Appendix

A.1 Unitarity

In this section, we will follow [1, 2] closely in the derivation of the partial wave unitarity bounds. Related to this is the helicity formalism [3, 4], which well suit to description of collider processes because helicity $h = \mathbf{S} \cdot \frac{\mathbf{p}}{|\mathbf{p}|}$ is invariant under rotation around and boost along $\mathbf{p}$. The helicity state is a quantum state where the spin component along the momentum is specified. For massive states, one has the rest frame $\mathring{p} := (m, 0, 0, 0)$ where:

$$J_z|\mathring{p}, \lambda\rangle = \lambda|\mathring{p}, \lambda\rangle. \tag{A.1}$$

To describe particles [5] not in the rest frame, one has to first boost along the z-axis to the particle's momentum $|\mathbf{p}|$ via the operator $Z_{|\mathbf{p}|}$, otherwise the spin will not be aligned with the momentum. The unitary operator describing rotation through the Euler angles α, β, γ is given by:

$$U(R(\alpha, \beta, \gamma)) := e^{-i\alpha\hat{J}_z} e^{-i\beta\hat{J}_y} e^{-i\gamma\hat{J}_z}. \tag{A.2}$$

One can then use $R(\phi, \theta, -\phi)$ to bring $\mathbf{e}_z$ to $\hat{\mathbf{p}} = (0, \sin\theta\cos\phi, \sin\theta\sin\phi, \cos\theta)$. The single particle state is then:

$$|p, \lambda\rangle = U(R(\phi, \theta, -\phi))U(Z_{|\mathbf{p}|})|\mathring{p}, \lambda\rangle. \tag{A.3}$$

Such construction gives an general momentum state but leaves the helicity eigenvalue preserved due to $U(Z_{|\mathbf{p}|})J_{\hat{p}} = J_{\hat{p}}U(Z_{|\mathbf{p}|})$ and $U(R(\phi, \theta, -\phi))J_z = J_{\hat{p}}U(R(\phi, \theta, -\phi))$.

The scattering matrix may be factorised into scattering part by removing the identity via $S = \mathbb{1} + iT$ so that:

$$\langle f|S|i\rangle = \langle f|i\rangle + i\langle f|T|i\rangle. \tag{A.4}$$

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Momentum conservation can be factored out from this so that the second term can be identified as:

$$\langle f|T|i\rangle = (2\pi)^4\delta^{(4)}(p_f - p_i)\mathcal{M}(i \to f). \tag{A.5}$$

The unitarity of the S-matrix then demands that:

$$(\mathbb{1} - iT^\dagger)(\mathbb{1} + iT) = \mathbb{1} \iff i(T - T^\dagger) = T^\dagger T.$$

Left operating this by $\langle f|$ and right operating by $|i\rangle$ gives:

$$i\langle f|(T - T^\dagger)|i\rangle = \sum_X \int d\Pi_X \langle f|T|X\rangle|\langle X|T|i\rangle, \tag{A.6}$$

where on the right hand side makes use of the completeness relation:

$$\mathbb{1} = \sum_X \int \Pi_X |X\rangle\langle X| := \sum_n \left(\prod_{i=1}^n \int \frac{d^4 p_i}{(2\pi)^4}\delta^{(4)}(p_i^2 - m_i^2)\theta(p_i^0)|\mathbf{p}_1 \dots \mathbf{p}_n\rangle\langle \mathbf{p}_1 \dots \mathbf{p}_n|\right). \tag{A.7}$$

The result is the *generalised optical theorem* which is described by the following relation:

$$i\left(\mathcal{M}_{i\to f} - \mathcal{M}^*_{f\to i}\right) = -\sum_X \mathcal{M}^*_{X\to f}\mathcal{M}_{i\to X}. \tag{A.8}$$

We now turn to two particle states because we are concerned with $2 \to 2$ scattering where working in centre of momentum frame:

$$\begin{aligned} |i\rangle &= |\mathbf{p}, \lambda_1; -\mathbf{p}, \lambda_2\rangle, \\ |f\rangle &= |\mathbf{p}', \lambda_1'; -\mathbf{p}', \lambda_2'\rangle. \end{aligned} \tag{A.9}$$

The two particle initial(final) state may be considered as a single particle state with helicity $\lambda_1^{(\prime)} - \lambda_2^{(\prime)}$. One can show that:

$$|J, M; \lambda_1, \lambda_2\rangle = \sqrt{\frac{2J+1}{4\pi}} \int d\Omega D^{(J)}_{\lambda_1-\lambda_2,M}\left(R^{-1}(\phi, \theta, -\phi)\right)|\mathbf{p}, \lambda_1; -\mathbf{p}, \lambda_2\rangle, \tag{A.10}$$

where with the Wigner D- and d-matrices defined with respect to the spherical harmonics functions:

$$\begin{aligned} D^{(j)}_{mm'}(R) &:= \langle jm|U(R(\alpha, \beta, \gamma))|jm'\rangle \\ &= e^{i\alpha m}e^{i\gamma m'}d^j_{mm'}(\beta). \end{aligned} \tag{A.11}$$

In elastic scattering where the integration over the intermediate state in (A.6) is:

$$\sum_{\lambda_1,\lambda_2}\int \frac{d^3p_1}{(2\pi)^3 2E_1}\frac{d^3p_2}{(2\pi)^3 2E_2}\delta^{(4)}(p_1+p_2-p_f) = \frac{1}{32\pi^2 s}\lambda^{1/2}(s,m_1^2,m_2^2)\int d\Omega_k, \tag{A.12}$$

where $\lambda(z_1,z_2,z_3) := (z_1^2+z_2^2+z_3^2)^2 - 2z_1z_2 - 2x_2x_3 - 2z_3z_1$. Using the Wigner-Eckart theorem, one can define:

$$\begin{aligned}\langle f|T|i\rangle &= \langle \mathbf{p}_1', \lambda_1'; \mathbf{p}_2', \lambda_2'|T|\mathbf{p}_1, \lambda_1; \mathbf{p}_2, \lambda_2\rangle \\ &= 16\pi\sum_J (2J+1)\langle\lambda_1', \lambda_2'|T^J(s)|\lambda_1,\lambda_2\rangle D^{J*}_{\lambda_1-\lambda_2,\lambda_1'-\lambda_2'}(\phi,\theta,-\phi),\end{aligned} \tag{A.13}$$

so that the LHS of (A.6) becomes:

$$16\pi\sum_J(2J+1)\left[\langle\lambda_1',\lambda_2'|T^J(s)|\lambda_1,\lambda_2\rangle - \langle\lambda_1',\lambda_2'|T^{J^\dagger}(s)|\lambda_1,\lambda_2\rangle\right]D^{J*}_{\lambda_1-\lambda_2,\lambda_1'-\lambda_2'}(R_{fi}), \tag{A.14}$$

and the RHS by:

$$i\frac{32\pi\lambda^{1/2}(s,m_1^2,m_2)}{s}\sum_J(2J+1)\left|\langle\lambda_1',\lambda_2'|T^J(s)|\lambda_1,\lambda_2\rangle\right|^2 D^{J*}_{\lambda_1-\lambda_2,\lambda_1'-\lambda_2'}(R_{fi}) \tag{A.15}$$

The J-th component can then be extracted using the orthogonality relation:

$$\int d\Omega D^{J*}_{m_1,am_2}(R))D^{J'}_{m_1',m_2'}(R) = \frac{4\pi}{2J+1}\delta_{J,J'}\delta_{m_1,m_1'}\delta_{m_2,m_2'}. \tag{A.16}$$

One arrives at amplitude relation:

$$\begin{aligned}&\langle\lambda_1',\lambda_2'|T^J(s)|\lambda_1,\lambda_2\rangle - \langle\lambda_1',\lambda_2'|T^{\dagger J}(s)|\lambda_1,\lambda_2\rangle \\ &= \frac{2i\lambda^{1/2}(s,m_1^2,m_2)}{s}\left|\langle\lambda_1',\lambda_2'|T^J(s)|\lambda_1,\lambda_2\rangle\right|^2\end{aligned} \tag{A.17}$$

which, in the high energy limit reduces to (suppressing the helicity indices):

$$\text{Im}\, T^J = |T^J|^2 = (\text{Re}\, T^J)^2 + (\text{Im}\, T^J)^2. \tag{A.18}$$

This is equivalent to a bound on the real part of T^J, consistent with [6]:

$$(\text{Re}\, T^J)^2 = \text{Im}\, T^J(1-\text{Im}\, T^J) \leq \frac{1}{4}. \tag{A.19}$$

The longitudinal modes of vector boson scatter behaves like a scalar which is spin-less $D^J_{00}(0, \theta, 0) = P_J(\cos\theta)$. Specifically, (A.13) becomes the Jacob-Wick expansion [7, 8]:

$$\langle f|T|i\rangle = 16\pi \sum_J (2J+1) a_J(s) P_J(\cos\theta). \tag{A.20}$$

In the Born approximation where a_0 is real, this becomes:

$$|\text{Re}\, a_0| \leq \frac{1}{2}. \tag{A.21}$$

A. 2 $hZ \to hZ$ Scattering

In this section we give matrix elements for $hZ \to hZ$ scattering which is presented in Sect. 2.3. The Mandelstam variables are defined by:

$$\begin{aligned} s &= (p_1 + k_1)^2 \approx 4p^2, \\ t &= (p_1 - k_2)^2 \approx 2p^2(\cos\theta - 1). \end{aligned} \tag{A.22}$$

A.2.1. s-channel

The matrix amplitude of the left panel of Fig. 2.1 may be evaluated as:

$$\begin{aligned} i\mathcal{M}_s = {} & i\frac{\kappa_Z}{2}\epsilon^{s,\mu\nu}(p_1)\epsilon^{\alpha}(k_1)\left[\left(m_Z^2 + k_1\cdot(p_1+k_1)\right)C_{\mu\nu,\alpha\beta} + D_{\mu\nu,\alpha\beta}(k_1, p_1+k_1)\right] \\ & \times\left(i\frac{-\eta^{\beta\gamma} + (p_1+k_1)^{\beta}(p_2+k_2)^{\gamma}/m_Z^2}{(p_1+k_1)^2 - m_Z^2}\right) \\ & \times\left[\left(m_Z^2 + k_2\cdot(p_2+k_2)\right)C_{\rho\sigma,\gamma\delta} + D_{\rho\sigma,\gamma\delta}(p_2+k_2, k_2)\right] i\frac{\kappa_Z}{2}\epsilon^{s',\rho\sigma}(p_2)\epsilon^{\delta}(k_2). \end{aligned} \tag{A.23}$$

where the diagonal term in the propagator is neglected because we are only concerned with the only the leading contributions.

A.2.2 t-channel

The corresponding matrix amplitude of the right panel of Fig. 2.1 is:

$$\begin{aligned}
i\mathcal{M}_t = {} & i\frac{\kappa}{2}\epsilon^{s,\rho\sigma}(p_2)\epsilon^{\alpha}(k_1)\left[\left(m_Z^2 + k_1\cdot(p_2-k_1)\right)C_{\rho\sigma,\alpha\beta} + D_{\rho\sigma,\alpha\beta}(k_1, p_2-k_1)\right] \\
& \times\left(i\frac{-\eta^{\beta\gamma} + (p_2-k_1)^{\beta}(p_1-k_2)^{\gamma}/m_Z^2}{t-m_Z^2}\right) \\
& \times\left[\left(m_Z^2 + k_2\cdot(p_1-k_2)\right)C_{\mu\nu,\gamma\delta} + D_{\mu\nu,\delta\gamma}(k_2, p_1-k_2)\right] i\frac{\kappa}{2}\epsilon^{s',\mu\nu}(p_1)\epsilon^{\delta}(k_2).
\end{aligned}$$

A.3 Decay Form Factors

The form factors in Sect. 3.2 are from [9] and are collected here:

$$\begin{aligned}
S^{\gamma}(m_h) &= 2\sum_{f=b,t,c,\tau} N_c Q_f^2 \kappa_f^S F_s(\tau_{h,f}) - \kappa_W F_1(\tau_{h,W}), \\
P^{\gamma}(m_h) &= 2\sum_{f=b,t,c,\tau} N_c Q_f^2 \kappa_f^P F_p(\tau_{h,f}), \\
S^{g}(m_h) &= \sum_{f=b,t,c} \kappa_f^S F_s(\tau_{h,f}), \\
P^{g} m_h) &= \sum_{f=b,t,c} \kappa_f^P F_p(\tau_{h,f}), \\
S^{Z\gamma}(m_h) &= 2\sum_{f=b,t,\tau} Q_f N_C^f m_f^2 \frac{I_3^f - 2s_W^2 Q_f}{s_W c_W}\kappa_f^S F_f^{(0)} + m_Z^2 \cot\theta_W \kappa_W F_W, \\
P^{Z\gamma}(m_h) &= -2\sum_{f=b,t,\tau} Q_f N_C^f m_f^2 \frac{I_3^f - 2s_W^2 Q_f}{s_W c_W}\kappa_f^P F_f^{(5)},
\end{aligned} \tag{A.24}$$

with $\tau_{h,i} := m_h^2/4m_i^2$, $\tau_{Z,i} := m_Z^2/4m_i^2$. The colour factor is $N_C = 3$ for quarks and $N_C = 1$ for τ's. The scaling functions are defined as follows:

$$\begin{aligned}
F_s(\tau) &= \tau^{-1}[1 + (1-\tau^{-1})f(\tau)], \\
F_p(\tau) &= \tau^{-1}f(\tau), \\
F_1(\tau) &= 2 + 3\tau^{-1} + 3\tau^{-1}(2-\tau^{-1})f(\tau),
\end{aligned} \tag{A.25}$$

where $f(\tau)$ is in terms defined in (A.28) and (A.29). We note that if Higgs boson mass is greater than twice the mass of the charged particle running in the loop, i.e. $m_h > 2m_x$ ($\tau > 1$) and imaginary parts arises in the form factors. We also have the isospin $I_3^u = 1/2$, $I_3^{d,\ell} = -1/2$. The remaining factors in the above expressions are defined in [10]:

$$\begin{aligned} F_f^{(0)} &= C_0(m_f^2) + 4C_2(m_f^2), \qquad F_f^{(5)} = C_0(m_f^2), \\ F_W &= 2\left[\frac{m_h^2}{m_W^2}(1-2c_W^2) + 2(1-6c_W^2)\right]C_2(m_W^2) + 4(1-4c_W^2)C_0(m_W^2), \end{aligned} \tag{A.26}$$

where the factors:

$$\begin{aligned} C_0(m_i^2) &= -\frac{1}{2m^2(\tau_{h,i}-\tau_{Z,i})}[f(\tau_{Z,i}) - f(\tau_{h,i})], \\ C_2(m_i^2) &= \frac{1}{8m^2(\tau_{h,i}-\tau_{Z,i})^2}\Big([f(\tau_{Z,i}) - f(\tau_{h,i})] + 2\tau_{Z,i}[g(\tau_{Z,i}) - g(\tau_{h,i})] + (\tau_{h,i}-\tau_{Z,i})\Big), \end{aligned} \tag{A.27}$$

are in turns defined using:

$$f(\tau) = -\frac{1}{2}\int_0^1 \frac{dy}{y}\ln[1-4\tau y(1-y)] = \begin{cases} [\sin^{-1}(\sqrt{\tau})]^2, & \tau \le 1, \\ -\frac{1}{4}\left[\ln\left(\frac{\sqrt{\tau}+\sqrt{\tau-1}}{\sqrt{\tau}-\sqrt{\tau-1}}\right) - i\pi\right]^2, & \tau \ge 1, \end{cases} \tag{A.28}$$

$$g(\tau) = \begin{cases} \sqrt{\tau^{-1}-1}\,\sin^{-1}(\sqrt{\tau}), & \tau \le 1, \\ \frac{1}{2}\sqrt{1-\tau^{-1}}\left[\ln\left(\frac{\sqrt{\tau}+\sqrt{\tau-1}}{\sqrt{\tau}-\sqrt{\tau-1}}\right) - i\pi\right], & \tau > 1. \end{cases} \tag{A.29}$$

A.4 Spinor-Helicity Formalism

The helicity operator in the **p** direction is given by:

$$J := \boldsymbol{\sigma}\cdot\mathbf{p} = \begin{pmatrix} \hat{p}^3 & \hat{p}^1 - i\hat{p}^2 \\ \hat{p}^1 + i\hat{p}^2 & -\hat{p}^3 \end{pmatrix} = \begin{pmatrix} \cos\theta & e^{-i\phi}\sin\theta \\ e^{i\phi}\sin\theta & -\cos\theta \end{pmatrix}. \tag{A.30}$$

The massless solutions to the Dirac equation:

$$\left(i\gamma_\mu\partial^\mu - m\right)\psi = 0, \tag{A.31}$$

have definite helicity, where the positive and negative states are respectively:

$$u_\pm(k) = P_\pm u(k), \quad v_\mp(k) = P_\pm v(k). \tag{A.32}$$

Similar relations can be obtained for their Dirac conjugates. The helicity eigenstates are then two component spinors satisfying:

$$J\chi_\lambda(k) = \lambda\chi_\lambda(k). \tag{A.33}$$

With the spin quantisation axis chosen to point in the $\hat{z}$ direction, one obtains:

$$\begin{aligned}
\chi_+(\mathbf{p}) &= \frac{1}{\sqrt{2|\mathbf{p}|\left(|\mathbf{p}| + p^3\right)}}\begin{pmatrix}|\mathbf{p}| + p^3 \\ p_+\end{pmatrix} = \begin{pmatrix}\cos\frac{\theta}{2} \\ e^{i\phi}\sin\frac{\theta}{2}\end{pmatrix}, \\
\chi_-(\mathbf{p}) &= \frac{1}{\sqrt{2|\mathbf{p}|\left(|\mathbf{p}| + p^3\right)}}\begin{pmatrix}-p_- \\ |\mathbf{p}| + p^3\end{pmatrix} = \begin{pmatrix}-e^{-i\phi}\sin\frac{\theta}{2} \\ \cos\frac{\theta}{2}\end{pmatrix}, \\
\chi_\lambda(-\mathbf{p}) &= -\lambda e^{i\lambda\phi}\chi_{-\lambda}(\mathbf{p}).
\end{aligned} \tag{A.34}$$

The four component Dirac spinor is then:

$$u_\lambda(p) = \begin{pmatrix}\sqrt{E - \lambda|\mathbf{p}|}\chi_\lambda(\mathbf{p}) \\ \sqrt{E + \lambda|\mathbf{p}|}\chi_\lambda(\mathbf{p})\end{pmatrix}, \qquad v_\lambda(p) = \begin{pmatrix}-\lambda\sqrt{E + \lambda|\mathbf{p}|}\chi_{-\lambda}(\mathbf{p}) \\ \lambda\sqrt{E - \lambda|\mathbf{p}|}\chi_{-\lambda}(\mathbf{p}),\end{pmatrix} \tag{A.35}$$

Whilst in the Weyl basis, this is:

$$\lambda_\alpha = \begin{pmatrix}\sqrt{p^+} \\ \sqrt{p^-}e^{i\phi}\end{pmatrix}, \qquad \tilde{\lambda}^{\dot{\alpha}} = \begin{pmatrix}-\sqrt{p^-}e^{-i\phi} \\ \sqrt{p^+}\end{pmatrix}, \tag{A.36}$$

$$u_+(p) = \begin{pmatrix}\sqrt{p^+} \\ \sqrt{p^-}e^{i\phi} \\ 0 \\ 0\end{pmatrix}, \qquad u_-(p) = \begin{pmatrix}0 \\ 0 \\ \sqrt{p^-}e^{-i\phi} \\ -\sqrt{p^+}\end{pmatrix}, \tag{A.37}$$

$$e^{\pm i\phi} = \frac{p^1 \pm ip^2}{\sqrt{p^+p^-}}, \qquad p^\pm = p^0 \pm p^3. \tag{A.38}$$

In the limit $p^+ = p^1 = p^2 = 0$, the spinors become:

$$u_+(p) = \begin{pmatrix}0 \\ \sqrt{2p^0} \\ 0 \\ 0\end{pmatrix}, \qquad u_-(p) = \begin{pmatrix}0 \\ 0 \\ \sqrt{2p^0} \\ 0\end{pmatrix}, \tag{A.39}$$

Table A.1 Spinor helicity formalism for massless momenta

Notation	Weyl Spinor	Positive energy Dirac spinor	Negative energy Dirac spinor
$i\rangle$	$\lambda_\alpha(k_i)$	$u_+(k_i)$	$v_-(k_i)$
$i]$	$\tilde{\lambda}^{\dot{\alpha}}(k_i)$	$u_-(k_i)$	$v_+(k_i)$
$\langle i$	$\lambda^\alpha(k_i)$	$\overline{u}_-(k_i)$	$\overline{v}_+(k_i)$
$[i$	$\tilde{\lambda}_{\dot{\alpha}}(k_i)$	$\overline{u}_+(k_i)$	$\overline{v}_-(k_i)$

The Dirac conjugate can then be evaluated as follows:

$$\overline{u}_+(p) = \left(0, 0, \sqrt{p^+}, \sqrt{p^-}e^{-i\phi}\right), \qquad \overline{u}_-(p) = \left(\sqrt{p^-}e^{i\phi}, -\sqrt{p^+}, 0, 0\right), \quad \text{(A.40)}$$

These definitions are summarised in Table A.1.
One may subsequently form binary products:

$$\langle ij\rangle = \sqrt{|s_{ij}|}e^{i\phi_{ij}}, \qquad [ij] = -\sqrt{|s_{ij}|}e^{-i\phi_{ij}}, \quad \text{(A.41)}$$

where:

$$s_{ij} = (k_i + k_j)^2,$$
$$\cos\phi_{ij} = \frac{p_i^1 p_j^+ - p_j^1 p_i^+}{\sqrt{|s_{ij}|}\sqrt{p_i^+ p_j^+}}, \qquad \sin\phi_{ij} = \frac{p_i^2 p_j^+ - p_j^2 p_i^+}{\sqrt{|s_{ij}|}\sqrt{p_i^+ p_j^+}}, \quad \text{(A.42)}$$

To facilitate the squaring the amplitude with real momenta, the relation $\langle ij\rangle = [ji]^*$ can be employed. The following expression for a massless momenta k^μ is often useful in the calculation of the amplitudes:

$$k^\mu = \frac{1}{2}\langle k\gamma^\mu k]. \quad \text{(A.43)}$$

We collect here some properties of the spinors:

1.
$$\langle ij\rangle[ji] = 2(k_i \cdot k_j). \quad \text{(A.44)}$$

2.
$$\frac{\langle i\gamma^\mu j]}{2}\frac{\langle l\gamma_\mu m]}{2} = \frac{1}{2}\langle il\rangle[mj]. \quad \text{(A.45)}$$

3. *Shouten identity*

$$\langle ij\rangle\langle lm\rangle + \langle il\rangle\langle mj\rangle + \langle im\rangle\langle jl\rangle = 0, \tag{A.46}$$

and similarly with $\langle \dots \rangle$ replaced by $[\dots]$.

4.

$$\langle ij\rangle[ji] = \mathrm{Tr}\left(\frac{1}{2}(1-\gamma^5)\not{k}_i\not{k}_j\right) = 2k_i \cdot k_j = s_{ij}. \tag{A.47}$$

5.

$$\langle ij\rangle[ji]\langle lm\rangle[mi] = \mathrm{Tr}\left(\frac{1}{2}(1-\gamma^5).\not{k}_i\not{k}_j\not{k}_l\not{k}_m\right) = \frac{1}{2}\left(s_{ij}s_{lm} - s_{il}s_{jm} + s_{im}s_{jl} - 4i\epsilon_{iljm}\right). \tag{A.48}$$

A.4.1 Massive Case

One defines a spin vector in the rest frame $s^\mu = (0, \mathbf{s})$ which is boosted to:

$$s^\mu = \left(\frac{1}{m}\mathbf{p}\cdot\mathbf{s}, \mathbf{s} + \frac{\mathbf{p}(\mathbf{p}\cdot\mathbf{s})}{m(E+m)}\right), \tag{A.49}$$

so as to split p^μ into two light-like momenta:

$$p_\pm := \frac{1}{2}\left(p \pm ms\right). \tag{A.50}$$

The massive spinors are then constructed [11]:

$$\begin{aligned}
u_+(p) &= |p_+\rangle + |p_-]\frac{\langle p_- p_+\rangle}{m},\\
u_-(p) &= |p_+] + |p_-\rangle\frac{[p_- p_+]}{m},\\
\overline{u}_+(p) &= [p_+| + \frac{[p_+ p_-]}{m}\langle p_-|,\\
\overline{u}_-(p) &= \langle p_+| + \frac{\langle p_+ p_-\rangle}{m}[p_-|,\\
v_+(p) &= |p_+] - |p_-\rangle\frac{[p_- p_+]}{m},\\
v_-(p) &= |p_+\rangle - |p_-]\frac{\langle p_- p_+\rangle}{m},\\
\overline{v}_+(p) &= \langle p_+| - \frac{\langle p_+ p_-\rangle}{m}[p_-|,\\
\overline{v}_-(p) &= [p_+| - \frac{[p_+ p_-]}{m}\langle p_-|.
\end{aligned} \tag{A.51}$$

There have been attempts to find a basis in order to maximise the polarised single top production cross section [12, 13] by picking a particular reference momenta to define the spin projection. This has been done in [14, 15] where the products of two spinors is defined as:

$$[ij] = \frac{(p_i \cdot k_0)(p_j \cdot k_1) - (p_j \cdot k_0)(p_i \cdot k_1) - i\epsilon\mu\nu\rho\sigma k_0^\mu p_i^\nu p_j^\rho k_1^\sigma}{\sqrt{(p_i \cdot k_0)(p_j \cdot k_0)}}, \tag{A.52}$$

where:

$$k_0 \cdot k_0 = 0, \qquad k_1 \cdot k_1 = -1, \qquad k_0 \cdot k_1 = 0. \tag{A.53}$$

A.5 Higgs Production Associated with Single Top

The amplitude radiating off the W-propagator (cf. Fig. 3.8) follows simply:

$$\begin{aligned}
-i\mathcal{M} =& \left(\overline{u}(k_4) V_{tb} \frac{-ig_2}{\sqrt{2}} \gamma^\mu P_- u(k_2)\right) \left(\frac{-i(g_{\mu\mu'} - k_{24}^\mu k_{24}^{\mu'}/m_W^2)}{k_{24}^2 - m_W^2 + im_W\Gamma_W}\right) \\
&\times \frac{i}{2} v g_2^2 g^{\mu'\nu'} \left(\frac{-i(g_{\nu\nu'} - k_{13}^\nu k_{13}^{\nu'}/m_W^2)}{k_{13}^2 - m_W^2 + im_W\Gamma_W}\right) \left(\overline{u}(k_3) V_{ud} \frac{-ig_2}{\sqrt{2}} \gamma^\nu P_- u(k_1)\right) \\
=& \frac{ig_2^4 v V_{tb} V_{ud}}{4} \left(\overline{u}(k_4)\gamma^\mu P_- u(k_2)\right) \left(\frac{(g_{\mu\nu} - k_{24}^\mu k_{24}^\nu/m_W^2)}{k_{24}^2 - m_W^2 + im_W\Gamma_W}\right) \\
&\times \left(\frac{1}{k_{13}^2 - m_W^2 + im_W\Gamma_W}\right) \langle 1\gamma^\nu 3],
\end{aligned} \tag{A.54}$$

where treating all quarks but the t-quark as massless gives $k_{13}^\mu \langle 3\gamma_\mu 1] = 0$, so that the above is:

$$\begin{aligned}
\mathcal{M}(ub \to t^{(\uparrow)} dh) =& \frac{ig_2^4 v V_{tb} V_{ud}}{4} \frac{1}{k_{13}^2 - m_W^2 + im_W\Gamma_W} \frac{1}{k_{24}^2 - m_W^2 + im_W\Gamma_W} \\
&\times \frac{[t_+ t_-]}{m_t} \langle t_- \gamma_\mu 2](g^{\mu\nu} - k_4^\mu k_{42}^\nu/m_W^2)\langle 1\gamma_\nu 3] \\
=& \frac{ig_2^4 v V_{tb} V_{ud}}{4} \frac{1}{k_{13}^2 - m_W^2 + im_W\Gamma_W} \frac{1}{k_{24}^2 - m_W^2 + im_W\Gamma_W} \\
&\times \frac{[t_+ t_-]}{m_t} \left[\langle t_- 1\rangle[32] - \frac{\langle t_- t_+\rangle[t_+ 2]}{m_W^2} (\langle 14\rangle[43] - \langle 12\rangle[23])\right].
\end{aligned} \tag{A.55}$$

Similarly for the other polarisation:

$$\begin{aligned}\mathcal{M}(ub \to t^{(\downarrow)}dh) =& \frac{i g_2^4 v V_{tb} V_{ud}}{4} \frac{1}{k_{13}^2 - m_W^2 + i m_W \Gamma_W} \frac{1}{k_{24}^2 - m_W^2 + i m_W \Gamma_W} \\ & \times \langle t_+ \gamma_\mu 2](g^{\mu\nu} - k_4^\mu k_{42}^\nu / m_W^2)\langle 1 \gamma_\nu 3] \\ =& \frac{i g_2^4 v V_{tb} V_{ud}}{4} \frac{1}{k_{13}^2 - m_W^2 + i m_W \Gamma_W} \frac{1}{k_{24}^2 - m_W^2 + i m_W \Gamma_W} \\ & \times \left[\langle t_+ 1\rangle[32] - \frac{\langle t_+ t_-\rangle[t_- 2]}{m_W^2} (\langle 14\rangle[43] - \langle 12\rangle[23]) \right].\end{aligned} \tag{A.56}$$

For the Higgs radiating off the top quark, one has an additional Yukawa vertex and t-propagator:

$$\begin{aligned}-i\mathcal{M} =& \overline{u}(k_4) \frac{-i y_t}{\sqrt{2}} \left(e^{i\xi} P_+ + e^{-i\xi} P_-\right) \frac{\gamma_\rho k_{45}^\rho - m_t}{k_{45}^2 - m_t^2 - i m_t \Gamma_t} V_{tb} \frac{-i g_2}{\sqrt{2}} \gamma_\mu P_- u(k_2) \\ & \times \left(\frac{-i(g^{\mu\nu} - k_{13}^\mu k_{13}^\nu / m_W^2)}{k_{13}^2 - m_W^2 + i m_W \Gamma_W} \right) \left(\overline{u}(k_3) V_{ud} \frac{-i g_2}{\sqrt{2}} \gamma_\nu P_- u(k_1) \right) \\ =& \frac{g_2^2 y_t V_{tb} V_{ud}}{2\sqrt{2}} \overline{u}(k_4) \frac{e^{-i\xi} \gamma_\rho k_{45}^\rho - e^{i\xi} m_t}{k_{45}^2 - m_t^2 + i m_t \Gamma_t} \gamma_\mu P_- u(k_2) \\ & \times \left(\frac{1}{k_{13}^2 - m_W^2 + i m_W \Gamma_W} \right) \langle 3 \gamma^\mu 1],\end{aligned} \tag{A.57}$$

where in the second line we have use the fact that $\gamma^\mu P\pm = P_\mp \gamma^\mu$. To make further progress, the completeness relation has to be used:

$$i\rangle[i + i]\langle i = \gamma_\mu k_i^\mu, \tag{A.58}$$

together with momentum conservation $k_{45} = k_1 + k_2 - k_3$ to give[1]:

[1]We have explicitly used the relation $[i\gamma^\mu j] = \langle i\gamma^\mu j\rangle = 0$.

$$\begin{aligned}\mathcal{M}(ub \to t^{(\uparrow)}dh) =& i\frac{g_2^2 y_t V_{tb} V_{ud}}{2\sqrt{2}}\left([t_+ + \frac{[t_+t_-]}{m_t}\langle t_-\right)\frac{e^{-i\xi}(1]\langle 1+2]\langle 2-3]\langle 3) - e^{i\xi}m_t}{k_{45}^2 - m_t^2 + im_t\Gamma_t}\gamma^\mu 2] \\ &\times\left(\frac{1}{k_{13}^2 - m_W^2 + im_W\Gamma_W}\right)\langle 3\gamma_\mu 1] \\ =& i\frac{g_2^2 y_t V_{tb} V_{ud}}{2\sqrt{2}}\frac{1}{k_{45}^2 - m_t^2 + im_t\Gamma_t}\frac{1}{k_{13}^2 - m_W^2 + im_W\Gamma_W} \\ &\times\left[e^{-i\xi}\left([t_+1]\langle 1\gamma^\mu 2] + [t_+2]\langle 2\gamma^\mu 2]\right) - e^{i\xi}[t_+t_-]\langle t_-\gamma^\mu 2]\right]\langle 3\gamma_\mu 1] \\ =& i\frac{g_2^2 y_t V_{tb} V_{ud}}{\sqrt{2}}\frac{1}{k_{45}^2 - m_t^2 + im_t\Gamma_t}\frac{1}{k_{13}^2 - m_W^2 + im_W\Gamma_W} \\ &\times[12]\left[e^{-i\xi}\left([t_+1]\langle 13\rangle + [t_+2]\langle 23\rangle\right) - [t_+t_-]e^{i\xi}\langle t_-3\rangle\right].\end{aligned} \tag{A.59}$$

Again we have:

$$\begin{aligned}\mathcal{M}(ub \to t^{(\downarrow)}dh) =& i\frac{g_2^2 y_t V_{tb} V_{ud}}{2\sqrt{2}}\left(\langle t_+ + \frac{\langle t_+t_-\rangle}{m_t}[t_-\right)\frac{e^{-i\xi}(1]\langle 1+2]\langle 2-3]\langle 3) - e^{i\xi}m_t}{k_{45}^2 - m_t^2 + im_t\Gamma_t}\gamma^\mu 2] \\ &\times\left(\frac{1}{k_{13}^2 - m_W^2 + im_W\Gamma_W}\right)\langle 3\gamma_\mu 1] \\ =& i\frac{g_2^2 y_t V_{tb} V_{ud}}{2\sqrt{2}}\frac{1}{k_{45}^2 - m_t^2 + im_t\Gamma_t}\frac{1}{k_{13}^2 - m_W^2 + im_W\Gamma_W} \\ &\times\left[e^{-i\xi}\frac{\langle t_+t_-\rangle}{m_t}\left([t_-1]\langle 1\gamma^\mu 2] + [t_-2]\langle 2\gamma^\mu 2]\right) - e^{i\xi}m_t\langle t_-\gamma^\mu 2]\right]\langle 3\gamma_\mu 1] \\ =& i\frac{g_2^2 y_t V_{tb} V_{ud}}{\sqrt{2}}\frac{1}{k_{45}^2 - m_t^2 + im_t\Gamma_t}\frac{1}{k_{13}^2 - m_W^2 + im_W\Gamma_W} \\ &\times[12]\left[e^{-i\xi}\frac{\langle t_+t_-\rangle}{m_t}\left([t_-1]\langle 13\rangle + [t_-2]\langle 23\rangle\right) - m_t e^{i\xi}\langle t_+3\rangle\right].\end{aligned} \tag{A.60}$$

We refer a reader who is interested in evaluating these amplitudes to [16] in order to parameterise the phase space for three-body decay.

A.6 Sphaleron Energy

When $\theta_W \to 0$, the sphaleron has an $O(3)$ symmetry and so one may take an *Ansatz* of the form:

$$\frac{i}{2}g_2\sigma^a W_i^a dx^i = f_W(\zeta)dU^\infty(U^\infty)^{-1}, \tag{A.61}$$

$$\phi = \frac{v}{\sqrt{2}}f_h(\zeta)U^\infty\begin{pmatrix}0\\1\end{pmatrix}, \tag{A.62}$$

where ζ is the dimensionless radial distance $\zeta := g_2|v|r$ and:

$$U^{\infty} := \frac{1}{r}\begin{pmatrix} z & x+iy \\ -x+iy & z \end{pmatrix}. \tag{A.63}$$

The functions f_W, f_h are to minimise the energy functional given by:

$$\begin{aligned} E_{sph} =& \frac{4\pi|v|}{g_2}\int_0^{\infty} d\zeta \left[4\left(\frac{df_W}{d\zeta}\right)^2 + \frac{8}{\zeta^2}[f_W(1-f_W)]^2 + \frac{1}{2}\zeta^2\left(\frac{df_h}{d\zeta}\right)^2\right. \\ &+ [f_h(1-f_W)]^2 \\ &\left.+ \frac{\zeta^2}{g_2^2 v^4}\left(\frac{\lambda}{4}(vf_h)^4 + \frac{\kappa}{3}(vf_h)^3 - \frac{\mu^2}{2}(vf_h)^2 + c\right)\right]. \end{aligned} \tag{A.64}$$

The resulting equations of motion is a system of coupled second order ordinary differential equations:

$$\begin{aligned} \zeta^2\frac{d^2 f_W}{d\zeta^2} &= 2f_W(1-f_W)(1-2f_W) - \frac{\zeta^2}{4}f_h^2(1-f_W), \\ \frac{d}{d\zeta}\left(\zeta^2\frac{df_h}{d\zeta}\right) &= 2f_h(1-f_W)^2 + \frac{\zeta^2}{g_2^2}\left(\lambda f_h^3 + \frac{\kappa}{v}f_h^2 - \frac{\mu^2}{v^2}f_h\right). \end{aligned} \tag{A.65}$$

This was solved numerically with the relaxation method as prescribed in appendix of [17], by compactifying the radial distance via $r := \frac{\xi}{\xi+1}$. The corresponding boundary conditions:

$$\lim_{\zeta\to 0}(f_W, f_h) = 0, \tag{A.66}$$

$$\lim_{\zeta\to\infty}(f_W, f_h) = 1, \tag{A.67}$$

can then be incorporated and guarantee the finiteness of the sphaleron energy. The sphaleron configuration are subsequently shown in Fig. A.1.

A.7 $\mathcal{CP}$-violating Source

In this section we outline the derivation of the source based on [18–22]. The top quark has a four momentum $(E, \mathbf{k})$ and a complex mass as indicated by $M(z) = m_t(z)e^{-i\theta_t(z)}$ (4.63) where z is the rest-frame coordinate transverse to the bubble wall. The t-quark emerges from the thermal ensemble and propagates for a mean-free time of τ_t between the scattering points z_0 and $z_0 + \tau v$, where $v = |\mathbf{k}_\perp|/E$ is the velocity perpendicular to the wall. We consider here the top-axial current whose divergence is non-zero due to the complex mass, and so this divergence will be taken as source

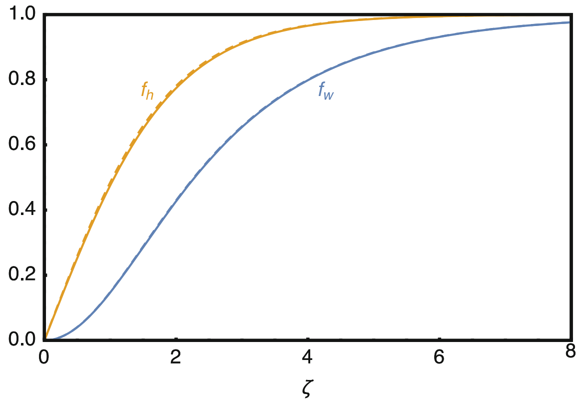

Fig. A.1 The sphaleron functions f_W and f_h as solutions to (A.65) for $\kappa = -0.5m_h^2/|v|$. The respective SM solutions $\kappa = 0$ are shown by the *dashed lines*

term. The currents J_+ will be defined as the contribution of particles moving from z_0 to τv, whereas J_- for those corresponding to the particles moving in the reverse direction. It then follows that:

$$\begin{aligned} J_+ &:= \int \frac{d^3k}{(2\pi)^3} \Big\langle \left[n(E,v) - n(E,-\tilde{v})\right] Q\,(z_0,\mathbf{k},\tau) \Big\rangle_z (1,0,0,\tilde{v}), \\ J_- &:= \int \frac{d^3k}{(2\pi)^3} \Big\langle \left[n(E,v) - n(E,-\tilde{v})\right] Q\,(z_0,\mathbf{k},\tau) \Big\rangle_z (1,0,0,-v), \end{aligned} \tag{A.68}$$

where $\tilde{v}^2 = v^2 + \left[M^2(z_0) - M^2(z_0 + v\tau)\right]/E^2$ is the squared velocity at $z_0 + v\tau$. The boosted Fermi-Dirac distribution is given by

$$n\,(E,v) = \frac{1}{e^{\beta\gamma_w E(1-vv_w)} + 1}, \tag{A.69}$$

and $\langle\cdot\rangle_z$ corresponds to averaging the coordinate z over $z \in [z_0 - \tau v/2, z_0 + \tau v/2]$. The chiral charge in this specific case is:

$$Q\,(z_0,\mathbf{k},\tau) = |T_L|^2 - |T_R|^2 - \left|T_{\overline{L}}\right|^2 + \left|T_{\overline{R}}\right|^2, \tag{A.70}$$

where $T_{L(R)}$ are the transmission amplitude or a left(right)-handed spinor over the distance $v\tau$ and $T_{\overline{L}(\overline{R})}$ are the anti-particle analogue. One solves the Dirac equation with a spatial dependent mass:

$$\begin{aligned} (-i\partial_z - P_L)\,t_L(z) &= -i\delta(z - z_0)t_L(z) + M(z)t_R, \\ (-i\partial_z - P_L)\,t_R(z) &= -M(z)^\dagger t_L. \end{aligned} \tag{A.71}$$

Due to hermiticity, one has $T_{\overline{L}} = T_L(M \to M^*)$. It also follows that $T_R = T_L(M \to -M^\dagger)$.

The transmission amplitude T_L is given by the sum of interference of paths of successive scattering and rescattering so it reads (up to a phase):

$$T_L(z_0, \tau) = 1 - \int_{z_0}^{z_0+v\tau} dz_1 \int_{z_0}^{z_1} dz_2 e^{2i|\mathbf{k}_\perp|(z_1-z_2)} M(z_2) M^\dagger(z_1) + \mathcal{O}\left(\frac{M^4}{|\mathbf{k}_\perp|^4}\right), \tag{A.72}$$

where the first term corresponds the path where the left-handed quark is straightforwardly transmitted, and the second term, where it is scattered backward into a right-handed quark at $z = z_1$ and forward again at $z = z_2$ into a left-handed quark.

One uses $\mathcal{CPT}$ symmetry to identify the amplitude of a particle transmitted from the left with that of its $\mathcal{CP}$-conjugate from the left.Similiarly, reflection amplitudes are related to transmission amplitudes through unitary condition $|T|^2 + |R|^2 = 1$. The result is:

$$\begin{aligned} Q(z_0, \mathbf{k}, \tau) =& 8 \int_{z_0}^{z_0+v\tau} dz_1 \int_{z_0}^{z_1} dz_2 \sin\left[2|\mathbf{k}_\perp|(z_1 - z_2)\right] \operatorname{Im}\left(M(z_2) M^\dagger(z_1)\right) \\ &+ \mathcal{O}\left(\frac{M^4}{|\mathbf{k}_\perp|^4}\right). \end{aligned} \tag{A.73}$$

The derivative expansion $M(z) = M(z_0) + (z_i - z_0)\partial_z M(z_0)$ can then be used in the thick wall limit $L_w \gg \tau$ to give:

$$\begin{aligned} Q(z_0, \mathbf{k}, \tau) &= 4 \frac{1}{|\mathbf{k}_\perp|^3} f(|\mathbf{k}_\perp| v\tau)\, m_t^2(z_0)\, \partial_z \theta_t(z_0), \\ f(\xi) &:= \sin\xi \left(\sin\xi - \xi \cos\xi\right). \end{aligned} \tag{A.74}$$

The source is computed from the current which deposits $(J_+ + J_-)^\mu$ per τ interval, over the typical relaxation time τ_R:

$$\begin{aligned} S(z, t) &= \partial_\mu \int_t^{t+t_R} \frac{1}{\tau} (J_+ + J_-)^\mu \\ &= -\frac{1}{\tau} \int \frac{d^3k}{(2\pi)^3} \Big\langle \left[n(E, v) - n(E, -\tilde{v})\right] Q\left(z_0, \mathbf{k}, \tau\right) \Big\rangle_z, \end{aligned} \tag{A.75}$$

since $(J_+ + J_-)^\mu$ is zero up to $\mathcal{O}(v_w)$ for $\mu = 3$ and may be neglected for $\mu = 0$ when the relaxation time is large. Using the fact that:

$$\begin{aligned} |\mathbf{k}_\perp|^2 &= v^2 E^2, \\ |\mathbf{k}_\parallel|^2 &= (1 - v^2)E^2 - m_t^2 \geq 0 \end{aligned} \quad \Longrightarrow \quad \begin{aligned} \int \frac{d^3k}{(2\pi)^3} &= \frac{1}{2\pi^2} \int_0^\infty |\mathbf{k}_\parallel| d|\mathbf{k}_\parallel| \int_0^\infty d|\mathbf{k}_\perp|, \\ &= \frac{1}{2\pi^2} \int_0^1 dv \int_{\gamma m_t}^\infty E^2 dE. \end{aligned} \tag{A.76}$$

and $\frac{dn}{dv}(E,0) = E\frac{dn}{dE}(E,0)$, one gets (up to first order in v_w and τ/L_w):

$$\begin{aligned} S(z,t) &= \frac{v_w}{2\pi^2}\int_0^1 dv \int_{z_0-\tau v/2}^{z_0+\tau v/2}\frac{dz'}{\tau v}\int_{\gamma m_t(z)}^{\infty} dE E^3 \frac{dn}{dE}(2v)\left.\frac{Q(z',\mathbf{k},\tau)}{\tau}\right|_{z'=\gamma_w(z-v_w t)} \\ &= T\gamma_w v_w\left[m^2(z_0)\partial_z\theta_t(z_0)\right]_{z_0=\gamma_w(z-v_w t)}\frac{2}{\pi^2}\mathcal{F}(\tau,m,T), \end{aligned} \tag{A.77}$$

where the form factor $\mathcal{F}(\tau,m,T)$ may be obtained by first reversing the E and v integrals and then defining:

$$y := \beta E, \qquad t := v^2 E\tau, \tag{A.78}$$

giving the result:

$$\mathcal{F}(\tau,m,T) = \frac{1}{\sqrt{\tau T}}\int_{m_t/T}^{\infty} y^{1/2}dy\int_0^{T\tau(y-\frac{m_t^2}{T^2}\frac{1}{y})}\frac{dt}{t^{3/2}}\frac{e^y}{(1+e^y)^2}. \tag{A.79}$$

Thermal corrections can then be considered by replacing:

$$m_t^2 \to \mathcal{M}_t^2 = m_t^2 + \frac{g_3^2T^2}{6}, \tag{A.80}$$

where the second term is the effect of substituting the particle for the quasiparticle to incorporate the self-energy thermal correction dominate by the gluons. Since the form factor takes the limit $\mathcal{F}(\tau,m,T) \to 0.25$ [18], (4.65) is recovered by setting the extra colour factor as $N_c = 3$.

A.8 Solving Diffusion Equations

In Sect. 4.5, it was mentioned that the system of diffusion equation (4.74) may be solved in terms of the Higgs density H in (4.75). The assumption is that the strong sphaleron and Yukawa rate is sufficiently larger than the Higgs rate and chirality flip rate:

$$\begin{aligned} \frac{Q}{k_Q}-\frac{H}{k_H}-\frac{T}{k_T} &= \mathcal{O}\left(\frac{1}{\Gamma_y}\right), \\ \frac{2Q}{k_Q}-\frac{T}{k_T}+\frac{9(Q+T)}{k_B} &= \mathcal{O}\left(\frac{1}{\Gamma_{ss}}\right) \end{aligned} \implies \begin{aligned} Q &= H\left(\frac{k_Q(9k_T-k_B)}{k_H(k_B+9k_Q+9k_T)}\right)+\mathcal{O}\left(\frac{1}{\Gamma_{ss}},\frac{1}{\Gamma_y}\right), \\ T &= -H\left(\frac{k_T(2k_B+9k_Q)}{k_H(k_B+9k_Q+9k_T)}\right)+\mathcal{O}\left(\frac{1}{\Gamma_{ss}},\frac{1}{\Gamma_y}\right). \end{aligned}$$

These are then substituted by into a linear combination of (4.74):

$$\left[D_Q\left(2T''+Q''\right)-D_H H''\right]-v_w(2T'+Q'-H)+\Gamma_h\left(\frac{H}{k_H}\right)-S_t^{CPV}=0, \tag{A.81}$$

which can be solved using the Greens function:

$$G_H(z|z_0):=\frac{\overline{D}^{-1}}{k_+-k_-}\begin{cases}e^{k_+(z-z_0)} & z\in(z_0,\infty),\\ e^{k_-(z-z_0)} & z\in(-\infty,z_0).\end{cases} \tag{A.82}$$

Similarly (4.82) can be solved with a Greens function:

$$G_B(z|z_0)=-\frac{1}{v_w}\begin{cases}1 & z\in(0,\infty),\\ e^{v_w(z-z_0)/D_q} & z\in(-\infty,0),\end{cases} \tag{A.83}$$

and Γ_{ws} is evaluated at $v^s_{T_c}$ according to (4.53) and the values in Table 4.1.

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