\alpha_2}(z)=f_{\alpha_1\alpha_2}(z-{\rm i}\epsilon)-f_{\alpha_1\alpha_2}(z+{\rm i}\epsilon). \tag{A2}$$

and

$$R_{\alpha_1\alpha_2}(z)=f_{\alpha_1\alpha_2}(z+{\rm i}\epsilon)$$

$$\times[\exp(-{\rm i}\pi\alpha_1)+1][\exp(-{\rm i}\pi\alpha_2)+1]. \tag{A3}$$

Because $z=s_1s_2/m_{\rm H}^2s$, $D_{\alpha_1\alpha_2}(z)$ is just the s-channel discontinuity of $F_{\alpha_1\alpha_2}(z)$ at fixed s_1 and s_2. When $\alpha_1=\alpha_2=1$, replacing $F_{\alpha_1\alpha_2}(z)$ with $D_{\alpha_1\alpha_2}(z)$ in (1.6) gives the s-channel discontinuity of the amplitude T. So, to demonstrate the result we need, we must show that $R_{\alpha_1\alpha_2}$ vanishes when $\alpha_1=\alpha_2=1$.

According to (3.2), $f_{\alpha_1\alpha_2}(z)$ consists of two terms. When we calculate its z-discontinuity, its first term gets multiplied by

$$1-\exp[-2\pi{\rm i}(\alpha_2-1)]. \tag{A4}$$

This factor vanishes when $\alpha_1=\alpha_2=1$, but there are the two Γ-functions which diverge. The properties of $\phi(u)$ must be such as to lead to the finite constant result (2.14) when the two terms in (3.2) are added together. That is, one of the Γ-function divergences in (3.2) cancels between the two terms, leaving the other to give a finite result when it is multiplied by (A4). In $R_{\alpha_1\alpha_2}$, however, we have the factor

$$[\exp(-{\rm i}\pi\alpha_1)+1][\exp(-{\rm i}\pi\alpha_2)+1], \tag{A5}$$

which is doubly zero when $\alpha_1=\alpha_2=1$. When this is multiplied by the single Γ-function divergence in (3.2), the result is 0.

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