Dissertation

Electroweak Corrections to Higgs boson Decays into Sfermions in the Minimal Supersymmetric Standard Model

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Univ. Prof. Walter Majerotto
E 136
Institut für Theoretische Physik
und
Institut für Hochenergiephysik der ÖAW

eingereicht an der Technischen Universität Wien
Technisch–naturwissenschaftliche Fakultät

von

Dipl. Ing. Christian Weber
E9207111
Schiffmühlenstraße 57/1/10, A-1220 Wien
weber@hephy.oeaw.ac.at

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Kurzfassung

Das Standardmodell der Elementarteilchenphysik beschreibt alle derzeit bekannten fundamentalen Materieteilchen sowie deren Wechselwirkungen, nämlich die starke, die schwache und die elektromagnetische Kraft. Obwohl das Standardmodell bei der korrekten Beschreibung experimenteller Ergebnisse im Rahmen der Messgenauigkeit große Erfolge verzeichnen kann, wird angenommen, dass es eine effektive Theorie ist, die nur für die derzeitig erreichbaren Energien Gültigkeit hat. Um physikalische Phänomene auch bei höheren Energien beschreiben zu können, muss das Standardmodell erweitert werden. Der aussichtsreichste Kandidat hierfür ist das Minimale Supersymmetrische Standard Model (MSSM), dem das Konzept der Supersymmetrie zugrunde liegt. Das MSSM sagt die Existenz von sogenannten supersymmetrischen Teilchen zu den bisher bekannten Teilchen voraus. Die Suche nach diesen supersymmetrischen Teilchen ist deshalb eines der wichtigsten Ziele des Large Hadron Colliders (LHC) am Kernforschungszentrum CERN, bei dem genügend hohe Energien zur Verfügung stehen sollen, um Kollisionen zu erzeugen, die die Existenz von Supersymmetrie bestätigen.

Unter diesen supersymmetrischen Teilchen befinden sich die Partnerteilchen der Fermionen, die Sfermionen, sowie fünf physikalische Higgsbosonen. Um diese Teilchen entdecken zu können, sind genaue Vorhersagen ihrer Zerfallskanäle und Verzweigungsverhältnisse erforderlich.

In dieser Arbeit werden die Zerfälle von Higgsbosonen in zwei Sfermionen und ihre gekreuzten Kanäle analysiert. Von großem Interesse sind dabei die Sfermionen der dritten Generation, weil man annimmt, dass sie wegen ihrer starken Yukawakopplungen sowie ihrer Links-Rechts Mischung sehr leicht sind. In der Berechnung der Zerfallsbreiten werden die vollständigen elektroschwachen Einschleifen-Korrekturen berücksichtigt. Aufgrund der Tatsache, dass in diesem Prozess fast alle Parameter des MSSM renormiert werden müssen und daher eine große Anzahl von Feynmangraphen berechnet werden muss, gestaltet sich diese Aufgabe als ziemlich schwierig. In bestimmten Zerfallskanälen, vor allem für große Werte des Parameters $\tan\beta$, führt das On-shell Renormierungverfahren zu inakzeptablen Ergebnissen, sodass eine Verbesserung nötig ist. Dieses Problem kann durch eine Redefinition des Tree-levels gelöst werden, indem die fermionischen Masses und die trilinearen Kopplungen als $\overline{\mathrm{DR}}$ Größen behandelt werden. Die Entwicklung um diesen redefinierten Tree-level führt daraufhin zu vernünftigen Ergebnissen. Um die benötigten $\overline{\mathrm{DR}}$ und On-shell Größen in konsistenter Weise zu erhalten, muss die Umrechnung von den $\overline{\mathrm{DR}}$ auf die On-shell Größen und umgekehrt sorgfältig behandelt werden. Die starke Verwicklung der in diesem Verfahren auftretenden Parameter machen die Entwicklung eines ausgeklügelten Iterationsverfahrens notwendig.

Abstract

The Standard Model of elementary particle physics describes all presently known fundamental particles that make up all matter as well as their interactions, i.e. the strong, weak, and electromagnetic forces. Despite its great success in explaining experimental results correctly within the scope of precision measurements at current particle accelerators, it is believed to be an effective theory valid only at energies accessible by today's particle accelerators. Therefore, the Standard Model has to be extended to describe physics also at higher energies. The most promising candidate is the minimal supersymmetric extension of the Standard Model (MSSM). Based on the concept of supersymmetry, it predicts the existence of supersymmetric particles to every fundamental known particle. The search for supersymmetry is one of the primary goals of the Large Hadron Collider (LHC) at the CERN laboratory which should be ready for use in 2007, producing collisions at sufficiently high energies to detect the superpartners many theorists expect to see.

Among these supersymmetric particles are the partners of the fermions, called sfermions, as well as the five supersymmetric counterparts of the Higgs boson in the Standard Model. For a discovery, precise predictions for their decay modes and branching ratios are necessary.

In this thesis, we study in detail the decays of Higgs bosons into two sfermions, as well as the corresponding crossed channels. In particular, the sfermions of the third generation are interesting because one expects them to be lighter than the other sfermions due to their large Yukawa couplings and left–right mixings. We will calculate the full electroweak one–loop corrections in the on–shell renormalization scheme. Owing to the fact that almost all parameters of the MSSM have to be renormalized in this process and hence a large number of graphs has to be computed, the calculation is very complex. Despite this complexity, we have performed the calculation in an analytic way. As we will see, in some cases the on–shell scheme will lead to unacceptable results in certain decay channels which makes an improvement necessary. Especially this is the case in the decay modes involving down–type sfermions for large values of the parameter $\tan\beta$. This problem can be solved by defining an appropriate tree level in terms of $\overline{\mathrm{DR}}$ running values for the fermion masses and the trilinear couplings. The expansion around this new tree level then no longer suffers from bad convergence. In order to get consistently all needed $\overline{\mathrm{DR}}$ running and on–shell masses, we have to pay special attention to the shifting from the $\overline{\mathrm{DR}}$ to the on–shell renormalization scheme and vice versa. Since the parameters involved in these calculations are very entangled, we have to perform a sophisticated iteration procedure.

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Chapter 1

Supersymmetry

The Standard Model (SM) of elementary particle physics [1, 2] is an impressively successful theory of quarks and leptons and their electroweak and strong interactions. It is a gauge theory described by the gauge group $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ in which the electroweak $SU(2)_L \otimes U(1)_Y$ symmetry is spontaneously broken to the $U(1)_{EM}$ electromagnetic symmetry. The mechanism of spontaneous electroweak symmetry breaking (EWSB) [3] is based on the nonvanishing vacuum expectation value v of the fundamental scalar field in the SM, the *Higgs* field, to give masses to all particles which couple to the Higgs boson, in particular to the $W^{\pm}$ and Z^0 weak vector bosons. Although the SM describes almost all phenomena presently known at energies up to $\simeq$ 100 GeV, there are several fundamental questions that remain unanswered:

- *Hierarchy problem*
  One of the main arguments to extend the SM is the solution of the hierarchy problem. The SM does not explain the scale of EWSB. Phenomenologically, the mass of the Higgs particle is expected to lie in the range of the EWSB scale, i.e. $v \sim 250$ GeV. In the SM, radiative corrections to the Higgs mass (squared) depend quadratically on the UV cut–off Λ where new physics should appear, since the masses of scalar particles are not protected by chiral symmetries. Therefore, the Higgs boson is unstable against quantum corrections, leading to a natural mass close to the high scale Λ.

- *Electroweak symmetry breaking*
  In the SM, the masses of fermions and gauge bosons are generated by the Higgs mechanism which is parameterized by the Higgs boson h and its potential $V(h) \propto \mu^2 h^2 + \lambda h^4$. However, this potential is introduced by hand and without any deeper justification for a negative squared mass parameter μ^2 which accounts for the typical 'mexican hat' potential with its minimum away from $h = 0$.

- *Gauge coupling unification*
  Despite its enormous success in confirming nearly all experimental data in high precision measurements, the SM cannot be the final truth in understanding nature but is rather an effective theory valid only at energies nowadays reachable at particle accelerators. Therefore, it is expected that there exists a Grand Unified Theory (GUT) at a high scale in which the fundamental forces are treated by means of one single gauge

group. However, the most recent measurements of the coupling show that unification within the SM is not possible.

- *Cosmological issues*
  One of the most fundamental open questions is the origin of the observed *baryon asymmetry* of the universe. Although the SM fulfills all the requirements for baryogenesis [5], the electroweak phase transition is too weak to preserve the generated baryon asymmetry. Therefore, baryon asymmetry generated at the electroweak phase transition claims for new physics at the electroweak scale.
  In addition, the cosmic microwave background data strongly indicate that only about 5% of the total matter density of the universe consist of quarks leptons of the SM, while there is about five to six times more mass in the form of invisible *cold dark matter*. Unfortunately, the SM does not have a reliable candidate with the right properties to form this cold dark matter.

Therefore, the Standard Model has to be extended to describe physics also at higher energies. In the early 70's, J. Wess and B. Zumino found an attractive symmetry relating the two fundamental species of elementary particles, bosons and fermions, by a *supersymmetry* transformation,

$$\mathcal{Q}|\,\text{Fermion}\rangle = |\,\text{Boson}\rangle \qquad \mathcal{Q}|\,\text{Boson}\rangle = |\,\text{Fermion}\rangle\,. \tag{1.1}$$

In such supersymmetric models each particle of a certain type gets a *superpartner* with equal mass and the same quantum numbers but differs in spin by 1/2. Due to this boson $\leftrightarrow$ fermion symmetry the scalar masses are protected from quadratically divergent loop corrections, as the masses of the fermions are protected by chiral symmetry, providing an elegant solution to the hierarchy problem.
Since none of the predicted supersymmetric particles have been observed, SUSY must be a broken symmetry. If this supersymmetry breaking is of a certain type known as *soft breaking* [6], it doesn't forfeit some of its advantages, e.g. it does not reintroduce quadratic divergences of scalar particle masses (squared).
Even though theories including SUSY have to explain why the masses of the predicted superpartners are that high and up to now there is no direct evidence that the fundamental structure of nature is supersymmetric, such theories provide many remarkable features:

- *Hierarchy problem*
  One of the main reasons for introducing SUSY theories is their ability to solve the hierarchy problem [7]. By grouping fermions and bosons together in supermultiplets, the putative quadratically divergent radiative fermionic corrections to the Higgs boson mass are cancelled by the corresponding bosonic loop contributions of opposite sign. Hence, SUSY stabilizes the hierarchy in the sense that the 'natural' mass of the Higgs boson lies in the range of electroweak symmetry breaking which is no longer in contradiction with a very high GUT scale.

- *Electroweak symmetry breaking*
  As already indicated, in the SM an effective Higgs potential $V(h) \propto \mu^2 h^2 + \lambda h^4$ with $\mu^2 < 0$ is introduced 'by hand' to achieve EWSB. In renormalizable (supersymmetric)

theories, however, the mass parameters which enter in the Lagrangian are also scale–dependent, and renormalization group equations (RGEs) can be used to evolve the parameters from the unification scale of the order of 10^{16} GeV down to the weak scale of order 10^2 GeV. In case of the Higgs mass parameter μ, the large top quark Yukawa coupling is responsible for a negative value of μ^2 of the correct order of magnitude at the electroweak scale, thus providing a plausible explanation of the origin of EWSB [8, 9].

- *Gauge coupling unification*
  It can further be shown that in the minimal supersymmetric extension of the SM, the extrapolation of the low energy values of the gauge couplings unify at a scale $M_{\text{GUT}} \simeq 3 \times 10^{16}$ GeV [10], well in agreement with the limits on the proton lifetime.

- *Cold dark matter*
  As we have seen, supersymmetric models can solve many problems which the SM suffers from. However, without any additional structure, they can give rise to baryon and lepton number violation at unacceptable levels, e.g. proton decay can be mediated by the superpartners of quarks, i.e. $p \to \pi^0 e^+$. The non-observation of such decays has lead to the introduction of a discrete symmetry known as R–parity [11], to forbid such decays and to ensure baryon and lepton number conservation in an elegant way. As a consequence, the lightest supersymmetric particle (LSP) is absolutely stable, and, if electrically neutral it serves as a nice cold dark matter candidate.

Chapter 2

The Minimal Supersymmetric Standard Model

The simplest and most attractive extension of the Standard Model is the *Minimal Supersymmetric Standard Model* (MSSM). As indicated by its name, 'minimal' means in this case that the number of superfields and interactions is kept as small as possible. In particular, the field content of the MSSM consists only of the SM fields and their supersymmetric partners, and an additional Higgs doublet.

- **Gauge fields**
  In order to respect the $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetry of the SM, the spin–1 gauge bosons are described by the corresponding vector superfields. In particular, the eight gluons of QCD, G^a_μ, get eight spin–$\frac{1}{2}$ partners $\tilde{G}^a$ called *gluinos*, the $SU(2)$ gauge bosons W^i_μ get three *winos* $\tilde{W}^i$ as partners and the $U(1)$ gauge boson B_μ gets a *bino* $\tilde{B}$. Note that since $SU(2)_L \times U(1)_Y$ is broken in the SM, the winos and the bino do not form mass eigenstates but mix with fields with the same charge but different $SU(2)_L \otimes U(1)_Y$ quantum numbers.

Superfield	spin–1	spin–1/2	$SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$	Names
$\hat{V}^a_s$	G^a_μ	$\tilde{G}^a$	$(\mathbf{8}, \mathbf{1}, 0)$	gluons, gluinos
$\hat{V}^i$	W^i_μ	$\tilde{W}^i$	$(\mathbf{1}, \mathbf{3}, 0)$	W-bosons, winos
$\hat{V}'$	B_μ	$\tilde{B}$	$(\mathbf{1}, \mathbf{1}, 0)$	B-boson, bino

Table 2.1: Gauge supermultiplet fields in the MSSM.

- **Matter fields**
  The matter content of the SM is described by three generations of leptons and quarks, i.e. for each generation two $SU(2)_L$ fermion doublets and three singlets for the right-

handed fermions,

$$L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}, \quad E = e_R^c, \quad Q = \begin{pmatrix} u_L \\ d_L \end{pmatrix}, \quad D = d_R^c, \quad U = u_R^c. \tag{2.1}$$

Therefore, one generation of the SM is represented by five left–chiral superfields which contain the leptons and quarks given above plus their supersymmetric partners, the *sleptons* and *squarks*:

$$\tilde{L} = \begin{pmatrix} \tilde{\nu}_L \\ \tilde{e}_L \end{pmatrix}, \quad \tilde{E} = \tilde{e}_R^*, \quad \tilde{Q} = \begin{pmatrix} \tilde{u}_L \\ \tilde{d}_L \end{pmatrix}, \quad \tilde{D} = \tilde{d}_R^*, \quad \tilde{U} = \tilde{u}_R^*. \tag{2.2}$$

Superfield	spin–1/2	spin–0	$SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$	Names
$\hat{Q}$	(u_L, d_L)	$(\tilde{u}_L, \tilde{d}_L)$	$(\mathbf{3}, \mathbf{2}, \frac{1}{3})$	quarks, squarks
$\hat{U}^c$	$\bar{u}_R$	$\tilde{u}_R^*$	$(\bar{\mathbf{3}}, \mathbf{1}, -\frac{4}{3})$	(× 3 families)
$\hat{D}^c$	$\bar{d}_R$	$\tilde{d}_R^*$	$(\bar{\mathbf{3}}, \mathbf{1}, \frac{2}{3})$	
$\hat{L}$	(ν_L, e_L)	$(\tilde{\nu}_L, \tilde{e}_L)$	$(\mathbf{1}, \mathbf{2}, \text{-}1)$	leptons, sleptons
$\hat{E}^c$	$\bar{e}_R$	$\tilde{e}_R^*$	$(\mathbf{1}, \mathbf{1}, 2)$	(× 3 families)
$\hat{H}_1$	$(\tilde{H}_1^1, \tilde{H}_1^2)$	(H_1^1, H_1^2)	$(\mathbf{1}, \mathbf{2}, \text{-}1)$	higgsinos, Higgs
$\hat{H}_2$	$\tilde{H}_2^1, \tilde{H}_2^2$	H_2^1, H_2^2	$(\mathbf{1}, \mathbf{2}, 1)$	

Table 2.2: Chiral supermultiplet fields in the MSSM.

- **Higgs sector**
  Contrary to the SM, two chiral superfield doublets with hypercharges ±1 are required to break $SU(2)_L \times U(1)_Y$ invariance and to give masses to both up– and down–type fermions. One reason is that if there was only one single chiral superfield doublet, the gauge symmetry would suffer from a fermion triangle gauge anomaly. This can easily be seen from the conditions for anomaly cancellation, $\mathrm{Tr}[Y^3] = \mathrm{Tr}[T_3^2 Y] = 0$, where T_3 and Y denote the third component of the isospin and the weak hypercharge, respectively, and the electric charge given by $Q = T_3 + Y/2$. In the SM these conditions are satisfied by a complete generation of SM fermions. To cancel the contribution from one superfield doublet, one needs a second doublet to get a consistent quantum theory. The two Higgs doublets and their superpartners, the *higgsinos* $\tilde{H}_i^j$, are given as follows:

$$H_1 = \begin{pmatrix} H_1^1 \\ H_1^2 \end{pmatrix}, \qquad H_2 = \begin{pmatrix} H_2^1 \\ H_2^2 \end{pmatrix},$$
$$\tilde{H}_1 = \begin{pmatrix} \tilde{H}_1^1 \\ \tilde{H}_1^2 \end{pmatrix}, \qquad \tilde{H}_2 = \begin{pmatrix} \tilde{H}_2^1 \\ \tilde{H}_2^2 \end{pmatrix}. \tag{2.3}$$

2.1 MSSM Lagrangian

The complete Lagrangian of the MSSM can be written as

$$\mathcal{L}_{\text{MSSM}} = \mathcal{L}_{\text{kinetic}} - V_Y - V_F - V_D - V_{\tilde{G}\psi\tilde{\psi}} + \mathcal{L}_{\text{soft}} , \tag{2.4}$$

where $\mathcal{L}_{\text{kinetic}}$ stands for both the standard kinetic terms for each particle and their interactions with the gauge bosons.
The interactions described by the potentials V_X are all restricted by supersymmetry, and the last term ($\mathcal{L}_{\text{soft}}$) includes the SUSY–breaking terms.
In general, interactions in the MSSM have two different sources:

- **Superpotential**
  The superpotential for the MSSM is given by

  $$W = \varepsilon_{ij} \left[h_e \hat{H}_1^i \hat{L}^j \hat{E}^c + h_d \hat{H}_1^i \hat{Q}^j \hat{D}^c + h_u \hat{H}_2^j \hat{Q}^i \hat{U}^c - \mu \hat{H}_1^i \hat{H}_2^j \right] , \tag{2.5}$$

  where the hatted quantities $\hat{H}_j^i, \hat{Q}^i, \hat{L}^j, \hat{U}, \hat{D}, \hat{E}$ are the chiral superfields given in Table 2.2. Due to better readability we have suppressed all colour, weak isospin and generation indices.

  The superpotential determines two kinds of interactions mentioned in eq. (2.4). Firstly, the Yukawa potential V_Y can be obtained by replacing two superfields in the superpotential by the corresponding fermionic fields and the remaining superfield by its scalar representative,

  $$\begin{aligned} V_Y &= \varepsilon_{ij} \left[h_e H_1^i L^j E^c + h_d H_1^i Q^j D^c + h_u H_2^j Q^i U^c - \mu \tilde{H}_1^i \tilde{H}_2^j \right] \\ &+ \varepsilon_{ij} \left[h_e \tilde{H}_1^i L^j \tilde{E}^c + h_d \tilde{H}_1^i Q^j \tilde{D}^c + h_u \tilde{H}_2^j Q^i \tilde{U}^c \right] \\ &+ \varepsilon_{ij} \left[h_e \tilde{H}_1^i \tilde{L}^j E^c + h_d \tilde{H}_1^i \tilde{Q}^j D^c + h_u \tilde{H}_2^j \tilde{Q}^i U^c \right] \\ &+ \text{h.c.} \end{aligned} \tag{2.6}$$

  The F–term potential V_F originates from using the equations of motion for the auxiliary fields F_i,

  $$V_F = \sum_i F^{*i} F_i = \sum_i \left| \frac{\delta W(\varphi)}{\delta \varphi_i} \right|^2 \tag{2.7}$$

  where the sum is taken over all scalar components φ_i of the superfields.

- **Gauge Symmetry**
  Apart from the usual SM–like gauge interactions, the MSSM also has terms that are related to gauge symmetry although they contain neither gauge bosons nor gauginos. These terms have their origin in eliminating the auxiliary fields D^a, therefore they are called D–terms. The corresponding D–term potential is given by

  $$V_D = \frac{1}{2} \sum D^a D^a , \tag{2.8}$$

with

$$D^a = g^a \varphi_i^* (T^a)_i^j \varphi_j \,, \tag{2.9}$$

where φ_i are the scalar components of the superfields and T^a denoting the generators of the gauge group satisfying $[T^a, T^b] = i f^{abc} T^c$.

In fact, there exists one additional kind of interaction allowed by gauge invariance involving the gaugino fields. The corresponding potential $V_{\tilde{G}\psi\tilde{\psi}}$ is given by

$$V_{\tilde{G}\psi\tilde{\psi}} = i\sqrt{2} g_a \varphi_k \bar{\lambda}^a (T^a)_{kl} \bar{\psi}_l + \text{h.c.} \,, \tag{2.10}$$

where (φ, ψ) are the spin–0 and spin spin–$\frac{1}{2}$ components of the chiral superfield, respectively, and λ^a denoting the gaugino field.

The last term in the full Lagrangian of the MSSM, $\mathcal{L}_{\text{soft}}$, involves the soft SUSY–breaking terms and can be explicitly written as

$$\begin{aligned} -\mathcal{L}_{\text{soft}} &= m_{H_1}^2 |H_1|^2 + m_{H_2}^2 |H_2|^2 - m_{12}^2 (H_1 H_2 + H_1^\dagger H_2^\dagger) + \frac{1}{2} m_{\tilde{g}}\, \tilde{g}^a \tilde{g}^a + \frac{1}{2} M\, \tilde{W}^i \tilde{W}^i \\ &+ \frac{1}{2} M'\, \tilde{B}\tilde{B} + M_{\tilde{Q}}^2 |\tilde{q}_L|^2 + M_{\tilde{U}}^2 |\tilde{u}_R^c|^2 + M_{\tilde{D}}^2 |\tilde{d}_R^c|^2 + M_{\tilde{L}}^2 |\tilde{l}_L|^2 + M_{\tilde{E}}^2 |\tilde{e}_R^c|^2 \\ &+ \varepsilon_{ij} \left(h_e A_e H_1^i \tilde{L}^j \tilde{E}^c + h_d A_d H_1^i \tilde{Q}^j \tilde{D}^c + h_u A_u H_2^j \tilde{Q}^i \tilde{U}^c + \text{h.c.} \right) , \end{aligned} \tag{2.11}$$

where we have introduced the SUSY–breaking mass parameters $m_{H_1}^2$, $m_{H_2}^2$, m_{12}^2, $m_{\tilde{g}}$, M, M', $M_{\tilde{Q}}^2$, $M_{\tilde{U}}^2$, $M_{\tilde{D}}^2$, $M_{\tilde{L}}^2$, $M_{\tilde{E}}^2$ as well as the trilinear scalar couplings A_e, A_u, A_d.

2.2 MSSM spectrum

2.2.1 Higgs sector

In the MSSM, two complex Higgs doublets or eight real scalar degrees of freedom (DOF) are required to describe electroweak symmetry breaking, i.e. the Higgs scalars acquire non-vanishing vacuum expectation values (VEVs). From these eight real scalar DOF three are massless and become the longitudinal modes of the massive vector bosons Z^0 and $W^\pm$. The masses of the remaining five DOF, representing the three neutral Higgs bosons h^0, H^0, A^0 and the two charged ones, $H^\pm$, can be obtained by expanding the Higgs potential around its minimum, up to second order in the fields. The scalar Higgs potential in the MSSM is given by

$$\begin{aligned} V &= m_1^2 |H_1|^2 + m_2^2 |H_2|^2 - m_{12}^2 (H_1 H_2 + H_1^\dagger H_2^\dagger) \\ &+ \frac{1}{8}(g^2 + g'^2)(|H_1|^2 - |H_2|^2)^2 + \frac{g^2}{2} |H_1^\dagger H_2|^2 \,, \end{aligned} \tag{2.12}$$

with $m_i^2 = m_{H_i}^2 + |\mu|^2$, where $m_{H_i}^2$ and m_{12}^2 are soft SUSY–breaking parameters. The quadratic terms $\propto |\mu|^2$ originate from F–terms, and the terms involving four scalar Higgs fields are

derived from the D–term potential. Note that in contrast to the SM, where the strength of the Higgs self–interaction is an unknown free parameter, the quartic interactions in the MSSM are completely determined by the gauge couplings g_i.
As a next step, electroweak symmetry should be broken down to electromagnetism $SU(2)_L \otimes U(1)_Y \rightarrow U(1)_{EM}$, when the Higgs fields

$$H_1 = \begin{pmatrix} H_1^1 \\ H_1^2 \end{pmatrix} = \begin{pmatrix} H_1^0 \\ H_1^- \end{pmatrix}, \qquad H_2 = \begin{pmatrix} H_2^1 \\ H_2^2 \end{pmatrix} = \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix},$$

get nonvanishing VEVs. Having the freedom to make $SU(2)_L$ gauge transformations, we can rotate away a possible VEV of one component of the scalar fields. Without loss of generality we can choose $\langle H_2^+ \rangle = 0$ at the minimum of the potential, implying that also the VEV of the negatively charged component of H_1 is vanishing, $\langle H_1^- \rangle = 0$. Since both charged components of the Higgs scalars remain unaffected, electromagnetism is not spontaneously broken, in agreement with experiment. Therefore, only the neutral Higgs boson fields acquire a nonvanishing VEV, i.e.

$$\langle H_1 \rangle = \begin{pmatrix} v_1 \\ 0 \end{pmatrix}, \qquad \langle H_2 \rangle = \begin{pmatrix} 0 \\ v_2 \end{pmatrix}. \tag{2.13}$$

In order to determine the minimum of the Higgs potential V, we can take m_{12}^2 in the term $\propto H_1 H_2$ to be real and positive, since this is the only piece which depends on the phases of the Higgs fields; any possible phase in m_{12}^2 can be absorbed into the phases of H_1 and H_2. Thus the product $H_1^0 H_2^0$ also has to be real and positive, and, concerning their opposite weak hypercharge, both phases can be made zero by a $U(1)_Y$ gauge transformation. As a consequence, CP invariance cannot be broken spontaneously in the MSSM, which means that the eigenstates of the Higgs boson are also eigenstates of CP.
For the two Higgs doublets we choose the following common parameterization:

$$H_1 \equiv \begin{pmatrix} H_1^0 \\ H_1^- \end{pmatrix} = \begin{pmatrix} v_1 + (\phi_1^0 + i\chi_1^0)/\sqrt{2} \\ \phi_1^- \end{pmatrix}, \qquad Y_{H_1} = -1 \tag{2.14}$$

$$H_2 \equiv \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix} = \begin{pmatrix} \phi_2^+ \\ v_2 + (\phi_2^0 + i\chi_2^0)/\sqrt{2} \end{pmatrix}, \qquad Y_{H_2} = +1 \tag{2.15}$$

The minimum of the Higgs potential can now be obtained easily by solving the equations

$$\left.\frac{\partial V}{\partial H_1^0}\right|_{\langle H_n^0 \rangle = v_n} = \left.\frac{\partial V}{\partial H_2^0}\right|_{\langle H_n^0 \rangle = v_n} = 0\,, \tag{2.16}$$

resulting in the two minimization conditions

$$m_1^2 v_1 = -m_{12}^2 v_2 - \frac{1}{4}(g^2 + g'^2)(v_1^2 - v_2^2)\,, \tag{2.17}$$

$$m_2^2 v_2 = -m_{12}^2 v_1 + \frac{1}{4}(g^2 + g'^2)(v_1^2 - v_2^2)\,. \tag{2.18}$$

Since one combination of the VEVs,

$$m_Z^2 = \frac{g^2 + g'^2}{2}(v_1^2 + v_2^2)\,, \qquad m_W^2 = \frac{g^2}{2}(v_1^2 + v_2^2)\,, \tag{2.19}$$

and hence

$$v^2 \equiv (v_1^2 + v_2^2) = \frac{2m_Z^2}{g^2 + g'^2} \approx (174 \text{ GeV})^2 , \tag{2.20}$$

is very well known from experiment, we can express both VEVs in terms of one single parameter,

$$\tan\beta \equiv \frac{v_2}{v_1} \geq 0 , \qquad 0 \leq \beta \leq \frac{\pi}{2} . \tag{2.21}$$

Eqs. (2.17) and (2.18) may now be written as

$$m_1^2 = -m_{12}^2 \tan\beta - \frac{1}{2} m_Z^2 \cos 2\beta , \tag{2.22}$$

$$m_2^2 = -m_{12}^2 \cot\beta + \frac{1}{2} m_Z^2 \cos 2\beta . \tag{2.23}$$

Therefore, the Higgs sector at tree level only depends on two free parameters.
The Higgs mass spectrum is obtained by evaluating the second derivatives of the Higgs potential, taken at its minimum,

$$M_{ij}^{2,\text{Higgs}} = \frac{1}{2} \frac{\partial^2 V}{\partial H_i \partial H_j}\bigg|_{\langle H_n^0 \rangle = v_n} . \tag{2.24}$$

At tree level, $M_{ij}^{2,\text{Higgs}}$ splits into four independent $2{\times}2$ mass matrices which can be separately diagonalized. In terms of the original gauge eigenstate fields, the mass eigenstates are given by

$$\begin{pmatrix} H^0 \\ h^0 \end{pmatrix} = \begin{pmatrix} \cos\alpha & \sin\alpha \\ -\sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} \phi_1^0 \\ \phi_2^0 \end{pmatrix} , \tag{2.25}$$

$$\begin{pmatrix} G^0 \\ A^0 \end{pmatrix} = \begin{pmatrix} -\cos\beta & \sin\beta \\ \sin\beta & \cos\beta \end{pmatrix} \begin{pmatrix} \chi_1^0 \\ \chi_2^0 \end{pmatrix} , \tag{2.26}$$

$$\begin{pmatrix} G^\pm \\ H^\pm \end{pmatrix} = \begin{pmatrix} -\cos\beta & \sin\beta \\ \sin\beta & \cos\beta \end{pmatrix} \begin{pmatrix} \phi_1^\pm \\ \phi_2^\pm \end{pmatrix} . \tag{2.27}$$

The Goldstone bosons G^0 and $G^\pm$ are 'eaten' by the longitudinal components of the massive vector bosons Z^0 and $W^\pm$, respectively. The five remaining physical Higgs bosons form two CP–even states (h^0, H^0), one CP–odd state A^0 and the two charged Higgs bosons $H^\pm$. As already mentioned above, the two free parameters in the Higgs sector are conventionally chosen to be the mass of the pseudo–scalar Higgs boson A^0 and the ratio of the two VEVs,

$$m_{A^0} , \qquad \text{and} \qquad \tan\beta . \tag{2.28}$$

The remaining parameters such as the masses and the mixing angle α can be expressed using these free parameters as

$$m_{h^0,H^0}^2 = \frac{1}{2}\left[m_{A^0}^2 + m_Z^2 \mp \sqrt{(m_{A^0}^2 + m_Z^2)^2 - 4m_{A^0}^2 m_Z^2 \cos^2\beta} \right] , \tag{2.29}$$

$$m_{H^\pm}^2 = m_{A^0}^2 + m_W^2 , \tag{2.30}$$

$$\tan 2\alpha = \tan 2\beta \frac{m_{A^0}^2 + m_Z^2}{m_{A^0}^2 - m_Z^2} , \qquad -\frac{\pi}{2} \leq \alpha \leq 0 . \tag{2.31}$$

2.2.2 Sfermion sector

The sfermion mixing is described by the sfermion mass matrix in the left–right basis $(\tilde{f}_L, \tilde{f}_R)$, and in the mass basis $(\tilde{f}_1, \tilde{f}_2)$, $\tilde{f} = \tilde{t}, \tilde{b}$ or $\tilde{\tau}$,

$$\mathcal{M}^2_{\tilde{f}} = \begin{pmatrix} m^2_{\tilde{f}_L} & a_f\, m_f \\ a_f\, m_f & m^2_{\tilde{f}_R} \end{pmatrix} = \left(R^{\tilde{f}}\right)^\dagger \begin{pmatrix} m^2_{\tilde{f}_1} & 0 \\ 0 & m^2_{\tilde{f}_2} \end{pmatrix} R^{\tilde{f}}\,, \tag{2.32}$$

where $R^{\tilde{f}}_{i\alpha}$ is a 2×2 rotation matrix with rotation angle $\theta_{\tilde{f}}$,

$$R^{\tilde{f}}_{ij} = \begin{pmatrix} \cos\theta_{\tilde{f}} & \sin\theta_{\tilde{f}} \\ -\sin\theta_{\tilde{f}} & \cos\theta_{\tilde{f}} \end{pmatrix}\,, \tag{2.33}$$

which relates the mass eigenstates $\tilde{f}_i$, $i = 1, 2$, $(m_{\tilde{f}_1} < m_{\tilde{f}_2})$ to the gauge eigenstates $\tilde{f}_\alpha$, $\alpha = L, R$, by $\tilde{f}_i = R^{\tilde{f}}_{i\alpha} \tilde{f}_\alpha$ and

$$m^2_{\tilde{f}_L} = M^2_{\{\tilde{Q},\tilde{L}\}} + (I^{3L}_f - e_f \sin^2\theta_W) \cos 2\beta\, m^2_Z + m^2_f\,, \tag{2.34}$$

$$m^2_{\tilde{f}_R} = M^2_{\{\tilde{U},\tilde{D},\tilde{E}\}} + e_f \sin^2\theta_W \cos 2\beta\, m^2_Z + m^2_f\,, \tag{2.35}$$

$$a_f = A_f - \mu\, (\tan\beta)^{-2I^{3L}_f}\,. \tag{2.36}$$

$M_{\tilde{Q}}$, $M_{\tilde{L}}$, $M_{\tilde{U}}$, $M_{\tilde{D}}$ and $M_{\tilde{E}}$ are soft SUSY–breaking masses, A_f is the trilinear scalar coupling parameter, μ the higgsino mass parameter, I^{3L}_f denotes the third component of the weak isospin of the fermion f, e_f the electric charge in terms of the elementary charge e_0, and θ_W is the Weinberg angle.
The mass eigenvalues and the mixing angle in terms of primary parameters are

$$m^2_{\tilde{f}_{1,2}} = \frac{1}{2}\left(m^2_{\tilde{f}_L} + m^2_{\tilde{f}_R} \mp \sqrt{(m^2_{\tilde{f}_L} - m^2_{\tilde{f}_R})^2 + 4a^2_f m^2_f}\right)\,, \tag{2.37}$$

$$\cos\theta_{\tilde{f}} = \frac{-a_f\, m_f}{\sqrt{(m^2_{\tilde{f}_L} - m^2_{\tilde{f}_1})^2 + a^2_f m^2_f}} \qquad (0 \le \theta_{\tilde{f}} < \pi)\,, \tag{2.38}$$

and the trilinear breaking parameter A_f can be written as

$$m_f A_f = \frac{1}{2}(m^2_{\tilde{f}_1} - m^2_{\tilde{f}_2}) \sin 2\theta_{\tilde{f}} + m_f\, \mu\, (\tan\beta)^{-2I^{3L}_f}\,. \tag{2.39}$$

The mass of the sneutrino $\tilde{\nu}_\tau$ is given by

$$m^2_{\tilde{\nu}_\tau} = M^2_{\tilde{L}} + \frac{1}{2}\, m^2_Z \cos 2\beta\,. \tag{2.40}$$

2.2.3 Chargino and Neutralino sector

The fermionic superpartners of the gauge bosons, the gauginos, and the superpartners of the Higgs bosons, the higgsinos, mix to form mass eigenstates called charginos and neutralinos.

The charginos are therefore the superpartners of the gauge bosons $W^\pm$ and the charged Higgs bosons $H^\pm$. In Weyl representation, the chargino fields [12]

$$\psi^+ = (-i\tilde{W}^+, \tilde{H}_2^+) \qquad\qquad \psi^- = (-i\tilde{W}^-, \tilde{H}_1^-)\,, \tag{2.41}$$

enter in the mass term of the Lagrangian in the following form:

$$\mathcal{L} = -\frac{1}{2}\left(\psi^+, \psi^-\right) \cdot \begin{pmatrix} 0 & X^T \\ X & 0 \end{pmatrix} \cdot \begin{pmatrix} \psi^+ \\ \psi^- \end{pmatrix} + \text{h.c.}\,, \tag{2.42}$$

with

$$X = \begin{pmatrix} M & \sqrt{2}m_W \sin\beta \\ \sqrt{2}m_W \cos\beta & \mu \end{pmatrix}\,. \tag{2.43}$$

Since we work in the CP–conserving MSSM, the mass matrix X can be diagonalized by two *real* 2×2 matrices U and V according to

$$UXV^{-1} = \begin{pmatrix} m_{\tilde{\chi}_1^\pm} & 0 \\ 0 & m_{\tilde{\chi}_2^\pm} \end{pmatrix}\,, \qquad |m_{\tilde{\chi}_1^\pm}| \leq |m_{\tilde{\chi}_2^\pm}|\,. \tag{2.44}$$

In Dirac representation, the mass eigenstates are related to the gauge eigenstates by

$$\tilde{\chi}_i^+ \;\equiv\; \begin{pmatrix} V_{ij}\,\psi_j^+ \\ U_{ij}\,\bar{\psi}_j^- \end{pmatrix}\,. \tag{2.45}$$

As these matrices are only of rank 2, the mass eigenvalues can be given analytically:

$$m^2_{\tilde{\chi}^\pm_{1,2}} \;=\; \frac{1}{2}\left[M^2 + \mu^2 + 2m_W^2 \mp \sqrt{(M^2+\mu^2+2m_W^2)^2 - 4(m_W^2 \sin 2\beta - \mu M)^2}\right] \tag{2.46}$$

The superpartners of the neutral gauge bosons, $\tilde{B}_\mu$ and $\tilde{W}_\mu^3$, and of the neutral Higgs bosons, $\tilde{H}_1^0$ and $\tilde{H}_2^0$, mix to form four neutral mass eigenstates called *neutralinos*. In the interaction base one can combine the four Weyl states as

$$\psi_j^0 = (-i\tilde{B}, -i\tilde{W}^3, \tilde{H}_1^0, \tilde{H}_2^0)\,. \tag{2.47}$$

In terms of the vector ψ^0 the neutralino mass terms in the Lagrangian are

$$\mathcal{L} \;=\; -\frac{1}{2}\left(\psi^0\right)^T Y\,\psi^0 + h.c.\,, \tag{2.48}$$

where we used the neutralino mass matrix defined as

$$Y = \begin{pmatrix} M' & 0 & -m_Z s_W \cos\beta & m_Z s_W \sin\beta \\ 0 & M & m_Z c_W \cos\beta & -m_Z c_W \sin\beta \\ -m_Z s_W \cos\beta & m_Z c_W \cos\beta & 0 & -\mu \\ m_Z s_W \sin\beta & -m_Z c_W \sin\beta & -\mu & 0 \end{pmatrix}\,. \tag{2.49}$$

We use the short forms s_W and c_W for the sine and the cosine of the Weinberg angle. Due to the Majorana nature of the neutralinos, the matrix can be diagonalized using only one single rotation matrix,

$$ZYZ^{-1} = \mathrm{diag}(m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}_2^0}, m_{\tilde{\chi}_3^0}, m_{\tilde{\chi}_4^0}), \qquad |m_{\tilde{\chi}_1^0}| \leq |m_{\tilde{\chi}_2^0}| \leq |m_{\tilde{\chi}_3^0}| \leq |m_{\tilde{\chi}_4^0}| \,, \tag{2.50}$$

where we again assume the mixing matrix to be real and allow the eigenvalues to be negative. The 4–component Majorana spinors for the neutralino fields can be constructed as

$$\tilde{\chi}_i^0 \equiv Z_{ij} \tilde{\psi}_j^0 \,, \tag{2.51}$$

with the corresponding mass term Lagrangian

$$\mathcal{L} = -\frac{1}{2} \sum_{i=1}^{4} m_{\tilde{\chi}_i^0} \, \bar{\tilde{\chi}}_i^0 \, \tilde{\chi}_i^0 \,. \tag{2.52}$$

Chapter 3

Renormalization

Nowadays, experiments at particle accelerators have reached such a high precision that models in elementary particle physics can be studied at the quantum level. Therefore, theoretical predictions of observables performed in the Born approximation are not sufficient anymore and have to be improved by the inclusion of higher order corrections in perturbation theory. Practically this means that one is confronted with the calculation of a large number of *Feynman diagrams* with loops illustrating integrals over indefinite momenta. In general, these integrals are divergent for large momenta and therefore have to be treated in a proper way. In order to give such expressions a physical meaning, the divergences have to be absorbed into the fields and parameters of the Lagrangian — this redefinition, which can be achieved in various ways, is called *renormalization procedure.*
In this thesis, we will make use of the so–called *multiplicative renormalization.* In this scheme all *bare* parameters and fields entering in the original Lagrangian are replaced by their corresponding *renormalized* ones, which are obtained by the multiplication with appropriate renormalization constants:

$$g_0 \;\rightarrow\; Z_g\, g \;=\; \Big(1 + \frac{\delta g}{g}\Big)\, g\,, \tag{3.1}$$

$$\phi_0 \;\rightarrow\; Z_\phi^{1/2} \phi \;=\; \Big(1 + \frac{1}{2}\delta Z_\phi\Big)\, \phi\,. \tag{3.2}$$

Expanding the renormalization constants Z_g and $Z_\phi^{1/2}$ around the value 1, the original Lagrangian splits into a renormalized Lagrangian and a part containing the *counter terms* δg and δZ_ϕ, i.e.

$$\mathcal{L}\big(g_0, \phi_0\big) \;=\; \mathcal{L}\big(g, \phi\big) + \delta\mathcal{L}\big(g, \phi, \delta g, \delta Z_\phi\big)\,. \tag{3.3}$$

In order to absorb the divergences mentioned above and to give the parameters a well–defined meaning, these counter terms have to fulfill several requirements depending on the chosen renormalization scheme.
In this thesis, we use the *on–shell renormalization scheme* [13] as far as is possible in the decay processes considered. In this scheme the finite parts of the renormalization constants are determined e.g. by the condition that the propagator of each particle is exactly at its physical mass. The main advantage of this on–shell approach is that it identifies the renormalized parameters with observable and, therefore, scale–independent quantities. However,

we will see that applying this scheme will lead to unacceptable results in certain decay channels which makes an improvement necessary.
In the following we will review some results from the renormalization of the Standard Model and discuss the renormalization of two–point–functions, one of the main ingredients when performing radiative correction within the on–shell scheme. Furthermore, we give all renormalization conditions of the parameters needed for the explicit calculation.

3.1 SM gauge sector

The gauge sector of the Standard Model is not affected by its minimal extension, the MSSM. Thus, the treatment of the electroweak gauge sector is identical to the renormalization procedure in the SM, which is discussed in detail in [14, 15]. For the gauge fields the renormalization constants are given by

$$W^{\pm}_{\mu} \;\rightarrow\; \left(1+\tfrac{1}{2}\delta Z_W\right) W^{\pm}_{\mu}\,, \tag{3.4}$$

$$\begin{pmatrix} A_\mu \\ Z_\mu \end{pmatrix} \;\rightarrow\; \begin{pmatrix} 1+\frac{1}{2}\delta Z_{AA} & \frac{1}{2}\delta Z_{AZ} \\ \frac{1}{2}\delta Z_{ZA} & 1+\frac{1}{2}\delta Z_{ZZ} \end{pmatrix} \begin{pmatrix} A_\mu \\ Z_\mu \end{pmatrix}. \tag{3.5}$$

Since the photon stays massless also after renormalization, only the weak gauge bosons Z^0 and $W^{\pm}$ receive mass corrections, i.e.

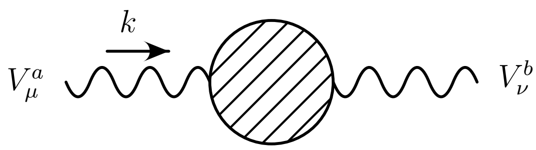
$$m_W^2 \rightarrow m_W^2 + \delta m_W^2\,, \qquad m_Z^2 \rightarrow m_Z^2 + \delta m_Z^2. \tag{3.6}$$

Decomposing the vector two–point–functions and the associated self–energies into their transverse and longitudinal parts,

$$\Gamma^W_{\mu\nu}(k) \;=\; -ig_{\mu\nu}\big(k^2-m_W^2\big) - i\Big(g_{\mu\nu}-\frac{k_\mu k_\nu}{k^2}\Big)\Pi^W_T(k^2) - i\frac{k_\mu k_\nu}{k^2}\Pi^W_L(k^2)\,, \tag{3.7}$$

$$\Gamma^{ab}_{\mu\nu}(k) \;=\; -ig_{\mu\nu}\big(k^2-m_a^2\big)\delta_{ab} - i\Big(g_{\mu\nu}-\frac{k_\mu k_\nu}{k^2}\Big)\Pi^{ab}_T(k^2) - i\frac{k_\mu k_\nu}{k^2}\Pi^{ab}_L(k^2)\,, \tag{3.8}$$

with $a,b=\{A,Z\}$, the corresponding renormalized self–energies in



$$\mathcal{M} = i\,\varepsilon^\mu(k)\,\hat{\Gamma}^{ab}_{\mu\nu}(k)\,\varepsilon^{*\nu}(k)$$

can be written as

$$\hat{\Pi}^W(k^2) \;=\; \Pi^W(k^2) + \left(k^2-m_W^2\right)\delta Z^W - \delta m_W^2\,, \tag{3.9}$$

$$\hat{\Pi}^{ab}(k^2) \;=\; \Pi^{ab}(k^2) + \tfrac{1}{2}\left(k^2-m_a^2\right)\delta Z^{ab} + \tfrac{1}{2}\left(k^2-m_b^2\right)\delta Z^{ba} - \delta_{ab}\delta m_a^2\,, \tag{3.10}$$

valid for both the transverse and longitudinal parts with $m_A^2=\delta m_A^2=0$. Applying the *on–shell renormalization conditions* to $\hat{\Gamma}^{ab}_{\mu\nu}(k)$,

$$\mathrm{Re}\,\hat{\Gamma}^{ab}_{\mu\nu}(k)\,\varepsilon^\nu(k)\Big|_{k^2=m_a^2} \;=\; 0, \qquad \lim_{k^2\to m_a^2}\frac{1}{k^2-m_a^2}\mathrm{Re}\,\hat{\Gamma}^{ab}_{\mu\nu}(k)\,\varepsilon^\nu(k) \;=\; -\,\varepsilon^\mu(k)\,, \tag{3.11}$$

which means that the poles of the propagators are determined by the physical (*pole*) masses and the residua are set to 1 on shell ($k^2 = m_a^2$), the counter terms are given by

$$\delta Z^{aa} = -\mathrm{Re}\,\dot{\Pi}_T^{aa}(m_a^2)\,, \qquad \delta Z^W = -\,\mathrm{Re}\,\dot{\Pi}_T^{W}(m_W^2)\,, \tag{3.12}$$

$$\delta Z^{ab} = \frac{2}{m_a^2 - m_b^2}\,\mathrm{Re}\,\Pi_T^{ab}(m_b^2)\,, \qquad a \neq b\,, \tag{3.13}$$

$$\delta m_Z^2 = \mathrm{Re}\,\Pi_T^{ZZ}(m_Z^2)\,, \qquad \delta m_W^2 = \mathrm{Re}\,\Pi_T^{WW}(m_W^2)\,, \tag{3.14}$$

with $\dot{\Pi}(m^2) = \frac{\partial}{\partial k^2}\Pi(k^2)\big|_{k^2=m^2}$.

Since the weak mixing angle is a derived quantity in the on–shell scheme, determined by the condition $m_W = m_Z\, c_W$ ($c_W \equiv \cos\theta_W$) [13], its renormalization constant can be expressed in terms of the mass counter terms of the weak gauge bosons,

$$\frac{\delta c_W^2}{c_W^2} = \frac{\delta m_W^2}{m_W^2} - \frac{\delta m_Z^2}{m_Z^2}\,, \qquad \frac{\delta s_W^2}{s_W^2} = -\,\frac{c_W^2}{s_W^2}\frac{\delta c_W^2}{c_W^2}\,. \tag{3.15}$$

3.2 Electric charge

For the renormalization of the electric charge one only has to renormalize one single vertex, for which usually the electron–positron–photon vertex is taken. In requiring for the renormalized elementary charge to describe the electromagnetic coupling in the Thomson limit, i.e. for on–shell external particles and vanishing photon momentum,

$$\bar{u}(p)\,\hat{\Gamma}_\mu^{ee\gamma}(p,p)\,u(p)\Big|_{p^2=m_e^2} = ie\,\bar{u}(p)\gamma_\mu u(p), \tag{3.16}$$

the counter term for the electric charge in $e_0 = e + \delta e$ is given by

$$\frac{\delta e}{e} = -\frac{1}{2}\delta Z^{AA} + \frac{s_W}{c_W}\frac{1}{2}\delta Z^{ZA} = \frac{1}{2}\,\dot{\Pi}_T^{AA}(0) + \frac{s_W}{c_W}\frac{\Pi_T^{AZ}(0)}{m_Z^2}\,. \tag{3.17}$$

However, the scale of high energy processes lies in the range of hundreds of GeV and thus far away from the Thomson limit. In addition, contributions of light hadrons in $\dot{\Pi}_T^{AA}(0)$ lead to large theoretical uncertainties [16, 15]. To avoid this problem, we use as input an effective $\overline{\mathrm{MS}}$ running coupling at $Q = m_Z$, where the contributions from light fermions are already absorbed [17, 18],

$$\alpha_{\overline{\mathrm{MS}}}^{\mathrm{eff}}(m_Z^2) = \frac{\alpha}{1 - \Delta\alpha_{\overline{\mathrm{MS}}}^{\mathrm{eff}}(m_Z^2)} \simeq \frac{1}{127.7}\,. \tag{3.18}$$

Here, α is the fine structure constant given in the Thomson limit, $\alpha = 1/137.036$, and

$$\Delta\alpha_{\overline{\mathrm{MS}}}^{\mathrm{eff}}(m_Z^2) = \frac{\alpha}{\pi}\left(\frac{5}{3} + \frac{55}{27}\Big(1 + \frac{\alpha}{\pi}\Big)\right) + \Delta\alpha_{\mathrm{lep}}(m_Z^2) + \Delta\alpha_{\mathrm{had}}^{(5)}(m_Z^2)\,, \tag{3.19}$$

where $\Delta\alpha_{\rm lep}(m_Z^2) \simeq 0.031497687$ are the leptonic and $\Delta\alpha^{(5)}_{\rm had}(m_Z^2) = 0.02769 \pm 0.00035$ are the hadronic contributions [19]. The counter term for the electric charge δe is then given by

$$\begin{aligned}\frac{\delta e}{e} &= \frac{1}{(4\pi)^2}\frac{e^2}{6}\Bigg[4\sum_f N_C^f\, e_f^2\Big(\Delta + \log\frac{Q^2}{x_f^2}\Big) + \sum_{\tilde f}\sum_{m=1}^{2} N_C^f\, e_f^2\Big(\Delta + \log\frac{Q^2}{m^2_{\tilde f_m}}\Big) \\ &\quad +4\sum_{k=1}^{2}\Big(\Delta + \log\frac{Q^2}{m^2_{\tilde\chi^+_k}}\Big) + \sum_{k=1}^{2}\Big(\Delta + \log\frac{Q^2}{m^2_{H^+_k}}\Big) - 22\Big(\Delta + \log\frac{Q^2}{m^2_W}\Big)\Bigg],\end{aligned} \tag{3.20}$$

with $x_f = m_Z\ \forall\ m_f < m_Z$ and $x_t = m_t$. N_C^f is the colour factor, $N_C^f = 1, 3$ for (s)leptons and (s)quarks, respectively. Δ denotes the UV divergence factor, $\Delta = 2/\epsilon - \gamma + \log 4\pi$, with γ being the Euler–Mascheroni constant $\gamma = \lim\limits_{m\to\infty}\left(\sum_{k=1}^m \frac{1}{k} - \log m\right) \sim 0.577216$.

3.3 Renormalization of two–point functions

Before we will turn to the renormalization of the remaining parameters and fields of the MSSM which we will need in our calculations, i.e. the ones of the Higgs and sfermion sector, we will have a short look on the subject of renormalizing two–point functions, as they are the basic building blocks for calculating higher order corrections.

3.3.1 Scalar particles with mixing

According to multiplicative renormalization, the unrenormalized fields $\phi_{0,i}$ and mass parameters $m_{0,i}$ in the bare Lagrangian

$$\mathcal{L}_0 = -\phi^*_{0,i}\,\delta_{ij}\left(\partial_\mu\partial^\mu + m^2_{0,i}\right)\phi_{0,j} \tag{3.21}$$

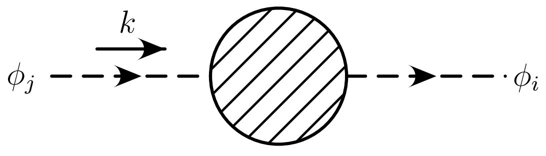
are replaced by the corresponding renormalized ones, i.e.

$$\mathcal{L} = -\phi^*_i\,\delta_{ij}\left(\partial_\mu\partial^\mu + m^2_i\right)\phi_j\,, \tag{3.22}$$

$$\phi_{0,j} = \sqrt{Z_{jk}}\,\phi_k = \big(\delta_{jk} + \tfrac{1}{2}\,\delta Z_{jk}\big)\phi_k + \mathcal{O}(\delta Z^2)\,, \tag{3.23}$$

$$(m_{0,i})^2 = m_i^2 + \delta m_i^2\,. \tag{3.24}$$

For the full renormalized two–point–function



$$\mathcal{M} = i\,\hat\Gamma_{ij}(k^2) = i\delta_{ij}(k^2 - m_i^2) + i\hat\Pi_{ij}(k^2)$$

we demand the on–shell renormalization conditions

$$\mathrm{Re}\,\hat\Gamma_{ij}(k^2)\Big|_{k^2=m_j^2} = 0\,, \qquad \lim_{k^2\to m_i^2}\frac{1}{k^2 - m_i^2}\,\mathrm{Re}\hat\Gamma_{ii}(k^2) = 1\,. \tag{3.25}$$

Inserting the renormalized self–energy $\hat{\Pi}_{ij}(k^2) = \Pi_{ij}(k^2) + \Pi^{(c)}_{ij}(k^2)$ with the self–energy counter–term

$$\Pi^{(c)}_{ij}(k^2) \;=\; -\delta_{ij}\, \delta m_i^2 + \tfrac{1}{2}\left(k^2 - m_i^2\right) \delta Z_{ij} + \tfrac{1}{2}\left(k^2 - m_j^2\right) \delta Z^*_{ji} \qquad (3.26)$$

into eq. (3.25) leads to

$$\delta m_i^2 \;=\; \mathrm{Re}\,\Pi_{ii}(m_i^2)\,, \qquad (3.27)$$

$$\delta Z_{ij} \;=\; \frac{2}{m_i^2 - m_j^2}\,\mathrm{Re}\,\Pi_{ij}(m_j^2) \qquad i \neq j\,, \qquad (3.28)$$

$$\delta Z_{ii} \;=\; \delta Z^*_{ii} = -\mathrm{Re}\,\dot{\Pi}_{ii}(m_i^2)\,. \qquad (3.29)$$

3.3.2 Fermionic particles with mixing

Like in the previous chapter we have the same structure for the physical as well as for the bare Lagrangian:

$$\mathcal{L} \;=\; \bar{\psi}_j\, \delta_{ij}\, (i\,\partial\!\!\!/ - m_i)\, \psi_i\,, \qquad (3.30)$$

$$\mathcal{L}_0 \;=\; \bar{\psi}_{0,j}\, \delta_{ij}\, (i\,\partial\!\!\!/ - m_{0,i})\, \psi_{0,i}\,. \qquad (3.31)$$

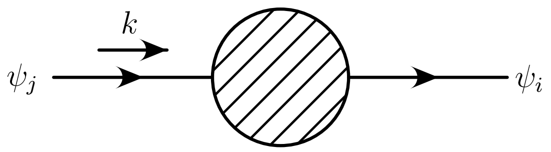
The relation between the unrenormalized and the renormalized quantities is given by attaching multiplicative renormalization constants to the unrenormalized fermion fields $\psi_{0,i}$ and the mass parameter $m_{0,i}$, i.e.

$$\psi_{0,j} \;=\; \left(\delta_{jk} + \tfrac{1}{2}\delta Z^L_{jk} P_L + \tfrac{1}{2}\delta Z^R_{jk} P_R\right)\psi_k\,, \qquad (3.32)$$

$$\bar{\psi}_{0,i} \;=\; \bar{\psi}_l\left(\delta_{il} + \tfrac{1}{2}\delta Z^{R\dagger}_{il} P_L + \tfrac{1}{2}\delta Z^{L\dagger}_{il} P_R\right), \qquad (3.33)$$

$$m_{0,i} \;=\; m_i + \delta m_i\,, \qquad (3.34)$$

where the 'dagger' † in $\delta Z^{L,R\dagger}_{il}$ indicates hermitian conjugation with regard to the spinor indices. For the renormalized one particle irreducible (1PI) two–point–function



$$\mathcal{M} = i\,\bar{u}_i(k)\hat{\Gamma}_{ij}(k)u_j(k)$$

$$\hat{\Gamma}_{ij}(k) = \delta_{ij}(k\!\!\!/ - m_i) + \hat{\Pi}_{ij}(k)$$

with the renormalized self–energy

$$\hat{\Pi}_{ij}(k) \;=\; k\!\!\!/\, P_L\, \hat{\Pi}^L_{ij}(k) + k\!\!\!/\, P_R\, \hat{\Pi}^R_{ij}(k) + \hat{\Pi}^{S,L}_{ij}(k) P_L + \hat{\Pi}^{S,R}_{ij}(k) P_R \qquad (3.35)$$

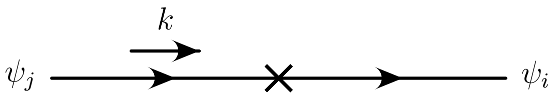
we require the on–shell renormalization conditions

$$\mathrm{Re}\,\hat{\Gamma}_{ij}(k)\, u_j(k)\Big|_{k^2=m_j^2} \;=\; 0\,, \qquad \lim_{k^2 \to m_i^2} \frac{1}{k\!\!\!/ - m_i}\,\mathrm{Re}\,\hat{\Gamma}_{ii}(k)u_i(k) \;=\; u_i(k)\,. \qquad (3.36)$$

Inserting the counter–term Lagrangian

$$\delta\mathcal{L} \;=\; \bar{\psi}_i\left(k\!\!\!/ P_L C^L_{ij} + k\!\!\!/ P_R C^R_{ij} - C^{S,L}_{ij} P_L - C^{S,R}_{ij} P_R\right)\psi_j\,, \qquad (3.37)$$

into $\mathcal{L}_0 = \mathcal{L} + \delta\mathcal{L}$,



$$\mathcal{M} = i\big(\not{k} P_L C^L_{ij} + \not{k} P_R C^R_{ij} - C^{S,L}_{ij} P_L - C^{S,R}_{ij} P_R \big)\,,$$

we get for the coefficients $C^{\cdots}_{ij}$

$$C^L_{ij} = \tfrac{1}{2}\big(\delta Z^L_{ij} + \delta Z^{L\dagger}_{ji}\big)\,, \tag{3.38}$$

$$C^R_{ij} = \tfrac{1}{2}\big(\delta Z^R_{ij} + \delta Z^{R\dagger}_{ji}\big)\,, \tag{3.39}$$

$$C^{S,L}_{ij} = \tfrac{1}{2}\big(m_i\,\delta Z^L_{ij} + m_j\,\delta Z^{R\dagger}_{ji}\big) + \delta_{ij}\,\delta m_i\,, \tag{3.40}$$

$$C^{S,R}_{ij} = \tfrac{1}{2}\big(m_i\,\delta Z^R_{ij} + m_j\,\delta Z^{L\dagger}_{ji}\big) + \delta_{ij}\,\delta m_i\,. \tag{3.41}$$

Thus the renormalized self–energies can be written as

$$\hat{\Pi}^L_{ij} = \Pi^L_{ij} + \tfrac{1}{2}\big(\delta Z^L_{ij} + \delta Z^{L\dagger}_{ji}\big)\,, \tag{3.42}$$

$$\hat{\Pi}^R_{ij} = \Pi^R_{ij} + \tfrac{1}{2}\big(\delta Z^R_{ij} + \delta Z^{R\dagger}_{ji}\big)\,, \tag{3.43}$$

$$\hat{\Pi}^{S,L}_{ij} = \Pi^{S,L}_{ij} - \tfrac{1}{2}\big(m_i\,\delta Z^L_{ij} + m_j\,\delta Z^{R\dagger}_{ji}\big) - \delta_{ij}\,\delta m_i\,, \tag{3.44}$$

$$\hat{\Pi}^{S,R}_{ij} = \Pi^{S,R}_{ij} - \tfrac{1}{2}\big(m_i\,\delta Z^R_{ij} + m_j\,\delta Z^{L\dagger}_{ji}\big) - \delta_{ij}\,\delta m_i\,. \tag{3.45}$$

Taking the renormalization conditions in eq. (3.36) into account, we obtain the counter terms for the mass parameter and the wave–function corrections

$$\delta m_i = \frac{1}{2}\,\mathrm{Re}\Big[m_i\Big(\Pi^L_{ii}(m_i) + \Pi^R_{ii}(m_i)\Big) + \Pi^{S,L}_{ii}(m_i) + \Pi^{S,R}_{ii}(m_i)\Big]\,. \tag{3.46}$$

$$\delta Z^L_{ij} = \frac{2}{m_i^2 - m_j^2}\,\mathrm{Re}\Big[m_j^2\,\Pi^L_{ij}(m_j) + m_i m_j\,\Pi^R_{ij}(m_j) + m_i\,\Pi^{S,L}_{ij} + m_j\,\Pi^{S,R}_{ij}\Big]\,, \tag{3.47}$$

$$\begin{aligned}\delta Z^L_{ii} = {} & -\Pi^L_{ii}(m_i) + \frac{1}{2m_i}\Big[\Pi^{S,L}_{ii}(m_i) - \Pi^{S,R}_{ii}(m_i)\Big] \\ & - m_i\frac{\partial}{\partial k^2}\Big[m_i\big(\Pi^L_{ii}(k) + \Pi^R_{ii}(k)\big) + \Pi^{S,L}_{ii}(k) + \Pi^{S,R}_{ii}(k)\Big]\Big|_{k^2=m_i^2}\,,\end{aligned} \tag{3.48}$$

and the corresponding right–handed terms, $\delta Z^{(S),R}_{ij} = \delta Z^{(S),L}_{ij}(L \leftrightarrow R)$.

3.4 Sfermion sector

According to the results of section 3.3, where we have derived the mass corrections and the wave–function renormalization constants in terms of self–energies, these counter terms are given in the sfermion sector by

$$\delta m^2_{\tilde{f}_i} = \mathrm{Re}\,\Pi^{\tilde{f}}_{ii}(m^2_{\tilde{f}_i}) \tag{3.49}$$

and

$$\delta Z_{ii}^{\tilde{f}} = -\operatorname{Re}\dot{\Pi}_{ii}^{\tilde{f}}(m_{\tilde{f}_i}^2)\,, \qquad i=(1,2)\,, \tag{3.50}$$

$$\delta Z_{ij}^{\tilde{f}} = \frac{2}{m_{\tilde{f}_i}^2 - m_{\tilde{f}_j}^2}\operatorname{Re}\Pi_{ij}^{\tilde{f}}(m_{\tilde{f}_j}^2)\,, \qquad i \neq j\,. \tag{3.51}$$

In order to renormalize the parameters entering in the sfermion mass matrix $\mathcal{M}_{\tilde{f}}^2$ (see eq. (2.32)), we have to look at its counter term $\delta\mathcal{M}_{\tilde{f}}^2$,

$$\delta\mathcal{M}_{\tilde{f}}^2 = \left(\delta R^{\tilde{f}}\right)^\dagger \begin{pmatrix} m_{\tilde{f}_1}^2 & 0 \\ 0 & m_{\tilde{f}_2}^2 \end{pmatrix} R^{\tilde{f}} + \left(R^{\tilde{f}}\right)^\dagger \begin{pmatrix} \delta m_{\tilde{f}_1}^2 & 0 \\ 0 & \delta m_{\tilde{f}_2}^2 \end{pmatrix} R^{\tilde{f}} + \left(R^{\tilde{f}}\right)^\dagger \begin{pmatrix} m_{\tilde{f}_1}^2 & 0 \\ 0 & m_{\tilde{f}_2}^2 \end{pmatrix} \delta R^{\tilde{f}}. \tag{3.52}$$

with

$$\delta R^{\tilde{f}} = -\begin{pmatrix} \sin\theta_{\tilde{f}} & -\cos\theta_{\tilde{f}} \\ \cos\theta_{\tilde{f}} & \sin\theta_{\tilde{f}} \end{pmatrix}\delta\theta_{\tilde{f}}\,. \tag{3.53}$$

The renormalization constant of the rotation matrix $R_{ij}^{\tilde{f}}$ is determined such as to cancel the anti–hermitian part of the sfermion wave–function corrections,

$$\delta R_{ij}^{\tilde{f}} = \sum_{k=1}^{2}\frac{1}{4}\left(\delta Z_{ik}^{\tilde{f}} - \delta Z_{ki}^{\tilde{f}}\right)R_{kj}^{\tilde{f}}\,. \tag{3.54}$$

Therefore, the counter term for the sfermion mixing angle $\theta_{\tilde{f}}$ is given by [22, 23]

$$\delta\theta_{\tilde{f}} = \frac{1}{4}\left(\delta Z_{12}^{\tilde{f}} - \delta Z_{21}^{\tilde{f}}\right) = \frac{1}{2\left(m_{\tilde{f}_1}^2 - m_{\tilde{f}_2}^2\right)}\operatorname{Re}\left(\Pi_{12}^{\tilde{f}}(m_{\tilde{f}_2}^2) + \Pi_{21}^{\tilde{f}}(m_{\tilde{f}_1}^2)\right)\,, \tag{3.55}$$

and thus

$$\delta\left(\mathcal{M}_{\tilde{f}}^2\right)_{ij} = \frac{1}{2}\sum_{k,l=1}^{2}\left(R_{ik}^{\tilde{f}}\right)^\dagger \operatorname{Re}\left[\Pi_{kl}^{\tilde{f}}(m_{\tilde{f}_l}^2) + \Pi_{lk}^{\tilde{f}}(m_{\tilde{f}_k}^2)\right]R_{lj}^{\tilde{f}}\,. \tag{3.56}$$

The counter terms of the remaining free parameters of the sfermion mass matrix are determined, if possible, through their tree–level relations in order to absorb all corrections from other parameters in the considered matrix element [24, 25]. In this way we get the counter terms for the trilinear couplings (see eq. (2.39))

$$\delta A_f = -\left(A_f - \mu\,(\tan\beta)^{-2I_f^{3L}}\right)\frac{\delta m_f}{m_f} + \left(\frac{\delta\mu}{\mu} - 2I_f^{3L}\,\frac{\delta\tan\beta}{\tan\beta}\right)\mu\,(\tan\beta)^{-2I_f^{3L}} + \frac{\delta m_{LR}^2}{m_f} \tag{3.57}$$

with $\delta m^2_{LR} = \frac{1}{2}(\delta m^2_{\tilde{f}_1} - \delta m^2_{\tilde{f}_2}) \sin 2\theta_{\tilde{f}} + (m^2_{\tilde{f}_1} - m^2_{\tilde{f}_2}) \cos 2\theta_{\tilde{f}}\, \delta\theta_{\tilde{f}}$ as well as the renormalization constants for the soft SUSY–breaking masses $M_{\tilde{U},\tilde{D},\tilde{E}}$

$$\begin{aligned}\delta M^2_{\tilde{U},\tilde{D},\tilde{E}}(\tilde{f}) &= \delta m^2_{\tilde{f}_1} \sin^2\theta_{\tilde{f}} + \delta m^2_{\tilde{f}_2} \cos^2\theta_{\tilde{f}} + (m^2_{\tilde{f}_1} - m^2_{\tilde{f}_2}) \sin 2\theta_{\tilde{f}}\, \delta\theta_{\tilde{f}} - 2 m_f\, \delta m_f \\ &\quad - \delta m_Z^2 \cos 2\beta\, e_f s_W^2 + 2 m_Z^2 \sin 2\beta\, \delta\beta\, e_f s_W^2 - m_Z^2 \cos 2\beta\, e_f \delta s_W^2 \,. \end{aligned} \tag{3.58}$$

However, in the on–shell scheme, there is a subtlety concerning the soft SUSY–breaking parameters in the 11–element of $\mathcal{M}_{\tilde{f}}$, i.e. $M_{\tilde{Q}}$ and $M_{\tilde{L}}$, where we additionally have to include UV–finite shifts [24, 25]. At the $\overline{\mathrm{DR}}$ scale the soft SUSY–breaking parameters $M_{\tilde{Q},\tilde{L}}$ enter both in the up– and down–sfermion sector. At one–loop level, however, $M_{\tilde{Q},\tilde{L}}$ obtain different shifts in the up– and down–sfermion sector. In this work, we define $M^2_{\tilde{Q},\tilde{L}} \equiv M^2_{\tilde{Q},\tilde{L}}(\tilde{d}) = m^2_{\tilde{f}_1} \cos^2\theta_{\tilde{f}} + m^2_{\tilde{f}_2} \sin^2\theta_{\tilde{f}} - m_f^2 - m_Z^2 \cos 2\beta (I_f^{3L} - e_f s_W^2)$ to be the on–shell parameter in the down–sfermion sector ($\tilde{d} = \tilde{b}, \tilde{\tau}$), therefore we get

$$M^2_{\tilde{Q},\tilde{L}}(\tilde{u}) = M^2_{\tilde{Q},\tilde{L}}(\tilde{d}) + \delta M^2_{\tilde{Q},\tilde{L}}(\tilde{d}) - \delta M^2_{\tilde{Q},\tilde{L}}(\tilde{u}) \tag{3.59}$$

with

$$\begin{aligned}\delta M^2_{\tilde{Q},\tilde{L}}(\tilde{f}) &= \delta m^2_{\tilde{f}_1} \cos^2\theta_{\tilde{f}} + \delta m^2_{\tilde{f}_2} \sin^2\theta_{\tilde{f}} - (m^2_{\tilde{f}_1} - m^2_{\tilde{f}_2}) \sin 2\theta_{\tilde{f}}\, \delta\theta_{\tilde{f}} - 2 m_f\, \delta m_f \\ &\quad - \delta m_Z^2 \cos 2\beta\, (I_f^{3L} - e_f s_W^2) + 2 m_Z^2 \sin 2\beta\, \delta\beta\, (I_f^{3L} - e_f s_W^2) \\ &\quad + m_Z^2 \cos 2\beta\, e_f \delta s_W^2 \,. \end{aligned} \tag{3.60}$$

3.5 Higgs sector

The renormalization of the Higgs mixing angle α is treated in a similar way as the sfermion mixing angle $\theta_{\tilde{f}}$. Consider the mass matrix of the CP–even Higgs bosons h^0 and H^0 (cf. section 2.2.1),

$$\begin{aligned} M^2(H^0, h^0) &= \begin{pmatrix} \sin^2\beta\, m^2_{A^0} + \cos^2\beta\, m_Z^2 & -\sin\beta\cos\beta (m^2_{A^0} + m_Z^2) \\ -\sin\beta\cos\beta (m^2_{A^0} + m_Z^2) & \cos^2\beta\, m^2_{A^0} + \sin^2\beta\, m_Z^2 \end{pmatrix} \\ &= \left(R^{H^0}\right)^T \cdot \begin{pmatrix} m^2_{H^0} & 0 \\ 0 & m^2_{h^0} \end{pmatrix} \cdot R^{H^0} \,, \qquad m_{h^0} < m_{H^0} \end{aligned} \tag{3.61}$$

with the rotation matrix

$$R^{H^0}_{ij} \equiv R_{ij}(\alpha) = \begin{pmatrix} \cos\alpha & \sin\alpha \\ -\sin\alpha & \cos\alpha \end{pmatrix}_{ij} \,. \tag{3.62}$$

Analogously to the case of the sfermion mixing angle, the renormalization constant of the rotation matrix $R^{H^0}_{ij}$ is determined such as to cancel the anti–hermitian part of the Higgs wave–function corrections,

$$\delta R^{H^0}_{ij} = \sum_{k=1}^{2} \frac{1}{4} \left(\delta Z^{H^0}_{ki} - \delta Z^{H^0}_{ik}\right) R^{H^0}_{kj} \,, \tag{3.63}$$

which leads to the counter term

$$\delta\alpha \;=\; \frac{1}{4}\left(\delta Z^H_{21} - \delta Z^H_{12}\right) = \frac{1}{2\left(m^2_{H^0} - m^2_{h^0}\right)}\,\mathrm{Re}\left(\Pi^H_{12}(m^2_{H^0}) + \Pi^H_{21}(m^2_{h^0})\right). \qquad (3.64)$$

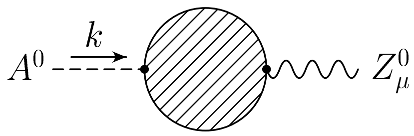
Note that the indices the wave–function renormalization constants $\delta Z^{H^0}_{ij}$ are interchanged due to the conventional nomenclature labelling the light Higgs boson by an index 1 and the heavy one by an index 2.

As stated in section 2.2.1, the parameter $\tan\beta$ plays a central role in the MSSM. Due to its close connection to spontaneous symmetry breaking, it enters in almost all sectors of the MSSM and, as a consequence, has a major effect on most MSSM observables. Though not directly connected to measurable quantities, the applied renormalization scheme determines its physical meaning and its relation to observables. In fact, several renormalization schemes for $\tan\beta$ suffer from specific disadvantages, leading either to gauge dependences or numerical instabilities [26]. However, in this thesis we apply the renormalization of $\tan\beta$ proposed by [20, 21]. The mixing angle β is fixed by the condition that the renormalized A^0–Z^0 transition vanishes at $p^2 = m^2_{A^0}$, i.e.

$$\mathrm{Im}\hat{\Pi}_{A^0Z^0}(m^2_{A^0}) \;=\; 0\,, \qquad (3.65)$$

which gives the counter term

$$\frac{\delta\tan\beta}{\tan\beta} \;=\; \frac{1}{m_Z \sin 2\beta}\,\mathrm{Im}\Pi_{A^0Z^0}(m^2_{A^0}). \qquad (3.66)$$



$$\mathcal{M} = -i\, k^\mu\, \Pi_{AZ}(k^2)\, \epsilon^*_\mu(k)$$

Figure 3.1: A^0Z^0 mixing self–energy relevant for the renormalization of $\tan\beta$.

Chapter 4

Higgs decays into sfermions in the MSSM

4.1 Introduction

The search for a Higgs boson is the primary goal of all present and future high energy experiments at the TEVATRON, LHC or an e^+e^- Linear Collider. Whereas the Standard Model (SM) predicts just one Higgs boson, with the present lower bound of its mass $m_H \geq 114.4$ GeV (at 95% confidence level) [27], extensions of the SM allow for more Higgs bosons. In particular, the Minimal Supersymmetric Standard Model (MSSM) contains five physical Higgs bosons: two neutral CP–even (h^0 and H^0), one neutral CP–odd (A^0), and two charged ones ($H^\pm$) [28, 29]. The existence of a charged Higgs boson or a CP–odd neutral one would give clear evidence for physics beyond the SM. For a discovery precise predictions for their decay modes and branching ratios are necessary. In case supersymmetric (SUSY) particles are not too heavy, the Higgs bosons can also decay into SUSY particles (neutralinos $\tilde{\chi}_i^0$, charginos $\tilde{\chi}_k^\pm$, sfermions $\tilde{f}_m$), $H^0, A^0 \to \tilde{\chi}_i^0 \tilde{\chi}_j^0$ $(i, j = 1 \ldots 4)$, $H^0, A^0 \to \tilde{\chi}_k^+ \tilde{\chi}_l^-$ $(k, l = 1, 2)$, $H^0, A^0 \to \tilde{f}_m \bar{\tilde{f}}_n$ $(m, n = 1, 2)$, $H^\pm \to \tilde{\chi}_k^\pm \tilde{\chi}_i^0$, $H^\pm \to \tilde{f}_m \bar{\tilde{f}}'_n$. At tree level, these decays were studied in [30, 31]. In particular, the branching ratios for the decays into sfermions, $H^0, A^0 \to \tilde{f}_m \bar{\tilde{f}}_n$, can be sizeable depending on the parameter space [32, 33]. The SUSY–QCD corrections to the decays into sfermions have also been calculated [34, 35].

In this thesis, we study in detail the decays of Higgs bosons into two sfermions, $\{h^0, H^0, A^0\} \to \tilde{f}_i \bar{\tilde{f}}_j$ and $H^+ \to \tilde{t}_i \bar{\tilde{b}}_j$ as well as the crossed channels $\tilde{f}_2 \to \tilde{f}_1 H_k^0$. In particular, the third generation sfermions $\tilde{t}_i, \tilde{b}_i$, and $\tilde{\tau}_i$ are interesting because one expects them to be lighter than the other sfermions due to their large Yukawa couplings and left–right mixings. Since A^0 only couples to $\tilde{f}_L$–$\tilde{f}_R$ (left–right states of $\tilde{f}$), and due to the CP nature of A^0, $A^0 \to \tilde{f}_i \bar{\tilde{f}}_i$ vanishes. (This is valid also beyond tree level for real parameters in the MSSM.) We will calculate the *full* electroweak one–loop corrections in the on–shell scheme. Owing to the fact that almost all parameters of the MSSM have to be renormalized in this process and hence a large number of graphs has to be computed, the calculation is very complex. Despite this complexity, we have performed the calculation in an analytic way.

4.2 Tree level

First, we review the tree–level results [32]. For the neutral Higgs fields we use the notation $H_k^0 = \{h^0, H^0, A^0, G^0\}$, $t/\tilde{t}$ stands for an up–type (s)fermion and $b/\tilde{b}$ for a down–type one. Following [28, 29] the Higgs–Sfermion–Sfermion couplings for neutral Higgs bosons, $G_{ijk}^{\tilde{f}}$, can be written as

$$G_{ijk}^{\tilde{f}} \equiv G\big(H_k^0 \tilde{f}_i^* \tilde{f}_j\big) = \Big[R^{\tilde{f}} G_{LR,k}^{\tilde{f}} (R^{\tilde{f}})^T\Big]_{ij} . \tag{4.1}$$

The 3rd generation left–right couplings $G_{LR,k}^{\tilde{f}}$ for up– and down–type sfermions are given by

$$G_{LR,1}^{\tilde{t}} = \begin{pmatrix} -\sqrt{2}h_t m_t c_\alpha + g_Z m_Z (I_t^{3L} - e_t s_W^2) s_{\alpha+\beta} & -\frac{h_t}{\sqrt{2}}(A_t\, c_\alpha + \mu s_\alpha) \\ -\frac{h_t}{\sqrt{2}}(A_t\, c_\alpha + \mu s_\alpha) & -\sqrt{2}h_t m_t c_\alpha + g_Z m_Z e_t s_W^2 s_{\alpha+\beta} \end{pmatrix} , \tag{4.2}$$

$$G_{LR,1}^{\tilde{b}} = \begin{pmatrix} \sqrt{2}h_b m_b s_\alpha + g_Z m_Z (I_b^{3L} - e_b s_W^2) s_{\alpha+\beta} & \frac{h_b}{\sqrt{2}}(A_b\, s_\alpha + \mu c_\alpha) \\ \frac{h_b}{\sqrt{2}}(A_b\, s_\alpha + \mu c_\alpha) & \sqrt{2}h_b m_b s_\alpha + g_Z m_Z e_b s_W^2 s_{\alpha+\beta} \end{pmatrix} , \tag{4.3}$$

$$G_{LR,2}^{\tilde{f}} = G_{LR,1}^{\tilde{f}} \quad \text{with } \alpha \to \alpha - \pi/2 , \tag{4.4}$$

$$G_{LR,3}^{\tilde{t}} = -\sqrt{2}h_t \begin{pmatrix} 0 & -\frac{i}{2}\Big(A_t\, c_\beta + \mu\, s_\beta\Big) \\ \frac{i}{2}\Big(A_t\, c_\beta + \mu\, s_\beta\Big) & 0 \end{pmatrix} , \tag{4.5}$$

$$G_{LR,3}^{\tilde{b}} = -\sqrt{2}h_b \begin{pmatrix} 0 & -\frac{i}{2}\Big(A_b\, s_\beta + \mu\, c_\beta\Big) \\ \frac{i}{2}\Big(A_b\, s_\beta + \mu\, c_\beta\Big) & 0 \end{pmatrix} , \tag{4.6}$$

where we have used the abbreviations $s_x \equiv \sin x$, $c_x \equiv \cos x$ and $s_W \equiv \sin\theta_W$. α denotes the mixing angle of the $\{h^0, H^0\}$–system, and h_t and h_b are the Yukawa couplings

$$h_t = \frac{g\, m_t}{\sqrt{2} m_W \sin\beta} , \qquad h_b = \frac{g\, m_b}{\sqrt{2} m_W \cos\beta} . \tag{4.7}$$

The couplings of the charged Higgs boson to two sfermions are given by

$$G_{ij1}^{\tilde{t}\tilde{b}} \equiv G\big(H^+ \tilde{t}_i^* \tilde{b}_j\big) = G_{ji1}^{\tilde{b}\tilde{t}} = \Big[R^{\tilde{t}}\, G_{LR,1}^{\tilde{t}\tilde{b}} (R^{\tilde{b}})^T\Big]_{ij} , \tag{4.8}$$

$$G_{LR,1}^{\tilde{t}\tilde{b}} = \begin{pmatrix} h_b m_b \sin\beta + h_t m_t \cos\beta - \frac{g m_W}{\sqrt{2}} \sin 2\beta & h_b (A_b \sin\beta + \mu\cos\beta) \\ h_t (A_t \cos\beta + \mu \sin\beta) & h_t m_b \cos\beta + h_b m_t \sin\beta \end{pmatrix} , \tag{4.9}$$

$$G_{LR,1}^{\tilde{b}\tilde{t}} = \begin{pmatrix} h_b m_b \sin\beta + h_t m_t \cos\beta - \frac{g m_W}{\sqrt{2}} \sin 2\beta & h_t (A_t \cos\beta + \mu\sin\beta) \\ h_b (A_b \sin\beta + \mu \cos\beta) & h_t m_b \cos\beta + h_b m_t \sin\beta \end{pmatrix} . \tag{4.10}$$

Starting with the tree–level interaction Lagrangian for the three neutral Higgs bosons,

$$\mathcal{L} = G_{ijk}^{\tilde{f}}\, H_k^0\, \tilde{f}_i^* \tilde{f}_j , \qquad (k = 1, 2, 3) , \tag{4.11}$$

the Feynman amplitude is simply given by

$$\mathcal{M}^{\rm tree} \;=\; iG^{\tilde f}_{ijk}\,. \tag{4.12}$$

The tree–level decay width for the process $H^0_k(p) \to \tilde f_i(k_1) + \bar{\tilde f}_j(k_2)$ then becomes

$$\begin{aligned}\Gamma^{\rm tree}(H^0_k \to \tilde f_i \bar{\tilde f}_j) &= \frac{N_C}{2m_{H^0_k}}\int\frac{d^3k_1}{(2\pi)^3}\frac{1}{2E_1}\int\frac{d^3k_2}{(2\pi)^3}\frac{1}{2E_2}\,(2\pi)^4\delta^4(p-k_1-k_2)|G^{\tilde f}_{ijk}|^2\\ &= \frac{N_C\,\kappa(m^2_{H^0_k}, m^2_{\tilde f_i}, m^2_{\tilde f_j})}{16\pi m^3_{H^0_k}}|G^{\tilde f}_{ijk}|^2 \end{aligned} \tag{4.13}$$

with the totally symmetric *Källen function* $\kappa(x,y,z) = \sqrt{(x-y-z)^2 - 4yz}$ and the colour factor $N^f_C = 3$ for squarks and $N^f_C = 1$ for sleptons, respectively.
Analogously, the decay width for the charged Higgs boson H^+ is given by

$$\Gamma^{\rm tree}(H^+ \to \tilde t_i\,\bar{\tilde b}_j) \;=\; \frac{N_C\,\kappa(m^2_{H^+}, m^2_{\tilde t_i}, m^2_{\tilde b_j})}{16\pi m^3_{H^+}}|G^{\tilde t\tilde b}_{ij1}|^2\,. \tag{4.14}$$

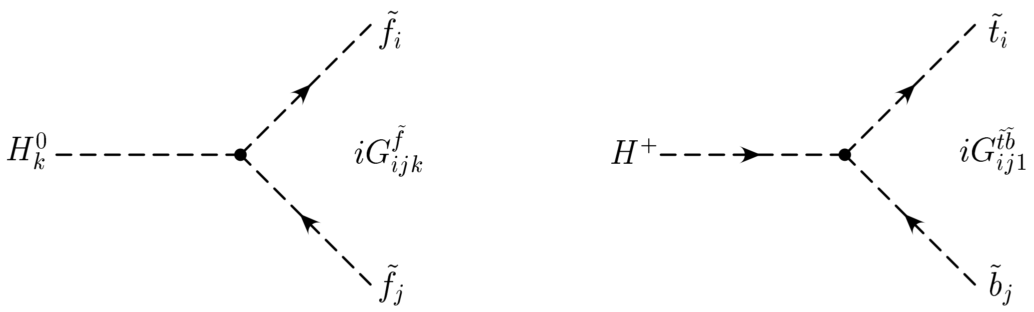


Figure 4.1: Tree level diagrams for $H^0_k \to \tilde f_i\,\bar{\tilde f}_j$ and $H^+ \to \tilde t_i\,\bar{\tilde b}_j$.

4.3 One–loop Corrections

Following the recipe of multiplicative renormalization, we replace the bare parameters and fields in the Lagrangian by the corresponding renormalized ones, i.e. we attach renormalization constants to each coupling and field. In case of the Higgs–Sfermion–Sfermion interaction Lagrangian

$$\mathcal{L}_0 \;=\; \big(G^{\tilde f}_{ijk}\big)^0\big(H^0_k\big)^0\tilde f^{0*}_i\tilde f^0_j + \big(G^{\tilde t\tilde b}_{ij1}\big)^0\big(H^+\big)^0\tilde t^{0*}_i\tilde b^0_j \tag{4.15}$$

the relations between the unrenormalized (bare) and renormalized (physical) fields and couplings are

$$\big(G^{\tilde{f}}_{ijk}\big)^0 = G^{\tilde{f}}_{ijk} + \delta G^{\tilde{f}(c)}_{ijk}\,, \qquad \big(G^{\tilde{t}\tilde{b}}_{ij1}\big)^0 = G^{\tilde{t}\tilde{b}}_{ij1} + \delta G^{\tilde{t}\tilde{b}(c)}_{ij1}\,, \tag{4.16}$$

$$\big(H^0_k\big)^0 = \big(\delta_{kl} + \frac{1}{2}\,\delta Z^{H^0}_{kl}\big) H^0_l\,, \qquad \big(H^+\big)^0 = \big(\delta_{1l} + \frac{1}{2}\,\delta Z^{H^+}_{1l}\big) H^+_l\,, \tag{4.17}$$

$$\tilde{f}^{0*}_i = \big(\delta_{ii'} + \frac{1}{2}\,\delta Z^{\tilde{f}}_{ii'}\big) \tilde{f}^*_{i'}\,, \qquad \tilde{t}^{0*}_i = \big(\delta_{ii'} + \frac{1}{2}\,\delta Z^{\tilde{t}}_{ii'}\big) \tilde{t}^*_{i'}\,, \tag{4.18}$$

$$\tilde{f}^0_j = \big(\delta_{jj'} + \frac{1}{2}\,\delta Z^{\tilde{f}}_{jj'}\big) \tilde{f}_{j'}\,, \qquad \tilde{b}^0_j = \big(\delta_{jj'} + \frac{1}{2}\,\delta Z^{\tilde{b}}_{jj'}\big) \tilde{b}_{j'}\,, \tag{4.19}$$

where we have used the notation $H^+_l = \{H^+, G^+\}$ for the charged Higgs and Goldstone boson, respectively. Note that due to the CP properties of the neutral Higgs bosons, the CP–even Higgs bosons (h^0 and H^0) don't mix with the CP–odd ones (A^0 and G^0), i.e. $\delta Z^{H^0}_{kl} = 0$ for $k = (1,2)$ and $l = (3,4)$. The bare Lagrangian $\mathcal{L}_0$ can then be written as a sum of the renormalized Lagrangian $\mathcal{L}^{\rm ren}$ plus its counter terms $\delta\mathcal{L}$,

$$\mathcal{L}_0 = \mathcal{L}^{\rm ren} + \delta\mathcal{L}\,, \tag{4.20}$$

$$\mathcal{L}^{\rm ren} = G^{\tilde{f}}_{ijk}\, H^0_k\, \tilde{f}^*_i \tilde{f}_j + G^{\tilde{t}\tilde{b}}_{ij1} H^+ \tilde{t}^*_i \tilde{b}_j\,, \tag{4.21}$$

$$\delta\mathcal{L} = -\delta G^{\tilde{f}(v)}_{ijk} H^0_k \tilde{f}^*_i \tilde{f}_j - \delta G^{\tilde{t}\tilde{b}(v)}_{ij1} H^+ \tilde{t}^*_i \tilde{b}_j\,, \tag{4.22}$$

which leads to the one–loop corrected (renormalized) couplings $G^{\tilde{f},\rm ren}_{ijk}$ and $G^{\tilde{t}\tilde{b},\rm ren}_{ij1}$:

$$G^{\tilde{f},\rm ren}_{ijk} = G^{\tilde{f}}_{ijk} + \Delta G^{\tilde{f}}_{ijk} = G^{\tilde{f}}_{ijk} + \delta G^{\tilde{f}(v)}_{ijk} + \delta G^{\tilde{f}(w)}_{ijk} + \delta G^{\tilde{f}(c)}_{ijk}\,, \tag{4.23}$$

$$G^{\tilde{t}\tilde{b},\rm ren}_{ij1} = G^{\tilde{t}\tilde{b}}_{ij1} + \Delta G^{\tilde{t}\tilde{b}}_{ij1} = G^{\tilde{t}\tilde{b}}_{ij1} + \delta G^{\tilde{t}\tilde{b}(v)}_{ij1} + \delta G^{\tilde{t}\tilde{b}(w)}_{ij1} + \delta G^{\tilde{t}\tilde{b}(c)}_{ij1}\,. \tag{4.24}$$

$\delta G^{\tilde{f}(v)}_{ijk}, \delta G^{\tilde{f}(w)}_{ijk}$ and $\delta G^{\tilde{f}(c)}_{ijk}$ and the corresponding terms for the couplings to the charged Higgs boson stand for the vertex corrections, the wave–function corrections and the coupling counter–term corrections, respectively. The full one–loop corrected decay widths for the neutral as well as the charged Higgs boson decays are then given by

$$\Gamma(H^0_k \to \tilde{f}_i \bar{\tilde{f}}_j) = \frac{N_C\,\kappa(m^2_{H^0_k}, m^2_{\tilde{f}_i}, m^2_{\tilde{f}_j})}{16\pi m^3_{H^0_k}} \Big[|G^{\tilde{f}}_{ijk}|^2 + 2{\rm Re}\big(G^{\tilde{f}}_{ijk}\cdot \Delta G^{\tilde{f}}_{ijk}\big)\Big]\,, \tag{4.25}$$

$$\Gamma(H^+ \to \tilde{t}_i\, \bar{\tilde{b}}_j) = \frac{N_C\,\kappa(m^2_{H^+}, m^2_{\tilde{t}_i}, m^2_{\tilde{b}_j})}{16\pi m^3_{H^+}} \Big[|G^{\tilde{t}\tilde{b}}_{ij1}|^2 + 2{\rm Re}\big(G^{\tilde{t}\tilde{b}}_{ij1}\cdot \Delta G^{\tilde{t}\tilde{b}}_{ij1}\big)\Big]\,. \tag{4.26}$$

Due to the lengthy formulae, we give the explicit form of the vertex corrections, $\delta G^{\tilde{f}(v)}_{ijk}$ and $\delta G^{\tilde{t}\tilde{b}(v)}_{ij1}$, in Appendix C.1.

For the CP–even Higgs bosons the wave–function corrections $\delta G^{\tilde{f}(w)}_{ijk}$ can be written as

$$\delta G^{\tilde{f}(w)}_{ijk} = \frac{1}{2}\Big[\delta Z^{\tilde{f}}_{i'i}\, G^{\tilde{f}}_{i'jk} + \delta Z^{\tilde{f}}_{j'j}\, G^{\tilde{f}}_{ij'k} + \delta Z^{H^0}_{lk}\, G^{\tilde{f}}_{ijl}\Big]\,, \tag{4.27}$$

with the implicit summations $i', j', l = (1,2)$. The wave–function renormalization constants are determined by imposing the on–shell renormalization conditions (see chapter 3),

$$\delta Z^{\tilde{f}}_{ii} = -\,\mathrm{Re}\,\dot{\Pi}^{\tilde{f}}_{ii}(m^2_{\tilde{f}_i})\,, \qquad i = (1,2)\,, \tag{4.28}$$

$$\delta Z^{\tilde{f}}_{ij} = \frac{2}{m^2_{\tilde{f}_i} - m^2_{\tilde{f}_j}}\,\mathrm{Re}\,\Pi^{\tilde{f}}_{ij}(m^2_{\tilde{f}_j})\,, \qquad i,j = (1,2),\ i \neq j,\ \tilde{f} \neq \tilde{\nu}_{e,\mu,\tau}\,, \tag{4.29}$$

$$\delta Z^{H^0}_{kk} = -\,\mathrm{Re}\,\dot{\Pi}^{H^0}_{kk}(m^2_{h^0_k})\,, \qquad k = 1,2\,, \tag{4.30}$$

$$\delta Z^{H^0}_{kl} = \frac{2}{m^2_{h^0_k} - m^2_{h^0_l}}\,\mathrm{Re}\,\Pi^{H^0}_{kl}(m^2_{h^0_l})\,, \qquad k,l = (1,2),\ k \neq l\,. \tag{4.31}$$

The explicit forms of the off–diagonal Higgs boson and sfermion self–energies, $\Pi^{H^0}_{kl}$ and $\Pi^{\tilde{f}}_{ij}$, as well as the derivatives of the diagonal ones, $\dot{\Pi}^{H^0}_{kk}$ and $\dot{\Pi}^{\tilde{f}}_{ii}$, are given in Appendix B.

Due to its CP–nature, the Higgs boson A^0 cannot only mix with its associated partner in the mass matrix, the neutral Goldstone boson G^0, but also with the weak vector boson Z^0. Therefore, we split the wave–function corrections into the diagonal ones,

$$\delta G^{\tilde{f}(w,\mathrm{diag})}_{123} = \frac{1}{2}\mathrm{Re}\Big[\delta Z^{\tilde{f}}_{11} + \delta Z^{\tilde{f}}_{22} + \delta Z^{H^0}_{33}\Big] G^{\tilde{f}}_{123}\,, \tag{4.32}$$

$$\delta G^{\tilde{f}(w)}_{123} = \delta G^{\tilde{f}(w,\mathrm{diag})}_{123} + \delta G^{\tilde{f}(w,AZ+AG)}_{123}\,, \tag{4.33}$$

and combine the amplitudes coming from A^0–G^0 and A^0–Z^0 mixing in the following convenient way. First we show that the sum of the parts coming from the propagators of Z^0 and G^0 outside the loops is independent of the gauge parameter $\xi = \xi_Z$. In a general R_ξ–gauge

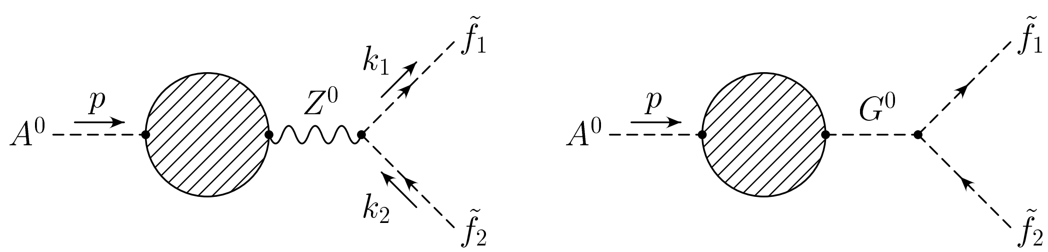


Figure 4.2: A^0–Z^0 contribution and A^0–G^0 wave–function correction

the amplitudes of the two graphs of Fig. 4.2 are given by

$$\mathcal{M}^Z = \Big(-ip^\mu \Pi_{AZ}(p^2)\Big)\frac{i}{p^2 - m^2_Z}\bigg(-g_{\mu\nu} + (1-\xi)\frac{p_\mu p_\nu}{p^2 - \xi m^2_Z}\bigg)\Big(-ig_Z\, z^{\tilde{f}}_{12}\Big)(k_1 + k_2)^\nu\,, \tag{4.34}$$

$$\mathcal{M}^G = \Big(i\Pi_{AG}(p^2)\Big)\frac{i}{p^2 - \xi m^2_Z}\, iG^{\tilde{f}}_{124}\,. \tag{4.35}$$

Contracting the Lorentz indices in $\mathcal{M}^Z$,

$$p^\mu\left(-g_{\mu\nu}+(1-\xi)\frac{p_\mu p_\nu}{p^2-\xi m_Z^2}\right)(k_1+k_2)^\nu \;=\; -\left(1-\frac{(1-\xi)p^2}{p^2-\xi m_Z^2}\right)\left(m_{\tilde f_1}^2-m_{\tilde f_2}^2\right), \quad (4.36)$$

and eliminating Π_{AG} in favor of Π_{AZ} by using the Slavnov–Taylor identity [20]

$$p^2\Pi_{AZ}(p^2)+im_Z\Pi_{AG}(p^2)=0\,, \quad (4.37)$$

we find the sum $\mathcal{M}^Z+\mathcal{M}^G$

$$\begin{aligned}\mathcal{M}^{Z+G} \;&=\; \frac{i}{p^2-m_Z^2}\Pi_{AZ}(p^2)\,g_Z\,z_{12}^{\tilde f}\big(m_{\tilde f_1}^2-m_{\tilde f_2}^2\big)\left(1-\frac{(1-\xi)p^2}{p^2-\xi m_Z^2}\right)\\ &+\;\frac{p^2}{p^2-\xi m_Z^2}\,\frac{\Pi_{AZ}(p^2)}{m_Z}\,G_{124}^{\tilde f}\,. \end{aligned} \quad (4.38)$$

Finally we use the identity

$$g_Z\,z_{ij}^{\tilde f}\left(m_{\tilde f_i}^2-m_{\tilde f_j}^2\right) \;=\; im_Z\,G_{ij4}^{\tilde f} \quad (4.39)$$

to obtain the result

$$\begin{aligned}\delta G_{123}^{\tilde f(w,AZ+AG)} \;&=\; -i\mathcal{M}^{Z+G}(p^2\to m_{A^0}^2) \;=\; -\frac{i\,\Pi_{AZ}(m_{A^0}^2)\,G_{124}^{\tilde f}}{m_Z\left(p^2-m_Z^2\right)\left(p^2-\xi m_Z^2\right)}\times\\ &\Big[-m_Z^2\Big(\big(p^2-\xi m_Z^2\big)-(1-\xi)p^2\Big)+p^2\big(p^2-m_Z^2\big)\Big]\\ &=\; -\frac{i}{m_Z}\,\Pi_{AZ}(m_{A^0}^2)\,G_{124}^{\tilde f}\,. \end{aligned} \quad (4.40)$$

The gauge dependence of the propagators of the Z^0 and G^0 in Fig. 4.2 is completely removed. However, there still remain gauge dependences from vector particles and Goldstone bosons in the loops of Π_{AZ} which cancel against their counter parts in the vertex, wave–function and counter–term corrections.

In a similar manner we can sum up the amplitudes stemming from H^+–G^+ and H^+–W^+ mixing. Using the Slavnov–Taylor identity

$$p^2\Pi_{HW}(p^2)-m_W\Pi_{HG}(p^2)=0 \quad (4.41)$$

(where H and G now certainly denote the charged Higgs and Goldstone bosons) as well as the relation between the $W^+\tilde t_i^*\tilde b_j$ and $G^+\tilde t_i^*\tilde b_j$ couplings,

$$\frac{g}{\sqrt2}\,R_{i1}^{\tilde t}R_{j1}^{\tilde b}\left(m_{\tilde t_i}^2-m_{\tilde b_j}^2\right) \;=\; m_W\,G_{ij2}^{\tilde t\tilde b}\,, \quad (4.42)$$

the wave–function corrections for the charged Higgs boson decays can be written as

$$\delta G_{ij1}^{\tilde t\tilde b(w)}=\frac12\left[\delta Z_{i'i}^{\tilde t}G_{i'jk}^{\tilde t\tilde b}+\delta Z_{j'j}^{\tilde b}G_{ij'k}^{\tilde t\tilde b}+\delta Z_{11}^{H^+}G_{ij1}^{\tilde t\tilde b}\right]+\delta G_{ij1}^{\tilde t\tilde b(w,HW+HG)} \quad (4.43)$$

with

$$\delta G_{ij1}^{\tilde{t}\tilde{b}(w,HW+HG)} = -\frac{1}{m_W}\,\Pi_{HW}(m_{H^+}^2)\,G_{ij2}^{\tilde{t}\tilde{b}}\,. \tag{4.44}$$

The coupling counter–term corrections originate from the shifting of the parameters in the Lagrangian. In the case of the CP–even Higgs bosons, h^0 and H^0, they can be expressed as $(k = 1, 2)$

$$\delta G_{ijk}^{\tilde{f}(c)} = \left[\delta R^{\tilde{f}} \cdot G_{LR,k}^{\tilde{f}} \cdot (R^{\tilde{f}})^T + R^{\tilde{f}} \cdot \delta G_{LR,k}^{\tilde{f}} \cdot (R^{\tilde{f}})^T + \; R^{\tilde{f}} \cdot G_{LR,k}^{\tilde{f}} \cdot (\delta R^{\tilde{f}})^T\right]_{ij}. \tag{4.45}$$

As we have fixed the counter terms for the sfermion and Higgs mixing angle by means of eqs. (3.55) and (3.64),

$$\delta\theta_{\tilde{f}} = \frac{1}{4}\left(\delta Z_{12}^{\tilde{f}} - \delta Z_{21}^{\tilde{f}}\right), \qquad \delta\alpha = \frac{1}{4}\left(\delta Z_{21}^{H^0} - \delta Z_{12}^{H^0}\right), \tag{4.46}$$

we can write the coupling counter–term corrections as

$$\delta G_{ijk}^{\tilde{f}(c)} = -\big(\varepsilon_{i'i}\,G_{i'jk}^{\tilde{f}} + \varepsilon_{j'j}\,G_{ij'k}^{\tilde{f}}\big)\,\delta\theta_{\tilde{f}} + \varepsilon_{lk}\,G_{ijl}^{\tilde{f}}\,\delta\alpha + \left[R^{\tilde{f}} \cdot \hat{\delta} G_{LR,k}^{\tilde{f}} \cdot (R^{\tilde{f}})^T\right]_{ij}, \tag{4.47}$$

where we have introduced the antisymmetric symbol ε_{ij} with $\varepsilon_{12} = 1$. The derivative $\hat{\delta}$ in eq. (4.47) indicates that the variation is taken with respect to all parameters except the mixing angle α. Using the relations

$$\frac{\delta G_{ij1}^{\tilde{f}}}{\delta\alpha} = -\,G_{ij2}^{\tilde{f}}\,, \qquad \frac{\delta G_{ij2}^{\tilde{f}}}{\delta\alpha} = G_{ij1}^{\tilde{f}}\,, \tag{4.48}$$

we can 'absorb' the counter terms for the mixing angles of the outer particles, $\delta\alpha$ and $\delta\theta_{\tilde{f}}$, into the wave–function corrections $\delta G_{ijk}^{\tilde{f}(w)}$ yielding the *symmetric* wave–function corrections in

$$\delta G_{ijk}^{\tilde{f}(w+c)} = \delta G_{ijk}^{\tilde{f}(w,\,\mathrm{symm.})} + \left[R^{\tilde{f}} \cdot \hat{\delta} G_{LR,k}^{\tilde{f}} \cdot (R^{\tilde{f}})^T\right]_{ij}, \tag{4.49}$$

given by

$$\delta G_{ijk}^{\tilde{f}(w,\,\mathrm{symm.})} = \frac{1}{4}\big(\delta Z_{ii'}^{\tilde{f}} + \delta Z_{i'i}^{\tilde{f}}\big)G_{i'jk}^{\tilde{f}} + \frac{1}{4}\big(\delta Z_{jj'}^{\tilde{f}} + \delta Z_{j'j}^{\tilde{f}}\big)G_{ij'k}^{\tilde{f}} + \frac{1}{4}\big(\delta Z_{kl}^{H} + \delta Z_{lk}^{H}\big)G_{ijl}^{\tilde{f}}\,. \tag{4.50}$$

Note that in this symmetrized form momentum–independent contributions from four–scalar couplings and tadpole shifts cancel out.

The explicit forms of the counter terms $\hat{\delta} G^{\tilde{f}}_{LR,k}$ for $k = 1, 2$ are given by

$$\begin{aligned}
\big(\hat{\delta} G^{\tilde{f}}_{LR,1}\big)_{11} &= -\sqrt{2}\, h_f\, m_f\, c_\alpha \left(\frac{\delta h_f}{h_f} + \frac{\delta m_f}{m_f}\right) - g_Z\, m_Z\, e_f\, \delta s_W^2\, s_{\alpha+\beta} \\
&\quad + g_Z\, m_Z (I_f^{3L} - e_f s_W^2) s_{\alpha+\beta} \left(\frac{\delta g_Z}{g_Z} + \frac{\delta m_Z}{m_Z} + \frac{\delta\beta}{t_{\alpha+\beta}}\right), \qquad (4.51)\\
\big(\hat{\delta} G^{\tilde{f}}_{LR,1}\big)_{12} &= \frac{\delta h_f}{h_f} \big(G^{\tilde{f}}_{LR,1}\big)_{12} - \frac{h_f}{\sqrt{2}} (\delta A_f\, c_\alpha + \delta\mu\, s_\alpha), \qquad (4.52)\\
\big(\hat{\delta} G^{\tilde{f}}_{LR,1}\big)_{22} &= -\sqrt{2}\, h_f\, m_f\, c_\alpha \left(\frac{\delta h_f}{h_f} + \frac{\delta m_f}{m_f}\right) \\
&\quad + g_Z\, m_Z\, e_f\, s_W^2\, s_{\alpha+\beta} \left(\frac{\delta g_Z}{g_Z} + \frac{\delta m_Z}{m_Z} + \frac{\delta s_W^2}{s_W^2} + \frac{\delta\beta}{t_{\alpha+\beta}}\right) \qquad (4.53)
\end{aligned}$$

for the sfermion couplings to the Higgs boson h^0 and

$$\begin{aligned}
\big(\hat{\delta} G^{\tilde{f}}_{LR,2}\big)_{11} &= -\sqrt{2}\, h_f\, m_f\, s_\alpha \left(\frac{\delta h_f}{h_f} + \frac{\delta m_f}{m_f}\right) + g_Z\, m_Z\, e_f\, \delta s_W^2\, c_{\alpha+\beta} \\
&\quad - g_Z\, m_Z (I_f^{3L} - e_f s_W^2) c_{\alpha+\beta} \left(\frac{\delta g_Z}{g_Z} + \frac{\delta m_Z}{m_Z} - t_{\alpha+\beta}\, \delta\beta\right), \qquad (4.54)\\
\big(\hat{\delta} G^{\tilde{f}}_{LR,2}\big)_{12} &= \frac{\delta h_f}{h_f} \big(G^{\tilde{f}}_{LR,2}\big)_{12} - \frac{h_f}{\sqrt{2}} (\delta A_f\, s_\alpha - \delta\mu\, c_\alpha), \qquad (4.55)\\
\big(\hat{\delta} G^{\tilde{f}}_{LR,2}\big)_{22} &= -\sqrt{2}\, h_f\, m_f\, s_\alpha \left(\frac{\delta h_f}{h_f} + \frac{\delta m_f}{m_f}\right) \\
&\quad - g_Z\, m_Z\, e_f\, s_W^2\, c_{\alpha+\beta} \left(\frac{\delta g_Z}{g_Z} + \frac{\delta m_Z}{m_Z} + \frac{\delta s_W^2}{s_W^2} - t_{\alpha+\beta}\, \delta\beta\right) \qquad (4.56)
\end{aligned}$$

for the couplings to H^0.
The higgsino mass parameter μ is fixed in the chargino sector, $\mu \equiv X_{22}$, where μ enters in the chargino mass matrix X [24, 25],

$$X = \begin{pmatrix} M & \sqrt{2} m_W \sin\beta \\ \sqrt{2} m_W \cos\beta & \mu \end{pmatrix} \qquad \rightarrow \delta\mu \;=\; (\delta X)_{22}\,. \qquad (4.57)$$

With the definition of the Yukawa couplings, eq. (4.7), the counter term δh_f can be further decomposed into

$$\frac{\delta h_f}{h_f} = \frac{\delta g}{g} + \frac{\delta m_f}{m_f} - \frac{\delta m_W}{m_W} + \left\{ \begin{matrix} -\cos^2\beta \\ \sin^2\beta \end{matrix} \right\} \frac{\delta \tan\beta}{\tan\beta} \qquad (4.58)$$

for $\left\{ \begin{matrix} \text{up} \\ \text{down} \end{matrix} \right\}$–type sfermions. The explicit forms of these counter terms and the corresponding self–energies are given in Appendix B.

The CP–odd Higgs boson A^0 only couples off–diagonally to the sfermions through terms proportional to Yukawa couplings. Therefore, the counter–term corrections simply read

$$\delta G_{123}^{\tilde{f}(c)} = \frac{\delta h_f}{h_f}\, G_{123}^{\tilde{f}} + \frac{i}{\sqrt{2}}\, h_f\, \delta\Big(A_f \begin{Bmatrix} \cos\beta \\ \sin\beta \end{Bmatrix} + \mu \begin{Bmatrix} \sin\beta \\ \cos\beta \end{Bmatrix} \Big) . \tag{4.59}$$

Analogously to the decays of the CP–even Higgs bosons, the sum of the wave–function and counter–term corrections of the charged Higgs boson can be expressed as

$$\delta G_{ij1}^{\tilde{t}\tilde{b}(w+c)} = \delta G_{ij1}^{\tilde{t}\tilde{b}(w,\,\mathrm{symm.})} + \Big[R^{\tilde{t}} \cdot \hat{\delta} G_{LR,1}^{\tilde{t}\tilde{b}} \cdot (R^{\tilde{b}})^T \Big]_{ij} + \delta G_{ij1}^{\tilde{t}\tilde{b}(w,HW+HG)} \tag{4.60}$$

with the symmetrized wave–function corrections

$$\delta G_{ij1}^{\tilde{t}\tilde{b}(w,\,\mathrm{symm.})} = \frac{1}{4}\big(\delta Z_{ii'}^{\tilde{t}} + \delta Z_{i'i}^{\tilde{t}}\big) G_{i'j1}^{\tilde{t}\tilde{b}} + \frac{1}{4}\big(\delta Z_{jj'}^{\tilde{b}} + \delta Z_{j'j}^{\tilde{b}}\big) G_{ij'1}^{\tilde{t}\tilde{b}} + \frac{1}{2}\delta Z_{11}^{H^+} G_{ij1}^{\tilde{t}\tilde{b}} . \tag{4.61}$$

The single elements of the matrix corresponding to the variation with respect to the couplings, $\hat{\delta} G_{LR,1}^{\tilde{t}\tilde{b}}$, are given explicitly as follows (cf. eq. (4.9)):

$$\begin{aligned} \big(\hat{\delta} G_{LR,1}^{\tilde{t}\tilde{b}}\big)_{11} &= h_b m_b s_\beta \Big(\frac{\delta h_b}{h_b} + \frac{\delta m_b}{m_b} + \frac{\delta s_\beta}{s_\beta} \Big) + h_t m_t c_\beta \Big(\frac{\delta h_t}{h_t} + \frac{\delta m_t}{m_t} + \frac{\delta c_\beta}{c_\beta} \Big) \\ &\quad - \frac{g m_W}{\sqrt{2}} \sin 2\beta \Big(\frac{\delta g}{g} + \frac{\delta m_W}{m_W} + \cos 2\beta\, \frac{\delta \tan\beta}{\tan\beta} \Big) \end{aligned} \tag{4.62}$$

$$\big(\hat{\delta} G_{LR,1}^{\tilde{t}\tilde{b}}\big)_{12} = \frac{\delta h_b}{h_b} \big(G_{LR,1}^{\tilde{t}\tilde{b}}\big)_{12} + h_b\big(\delta A_b s_\beta + A_b \delta s_\beta + \delta\mu\, c_\beta + \mu\, \delta c_\beta\big) \tag{4.63}$$

$$\big(\hat{\delta} G_{LR,1}^{\tilde{t}\tilde{b}}\big)_{21} = \frac{\delta h_t}{h_t} \big(G_{LR,1}^{\tilde{t}\tilde{b}}\big)_{21} + h_t\big(\delta A_t c_\beta + A_t \delta c_\beta + \delta\mu\, s_\beta + \mu\, \delta s_\beta\big) \tag{4.64}$$

$$\big(\hat{\delta} G_{LR,1}^{\tilde{t}\tilde{b}}\big)_{22} = h_t m_b c_\beta \Big(\frac{\delta h_t}{h_t} + \frac{\delta m_b}{m_b} + \frac{\delta c_\beta}{c_\beta} \Big) + h_b m_t s_\beta \Big(\frac{\delta h_b}{h_b} + \frac{\delta m_t}{m_t} + \frac{\delta s_\beta}{s_\beta} \Big) \tag{4.65}$$

4.4 Infrared divergences

Analyzing the virtual corrections in more detail, i.e. the vertex and wave–function corrections with one single photon in the loop, one finds that the corresponding amplitudes are divergent. These *infrared divergences* (IR) originate from massless photons which, for small momenta, lead to divergent integrals. In order to regularize the divergent expressions, we introduce a small photon mass λ in e.g.

$$\frac{\partial}{\partial p^2} \frac{1}{i\pi^2} \int d^d q \frac{1}{(q^2-\lambda^2)[(q+p)^2-m^2]}\bigg|_{p^2=m^2} = -\frac{1}{2m^2}\Big(2 - \log\frac{m^2}{\lambda^2}\Big) , \tag{4.66}$$

which diverges for $\lambda \to 0$. Following a theorem of Bloch and Nordsieck [36], the IR–divergences can be cancelled by the inclusion of *real Bremsstrahlung* processes which contain one additional single photon in the final state. In the case of the decay of the charged Higgs

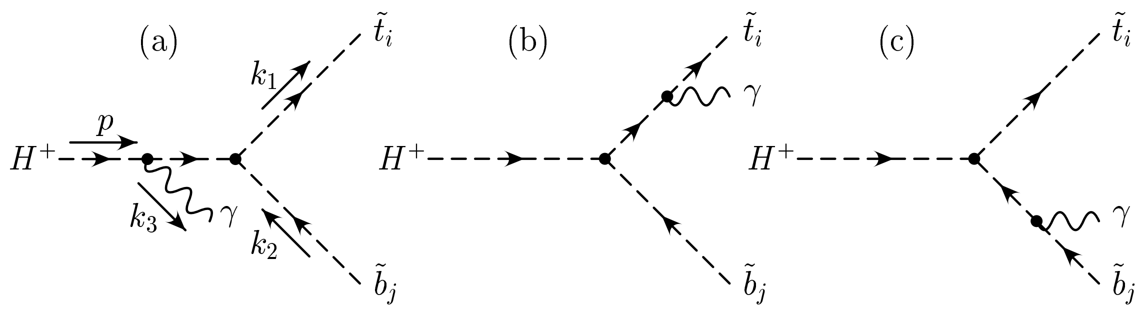


Figure 4.3: Real Bremsstrahlung diagrams relevant to cancel the IR–divergences in $H^+(p) \to \tilde{t}_i(k_1) + \bar{\tilde{b}}_j(-k_2)$.

boson, this means that the real photon emission process $H^+(p) \to \tilde{t}_i(k_1) + \bar{\tilde{b}}_j(-k_2) + \gamma(k_3)$ has to be calculated. The Feynman amplitudes of the diagrams in Fig. 4.3 are given by

$$
\begin{aligned}
\mathcal{M}^a &= iG^{\tilde{t}\tilde{b}}_{ij1}\big(-ie_0(2p-k_3)_\mu\big)\frac{i}{(p-k_3)^2-m_H^2}\,\varepsilon^*_\mu(k_3) &&= iG^{\tilde{t}\tilde{b}}_{ij1}(-e_0)\Big(-\frac{p\cdot\varepsilon^*}{-p\cdot k_3}\Big)\,, \\
\mathcal{M}^b &= iG^{\tilde{t}\tilde{b}}_{ij1}\big(-ie_0e_t(2k_1+k_3)_\mu\big)\frac{i}{(k_1+k_3)^2-m_i^2}\,\varepsilon^*_\mu(k_3) &&= iG^{\tilde{t}\tilde{b}}_{ij1}(-e_0)\Big(-e_t\frac{k_1\cdot\varepsilon^*}{k_1\cdot k_3}\Big)\,, \\
\mathcal{M}^c &= iG^{\tilde{t}\tilde{b}}_{ij1}\big(-ie_0e_b(2k_2-k_3)_\mu\big)\frac{i}{(k_2-k_3)^2-m_j^2}\,\varepsilon^*_\mu(k_3) &&= iG^{\tilde{t}\tilde{b}}_{ij1}(-e_0)\Big(-e_b\frac{k_2\cdot\varepsilon^*}{-k_2\cdot k_3}\Big)\,,
\end{aligned}
$$

which leads to the decay width

$$
\begin{aligned}
\Gamma^{\text{real}} &= \frac{N_C}{2m_{H^+}}\int\frac{d^3k_1}{(2\pi)^3 2E_1}\int\frac{d^3k_2}{(2\pi)^3 2E_2}\int\frac{d^3k_3}{(2\pi)^3 2E_3}(2\pi)^4\delta^4(p-k_1+k_2-k_3)|\mathcal{M}^{\text{hard}}|^2 \\
&= \frac{N_C}{16\pi^3 m_{H^+}}|G^{\tilde{t}\tilde{b}}_{ij1}|^2(-e_0^2)\Bigg[\frac{1}{\pi^2}\int\frac{d^3k_1}{2E_1}\int\frac{d^3k_2}{2E_2}\int\frac{d^3k_3}{2E_3}\,\delta^4(p-k_1+k_2-k_3) \\
&\quad\times\Bigg(\frac{m_H^2}{(-2p\cdot k_3)^2}+e_t^2\frac{m_i^2}{(2k_1\cdot k_3)^2}+e_b^2\frac{m_j^2}{(-2k_2\cdot k_3)^2}-e_te_b\frac{(-2k_1\cdot k_2)}{(2k_1\cdot k_3)(-2k_2\cdot k_3)} \\
&\qquad -e_t\frac{(-2p\cdot k_1)}{(-2p\cdot k_3)(2k_1\cdot k_3)}+e_b\frac{(2p\cdot k_2)}{(-2p\cdot k_3)(-2k_2\cdot k_3)}\Bigg)\Bigg]\,.
\end{aligned}
\tag{4.67}
$$

Using the phase–space integrals I_n and I_{mn} in the convention of [15],

$$
I_{i_1\ldots i_n} = \frac{1}{\pi^2}\int\frac{d^3k_1}{2E_1}\frac{d^3k_2}{2E_2}\frac{d^3k_3}{2E_3}\,\delta^4(p-k_1+k_2-k_3)\frac{1}{(\pm 2k_{i_1}\cdot k_3)\ldots(\pm 2k_{i_n}\cdot k_3)}\,, \tag{4.68}
$$

where the plus signs belong to the momenta of the outgoing particles k_1 and k_3 and the minus signs to the momenta p and k_2, the real Bremsstrahlung decay width can be expressed as

$$\begin{aligned}\Gamma(H^+ \to \tilde{t}_i \bar{\tilde{b}}_j \gamma) &= \frac{N_C}{16\pi^3 m_{H^+}} |G_{ij1}^{\tilde{t}\tilde{b}}|^2 (-e_0)^2 \Big[m_{H^+}^2 I_{00} + e_t^2 m_i^2 I_{11} + e_b^2 m_j^2 I_{22} \\ &\quad - e_t e_b \Big((m_{H^+}^2 - m_i^2 - m_j^2) I_{12} - I_2 - I_1 \Big) \\ &\quad - e_t \Big((m_j^2 - m_{H^+}^2 - m_i^2) I_{01} - I_1 - I_0 \Big) \\ &\quad + e_b \Big((m_i^2 - m_{H^+}^2 - m_j^2) I_{02} - I_2 - I_0 \Big) \Big] \,. \end{aligned} \tag{4.69}$$

Analogously, the real Bremsstrahlung contributions to the neutral Higgs boson decays are

$$\begin{aligned}\Gamma(H_k^0 \to \tilde{f}_i \bar{\tilde{f}}_j \gamma) &= \frac{N_C}{16\pi^3 m_{H_k^0}} |G_{ijk}^{\tilde{f}}|^2 (-e_0 e_f)^2 \Big(m_i^2 I_{11} + m_j^2 I_{22} - I_2 - I_1 \\ &\quad + (m_{H_k^0}^2 - m_i^2 - m_j^2) I_{12} \Big) \,. \end{aligned} \tag{4.70}$$

The corrected (UV– and IR–convergent) decay widths for the physical value $\lambda = 0$ are then given by

$$\Gamma^{\text{corr}}(H_k^0 \to \tilde{f}_i \bar{\tilde{f}}_j) \equiv \Gamma(H_k^0 \to \tilde{f}_i \bar{\tilde{f}}_j) + \Gamma(H_k^0 \to \tilde{f}_i \bar{\tilde{f}}_j \gamma) \,, \tag{4.71}$$

$$\Gamma^{\text{corr}}(H^+ \to \tilde{t}_i \bar{\tilde{b}}_j) \equiv \Gamma(H^+ \to \tilde{t}_i \bar{\tilde{b}}_j) + \Gamma(H^+ \to \tilde{t}_i \bar{\tilde{b}}_j \gamma) \,. \tag{4.72}$$

4.5 Improvement of One–loop Corrections

Having collected all pieces which are necessary for a reliable calculation of the decay widths, i.e. the vertex, wave–function and counter–term corrections to make the result UV–convergent as well as the real Bremsstrahlung processes to cancel the IR divergences, we could perform a numerical analysis which shows the contributions of the higher–order corrections. However, in the case of bottom squarks and tau sleptons, especially for large $\tan\beta$ the corrections to the decay widths can be very large in the on–shell renormalization scheme. If the corrections are negative, the one–loop corrected width can even become negative and therefore unphysical. Hence the perturbation expansion around the on–shell tree level is no longer reliable. It has been shown in [37, 38] that, in the case of the decays into bottom squarks, the source of these large corrections are mainly the counter terms for m_b and the trilinear coupling A_b, in particular the SUSY–QCD corrections. However, despite the absence of strong interactions for the decay into tau sleptons, the corrections become extremely large. This problem can be solved by defining an appropriate tree level in terms of $\overline{\text{DR}}$ running values for m_f and A_f. The expansion around this new tree level then no longer suffers from bad convergence.

Correction to m_b

First we review the improvement of the perturbation expansion by using $\overline{\rm DR}$ running bottom quark masses, following [38, 39, 40].

If the Yukawa coupling h_b is given at tree level in terms of the pole mass m_b, the one–loop corrections become very large due to gluon and gluino exchange contributions to the counter term δm_b. The large counter term caused by the gluon loop is absorbed by using SM two–loop RGEs in the $\overline{\rm MS}$ scheme [38, 39, 40]. Thus we obtain the SM running bottom quark mass $\hat{m}_b(Q)_{\rm SM}$:

$$\hat{m}_b(Q)^{\overline{\rm MS}}_{\rm SM} = \left(\frac{\hat{m}_b(Q)^{\overline{\rm MS}}_{\rm SM}}{\hat{m}_b(m_b)^{\overline{\rm MS}}_{\rm SM}}\right) \hat{m}_b(m_b)^{\overline{\rm MS}}_{\rm SM} \tag{4.73}$$

The ratio $\left(\hat{m}_b(Q)^{\overline{\rm MS}}_{\rm SM}/\hat{m}_b(m_b)^{\overline{\rm MS}}_{\rm SM}\right)$ can be expressed as

$$\frac{\hat{m}_b(Q)^{\overline{\rm MS}}_{\rm SM}}{\hat{m}_b(m_b)^{\overline{\rm MS}}_{\rm SM}} = \begin{cases} \dfrac{c_5(\alpha_s^{(2)}(Q)/\pi)}{c_5(\alpha_s^{(2)}(m_b)/\pi)} & (m_b < Q \le m_t)\,, \\[2ex] \dfrac{c_6(\alpha_s^{(2)}(Q)/\pi)}{c_6(\alpha_s^{(2)}(m_t)/\pi)}\,\dfrac{c_5(\alpha_s^{(2)}(m_t)/\pi)}{c_5(\alpha_s^{(2)}(m_b)/\pi)} & (Q > m_t)\,, \end{cases}$$

where we have used the functions

$$\begin{aligned} c_5(x) &= \left(\frac{23}{6}x\right)^{\frac{12}{23}} (1 + 1.175x) \qquad && (m_b < Q \le m_t)\,, \\ c_6(x) &= \left(\frac{7}{2}x\right)^{\frac{4}{7}} (1 + 1.398x) \qquad && (Q > m_t)\,, \end{aligned}$$

and the two–loop RGEs for α_s [40],

$$\alpha_s^{(2)}(Q) = \frac{12\pi}{(33 - 2\,n_f)\ln\frac{Q^2}{\Lambda^2_{n_f}}} \left(1 - \frac{6(153 - 19\,n_f)}{(33 - 2\,n_f)^2} \frac{\ln\ln\frac{Q^2}{\Lambda^2_{n_f}}}{\ln\frac{Q^2}{\Lambda^2_{n_f}}}\right)\,, \tag{4.74}$$

with $n_f = 5$ or 6 for $m_b < Q \le m_t$ or $Q > m_t$, respectively. For the SM $\overline{\rm DR}$ running bottom quark mass at the scale $Q = m_b$ we use the $\overline{\rm MS}$ equation

$$\hat{m}_b(m_b)^{\overline{\rm MS}}_{\rm SM} = m_b \left[1 + \frac{4}{3}\frac{\alpha_s^{(2)}(m_b)}{\pi} + K_q\left(\frac{\alpha_s^{(2)}(m_b)}{\pi}\right)^2\right]^{-1}\,, \tag{4.75}$$

with $K_q = 12.4$ and then convert to $\overline{\rm DR}$ using one–loop running $\alpha_s(Q)$:

$$\hat{m}_b(Q)_{\rm SM} \;=\; \left(\frac{\hat{m}_b(Q)^{\overline{\rm MS}}_{\rm SM}}{\hat{m}_b(m_b)^{\overline{\rm MS}}_{\rm SM}}\right) \hat{m}_b(m_b)^{\overline{\rm MS}}_{\rm SM} \;-\; \frac{\alpha_s(Q)}{3\pi}\,m_b \tag{4.76}$$

In the MSSM, for large $\tan\beta$ the counter term to m_b can be very large due to the gluino–mediated graph [37, 41, 42]. Here we absorb the gluino contribution as well as the sizeable contributions from neutralino and chargino loops and the remaining electroweak self–energies into the Higgs–sfermion–sfermion tree–level coupling. In such a way we obtain the full $\overline{\rm DR}$ running bottom quark mass

$$\hat{m}_b(Q)_{\rm MSSM} \;=\; \hat{m}_b(Q)_{\rm SM} + \delta m_b(Q)\,. \tag{4.77}$$

The explicit form of the electroweak contribution to the counter term $\delta m_b(Q)$ is given in Appendix B.9.

Correction to $\boldsymbol{A_{b,\tau}}$

The second source of a very large correction (in the on–shell scheme) are the counter terms for the trilinear coupling $A_{b,\tau}$ (see eq. (2.39)),

$$\delta A_{b,\tau} \;=\; \frac{\delta m_{LR}^2}{m_{b,\tau}} \;-\; \frac{m_{LR}^2}{m_{b,\tau}}\,\frac{\delta m_{b,\tau}}{m_{b,\tau}} \;+\; \delta\mu\tan\beta \;+\; \mu\,\delta\tan\beta\,. \tag{4.78}$$

Again, the big bottom mass correction δm_b contributes to δA_b, but also the counter term of the left–right mixing elements of the sfermion mass matrix, δm_{LR}^2, gives a very large correction for higher values of $\tan\beta$. In particular, in the case of the decay into staus, this is the main source for the bad convergence of the tree–level expansion. As in the case of the large correction to m_b, we redefine the Higgs–sfermion–sfermion tree–level coupling in terms of $\overline{\rm DR}$ running $\hat{A}_{b,\tau}(m_{A^0})$. Because of the fact that the counter terms $\delta A_{b,\tau}$ (for large $\tan\beta$) can become several orders of magnitude larger than the on–shell $A_{b,\tau}$ we use $\hat{A}_{b,\tau}(m_{A^0})$ as input [38]. In order to be consistent we have to perform an iteration procedure to get all the correct running and on–shell masses, mixing angles and other parameters. This procedure is described below.

4.6 Method of improvement

In this section we will explain in detail how we can improve the perturbation calculation for the sbottom and stau case by using $\overline{\rm DR}$ running values for m_b and $A_{b,\tau}$ in the Higgs–sfermion–sfermion tree–level couplings. Since we take $\overline{\rm DR}$ running values for $\hat{A}_b$ and $\hat{A}_\tau$ as input and all other parameters on–shell we will have to pay attention to the sbottom and stau sector in order to get consistently all needed running and on–shell masses, mixing angles and other parameters. Here we adopt the procedure developed in [38] and also extend it to the electroweak case.

Stop sector:

We start our calculation in the stop sector. Because all input parameters in the stop sector are on–shell we obtain the on–shell masses $m_{\tilde{t}_1}$, $m_{\tilde{t}_2}$ and the stop mixing angle $\theta_{\tilde{t}}$ by diagonalizing the stop mass matrix in the $\tilde{t}_L$–$\tilde{t}_R$ basis, see chapter 4.2. The running stop masses $\hat{m}_{\tilde{t}_i}$ and the mixing angle $\hat{\theta}_{\tilde{t}}$ are calculated at the scale $Q = Q_{\tilde{t}} = \sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}}$ by adding the appropriate counter terms to the on–shell values in

$$\hat{m}^2_{\tilde{t}_i}(Q_{\tilde{t}}) = m^2_{\tilde{t}_i} + \delta m^2_{\tilde{t}_i}\,, \tag{4.79}$$

$$\delta m^2_{\tilde{t}_i} = \mathrm{Re}\,\Pi^{\tilde{t}}_{ii}(m^2_{\tilde{t}_i})\,, \tag{4.80}$$

$$\hat{\theta}_{\tilde{t}}(Q_{\tilde{t}}) = \theta_{\tilde{t}} + \delta\theta_{\tilde{t}}\,. \tag{4.81}$$

The electroweak parts of the sfermion self–energies $\Pi^{\tilde{f}}_{ii}(m^2_{\tilde{f}_i})$ are given in Appendix B.4 and the SUSY–QCD contributions, $\Pi^{\mathrm{SUSY-QCD}}_{ii}(m^2_{\tilde{f}_i})$ are given in eqs. (25)–(27) in [34]. Here and in the following all running parameters $\hat{X}(Q)$ are related to their on–shell values X by $\hat{X}(Q) = X + \delta X$, with δX being the full one–loop counter term — also including the SUSY–QCD parts. According to eq. (3.55) we fix the sfermion mixing angle by

$$\delta\theta_{\tilde{f}} = \frac{1}{4}\left(\delta Z^{\tilde{f}}_{12} - \delta Z^{\tilde{f}}_{21}\right) = \frac{1}{2\left(m^2_{\tilde{f}_1} - m^2_{\tilde{f}_2}\right)}\,\mathrm{Re}\left(\Pi^{\tilde{f}}_{12}(m^2_{\tilde{f}_2}) + \Pi^{\tilde{f}}_{21}(m^2_{\tilde{f}_1})\right)\,. \tag{4.82}$$

For $\overline{\mathrm{DR}}$ running $\hat{m}_t$ we use the formulae from section 4.5 with the obvious substitutions $m_b \to m_t$ and $K_q = 10.9$ for the top–case. Next we evaluate the running parameters $\hat{M}_{\tilde{Q}}(Q)$ and $\hat{M}_{\tilde{U}}(Q)$ by inserting the running values $\hat{m}^2_{\tilde{t}_i}(Q), \hat{\theta}_{\tilde{t}}(Q), \hat{m}_t(Q)_{\mathrm{MSSM}}$, $\hat{m}_Z(Q) = m_Z + \delta m_Z$, $\hat{\beta}(Q) = \beta + \delta\beta$ and $\hat{\theta}_W = \theta_W - \frac{1}{\sin\theta_W}\left(\frac{\delta m_W}{m_W} - \frac{\delta m_Z}{m_Z}\right)$ into the equations

$$M^2_{\tilde{Q}} = m^2_{\tilde{t}_1}\cos^2\theta_{\tilde{t}} + m^2_{\tilde{t}_2}\sin^2\theta_{\tilde{t}} - m^2_t - m^2_Z\cos 2\beta\left(I^{3L}_t - e_t\sin^2\theta_W\right)\,, \tag{4.83}$$

$$M^2_{\tilde{U}} = m^2_{\tilde{t}_1}\sin^2\theta_{\tilde{t}} + m^2_{\tilde{t}_2}\cos^2\theta_{\tilde{t}} - m^2_t - m^2_Z\cos 2\beta\, e_t\sin^2\theta_W\,. \tag{4.84}$$

For the running value of A_t we use

$$\hat{A}_t = (\hat{m}^2_{\tilde{t}_1} - \hat{m}^2_{\tilde{t}_2})\,\frac{\sin 2\hat{\theta}_{\tilde{t}}}{\hat{m}_t} + \hat{\mu}\,\cot\hat{\beta}\,, \tag{4.85}$$

where we have taken running $\hat{\mu}(Q) = \mu + (\delta X)_{22}$ (cf. eq. (4.57)).

Sbottom sector:
In the sbottom sector we have given all parameters on–shell except the parameter for the trilinear coupling, $\hat{A}_b(Q)$, which is running. First we calculate $\hat{m}_b(Q_{\tilde{b}})_{\mathrm{MSSM}}$ from eq. (4.77) at the scale $Q_{\tilde{b}} = \sqrt{m_{\tilde{b}_1} m_{\tilde{b}_2}}$. From the stop sector we already know the running values of $M_{\tilde{Q}}, \tan\beta$ and μ. Then we diagonalize the sbottom mass matrix using $\hat{m}_b(Q_{\tilde{b}})_{\mathrm{MSSM}}, \hat{M}_{\tilde{Q}}, \tan\hat{\beta}, \hat{\mu}$ and on–shell $M_{\tilde{D}}$, which is near its running value $\hat{M}_{\tilde{D}}$, to obtain the starting values for $\hat{m}_{\tilde{b}_i}$ and $\hat{\theta}_{\tilde{b}}$. The on–shell sbottom masses $m_{\tilde{b}_i}$ and the mixing angle $\theta_{\tilde{b}}$ are calculated from their running values by subtracting the appropriate counter terms, i.e. $m^2_{\tilde{b}_i} = \hat{m}^2_{\tilde{b}_i}(Q) - \delta m^2_{\tilde{b}_i}$, $\theta_{\tilde{b}} = \hat{\theta}_{\tilde{b}}(Q) - \delta\theta_{\tilde{b}}$. Now we can compute the running value for $M_{\tilde{D}}$. Using the relation

$$M^2_{\tilde{D}} = m^2_{\tilde{b}_1}\sin^2\theta_{\tilde{b}} + m^2_{\tilde{b}_2}\cos^2\theta_{\tilde{b}} - m^2_b - m^2_Z\cos 2\beta\, e_b\sin^2\theta_W \tag{4.86}$$

we get $\hat{M}_{\tilde{D}} = (M^2_{\tilde{D}} + \delta M^2_{\tilde{D}})^{1/2} \approx M_{\tilde{D}} + \delta M^2_{\tilde{D}}/(2M_{\tilde{D}})$ with

$$\begin{aligned} \delta M^2_{\tilde{D}} &= \delta m^2_{\tilde{b}_1} \sin^2\theta_{\tilde{b}} + \delta m^2_{\tilde{b}_2} \cos^2\theta_{\tilde{b}} + \left(m^2_{\tilde{b}_1} - m^2_{\tilde{b}_2}\right) \sin 2\theta_{\tilde{b}}\, \delta\theta_{\tilde{b}} - 2m_b\, \delta m_b \\ &\quad -\delta m_Z^2 \cos 2\beta\, e_b \sin^2\theta_W + 2m_Z^2 \sin 2\beta\, \delta\beta\, e_b \sin^2\theta_W \\ &\quad -m_Z^2 \cos 2\beta\, e_b\, \delta\sin^2\theta_W \end{aligned} \tag{4.87}$$

and $\delta m_b = \hat{m}_b(Q_{\tilde{b}})_{\rm MSSM} - m_b$. Because the parameters involved in these calculations are very entangled, e.g. $\hat{M}_{\tilde{D}}$ depends on the $\delta m_{\tilde{b}_i}$ which themselves depend on $M_{\tilde{D}}$, we have to perform an iteration procedure.

Iteration procedure:
Here we will describe in detail the procedure how we obtain all necessary on–shell and running parameters. For convenience, we shortly denote all masses, parameters, couplings etc. for a certain $n \geq 1$ in the iteration by $\hat{\mathcal{X}}^{(n)}$. As starting values $\hat{\mathcal{X}}^{(0)}$ we take on–shell masses and parameters (except $\hat{A}_b$ which is running) and the couplings derived from these quantities. The only exceptions are the standard model running fermion masses $\hat{m}_f^{(0)} = \hat{m}_f(Q)_{\rm SM}$. $\hat{m}_f$ shortly stands for the full $\overline{\rm DR}$ running fermion masses, $\hat{m}_f(Q)_{\rm MSSM}$.
The single steps of the iteration procedure are the following:

1. The running stop masses and the stop mixing angle are calculated as explained above by $\hat{m}^{2\,(n)}_{\tilde{t}_i} = m^2_{\tilde{t}_i} + \delta m^{2\,(n)}_{\tilde{t}_i}\big(\hat{\mathcal{X}}^{(n-1)}\big)$ and $\hat{\theta}^{(n)}_{\tilde{t}} = \theta_{\tilde{t}} + \delta\theta^{(n)}_{\tilde{t}}\big(\hat{\mathcal{X}}^{(n-1)}\big)$.

2. $\hat{m}_t^{(n)} = \hat{m}_{t,\rm SM} + \delta m_t^{(n)}\big(\hat{\mathcal{X}}^{(n-1)}\big)$

3. $\hat{m}_Z^{(n)} = m_Z + \delta m_Z^{(n)}\big(\hat{\mathcal{X}}^{(n-1)}\big)$ and
   $$\sin^2\hat{\theta}_W^{(n)} = \sin^2\theta_W + \delta\sin^2\theta_W{}^{(n)} \text{ with } \delta\sin^2\theta_W{}^{(n)} = -\cos^2\theta_W \left(\frac{\delta m_W}{m_W} - \frac{\delta m_Z}{m_Z}\right)\big(\hat{\mathcal{X}}^{(n-1)}\big)$$

4. The running value of $\tan\beta$, $\tan\hat{\beta}^{(n)} = \tan\beta + \delta\tan\beta^{(n)}$,
   with $\delta\tan\beta^{(n)} = \dfrac{1}{m_Z \sin 2\beta}\, {\rm Im}\Pi_{A^0 Z^0}\big(\hat{\mathcal{X}}^{(n-1)}\big) \tan\beta$ [21].

5. $\hat{\mu}^{(n)} = \mu + \delta\mu^{(n)}$ with $\delta\mu^{(n)} = \delta X_{22}\big(\hat{\mathcal{X}}^{(n-1)}\big)$.

6. The soft SUSY–breaking masses $\hat{M}^{(n)}_{\tilde{Q},\tilde{U}}$ are calculated from $\hat{m}^{(n)}_{\tilde{t}_i}$, $\hat{\theta}^{(n)}_{\tilde{t}}$, $\hat{m}_t(Q_{\tilde{t}})^{(n)}$, $\hat{m}_Z^{(n)}$, $\sin^2\hat{\theta}_W^{(n)}$ and $\tan\hat{\beta}^{(n)}$.

7. We compute the running $\hat{A}_t$ by using running values in eq. (4.85):
   $$\hat{A}_t^{(n)} = \left(\hat{m}^{2\,(n)}_{\tilde{t}_1} - m^{2\,(n)}_{\tilde{t}_2}\right) \frac{\sin 2\hat{\theta}^{(n)}_{\tilde{t}}}{\hat{m}_t^{(n)}} + \hat{\mu}^{(n)} \cot\hat{\beta}^{(n)}$$

8. In the sbottom sector we obtain $\delta m_b^{(n)}$ from the running values already calculated in steps 1.–7., like $\hat{m}_{\tilde{b}_i}^{(n)}$, $\hat{\theta}_{\tilde{b}}^{(n)}$ or $\hat{m}_f^{(n)}$, and the remaining masses, couplings etc. from $\hat{\mathcal{X}}^{(n-1)}$.

9. $\hat{m}_b^{(n)} = \hat{m}_{b,\mathrm{SM}} + \delta m_b^{(n)}$.

10. We receive the running sbottom masses, $\hat{m}_{\tilde{b}_i}^{(n)}$, and the mixing angle, $\hat{\theta}_{\tilde{b}}^{(n)}$, by solving the mass eigenvalue problem with the running values of $\hat{M}_{\tilde{Q}}^{(n)}$, $\hat{M}_{\tilde{D}}^{(n-1)}$, $\hat{m}_b^{(n)}$, $\hat{A}_b$, $\hat{\mu}^{(n)}$ and $\tan\hat{\beta}^{(n)}$.

11. The on–shell sbottom masses $m_{\tilde{b}_i}^{2\,(n)} = \hat{m}_{\tilde{b}_i}^{2\,(n)} - \delta m_{\tilde{b}_i}^{2\,(n)}(Q_{\tilde{b}}^{(n)})$ at the scale $Q_{\tilde{b}}^{(n)} = \sqrt{\hat{m}_{\tilde{b}_1}^{(n)} \hat{m}_{\tilde{b}_2}^{(n)}}$, and $\theta_{\tilde{b}}^{(n)} = \hat{\theta}_{\tilde{b}}^{(n)} - \delta\theta_{\tilde{b}}^{(n)}$.

12. $$\begin{aligned} \delta M_{\tilde{D}}^{2\,(n)} &= \delta m_{\tilde{b}_1}^{2\,(n)} \sin^2\theta_{\tilde{b}} + \delta m_{\tilde{b}_2}^{2\,(n)} \cos^2\theta_{\tilde{b}} + \left(m_{\tilde{b}_1}^2 - m_{\tilde{b}_2}^2\right) \sin 2\theta_{\tilde{b}}\, \delta\theta_{\tilde{b}}^{(n)} - 2m_b \left(\hat{m}_b^{(n)} - m_b\right) \\ &\quad - \delta m_Z^{2\,(n)} \cos 2\beta\, e_b \sin^2\theta_W + 2m_Z^2 \sin 2\beta\, \delta\beta^{(n)}\, e_b \sin^2\theta_W \\ &\quad - m_Z^2 \cos 2\beta\, e_b\, \delta\sin^2\theta_W{}^{(n)} . \end{aligned}$$
(Remember that the values without a hat (ˆ) are on–shell ones!)

13. $\hat{M}_{\tilde{D}}^{(n)} = M_{\tilde{D}} + \dfrac{1}{2} \dfrac{\delta M_{\tilde{D}}^{2\,(n)}}{M_{\tilde{D}}}$.

14. In the sneutrino sector we calculate the running sneutrino mass $\hat{m}_{\tilde{\nu}_\tau}^{2\,(n)} = m_{\tilde{\nu}_\tau}^2 + \delta m_{\tilde{\nu}_\tau}^{2\,(n)}\left(\hat{\mathcal{X}}^{(n-1)}\right)$ and $\hat{M}_{\tilde{L}}^{2\,(n)} = \hat{m}_{\tilde{\nu}_\tau}^{2\,(n)} - \frac{1}{2}\,\hat{m}_Z^{(n)} \cos 2\hat{\beta}^{(n)}$, see also eq (4.83).

15. In the stau sector the values for running $\hat{m}_\tau^{(n)}, \hat{m}_{\tilde{\tau}_i}^{(n)}$ etc. are calculated like in the steps 8–13 in the sbottom sector with the evident substitution $\tilde{b} \to \tilde{\tau}$ for the corresponding parameters and $M_{\tilde{Q}} \to M_{\tilde{L}}, M_{\tilde{D}} \to M_{\tilde{E}}$.

16. All couplings are recalculated with the new running parameters $\to \hat{\mathcal{X}}^n$.

The iteration starts with $n = 1$ and ends, when certain parameters are calculated precisely enough for a given accuracy, i.e. $\left|1 - \frac{\hat{x}^{(n)}}{\hat{x}^{(n-1)}}\right| < \varepsilon$ for $\hat{x} = \{\hat{m}_b, \hat{M}_{\tilde{D}}, \hat{m}_\tau, \hat{M}_{\tilde{E}}\}$. For ε we choose $\varepsilon = 10^{-8}$. We have checked the consistency of this procedure by computing the on–shell $M_{\tilde{D}}$ and running $M_{\tilde{Q}}$ from the sbottom sector by using

$$M_{\tilde{D}}^2 = m_{\tilde{b}_1}^2 \sin^2\theta_{\tilde{b}} + m_{\tilde{b}_2}^2 \cos^2\theta_{\tilde{b}} - m_b^2 - m_Z^2 \cos 2\beta\, e_b \sin^2\theta_W \,, \tag{4.88}$$

$$\hat{M}_{\tilde{Q}}^2 = \hat{m}_{\tilde{b}_1}^2 \cos^2\hat{\theta}_{\tilde{b}} + \hat{m}_{\tilde{b}_2}^2 \sin^2\hat{\theta}_{\tilde{b}} - \hat{m}_b^2 - \hat{m}_Z^2 \cos 2\hat{\beta} \left(I_b^{3L} - e_b \sin^2\hat{\theta}_W\right) , \tag{4.89}$$

which are equal (up to higher–order corrections) to the on–shell input $M_{\tilde{D}}$ and running $M_{\tilde{Q}}$ from the stop sector.
For easier reading the single steps of the iteration procedure of the stop and sbottom sector are depicted in the flowchart in Fig. 4.4.

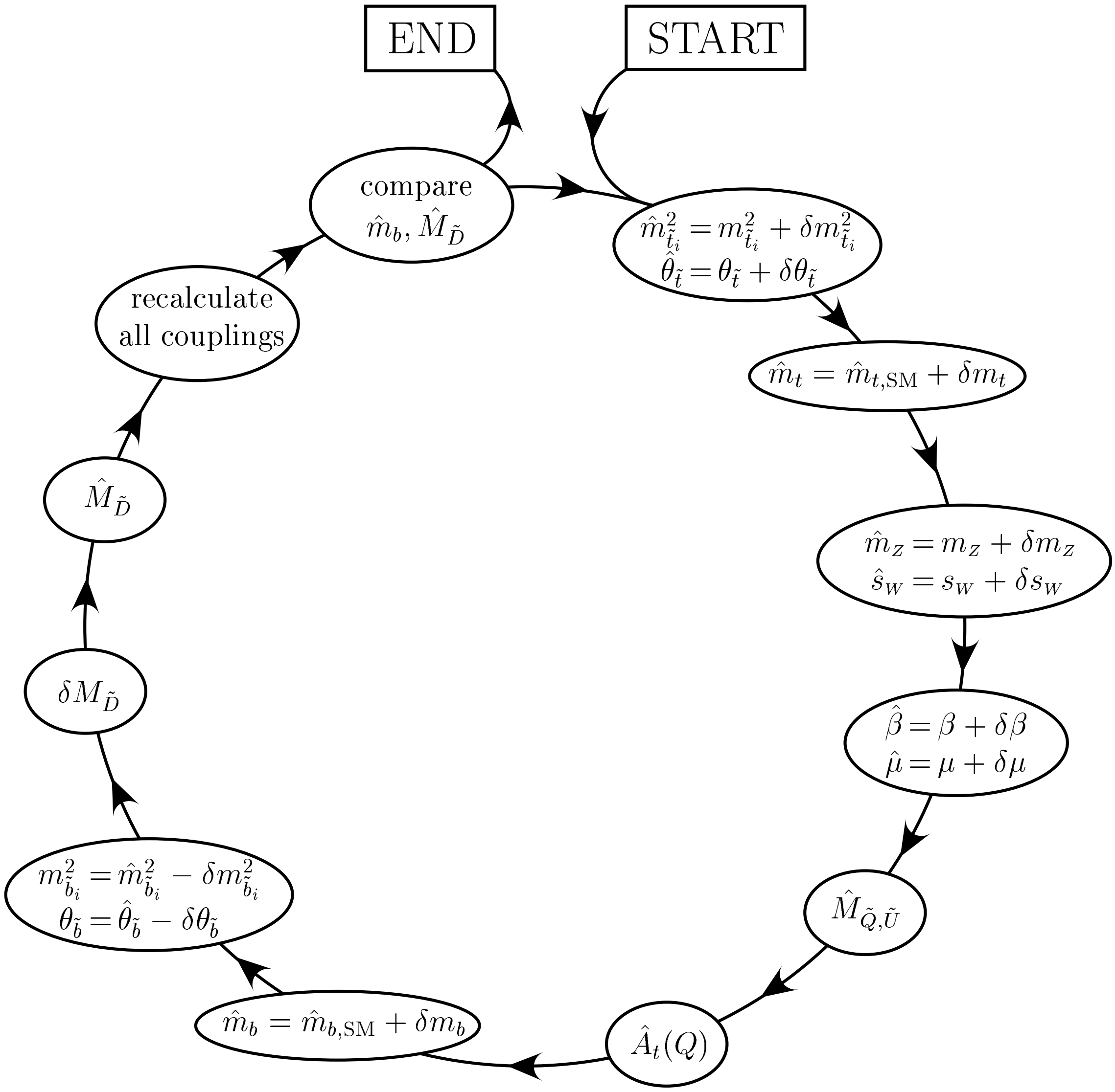


Figure 4.4: Simplified flowchart for the iteration procedure. For details see section 4.6.

4.7 Numerical results

In the following numerical examples, we take for the standard model parameters m_Z = 91.1875 GeV, m_W = 80.45 GeV, $\sin^2\theta_W = 1 - m_W^2/m_Z^2$, $\alpha(m_Z) = 1/127.7$, m_t = 178 GeV, m_b = 4.7 GeV, m_τ = 1.777 GeV and $\{m_u, m_d, m_e, m_c, m_s, m_\mu\} = \{5.38, 5.38, 0.511, 1500, 150, 106\}$ MeV for 1st and 2nd generation fermions. M' is fixed by the gaugino unification relation $M' = \frac{5}{3}\tan^2\theta_W M$, therefore the gluino mass is related to M by $m_{\tilde{g}} = (\alpha_s(m_{\tilde{g}})/\alpha)\sin^2\theta_W M$. In order to reduce the number of parameters in the input parameter set, we assume $M_{\tilde{Q}} \equiv M_{\tilde{Q}_3} = \frac{10}{9}M_{\tilde{U}_3} = \frac{10}{11}M_{\tilde{D}_3} = M_{\tilde{L}_3} = M_{\tilde{E}_3} = M_{\tilde{Q}_{1,2}} = M_{\tilde{U}_{1,2}} = M_{\tilde{D}_{1,2}} = M_{\tilde{L}_{1,2}} = M_{\tilde{E}_{1,2}}$ for the first, second and third generation soft SUSY–breaking masses as well as $A \equiv A_t = A_b = A_\tau$ for all (s)fermion generations, if not stated otherwise.

4.7.1 Decay processes into the light Higgs boson

Fig. 4.5 shows the tree–level, the full electroweak and the full one–loop corrected (electroweak and SUSY–QCD) decay width of $\tilde{t}_2 \to \tilde{t}_1 h^0$ as a function of the lighter stop mass, $m_{\tilde{t}_1}$, where $M_{\tilde{Q}_3}$ is varied from 300 to 700 GeV. To get a larger mass splitting for the top squarks, we relax the conditions for 3rd generation squarks and take $M_{\tilde{U}_3}$ = 500 GeV. All other SUSY–breaking masses are fixed at 300 GeV. For the remaining input parameters we choose $\{m_{A^0}, \mu, M, A\} = \{120, 1000, 250, -800\}$ GeV and $\tan\beta = 20$. At $m_{\tilde{t}_1}$ = 203 GeV and $m_{\tilde{t}_1}$ = 429 GeV one can identify two pseudo–thresholds originating from $\tilde{t}_2 \to \tilde{b}_2 H^+$ and $\tilde{t}_2 \to \tilde{b}_1 H^+$ in the wave–function correction, respectively.

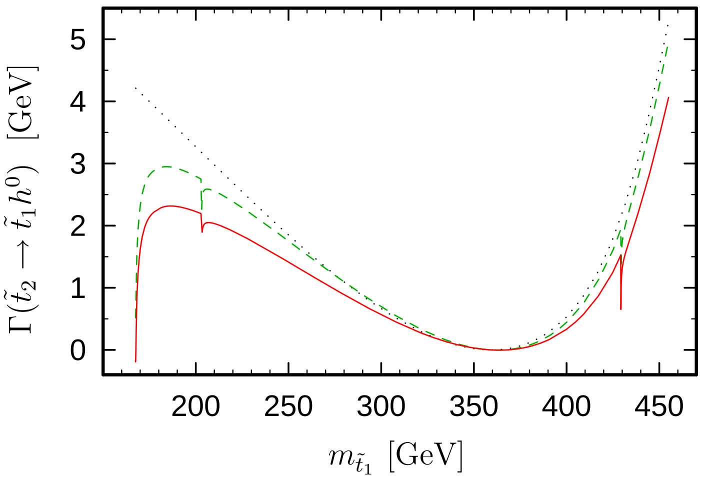


Figure 4.5: Tree–level (dotted line), full electroweak corrected decay width (dashed line) and full one–loop (electroweak and SUSY–QCD) corrected width (solid line) of $\tilde{t}_2 \to \tilde{t}_1 h^0$ as a function of $m_{\tilde{t}_1}$.

In Fig. 4.6 the radiative corrections to the decay width of $\tilde{t}_2 \to \tilde{t}_1 h^0$ as a function of $\tan\beta$ are quite constant in the considered region of the parameter space and make the perturbation expansion reliable also for large $\tan\beta$. The dotted, dashed and solid lines correspond to the tree–level, full electroweak corrected and full one–loop corrected decay widths, respectively.

Whereas the SUSY–QCD corrections reduce the decay width up to -10%, the electroweak corrections stay well below 3%. At $\tan\beta \approx 20.5$ one can see a pseudo–threshold $\tilde{t}_2 \to \tilde{\chi}^0_3 t$ entering in the sfermion wave–function corrections. As input parameters we take the values $\{m_{A^0}, \mu, A, M, M_{\tilde{Q}}\} = \{120, 450, -300, 300, 250\}$ GeV as well as $M_{\tilde{U}_3} = 600$ GeV for kinematical reasons.

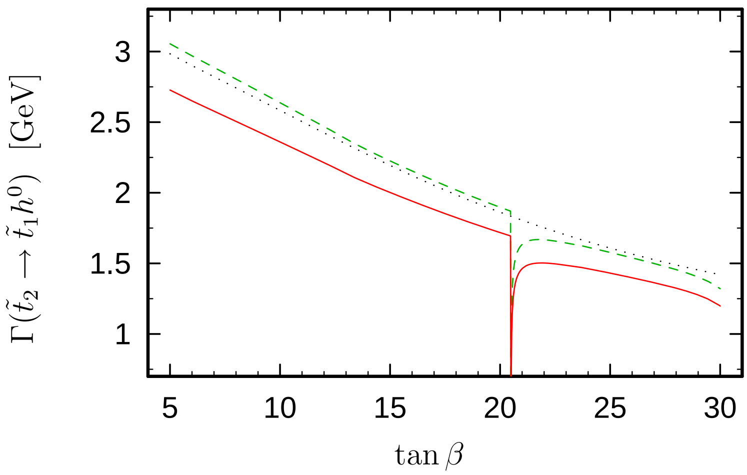


Figure 4.6: Tree–level (dotted line), full electroweak corrected decay width (dashed line) and full one–loop (electroweak and SUSY–QCD) corrected width (solid line) of $\tilde{t}_2 \to \tilde{t}_1 h^0$ as a function of $\tan\beta$.

In Fig. 4.7 we show two kinds of perturbation expansion for $\Gamma(\tilde{b}_2 \to \tilde{b}_1 h^0)$ with $\{m_{A^0}, \mu, A, M, M_{\tilde{Q}}\} = \{200, 550, 450, 200, 300\}$ GeV and $M_{\tilde{D}_3} = 600$ GeV: First we show the tree–level width, given in terms of on–shell input parameters (dotted line). The dashed and dash–dot–dotted lines correspond to the on–shell electroweak and SUSY–QCD corrected one–loop widths, respectively. For both corrections one can clearly see the invalidity of the on–shell perturbation expansion, which leads to an improper high decay width. The second way of perturbation expansion is given by the dash–dotted and the solid line which correspond to the improved tree–level and improved full one–loop decay width, respectively. Here we take the same input parameters as in the first case but with running $A = 450$ GeV. It is clearly seen that one needs the improvement as described in section 4.5 and 4.6.

4.7.2 Decays of the CP–even heavy Higgs boson

In Fig. 4.8 we show the tree level and the corrected widths to $H^0 \to \tilde{t}_1\tilde{t}_1$ for $\tan\beta = 30$ and $\{M_{\tilde{Q}}, A, M, \mu\} = \{250, 300, 120, 300\}$ GeV as a function of the mass of the CP–odd Higgs boson, m_{A^0}. As can be seen for larger values of m_{A^0}, the radiative corrections decrease and make the electroweak corrections comparable to the SUSY–QCD ones. Due to top and bottom squark loops in the Higgs wave–function corrections, two pseudo–thresholds $H^0 \to \tilde{b}_2\tilde{b}_2$, $H^0 \to \tilde{t}_2\tilde{t}_2$ appear at $m_{A^0} \approx 671$ and $m_{A^0} \approx 745$ GeV, respectively.
In Fig. 4.9 the dependence of the decay width $\Gamma(H^0 \to \tilde{t}_1\tilde{t}_1)$ as a function of $\tan\beta$ is given. In the considered region the electroweak corrections have different sign compared to the SUSY–QCD ones and, for larger values of $\tan\beta$, are the dominant contributions. Again,

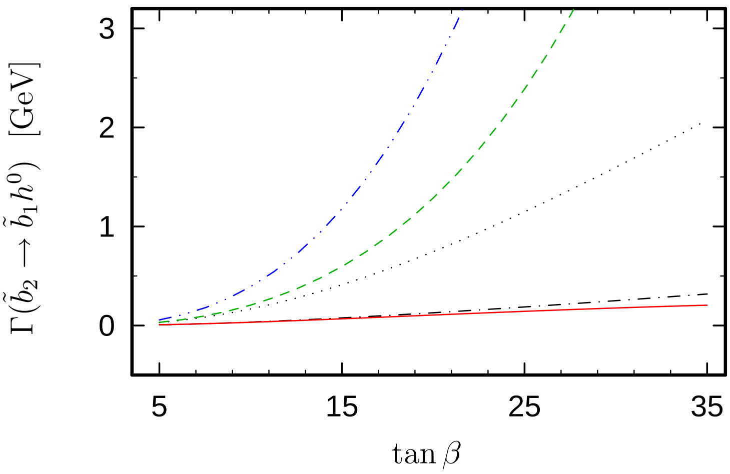


Figure 4.7: Two kinds of perturbation expansion: the dotted line corresponds to the on–shell tree–level width, the dashed and dash–dot–dotted lines correspond to electroweak and SUSY–QCD corrected on–shell one–loop width, respectively. The dash–dotted line corresponds to the improved tree–level, and the solid line to the (full) improved one–loop width.

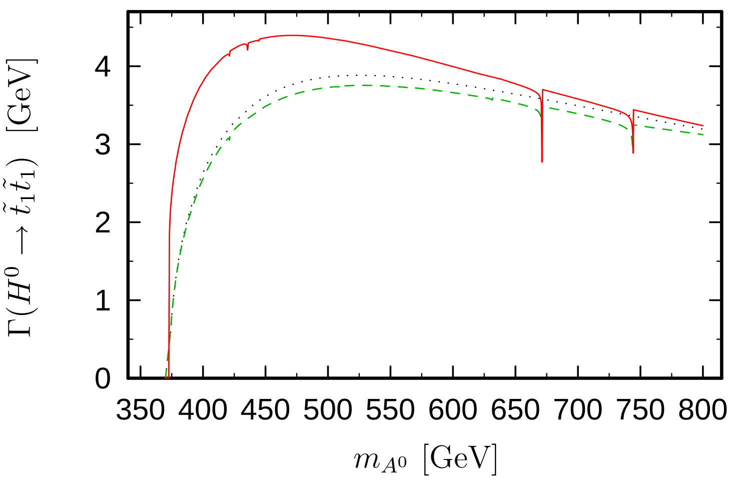


Figure 4.8: Tree–level (dotted line), full electroweak corrected decay width (dashed line) and full one–loop (electroweak and SUSY–QCD) corrected width (solid line) of $H^0 \to \tilde{t}_1 \tilde{t}_1$ as a function of m_{A^0}.

sbottom loops entering in the Higgs wave–function corrections cause a pseudo–threshold $H^0 \to \tilde{b}_2 \tilde{b}_2$ for $\tan\beta \approx 25.6$. As input parameters we have chosen $\{m_{A^0}, \mu, A, M, M_{\tilde{Q}}\} = \{800., 500, 550, 120, 300\}$ GeV.
As in the case of $\Gamma(\tilde{b}_2 \to \tilde{b}_1 h^0)$, we show two kinds of perturbation expansion for $\Gamma(H^0 \to \tilde{b}_1 \tilde{b}_1)$ in Fig. 4.10. Owing to the extremely large counter–term corrections δm_b and δA_b in the on–shell scheme, the perturbation expansion around the on–shell tree level leads to unacceptably big corrections even for lower values of $\tan\beta$ and, moreover, to improper negative decay widths. The dotted, dashed and dash–dot–dotted lines, which correspond to the on–

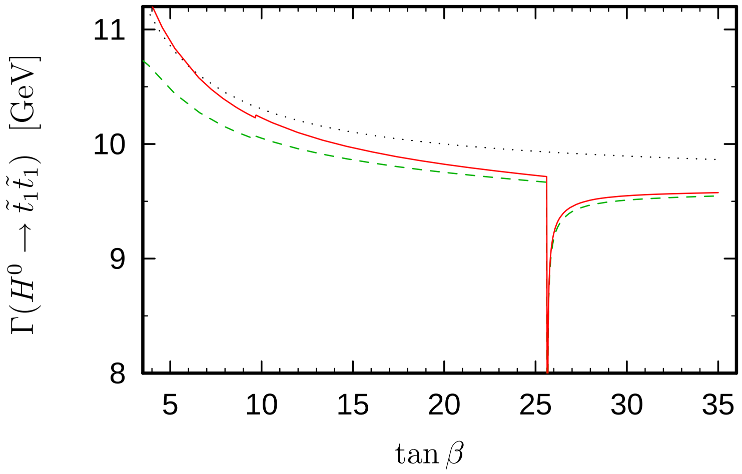


Figure 4.9: Tree–level (dotted line), full electroweak corrected decay width (dashed line) and full one–loop (electroweak and SUSY–QCD) corrected width (solid line) of $H^0 \to \tilde{t}_1\tilde{t}_1$ as a function of $\tan\beta$.

shell tree–level, electroweak and SUSY–QCD corrected widths, respectively, show this first expansion using the input parameters $\{m_{A^0}, \mu, A, M, M_{\tilde{Q}}\} = \{800, -300, 300, 120, 300\}$ GeV. The second perturbation expansion, where these large counter–term corrections are absorbed into the tree level, is given by the dash–dotted and the solid line which correspond to the improved tree–level and improved full one–loop decay width, respectively. The input parameters are the same as above, but with $\overline{\text{DR}}$ running $A = 300$ GeV. The corrections are now relatively small indicating that the (improved) tree level is already a good approximation for the full one–loop corrected decay width.

4.7.3 Decays of the CP–odd neutral Higgs boson

In Fig. 4.11 we show the tree–level and the corrected widths to $A^0 \to \tilde{t}_1\tilde{t}_2$ for $\tan\beta = 15$ and $\{m_{A^0}, A, M, M_{\tilde{Q}}\} = \{700, -500, 120, 300\}$ GeV as a function of the higgsino mass parameter μ. The electroweak corrections are about 10–15% an thus comparable to the SUSY–QCD ones. At $\mu \approx -242$ GeV one can identify the pseudo–threshold $\tilde{t}_2 \to \tilde{\chi}^0_4 t$ coming from the sfermion wave–function corrections.
Fig. 4.12 shows the tree–level, the full electroweak and the full one–loop corrected decay width of $A^0 \to \tilde{t}_1\tilde{t}_2$ as a function of the lighter stop mass, $m_{\tilde{t}_1}$, where $M_{\tilde{Q}}$ is varied from 140 to 410 GeV. As input parameters we choose $\{m_{A^0}, \mu, A, M\} = \{900, 250, 300, 120\}$ GeV and $\tan\beta = 7$. Again, in a large region of the parameter space the electroweak corrections are comparable to the SUSY–QCD ones. The pseudo–threshold at $m_{\tilde{t}_1} \approx 306$ GeV originates from $\tilde{t}_2 \to t\,\tilde{\chi}^0_3$ in the wave–function correction.
Fig. 4.13 displays the decay widths of the crossed channel $\tilde{t}_2 \to \tilde{t}_1 A^0$ as a function of A_t. As can be seen, the electroweak corrections are as large as the SUSY–QCD ones in the considered region. The values of the input parameters are $\tan\beta = 35$, $\{m_{A^0}, \mu, M, M_{\tilde{Q}}\} = \{150, 240, 300, 300, \}$ GeV, and $A_f = 700$ GeV for all trilinear couplings except A_t. With the relations for the SUSY–breaking masses given at the top of this section but with $M_{\tilde{U}_3} =$

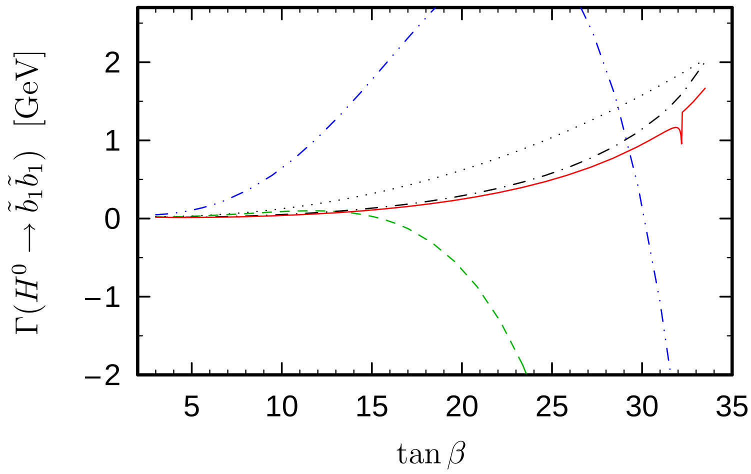


Figure 4.10: Two kinds of perturbation expansion: the dotted line corresponds to the on–shell tree–level width, the dashed and dash–dot–dotted lines correspond to electroweak and SUSY–QCD corrected on–shell one–loop width, respectively. The dash–dotted line corresponds to the improved tree–level, and the solid line to the (full) improved one–loop width.

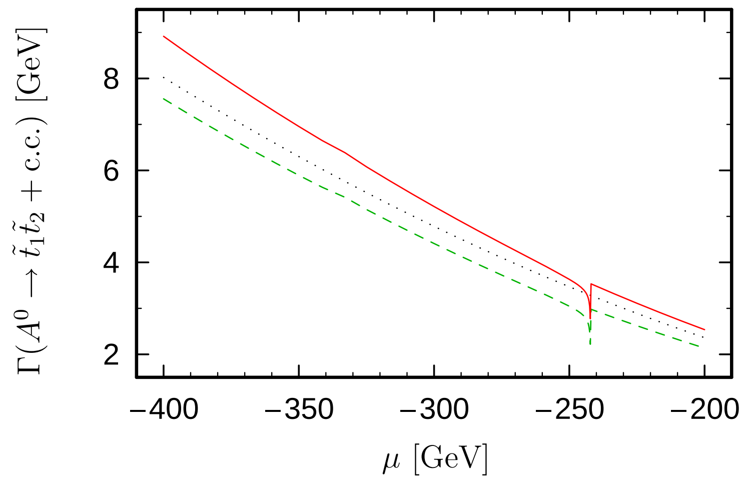


Figure 4.11: Tree–level (dotted line), full electroweak corrected decay width (dashed line) and full one–loop (electroweak and SUSY–QCD) corrected width (solid line) of $A^0 \to \tilde{t}_1\tilde{t}_2$ as a function of μ.

500 GeV we provide a quite acceptable mass splitting in the stop sector.
In Fig. 4.14 we show the decay width $\Gamma(A^0 \to \tilde{b}_1\tilde{b}_2)$ as a function of the trilinear couplings A for $\{m_{A^0}, \mu, M, M_{\tilde{Q}}\} = \{800, -300, 300, 300\}$ GeV and a large value for $\tan\beta$, $\tan\beta = 30$. As can be seen, the on–shell expansion dramatically suffers from bad convergence; the dotted line corresponds to the on–shell tree–level width, the dashed line to the electroweak and the dash–dot–dotted line to the SUSY–QCD one–loop width. After a redefinition of the tree level in terms of $\overline{\text{DR}}$ running m_b and A_b, the corrections lie in an acceptable range.
Fig. 4.15 shows the behaviour of the decay width $\Gamma(\tilde{b}_2 \to \tilde{b}_1 A^0)$ for large $\tan\beta$ for two kinds of

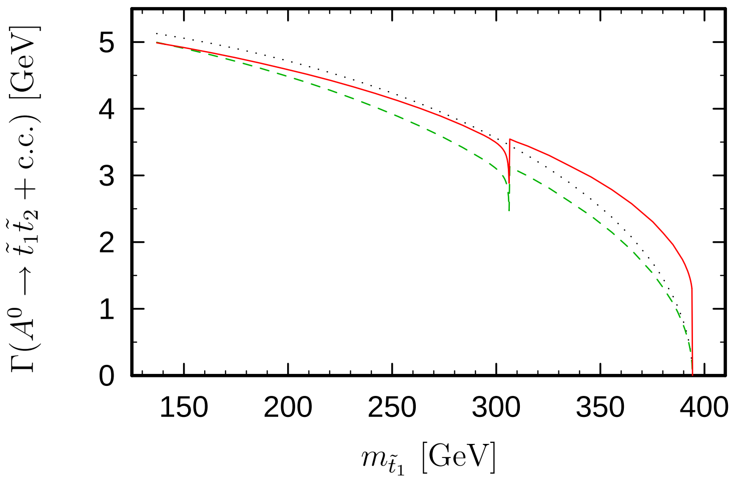


Figure 4.12: Tree–level (dotted line), full electroweak corrected decay width (dashed line) and full one–loop (electroweak and SUSY–QCD) corrected width (solid line) of $A^0 \to \tilde{t}_1\tilde{t}_2$ as a function of $m_{\tilde{t}_1}$.

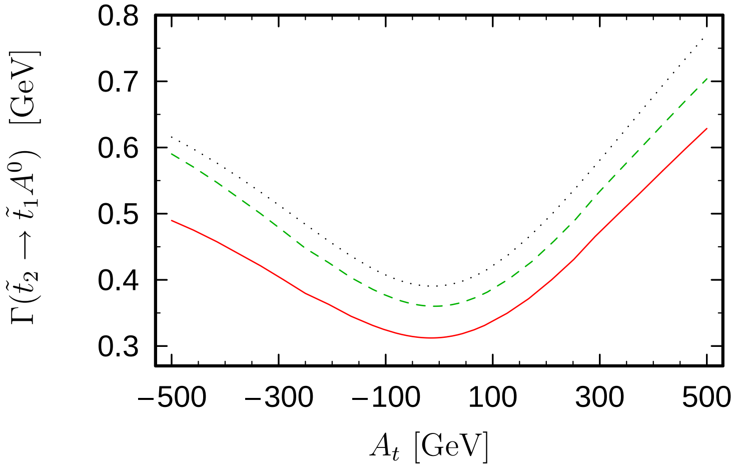


Figure 4.13: A_t–dependence of the tree–level (dotted line), full electroweak corrected (dashed line) and full one–loop corrected (solid line) decay widths of $\tilde{t}_2 \to \tilde{t}_1 A^0$.

perturbation expansion. The dotted and dash–dot–dotted lines correspond to the tree–level and full one–loop decay widths in the pure on–shell scheme. For large $\tan\beta$ one can clearly see the invalidity of the perturbation series, leading to a negative decay width. In the second case we show the expansion around the tree–level decay width, given in terms of $\overline{\text{DR}}$ running A_b and m_b. The dash–dotted line corresponds to the improved tree–level and the solid one to the one–loop decay width. Up to $\tan\beta \sim 35$ the corrections stay relatively small which indicates that already the (improved) tree level is a good approximation for $\Gamma(\tilde{b}_2 \to \tilde{b}_1 A^0)$. As input parameters we take the values $\{m_{A^0}, \mu, A, M, M_{\tilde{Q}}\} = \{150, -220, 500, 200, 300\}$ GeV as well as $M_{\tilde{D}_3} = 500$ GeV for kinematical reasons.

In Fig. 4.16 the A^0 decay into two staus is given as a function of $\tan\beta$. Despite the absence

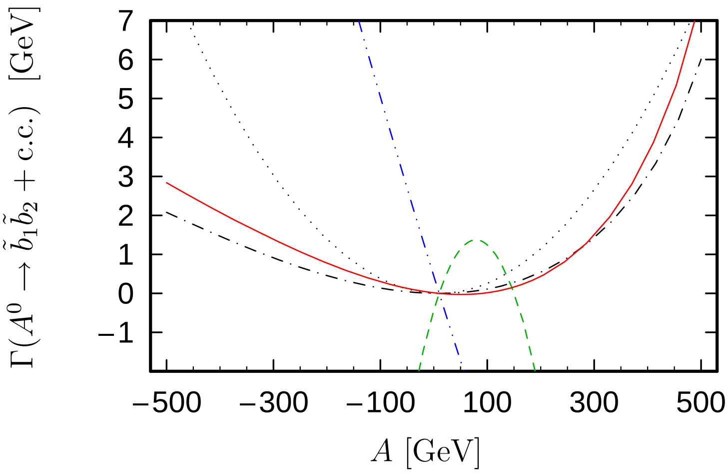


Figure 4.14: Full one–loop corrections to $A^0 \to \tilde{b}_1\tilde{b}_2$ for two kinds of perturbation expansion depending on the trilinear couplings A. The dotted, dashed and dash–dot–dotted lines correspond to the on–shell tree–level, electroweak and SUSY–QCD corrected decay widths, respectively. The dash–dotted line corresponds to the improved tree–level, and the solid line to the improved one–loop corrected width.

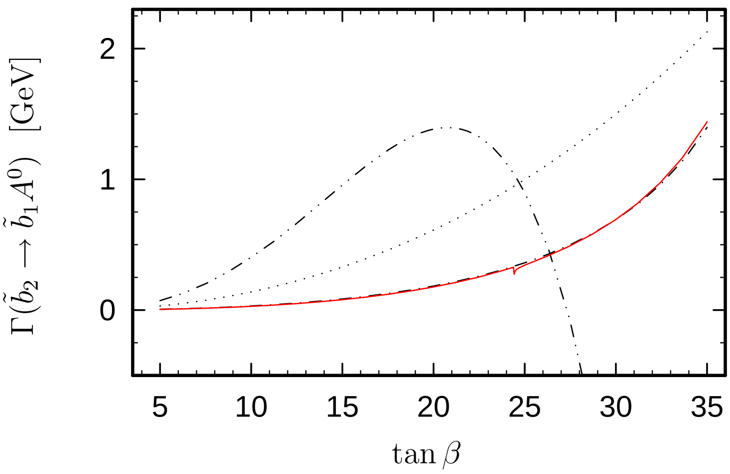


Figure 4.15: $\tan\beta$–dependence of $\Gamma(\tilde{b}_2 \to \tilde{b}_1 A^0)$ for two kinds of perturbation expansion. The dotted and dash–dot–dotted lines corresponds to on–shell tree–level and full one–loop width, respectively, the dash–dotted line corresponds to improved tree–level and the solid line shows the full improved one–loop width.

of SUSY–QCD corrections the perturbation expansion around the on–shell tree level (dotted line) leads to an improper negative decay width (dashed line) coming from large $\mathcal{O}(h_b^2)$ corrections. As input parameters we take $\{m_{A^0}, \mu, A, M, M_{\tilde{Q}}\} = \{800, 400, -500, 120, 300\}$ GeV. The dash–dotted line corresponds to the improved tree–level and the solid line shows the improved one–loop width for the same input parameters as above and running $A = -500$ GeV.

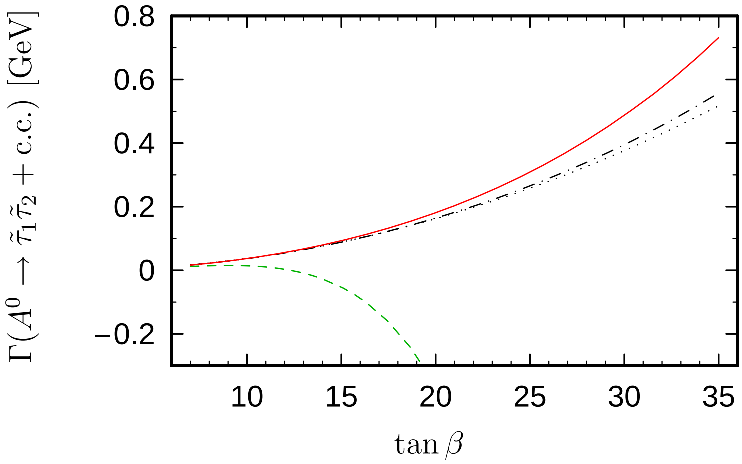


Figure 4.16: On–shell tree–level (dotted line) and full electroweak on–shell corrected decay width (dashed line) of $A^0 \to \tilde{\tau}_1\tilde{\tau}_2$ as a function of $\tan\beta$. The dash–dotted and solid lines correspond to improved tree–level and full improved one–loop decay widths, respectively.

4.7.4 Decays of the charged Higgs boson

Finally, we show two plots of the decay widths of the charged Higgs boson H^+ into a top and a bottom squark. Again, we show two kinds of perturbation expansion for $\Gamma(H^+ \to \tilde{t}_1\tilde{b}_1)$ in Fig. 4.17, showing the invalidity of the on–shell scheme for large values of $\tan\beta$ as well as the expansion around the improved tree level. The dotted, dashed and dash–dot–dotted lines, which correspond to the on–shell tree–level, electroweak and SUSY–QCD corrected widths, respectively, show this first expansion using the input parameters $\{m_{A^0}, \mu, A, M, M_{\tilde{Q}}\} = \{800, -260, 150, 120, 300\}$ GeV and $A_b = -700$ GeV. The second perturbation expansion is given by the dash–dotted and the solid line which correspond to the improved tree–level and improved full one–loop decay width, respectively.
In Fig. 4.18 we show the decay width $\Gamma(H^+ \to \tilde{t}_1\tilde{b}_2)$ as a function of the trilinear couplings A for $\{m_{A^0}, \mu, M, M_{\tilde{Q}}\} = \{800, 250, 300, 300\}$ GeV and $\tan\beta = 20$. As in the case of the decay of CP–odd Higgs boson A^0, see Fig. 4.14, the on–shell expansion dramatically suffers from bad convergence; the dotted line corresponds to the on–shell tree–level width, the dashed line to the electroweak and the dash–dot–dotted line to the SUSY–QCD one–loop width. After a redefinition of the tree level in terms of $\overline{\text{DR}}$ running m_b and A_b, the corrections are visibly smaller.

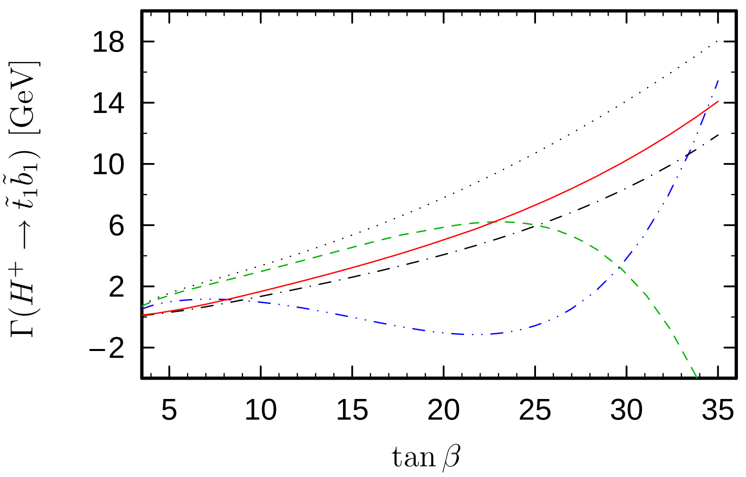


Figure 4.17: Two kinds of perturbation expansion: the dotted line corresponds to the on–shell tree–level width, the dashed and dash–dot–dotted lines correspond to electroweak and SUSY–QCD corrected on–shell one–loop width, respectively. The dash–dotted line corresponds to the improved tree–level, and the solid line to the (full) improved one–loop width.

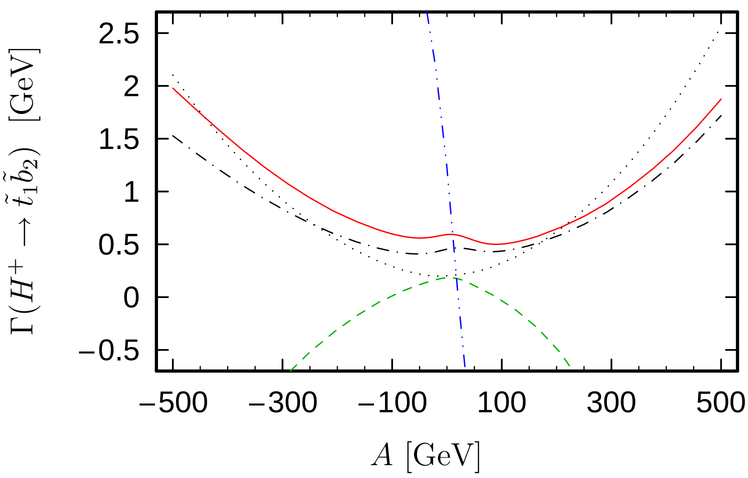


Figure 4.18: Full one–loop corrections to $H^+ \to \tilde{t}_1\tilde{b}_2$ for two kinds of perturbation expansion depending on the trilinear couplings A. The dotted, dashed and dash–dot–dotted lines correspond to the on–shell tree–level, electroweak and SUSY–QCD corrected decay widths, respectively. The dash–dotted line corresponds to the improved tree–level, and the solid line to the improved one–loop corrected width.

Appendix A

Electroweak Interactions in the MSSM

In this chapter we give all couplings which are necessary for the calculation of electroweak corrections to Higgs decay processes into sfermions. For the whole set of Feynman rules and a complete list of all terms of the MSSM Lagrangian we refer to [43].

A.1 Higgs–Sfermion–Sfermion couplings

For the neutral and charged Higgs fields we use the notation $H^0_k = \{h^0, H^0, A^0, G^0\}$, $H^+_k = \{H^+, G^+, H^-, G^-\}$ and $H^-_k \equiv (H^+_k)^* = \{H^-, G^-, H^+, G^+\}$. $t/\tilde{t}$ stands for an up–type (s)fermion and $b/\tilde{b}$ for a down–type one. Following [28, 29] the Higgs–sfermion–sfermion couplings for neutral Higgs bosons, $G^{\tilde{f}}_{ijk}$, can be written as

$$G^{\tilde{f}}_{ijk} \equiv G\big(H^0_k \tilde{f}^*_i \tilde{f}_j\big) = \Big[R^{\tilde{f}} G^{\tilde{f}}_{LR,k} (R^{\tilde{f}})^T\Big]_{ij} . \tag{A.1}$$

The 3rd generation left–right couplings $G^{\tilde{f}}_{LR,k}$ for up– and down–type sfermions are given by

$$G^{\tilde{t}}_{LR,1} = \begin{pmatrix} -\sqrt{2}h_t m_t c_\alpha + g_Z m_Z (I^{3L}_t - e_t s^2_W) s_{\alpha+\beta} & -\frac{h_t}{\sqrt{2}}(A_t\, c_\alpha + \mu s_\alpha) \\ -\frac{h_t}{\sqrt{2}}(A_t\, c_\alpha + \mu s_\alpha) & -\sqrt{2}h_t m_t c_\alpha + g_Z m_Z e_t s^2_W s_{\alpha+\beta} \end{pmatrix} , \tag{A.2}$$

$$G^{\tilde{b}}_{LR,1} = \begin{pmatrix} \sqrt{2}h_b m_b s_\alpha + g_Z m_Z (I^{3L}_b - e_b s^2_W) s_{\alpha+\beta} & \frac{h_b}{\sqrt{2}}(A_b\, s_\alpha + \mu c_\alpha) \\ \frac{h_b}{\sqrt{2}}(A_b\, s_\alpha + \mu c_\alpha) & \sqrt{2}h_b m_b s_\alpha + g_Z m_Z e_b s^2_W s_{\alpha+\beta} \end{pmatrix} , \tag{A.3}$$

$$G^{\tilde{f}}_{LR,2} = G^{\tilde{f}}_{LR,1} \qquad \text{with } \alpha \to \alpha - \pi/2 , \tag{A.4}$$

$$G^{\tilde{t}}_{LR,3} = -\sqrt{2}h_t \begin{pmatrix} 0 & -\frac{i}{2}\Big(A_t\, c_\beta + \mu\, s_\beta\Big) \\ \frac{i}{2}\Big(A_t\, c_\beta + \mu\, s_\beta\Big) & 0 \end{pmatrix} , \tag{A.5}$$

$$G^{\tilde{b}}_{LR,3} = -\sqrt{2}h_b \begin{pmatrix} 0 & -\frac{i}{2}\Big(A_b\, s_\beta + \mu\, c_\beta\Big) \\ \frac{i}{2}\Big(A_b\, s_\beta + \mu\, c_\beta\Big) & 0 \end{pmatrix} , \tag{A.6}$$

$$G^{\tilde{f}}_{LR,4} \;=\; G^{\tilde{f}}_{LR,3} \qquad \text{with } \beta \to \beta - \pi/2\,, \tag{A.7}$$

where we have used the abbreviations $s_x \equiv \sin x$, $c_x \equiv \cos x$ and $s_W \equiv \sin\theta_W$. α denotes the mixing angle of the $\{h^0, H^0\}$–system, and h_t and h_b are the Yukawa couplings

$$h_t \;=\; \frac{g\, m_t}{\sqrt{2} m_W \sin\beta}\,, \qquad h_b \;=\; \frac{g\, m_b}{\sqrt{2} m_W \cos\beta}\,. \tag{A.8}$$

The couplings of charged Higgs bosons to two sfermions are given by ($l = 1, 2$)

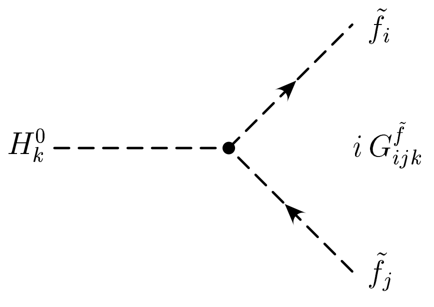
$$G^{\tilde{f}\tilde{f}'}_{ijl} \;\equiv\; G\Big(H^{\pm}_l \tilde{f}^*_i \tilde{f}'_j\Big) \;=\; G^{\tilde{f}'\tilde{f}}_{jil} \;=\; \left(R^{\tilde{f}}\, G^{\tilde{f}\tilde{f}'}_{LR,l} \left(R^{\tilde{f}'}\right)^T\right)_{ij}\,, \tag{A.9}$$

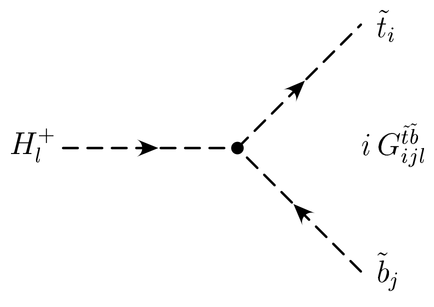
$$G^{\tilde{t}\tilde{b}}_{LR,1} \;=\; \begin{pmatrix} h_b m_b \sin\beta + h_t m_t \cos\beta - \frac{g m_W}{\sqrt{2}} \sin 2\beta & h_b(A_b \sin\beta + \mu\cos\beta) \\ h_t(A_t\cos\beta + \mu \sin\beta) & h_t m_b \cos\beta + h_b m_t \sin\beta \end{pmatrix}, \tag{A.10}$$

$$G^{\tilde{b}\tilde{t}}_{LR,1} \;=\; \begin{pmatrix} h_b m_b \sin\beta + h_t m_t \cos\beta - \frac{g m_W}{\sqrt{2}} \sin 2\beta & h_t(A_t \cos\beta + \mu\sin\beta) \\ h_b(A_b\sin\beta + \mu \cos\beta) & h_t m_b \cos\beta + h_b m_t \sin\beta \end{pmatrix}, \tag{A.11}$$

$$G^{\tilde{f}\tilde{f}'}_{LR,2} \;=\; G^{\tilde{f}\tilde{f}'}_{LR,1} \qquad \text{with} \quad \beta \;\to\; \beta - \frac{\pi}{2}\,. \tag{A.12}$$

f' denotes the isospin partner of the fermion f, i.e. $t' = b$, $\tilde{b}'_i = \tilde{t}_i$ etc. Note that only the angle β explicitly given in the matrices above has to be substituted; the dependence of β in the Yukawa couplings has to remain the same.
The Feynman rules for the Higgs–sfermion–sfermion couplings are (for $k = 1, \ldots, 4$; $l = 1, 2$)





A.2 Higgs–Fermion–Fermion couplings

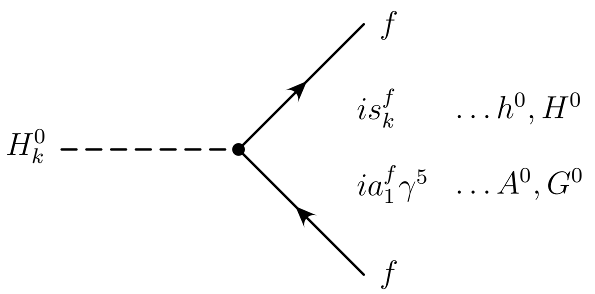
For the Higgs–fermion–fermion couplings the interaction Lagrangian reads

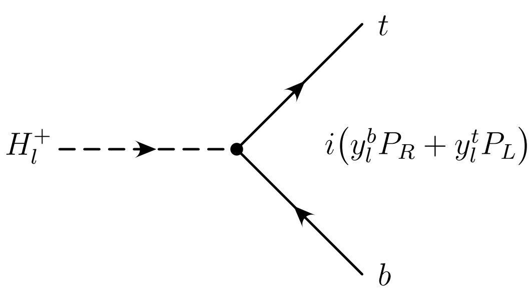
$$\mathcal{L} = \sum_{k=1}^{2} s^f_k\, H^0_k \bar{f} f + \sum_{k=3}^{4} s^f_k\, H^0_k \bar{f}\gamma^5 f + \sum_{l=1}^{2} \Big[H^+_l \bar{t}\, \big(y^b_l P_R + y^t_l P_L\big)\, b + \text{h.c.}\Big] \tag{A.13}$$

with the couplings

$$
\begin{aligned}
s_1^t &= -\,g\tfrac{m_t\cos\alpha}{2m_W\sin\beta} = -\,\tfrac{h_t}{\sqrt{2}}\cos\alpha, & s_1^b &= g\tfrac{m_b\sin\alpha}{2m_W\cos\beta} = \tfrac{h_b}{\sqrt{2}}\sin\alpha\,,\\
s_2^t &= -\,g\tfrac{m_t\sin\alpha}{2m_W\sin\beta} = -\,\tfrac{h_t}{\sqrt{2}}\sin\alpha, & s_2^b &= -\,g\tfrac{m_b\cos\alpha}{2m_W\cos\beta} = -\,\tfrac{h_b}{\sqrt{2}}\cos\alpha\,,\\
s_3^t &= ig\tfrac{m_t\cot\beta}{2m_W} = i\tfrac{h_t}{\sqrt{2}}\cos\beta, & s_3^b &= ig\tfrac{m_b\tan\beta}{2m_W} = i\tfrac{h_b}{\sqrt{2}}\sin\beta\,,\\
s_4^t &= ig\tfrac{m_t}{2m_W} = i\tfrac{h_t}{\sqrt{2}}\sin\beta, & s_4^b &= -\,ig\tfrac{m_b}{2m_W} = -\,i\tfrac{h_b}{\sqrt{2}}\cos\beta\,,\\
y_1^t &= g\tfrac{m_t\cot\beta}{\sqrt{2}m_W} = h_t\cos\beta, & y_1^b &= g\tfrac{m_b\tan\beta}{\sqrt{2}m_W} = h_b\sin\beta\,,\\
y_2^t &= g\tfrac{m_t}{\sqrt{2}m_W} = h_t\sin\beta, & y_2^b &= -\,g\tfrac{m_b}{\sqrt{2}m_W} = -\,h_b\cos\beta\,.
\end{aligned}
\tag{A.14}
$$

The Feynman rules for the couplings to the Higgs bosons are





A.3 Higgs–Gaugino–Gaugino couplings

The interaction Lagrangian for Higgs bosons and gauginos is given by

$$
\begin{aligned}
\mathcal{L} = & -\frac{g}{2}\sum_{k=1}^{2} H_k^0\,\bar{\tilde{\chi}}_l^0 F_{lmk}^0\,\tilde{\chi}_m^0 - i\frac{g}{2}\sum_{k=3}^{4} H_k^0\,\bar{\tilde{\chi}}_l^0\,F_{lmk}^0\gamma_5\,\tilde{\chi}_m^0\\
& -g\sum_{k=1}^{2} H_k^0\,\bar{\tilde{\chi}}_i^+\left(F_{ijk}^+P_R + F_{jik}^+P_L\right)\tilde{\chi}_j^+ + ig\sum_{k=3}^{4} H_k^0\,\bar{\tilde{\chi}}_i^+\left(F_{ijk}^+P_R + F_{jik}^+P_L\right)\tilde{\chi}_j^+\\
& -g\sum_{k=1}^{2}\Big[H_k^+\,\bar{\tilde{\chi}}_i^+\left(F_{ilk}^R P_R + F_{ilk}^L P_L\right)\tilde{\chi}_l^0 + \text{h.c.}\Big]\,.
\end{aligned}
\tag{A.15}
$$

with

$$
\begin{aligned}
F_{lmk}^0 = & \frac{e_k}{2}\Big[Z_{l3}Z_{m2} + Z_{m3}Z_{l2} - \tan\theta_W\left(Z_{l3}Z_{m1} + Z_{m3}Z_{l1}\right)\Big]\\
& + \frac{d_k}{2}\Big[Z_{l4}Z_{m2} + Z_{m4}Z_{l2} - \tan\theta_W\left(Z_{l4}Z_{m1} + Z_{m4}Z_{l1}\right)\Big] = F_{mlk}^0\,,
\end{aligned}
\tag{A.16}
$$

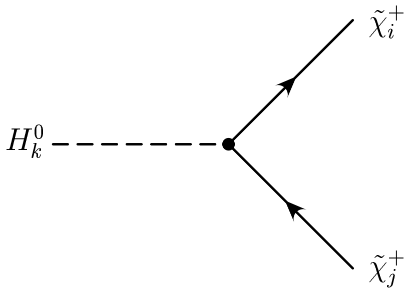
$$
F_{ijk}^+ = \frac{1}{\sqrt{2}}\left(e_k V_{i1}U_{j2} - d_k V_{i2}U_{j1}\right)\,, \tag{A.17}
$$

and

$$F^R_{ilk} = d_{k+2}\Big[V_{i1}Z_{l4} + \frac{1}{\sqrt{2}}(Z_{l2} + Z_{l1}\tan\theta_W)V_{i2}\Big]\,,$$
$$F^L_{ilk} = -e_{k+2}\Big[U_{i1}Z_{l3} - \frac{1}{\sqrt{2}}(Z_{l2} + Z_{l1}\tan\theta_W)U_{i2}\Big]\,. \qquad \text{(A.18)}$$

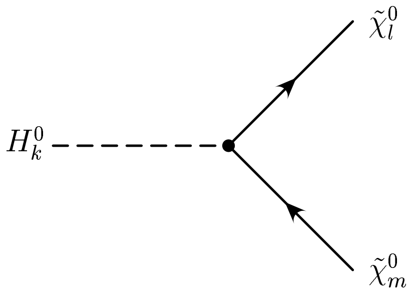
U, V and Z are rotation matrices which diagonalize the chargino and neutralino mass matrices (see chapter 2.2.3), and d_k and e_k take the values

$$d_k = \{-\cos\alpha, -\sin\alpha, \cos\beta, \sin\beta\}\,, \qquad e_k = \{-\sin\alpha, \cos\alpha, -\sin\beta, \cos\beta\}\,.$$



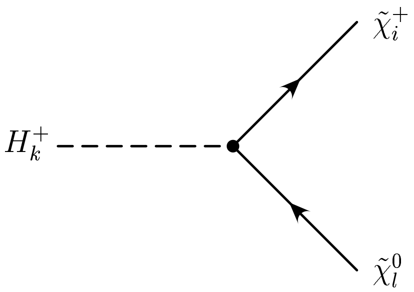
$-ig\left(F^+_{ijk}P_R + F^+_{jik}P_L\right)$ $\ldots h^0, H^0$

$-g\left(F^+_{ijk}P_R + F^+_{jik}P_L\right)$ $\ldots A^0, G^0$



$-igF^0_{lmk}$ $\ldots h^0, H^0$

$-gF^0_{lmk}(P_R - P_L)$ $\ldots A^0, G^0$



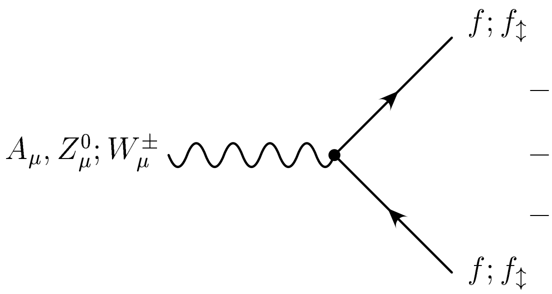
$-ig\left(F^R_{ilk}P_R + F^L_{ilk}P_L\right)$

A.4 Vector boson–Fermion–Fermion couplings

The Lagrangian describing interactions of vector bosons to fermions in the MSSM is given by

$$\mathcal{L} = -e\,e_f\,A_\mu\bar{f}\gamma^\mu f - g_Z Z^0_\mu\,\bar{f}\gamma^\mu(C^f_L P_L + C^f_R P_R)f$$
$$-\frac{g}{\sqrt{2}}\Big(W^+_\mu\bar{f}_\uparrow\gamma^\mu P_L f_\downarrow + W^-_\mu\bar{f}_\downarrow\gamma^\mu P_L f_\uparrow\Big)\,, \qquad \text{(A.19)}$$

where C_L^f and C_R^f are defined as $C_L^f = I_f^{3L} - e_f s_W^2$ and $C_R^f = -e_f s_W^2$. Here and in the following the arrows $\uparrow$ and $\downarrow$ attached at (s)fermions denote up– and down–type (s)fermions, respectively.



$$
\begin{array}{ll}
-ie\, e_f \gamma_\mu & \ldots A_\mu \\
-ig_Z \gamma_\mu (C_L^f P_L + C_R^f P_R) & \ldots Z_\mu^0 \\
-i\dfrac{g}{\sqrt{2}} \gamma_\mu P_L & \ldots W_\mu^\pm
\end{array}
$$

A.5 Vector boson–Gaugino–Gaugino couplings

The interaction of a vector boson with two gauginos is described by the Lagrangian

$$
\begin{aligned}
\mathcal{L} \;=\;& g\, \bar{\tilde{\chi}}_j^+ \gamma^\mu \Big(O_{ij}^L P_L + O_{ij}^R P_R \Big) \tilde{\chi}_i^0\, W_\mu^+ + g\, \bar{\tilde{\chi}}_i^0 \Big(O_{ij}^L P_L + O_{ij}^R P_R \Big) \tilde{\chi}_j^+\, W_\mu^- \\
& - e\, A_\mu\, \bar{\tilde{\chi}}_i^+ \gamma^\mu \tilde{\chi}_i^+ + g_Z Z_\mu^0\, \bar{\tilde{\chi}}_i^+ \gamma^\mu \Big(O_{ij}^{\prime L} P_L + O_{ij}^{\prime R} P_R \Big) \tilde{\chi}_j^+ \\
& + \frac{g_Z}{2} Z_\mu^0\, \bar{\tilde{\chi}}_i^0 \gamma^\mu \Big(O_{ij}^{\prime\prime L} P_L + O_{ij}^{\prime\prime R} P_R \Big) \tilde{\chi}_j^0 \,,
\end{aligned}
\tag{A.20}
$$

with the 4×2 coupling matrices for the $W^\pm$–chargino–neutralino vertex

$$
O_{ij}^L = Z_{i2} V_{j1} - \tfrac{1}{\sqrt{2}} Z_{i4} V_{j2} \,, \qquad O_{ij}^R = Z_{i2} U_{j1} + \tfrac{1}{\sqrt{2}} Z_{i3} U_{j2} \,, \tag{A.21}
$$

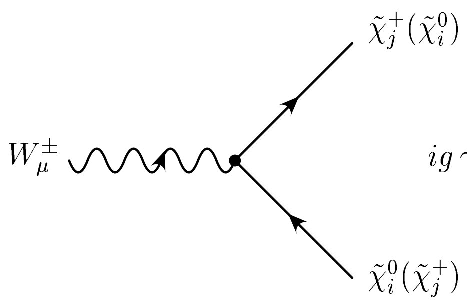
and the symmetric 2×2 and 4×4 coupling matrices

$$
O_{ij}^{\prime L} = -V_{i1} V_{j1} - \tfrac{1}{2} V_{i2} V_{j2} + \delta_{ij} s_W^2 \,, \tag{A.22}
$$

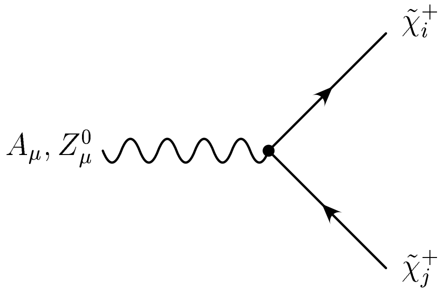
$$
O_{ij}^{\prime R} = -U_{i1} U_{j1} - \tfrac{1}{2} U_{i2} U_{j2} + \delta_{ij} s_W^2 \,, \tag{A.23}
$$

$$
O_{ij}^{\prime\prime L} = -\tfrac{1}{2} Z_{i3} Z_{j3} + \tfrac{1}{2} Z_{i4} Z_{j4} = -\, O_{ij}^{\prime\prime R} \,. \tag{A.24}
$$

Z_{ij}, U_{ij}, V_{ij} are the neutralino and chargino mixing matrices, respectively.

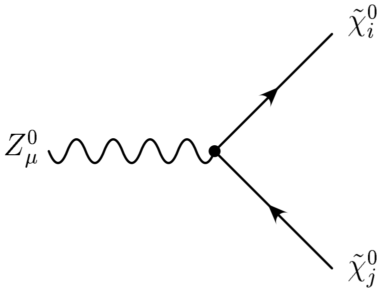


$$
ig\, \gamma_\mu \Big(O_{ij}^L P_L + O_{ij}^R P_R \Big)
$$



$-ie\,\delta_{ij}\gamma_\mu \qquad \ldots A_\mu$

$ig_Z\gamma_\mu\Big(O'^{L}_{ij}P_L + O'^{R}_{ij}P_R\Big) \qquad \ldots Z^0_\mu$



$-ig_Z\,\gamma_\mu\,O''^{L}_{ij}(P_R - P_L)$

A.6 Four–Scalar couplings ($F-$ and D–terms)

In this section we will derive all couplings involving four scalar particles in the MSSM. These interactions have two different sources originating from 'F–term' and 'D–term' contributions, which build up the complete scalar potential $V = V_F + V_D$ discussed in the following.

A.6.1 F–term potential

We start with the superpotential

$$W \;=\; h_t\,\tilde{t}^*_R\tilde{t}_L H^0_2 + h_b\,\tilde{b}^*_R\tilde{b}_L H^0_1 + h_\tau\,\tilde{\tau}^*_R\tilde{\tau}_L H^0_1 - h_t\,\tilde{t}^*_R\tilde{b}_L H^1_2 - h_b\,\tilde{b}^*_R\tilde{t}_L H^2_1 - h_\tau\,\tilde{\tau}^*_R\tilde{\nu}_\tau H^2_1$$

from which we can derive the F–term interaction potential $V_F = -\mathcal{L}_{int} = F_i^\dagger F_i$ with $F_i = \frac{\partial W}{\partial A_i}$, where A_i denotes all scalar (super)fields numbered by the index i, $A_i = \{\tilde{t}_L, \tilde{b}_L, \tilde{\nu}_\tau, \tilde{\tau}_L, \tilde{t}^*_R, \tilde{b}^*_R, \tilde{\tau}^*_R, H^0_1, H^0_2,\, H^2_1, H^1_2\}$.
The potential V_F then reads

$$\begin{aligned}
V_F \;=\; & \big(h_t\,\tilde{t}_R H^{0*}_2 - h_b\,\tilde{b}_R H^{2*}_1\big)\big(h_t\,\tilde{t}^*_R H^0_2 - h_b\,\tilde{b}^*_R H^2_1\big) + \big(h_b\,\tilde{b}_R H^{0*}_1 - h_t\,\tilde{t}_R H^{1*}_2\big)\big(h_b\,\tilde{b}^*_R H^0_1 - h_t\,\tilde{t}^*_R H^1_2\big) \\
& + h_t^2\big(\tilde{t}^*_L H^{0*}_2 - \tilde{b}^*_L H^{1*}_2\big)\big(\tilde{t}_L H^0_2 - \tilde{b}_L H^1_2\big) + h_b^2\big(\tilde{b}^*_L H^{0*}_1 - \tilde{t}^*_L H^{2*}_1\big)\big(\tilde{b}_L H^0_1 - \tilde{t}_L H^2_1\big) \\
& + h_\tau^2\,\tilde{\tau}_R H^{2*}_1 \tilde{\tau}^*_R H^2_1 + h_\tau^2\,\tilde{\tau}_R H^{0*}_1 \tilde{\tau}^*_R H^0_1 + h_\tau^2\big(\tilde{\tau}^*_L H^{0*}_1 - \tilde{\nu}^*_\tau H^{2*}_1\big)\big(\tilde{\tau}_L H^0_1 - \tilde{\nu}_\tau H^2_1\big) \\
& + \big(h_b\tilde{b}_R\tilde{b}^*_L + h_\tau\tilde{\tau}_R\tilde{\tau}^*_L\big)\big(h_b\tilde{b}^*_R\tilde{b}_L + h_\tau\tilde{\tau}^*_R\tilde{\tau}_L\big) + h_t^2\,\big(\tilde{t}_R\tilde{t}^*_L\big)\big(\tilde{t}^*_R\tilde{t}_L\big) \\
& + \big(h_b\tilde{b}_R\tilde{t}^*_L + h_\tau\tilde{\tau}_R\tilde{\nu}^*_\tau\big)\big(h_b\tilde{b}^*_R\tilde{t}_L + h_\tau\tilde{\tau}^*_R\tilde{\nu}_\tau\big) + h_t^2\,\big(\tilde{t}_R\tilde{b}^*_L\big)\big(\tilde{t}^*_R\tilde{b}_L\big) \\
\;=\; & h_t^2\,\big|H^0_2\big|^2\,\big(\tilde{t}^*_R\tilde{t}_R + \tilde{t}^*_L\tilde{t}_L\big) + h_b^2\,\big|H^0_1\big|^2\,\big(\tilde{b}^*_R\tilde{b}_R + \tilde{b}^*_L\tilde{b}_L\big) + h_\tau^2\,\big|H^0_1\big|^2\,\big(\tilde{\tau}^*_R\tilde{\tau}_R + \tilde{\tau}^*_L\tilde{\tau}_L\big)
\end{aligned}$$

$$
\begin{aligned}
&+ h_t^2 \left|H_2^1\right|^2 \left(\tilde{t}_R^* \tilde{t}_R + \tilde{b}_L^* \tilde{b}_L\right) + h_b^2 \left|H_1^2\right|^2 \left(\tilde{b}_R^* \tilde{b}_R + \tilde{t}_L^* \tilde{t}_L\right) + h_\tau^2 \left|H_1^2\right|^2 \left(\tilde{\tau}_R^* \tilde{\tau}_R + \tilde{\nu}_\tau^* \tilde{\nu}_\tau\right) \\
&- h_t^2 \left(\tilde{t}_L^* \tilde{b}_L H_2^{0*} H_2^1 + \tilde{b}_L^* \tilde{t}_L H_2^{1*} H_2^0\right) - h_b^2 \left(\tilde{b}_L^* \tilde{t}_L H_1^{0*} H_1^2 + \tilde{t}_L^* \tilde{b}_L H_1^{2*} H_1^0\right) \\
&- h_\tau^2 \left(\tilde{\tau}_L^* \tilde{\nu}_\tau H_1^{0*} H_1^2 + \tilde{\nu}_\tau^* \tilde{\tau}_L H_1^{2*} H_1^0\right) \\
&- h_t h_b \left(\tilde{t}_R^* \tilde{b}_R H_1^{2*} H_2^0 + \tilde{b}_R^* \tilde{t}_R H_2^{0*} H_1^2 + \tilde{b}_R^* \tilde{t}_R H_2^{1*} H_1^0 + \tilde{t}_R^* \tilde{b}_R H_1^{0*} H_2^1\right) \\
&+ h_t^2 \Big(\left(\tilde{t}_R^* \tilde{t}_L\right)\left(\tilde{t}_L^* \tilde{t}_R\right) + \left(\tilde{t}_R^* \tilde{b}_L\right)\left(\tilde{b}_L^* \tilde{t}_R\right)\Big) + h_b^2 \Big(\left(\tilde{b}_R^* \tilde{b}_L\right)\left(\tilde{b}_L^* \tilde{b}_R\right) + \left(\tilde{b}_R^* \tilde{t}_L\right)\left(\tilde{t}_L^* \tilde{b}_R\right)\Big) \\
&+ h_\tau^2 \Big(\tilde{\tau}_R^* \tilde{\tau}_L \tilde{\tau}_L^* \tilde{\tau}_R + \tilde{\tau}_R^* \tilde{\nu}_\tau \tilde{\nu}_\tau^* \tilde{\tau}_R\Big) \\
&+ h_b h_\tau \Big(\tilde{b}_R^* \tilde{b}_L \tilde{\tau}_L^* \tilde{\tau}_R + \tilde{b}_L^* \tilde{b}_R \tilde{\tau}_R^* \tilde{\tau}_L + \tilde{t}_L^* \tilde{b}_R \tilde{\tau}_R^* \tilde{\nu}_\tau + \tilde{b}_R^* \tilde{t}_L \tilde{\nu}_\tau^* \tilde{\tau}_R\Big)
\end{aligned} \tag{A.25}
$$

with the couplings of the neutral Higgs bosons and sfermions in the first line, those of the charged Higgs bosons in the second, the couplings of a neutral Higgs boson with a charged one and two sfermions and the four–sfermion couplings. Note that in the detailed calculation of the Feynman rules for the four–sfermion couplings we have to take care about the colour flow, see section A.6.5.

Transforming the interaction fields into the mass eigenstate fields,

$$
\begin{aligned}
H_1^0 &= v_1 + \tfrac{1}{\sqrt{2}} \left[\cos\alpha\, H^0 - \sin\alpha\, h^0 + i\left(-\cos\beta\, G^0 + \sin\beta\, A^0\right)\right] , \\
H_1^2 &= -\cos\beta\, G^- + \sin\beta\, H^- , \\
H_2^1 &= \sin\beta\, G^+ + \cos\beta\, H^+ , \\
H_2^0 &= v_2 + \tfrac{1}{\sqrt{2}} \left[\sin\alpha\, H^0 + \cos\alpha\, h^0 + i\left(\sin\beta\, G^0 + \cos\beta\, A^0\right)\right] ,
\end{aligned} \tag{A.26}
$$

$$
\begin{aligned}
\tilde{f}_L &= R_{i1}^{\tilde{f}} \tilde{f}_i = \cos\theta_{\tilde{f}}\, \tilde{f}_1 - \sin\theta_{\tilde{f}}\, \tilde{f}_2 , \\
\tilde{f}_R &= R_{i2}^{\tilde{f}} \tilde{f}_i = \sin\theta_{\tilde{f}}\, \tilde{f}_1 + \cos\theta_{\tilde{f}}\, \tilde{f}_2 ,
\end{aligned} \tag{A.27}
$$

we simplify our notations as follows:

$$
\begin{aligned}
\left|H_1^0\right|^2 &= \frac{1}{2}\Big[\sin^2\!\alpha\,(h^0)^2 - \sin 2\alpha\, h^0 H^0 + \cos^2\!\alpha\,(H^0)^2 + \sin^2\!\beta\,(A^0)^2 - \sin 2\beta\, A^0 G^0 \\
&\qquad + \cos^2\!\beta\,(G^0)^2\Big] = \frac{1}{2}\, H_k^0\, c_{kl}^{\tilde{b}}\, H_l^0 ,
\end{aligned} \tag{A.28}
$$

$$
\begin{aligned}
\left|H_2^0\right|^2 &= \frac{1}{2}\Big[\cos^2\!\alpha\,(h^0)^2 + \sin 2\alpha\, h^0 H^0 + \sin^2\!\alpha\,(H^0)^2 + \cos^2\!\beta\,(A^0)^2 + \sin 2\beta\, A^0 G^0 \\
&\qquad + \sin^2\!\beta\,(G^0)^2\Big] = \frac{1}{2}\, H_k^0\, c_{kl}^{\tilde{t}}\, H_l^0 ,
\end{aligned} \tag{A.29}
$$

$$\left|H_1^2\right|^2 = \sin^2\beta\, H^+H^- - \frac{1}{2}\sin 2\beta\, H^+G^- - \frac{1}{2}\sin 2\beta\, H^-G^+ + \cos^2\beta G^+G^- = \frac{1}{2}\, H_k^+\, d_{kl}^{\tilde{b}}\, H_l^- \,, \tag{A.30}$$

$$\left|H_2^1\right|^2 = \cos^2\beta\, H^+H^- + \frac{1}{2}\sin 2\beta\, H^+G^- + \frac{1}{2}\sin 2\beta\, H^-G^+ + \sin^2\beta G^+G^- = \frac{1}{2}\, H_k^+\, d_{kl}^{\tilde{t}}\, H_l^- \,, \tag{A.31}$$

with $H_k^0 = \{h^0, H^0, A^0, G^0\}$, $H_k^+ = \{H^+, G^+, H^-, G^-\}$, $H_k^- \equiv (H_k^+)^\dagger = \{H^-, G^-, H^+, G^+\}$ and

$$c_{kl}^{\tilde{b}} = \begin{pmatrix} \sin^2\alpha & -\frac{1}{2}\sin 2\alpha & 0 & 0 \\ -\frac{1}{2}\sin 2\alpha & \cos^2\alpha & 0 & 0 \\ 0 & 0 & \sin^2\beta & -\frac{1}{2}\sin 2\beta \\ 0 & 0 & -\frac{1}{2}\sin 2\beta & \cos^2\beta \end{pmatrix}, \tag{A.32}$$

$$c_{kl}^{\tilde{t}} = \begin{pmatrix} \cos^2\alpha & \frac{1}{2}\sin 2\alpha & 0 & 0 \\ \frac{1}{2}\sin 2\alpha & \sin^2\alpha & 0 & 0 \\ 0 & 0 & \cos^2\beta & \frac{1}{2}\sin 2\beta \\ 0 & 0 & \frac{1}{2}\sin 2\beta & \sin^2\beta \end{pmatrix}, \tag{A.33}$$

$$d_{kl}^{\tilde{b}} = \begin{pmatrix} \sin^2\beta & -\frac{1}{2}\sin 2\beta & 0 & 0 \\ -\frac{1}{2}\sin 2\beta & \cos^2\beta & 0 & 0 \\ 0 & 0 & \sin^2\beta & -\frac{1}{2}\sin 2\beta \\ 0 & 0 & -\frac{1}{2}\sin 2\beta & \cos^2\beta \end{pmatrix}, \tag{A.34}$$

$$d_{kl}^{\tilde{t}} = \begin{pmatrix} \cos^2\beta & \frac{1}{2}\sin 2\beta & 0 & 0 \\ \frac{1}{2}\sin 2\beta & \sin^2\beta & 0 & 0 \\ 0 & 0 & \cos^2\beta & \frac{1}{2}\sin 2\beta \\ 0 & 0 & \frac{1}{2}\sin 2\beta & \sin^2\beta \end{pmatrix}. \tag{A.35}$$

A.6.2 D–term potential

The D–term potential reads

$$V_D = \frac{1}{2}\left(D'D' + \sum_{i=1}^{3} D^iD^i + \sum_{a=1}^{8} D^aD^a\right) \tag{A.36}$$

with $D' = g' A_i^* \frac{Y_i}{2}\,\delta_{ij} A_j$, $D^i = g A_k^* \frac{\sigma_{kl}^i}{2}\, A_l$ and $D^a = g_s A_i^* \frac{\lambda_{ij}^a}{2}\, A_j$ being the terms according to $U(1)$–hypercharge, $SU(2)$–weak isospin and $SU(3)$–strong interaction. The matrices σ_{kl}^i and λ_{ij}^a are the well–known Pauli and Gell–Mann matrices. A_i stands for the scalar (super)fields, $A_i = \{\tilde{Q}, \tilde{L}, \tilde{U}, \tilde{D}, \tilde{E},\, H_1, H_2\}$,

$$\tilde{Q} = \begin{pmatrix} \tilde{t}_L \\ \tilde{b}_L \end{pmatrix}, \qquad \tilde{L} = \begin{pmatrix} \tilde{\nu}_\tau \\ \tilde{\tau}_L \end{pmatrix}, \qquad \tilde{U} = \tilde{t}_R^*\,, \qquad \tilde{D} = \tilde{b}_R^*\,, \qquad \tilde{E} = \tilde{\tau}_R^*\,,$$

$$H_1 = \begin{pmatrix} H_1^1 \\ H_1^2 \end{pmatrix} = \begin{pmatrix} H_1^0 \\ H_1^2 \end{pmatrix}, \qquad H_2 = \begin{pmatrix} H_2^1 \\ H_2^2 \end{pmatrix} = \begin{pmatrix} H_2^1 \\ H_2^0 \end{pmatrix}.$$

The $U(1)$–hypercharge term reads with

$$\frac{Y}{2} \;=\; Q - I_3 \qquad \rightarrow \qquad Y_{H_1} = -1,\; Y_{H_2} = 1 \tag{A.37}$$

$$D' = \frac{g'}{2}\left[\sum_f \left(Y_{\tilde f_L}\tilde f_L^* \tilde f_L - Y_{\tilde f_R}\tilde f_R^* \tilde f_R\right) - \left|H_1^0\right|^2 + \left|H_2^0\right|^2 - \left|H_1^2\right|^2 + \left|H_2^1\right|^2\right]. \tag{A.38}$$

Inserting the Pauli matrices

$$\sigma^1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \qquad \sigma^2 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \qquad \sigma^3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \tag{A.39}$$

we get for the electroweak parts

$$\begin{aligned}
D^1 &= \frac{g}{2}\left(\tilde t_L^* \tilde b_L + \tilde b_L^* \tilde t_L + \tilde\nu_\tau^* \tilde\tau_L + \tilde\tau_L^* \tilde\nu_\tau + H_1^{0*}H_1^2 + H_1^{2*}H_1^0 + H_2^{1*}H_2^0 + H_2^{0*}H_2^1\right), \\
D^2 &= -i\frac{g}{2}\left(\tilde t_L^* \tilde b_L - \tilde b_L^* \tilde t_L + \tilde\nu_\tau^* \tilde\tau_L - \tilde\tau_L^* \tilde\nu_\tau + H_1^{0*}H_1^2 - H_1^{2*}H_1^0 + H_2^{1*}H_2^0 - H_2^{0*}H_2^1\right), \\
D^3 &= \frac{g}{2}\left(\tilde t_L^* \tilde t_L - \tilde b_L^* \tilde b_L + \tilde\nu_\tau^* \tilde\nu_\tau - \tilde\tau_L^* \tilde\tau_L + \left|H_1^0\right|^2 - \left|H_1^2\right|^2 + \left|H_2^1\right|^2 - \left|H_2^0\right|^2\right) \\
&= g\sum_f I_f^{3L}\tilde f_L^* \tilde f_L + \sum_{i=1,2}\left(I^3_{H_1^i}\left|H_1^i\right|^2 + I^3_{H_2^i}\left|H_2^i\right|^2\right).
\end{aligned}$$

(The meaning of $I^3_{H_j^i}$ should be clear.)
When we take the square of the single terms we have to take care about the colours of the sfermion fields. The result is

$$\begin{aligned}
D'D' \;=\; \frac{g'^2}{4}\Bigg\{&\left[\sum_f\left(Y_{\tilde f_L}(\tilde f_L^*\tilde f_L) - Y_{\tilde f_R}(\tilde f_R^*\tilde f_R)\right)\right]^2 + \left[-\left|H_1^0\right|^2 + \left|H_2^0\right|^2 - \left|H_1^2\right|^2 + \left|H_2^1\right|^2\right]^2 \\
&+ 2\sum_f\left(Y_{\tilde f_L}(\tilde f_L^*\tilde f_L) - Y_{\tilde f_R}(\tilde f_R^*\tilde f_R)\right)\left(-\left|H_1^0\right|^2 + \left|H_2^0\right|^2 - \left|H_1^2\right|^2 + \left|H_2^1\right|^2\right)\Bigg\},
\end{aligned} \tag{A.40}$$

$$\begin{aligned}
D^1D^1 + D^2D^2 \;=\; g^2\Bigg[&\left((\tilde t_L^*\tilde b_L) + \tilde\nu_\tau^*\tilde\tau_L\right)\left((\tilde b_L^*\tilde t_L) + \tilde\tau_L^*\tilde\nu_\tau\right) + \left|H_1^0\right|^2\left|H_1^2\right|^2 + \left|H_2^0\right|^2\left|H_2^1\right|^2 \\
&+ H_1^{0*}H_1^2H_2^{0*}H_2^1 + H_2^{1*}H_2^0H_1^{2*}H_1^0 + \left((\tilde t_L^*\tilde b_L) + \tilde\nu_\tau^*\tilde\tau_L\right)\left(H_1^{2*}H_1^0 + H_2^{0*}H_2^1\right) \\
&+ \left((\tilde b_L^*\tilde t_L) + \tilde\tau_L^*\tilde\nu_\tau\right)\left(H_1^{0*}H_1^2 + H_2^{1*}H_2^0\right)\Bigg],
\end{aligned} \tag{A.41}$$

$$
\begin{aligned}
D^3D^3 \;=\; \frac{g^2}{4}\Bigg\{ &\sum_f (\tilde{f}_L^*\tilde{f}_L)\Big[(\tilde{f}_L^*\tilde{f}_L) - (\tilde{f}'^*_L\tilde{f}'_L) + (\hat{\tilde{f}}{}_L^*\hat{\tilde{f}}_L) - (\hat{\tilde{f}}{}'^*_L\hat{\tilde{f}}{}'_L)\Big] \\
&+ \Big[\left|H_1^0\right|^2 - \left|H_1^2\right|^2 + \left|H_2^1\right|^2 - \left|H_2^0\right|^2\Big]^2 \\
&+ 4\sum_f I_f^{3L}(\tilde{f}_L^*\tilde{f}_L)\Big(\left|H_1^0\right|^2 - \left|H_1^2\right|^2 + \left|H_2^1\right|^2 - \left|H_2^0\right|^2\Big)\Bigg\}, \qquad \text{(A.42)}
\end{aligned}
$$

where we have used $\left(I_f^{3L}\right)^2 = \frac{1}{4}$ and $I_f^{3L} I_{f'}^{3L} = -\frac{1}{4}$. Now we are able to calculate the Feynman rules, beginning with the easiest ones, those of two sfermions and two neutral Higgs bosons.

A.6.3 Higgs–Higgs–Sfermion–Sfermion

$H_k^0 H_l^0 \tilde{f}_i \tilde{f}_j$ couplings

The interesting part in the F–term potential for this coupling is (see eq. (A.25))

$$
\begin{aligned}
V_F \;\supset\; & h_t^2\left|H_2^0\right|^2\left(\tilde{t}_R^*\tilde{t}_R + \tilde{t}_L^*\tilde{t}_L\right) + h_b^2\left|H_1^0\right|^2\left(\tilde{b}_R^*\tilde{b}_R + \tilde{b}_L^*\tilde{b}_L\right) + h_\tau^2\left|H_1^0\right|^2\left(\tilde{\tau}_R^*\tilde{\tau}_R + \tilde{\tau}_L^*\tilde{\tau}_L\right) \\
=\; & \frac{h_t^2}{2} H_k^0\, c_{kl}^{\tilde{t}}\, H_l^0\big(\tilde{t}_R^*\tilde{t}_R + \tilde{t}_L^*\tilde{t}_L\big) + \frac{h_b^2}{2} H_k^0\, c_{kl}^{\tilde{b}}\, H_l^0\big(\tilde{b}_R^*\tilde{b}_R + \tilde{b}_L^*\tilde{b}_L\big) \\
& + \frac{h_\tau^2}{2} H_k^0\, c_{kl}^{\tilde{b}}\, H_l^0\big(\tilde{\tau}_R^*\tilde{\tau}_R + \tilde{\tau}_L^*\tilde{\tau}_L\big) \\
=\; & \sum_f \frac{h_f^2}{2} H_k^0\, c_{kl}^{\tilde{f}}\, H_l^0\left(\tilde{f}_R^*\tilde{f}_R + \tilde{f}_L^*\tilde{f}_L\right) \;=\; \sum_f \frac{h_f^2}{2} H_k^0\, c_{kl}^{\tilde{f}}\, H_l^0\, \delta_{ij}\, \tilde{f}_i^*\tilde{f}_j\,, \qquad \text{(A.43)}
\end{aligned}
$$

where in the last step we have transformed the sfermion interaction fields to the mass eigenstate fields (see eq. (A.27)) and made use of the relation $R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} + R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}} = \delta_{ij}$.

From the D–term potential we need the terms $\propto H_k^0 H_l^0 \tilde{f}_i \tilde{f}_j$ of

$$
\begin{aligned}
V_D \;\supset\; \frac{1}{2}\left(D'D' + D^3D^3\right) \;\sim\; & -\frac{g'^2}{4}\sum_f\Big(Y_{\tilde{f}_L}\tilde{f}_L^*\tilde{f}_L - Y_{\tilde{f}_R}\tilde{f}_R^*\tilde{f}_R\Big)\Big(\left|H_1^0\right|^2 - \left|H_2^0\right|^2\Big) \\
& + \frac{g^2}{2}\sum_f I_f^{3L}(\tilde{f}_L^*\tilde{f}_L)\Big(\left|H_1^0\right|^2 - \left|H_2^0\right|^2\Big)\,.
\end{aligned}
$$

Using the abbreviations defined in eqs. (A.28) – (A.35), we get

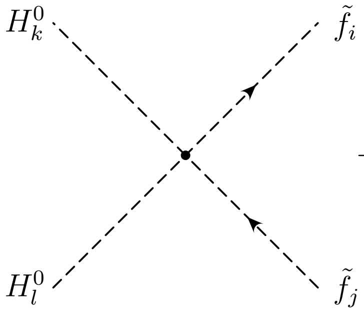
$$
\begin{aligned}
V_D \;\supset\; & \frac{g^2}{4}\sum_f\Big[\Big(I_f^{3L} + (I_f^{3L} - e_f)\, t_W^2\Big)\tilde{f}_L^*\tilde{f}_L + e_f\, t_W^2\, \tilde{f}_R^*\tilde{f}_R\Big] H_k^0\left(c_{kl}^{\tilde{b}} - c_{kl}^{\tilde{t}}\right) H_l^0 \\
=\; & \frac{g^2}{4c_W^2}\sum_f\Big[(I_f^{3L} - e_f s_W^2) R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} + e_f s_W^2 R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}}\Big] H_k^0\left(c_{kl}^{\tilde{b}} - c_{kl}^{\tilde{t}}\right) H_l^0\, \tilde{f}_i^*\tilde{f}_j
\end{aligned}
$$

$$= \sum_f \frac{g^2}{2} H_k^0 \left(c_{kl}^{\tilde{b}} - c_{kl}^{\tilde{t}}\right) H_l^0 \, e_{ij}^{\tilde{f}} \, \tilde{f}_i^* \tilde{f}_j \, . \tag{A.44}$$

with

$$e_{ij}^{\tilde{f}} = \frac{1}{2c_W^2}\left[(I_f^{3L} - e_f s_W^2) R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} + e_f s_W^2 R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}}\right] . \tag{A.45}$$

Therefore the Feynman rule for this coupling becomes



$$-i\left[h_f^2 \, c_{kl}^{\tilde{f}} \, \delta_{ij} + g^2 \left(c_{kl}^{\tilde{b}} - c_{kl}^{\tilde{t}}\right) e_{ij}^{\tilde{f}}\right] .$$

$H_k^+ H_l^+ \tilde{f}_i \tilde{f}_j$ couplings

For the couplings of the charged Higgs bosons and two sfermions we start with the interaction potential

$$\begin{aligned} V_F \supset\ & h_t^2 \left|H_2^1\right|^2 \left(\tilde{t}_R^* \tilde{t}_R + \tilde{b}_L^* \tilde{b}_L\right) + h_b^2 \left|H_1^2\right|^2 \left(\tilde{b}_R^* \tilde{b}_R + \tilde{t}_L^* \tilde{t}_L\right) + h_\tau^2 \left|H_1^2\right|^2 \left(\tilde{\tau}_R^* \tilde{\tau}_R + \tilde{\nu}_\tau^* \tilde{\nu}_\tau\right) \\ =\ & \frac{h_t^2}{2} H_k^+ \, d_{kl}^{\tilde{t}} \, H_l^- \left(\tilde{t}_R^* \tilde{t}_R + \tilde{b}_L^* \tilde{b}_L\right) + \frac{h_b^2}{2} H_k^+ \, d_{kl}^{\tilde{b}} \, H_l^- \left(\tilde{b}_R^* \tilde{b}_R + \tilde{t}_L^* \tilde{t}_L\right) \\ & + \frac{h_\tau^2}{2} H_k^+ \, d_{kl}^{\tilde{b}} \, H_l^- \left(\tilde{\tau}_R^* \tilde{\tau}_R + \tilde{\nu}_\tau^* \tilde{\nu}_\tau\right) = \sum_f \frac{h_f^2}{2} H_k^+ \, d_{kl}^{\tilde{f}} \, H_l^- \left(\tilde{f}_R^* \tilde{f}_R + \tilde{f}_L'^* \tilde{f}_L'\right) . \end{aligned}$$

To get the Feynman rule for this coupling we have to calculate the first (nontrivial) term of the S–matrix;

$$S_{fi}^{(1)} = -i \sum_f \frac{h_f^2}{2} \int d^4x \, \langle f| \, d_{mn}^{\tilde{f}} : H_m^+ H_n^- \left(\tilde{f}_R^* \tilde{f}_R + \tilde{f}_L'^* \tilde{f}_L'\right) : |i\rangle \tag{A.46}$$

with

$$|\, i \,\rangle = a_{H_k^+}^\dagger \, a_{H_l^-}^\dagger |\, 0 \,\rangle \qquad \text{and} \qquad \langle\, f \,| = \langle\, 0 \,|\, a_i \, b_j$$

for two incoming Higgs bosons (H_k^+ and H_l^-) and two outgoing sfermions (sfermion $\tilde{f}_i$ and anti–sfermion $\tilde{f}_j$).

Contracting the sfermions gives (here we use a short notation neglecting all space coordinates and momenta, p_i belongs to the particle with index i)

$$S_{fi}^{(1)} = -i \sum_f \frac{h_f^2}{2} \, d_{mn}^{\tilde{f}} \int d^4x \, \langle 0| \, a_i \, b_j : H_m^+ H_n^- \left(R_{i'2}^{\tilde{f}} R_{j'2}^{\tilde{f}} \, \tilde{f}_{i'}^* \tilde{f}_{j'} + R_{i'1}^{\tilde{f}'} R_{j'1}^{\tilde{f}'} \, \tilde{f}_{i'}'^* \tilde{f}_{j'}'\right) : a_{H_k^+}^\dagger \, a_{H_l^-}^\dagger |\, 0 \,\rangle$$

$$= \; -i\sum_f \frac{h_f^2}{2}\, d^{\tilde f}_{mn}\left(R^{\tilde f}_{i2}R^{\tilde f}_{j2} + R^{\tilde f'}_{i1}R^{\tilde f'}_{j1}\right)\int\! d^4x\, e^{i(p_i-p_j)}\,\langle\, 0|:H^+_m H^-_n:\, a^\dagger_{H^+_k}\, a^\dagger_{H^-_l}|\, 0\,\rangle$$

where we have used

$$\langle\, 0\,|a_i(p_i)\tilde f^*_{i'}(x)|\, 0\rangle \;=\; \underbrace{a_i(p_i)\,\tilde f^*_{i'}}(x) \;=\; \delta_{ii'}\, e^{ip_ix} \tag{A.47}$$

$$\langle\, 0\,|b_j(p_j)\tilde f_{j'}(x)|\, 0\rangle \;=\; \underbrace{b_j(p_j)\,\tilde f_{j'}}(x) \;=\; \delta_{jj'}\, e^{-ip_jx} \tag{A.48}$$

In the contractions of the Higgs fields we have to take care of the Higgs creation and annihilation operators e.g. $a^\dagger_{H^+}$ which can create a Higgs boson H^+ and therefore gives **two** contributions from H^+_1 **and** H^-_3:

$$\underbrace{H^+_k(x)\, a_{H^+_{k'}}}(p_{k'}) \;=\; \delta_{kk'}\, e^{ip_kx} \tag{A.49}$$

$$\underbrace{H^-_k(x)\, a_{H^+_{k'}}}(p_{k'}) \;=\; \begin{pmatrix} 0&0&1&0\\ 0&0&0&1\\ 1&0&0&0\\ 0&1&0&0 \end{pmatrix}_{kk'} e^{ip_kx} \tag{A.50}$$

$$\begin{aligned}
\langle\, 0|:H^+_m H^-_n:\, a^\dagger_{H^+_k}\, a^\dagger_{H^-_l}|\, 0\,\rangle \;&=\; H^+_m H^-_n a^\dagger_{H^+_k}\, a^\dagger_{H^-_l} + H^+_m H^-_n a^\dagger_{H^+_k}\, a^\dagger_{H^-_l}\\
&\propto\; \delta_{mk}\,\delta_{nl} \;+\; (\delta_{n1}\,\delta_{k3} + \delta_{n2}\,\delta_{k4} + \delta_{n3}\,\delta_{k1} + \delta_{n4}\,\delta_{k2})\\
&\quad\;\times\; (\delta_{m1}\,\delta_{l3} + \delta_{m2}\,\delta_{l4} + \delta_{m3}\,\delta_{l1} + \delta_{m4}\,\delta_{l2})\\
&=\; \delta_{mk}\,\delta_{nl} \;+\; \begin{pmatrix} 0&0&1&0\\ 0&0&0&1\\ 1&0&0&0\\ 0&1&0&0 \end{pmatrix}_{nk} \begin{pmatrix} 0&0&1&0\\ 0&0&0&1\\ 1&0&0&0\\ 0&1&0&0 \end{pmatrix}_{ml}
\end{aligned} \tag{A.51}$$

For the Feynman amplitude $\mathcal{M}$ we finally get

$$\begin{aligned}
\mathcal{M} \;&=\; -i\sum_f \frac{h_f^2}{2}\left(R^{\tilde f}_{i2}R^{\tilde f}_{j2}+R^{\tilde f'}_{i1}R^{\tilde f'}_{j1}\right)\left[d^{\tilde f}_{kl} + \begin{pmatrix} 0&0&1&0\\ 0&0&0&1\\ 1&0&0&0\\ 0&1&0&0 \end{pmatrix}_{nk} d^{\tilde f}_{mn} \begin{pmatrix} 0&0&1&0\\ 0&0&0&1\\ 1&0&0&0\\ 0&1&0&0 \end{pmatrix}_{ml}\right]\\
&=\; -i\sum_f h_f^2\left(R^{\tilde f}_{i2}R^{\tilde f}_{j2} + R^{\tilde f'}_{i1}R^{\tilde f'}_{j1}\right) d^{\tilde f}_{kl}\,.
\end{aligned}$$

The D–terms read with eqs. (A.40), (A.42) and (A.28)–(A.35)

$$V_D \;\supset\; -\frac{g'^2}{4}\sum_f\left(Y_{\tilde f_L}\tilde f^*_{L\alpha}\tilde f_{L\alpha} - Y_{\tilde f_R}\tilde f^*_{R\alpha}\tilde f_{R\alpha}\right)\left(\left|H^2_1\right|^2 - \left|H^1_2\right|^2\right)$$

$$
\begin{aligned}
& - \frac{g^2}{2} \sum_f I_f^{3L} \big(\tilde{f}_L^* \tilde{f}_L \big) \left(\left| H_1^2 \right|^2 - \left| H_2^1 \right|^2 \right) \\
= \; & \frac{g^2}{4c_W^2} \sum_f \left[\big(- I_f^{3L} \cos 2\theta_W - e_f s_W^2 \big) \, R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} + e_f s_W^2 R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}} \right] H_k^+ \big(d_{kl}^{\tilde{b}} - d_{kl}^{\tilde{t}} \big) H_l^- \; \tilde{f}_i^* \tilde{f}_j \\
\equiv \; & \frac{g^2}{2} \sum_f H_k^+ \big(d_{kl}^{\tilde{b}} - d_{kl}^{\tilde{t}} \big) H_l^- \, f_{ij}^{\tilde{f}} \; \tilde{f}_i^* \tilde{f}_j \, . \qquad \text{(A.52)}
\end{aligned}
$$

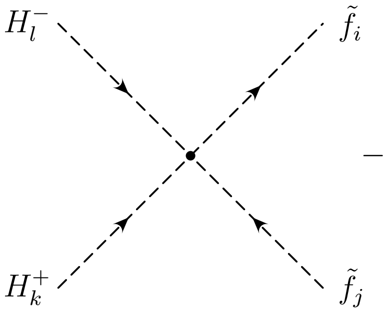
with

$$
f_{ij}^{\tilde{f}} \;=\; \frac{1}{2c_W^2} \left[\big(- I_f^{3L} \cos 2\theta_W - e_f s_W^2 \big) \, R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} + e_f s_W^2 R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}} \right] \qquad \text{(A.53)}
$$

Analogously to the previous calculation we get for the Feynman amplitude

$$
\mathcal{M} \;=\; -i \sum_f g^2 f_{ij}^{\tilde{f}} \left(d_{kl}^{\tilde{b}} - d_{kl}^{\tilde{t}} \right) . \qquad \text{(A.54)}
$$

In the Feynman rule we have to take into consideration that only the terms $\propto R_{ij}^{\tilde{f}}$ are valid for a coupling with sfermions and not the terms $\propto R_{ij}^{\tilde{f}'}$!



$$
-i \left[h_{f'}^2 \, d_{kl}^{\tilde{f}'} \, R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} + h_f^2 \, d_{kl}^{\tilde{f}} \, R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}} + g^2 f_{ij}^{\tilde{f}} \left(d_{kl}^{\tilde{b}} - d_{kl}^{\tilde{t}} \right) \right]
$$

$H_k^0 H_l^+ \tilde{f}_i \tilde{f}_j'$ couplings

The terms of the superpotential V_F which are necessary for calculating these couplings can be picked out of eq. (A.25):

$$
\begin{aligned}
V_F \;\supset\; & - h_t^2 \big(\tilde{t}_L^* \tilde{b}_L H_2^{0*} H_2^1 + \tilde{b}_L^* \tilde{t}_L H_2^{1*} H_2^0 \big) - h_b^2 \big(\tilde{b}_L^* \tilde{t}_L H_1^{0*} H_1^2 + \tilde{t}_L^* \tilde{b}_L H_1^{2*} H_1^0 \big) \\
& - h_\tau^2 \big(\tilde{\tau}_L^* \tilde{\nu}_\tau H_1^{0*} H_1^2 + \tilde{\nu}_\tau^* \tilde{\tau}_L H_1^{2*} H_1^0 \big) \\
& - h_t h_b \left(\tilde{t}_R^* \tilde{b}_R H_1^{2*} H_2^0 + \tilde{b}_R^* \tilde{t}_R H_2^{0*} H_1^2 + \tilde{b}_R^* \tilde{t}_R H_2^{1*} H_1^0 + \tilde{t}_R^* \tilde{b}_R H_1^{0*} H_2^1 \right) \qquad \text{(A.55)}
\end{aligned}
$$

Accordingly to the abbreviations before we introduce a few more coupling matrices for better reading:

$$
H_2^{0*} H_2^1 \;=\; \frac{1}{\sqrt{2}} \left[\cos\alpha \, h^0 + \sin\alpha \, H^0 - i \left(\cos\beta \, A^0 + \sin\beta \, G^0 \right) \right] \left[\cos\beta \, H^+ + \sin\beta \, G^+ \right]
$$

$$= \frac{1}{\sqrt{2}}\left(h^0, H^0, A^0, G^0\right)\begin{pmatrix} \cos\alpha\cos\beta & \cos\alpha\sin\beta & 0 & 0 \\ \sin\alpha\cos\beta & \sin\alpha\sin\beta & 0 & 0 \\ -i\cos^2\beta & -\frac{i}{2}\sin 2\beta & 0 & 0 \\ -\frac{i}{2}\sin 2\beta & -i\sin^2\beta & 0 & 0 \end{pmatrix}\begin{pmatrix} H^+ \\ G^+ \\ H^- \\ G^- \end{pmatrix}$$

$$= \frac{1}{\sqrt{2}} H_k^0\, c_{kl}^{\tilde{t},0+} H_l^+ \tag{A.56}$$

$$H_2^{1*} H_2^0 = \left(H_2^{0*} H_2^1\right)^* = \frac{1}{\sqrt{2}}\left(H_k^0\, c_{kl}^{\tilde{t},0+} H_l^+\right)^* \tag{A.57}$$

$$H_1^{0*} H_1^2 = \frac{1}{\sqrt{2}}\left[-\sin\alpha\, h^0 + \cos\alpha\, H^0 - i\left(\sin\beta\, A^0 - \cos\beta\, G^0\right)\right]\left[\sin\beta\, H^- - \cos\beta\, G^-\right]$$

$$= \frac{1}{\sqrt{2}}\left(h^0, H^0, A^0, G^0\right)\begin{pmatrix} 0 & 0 & -\sin\alpha\sin\beta & \sin\alpha\cos\beta \\ 0 & 0 & \cos\alpha\sin\beta & -\cos\alpha\cos\beta \\ 0 & 0 & -i\sin^2\beta & \frac{i}{2}\sin 2\beta \\ 0 & 0 & \frac{i}{2}\sin 2\beta & -i\cos^2\beta \end{pmatrix}\begin{pmatrix} H^+ \\ G^+ \\ H^- \\ G^- \end{pmatrix}$$

$$= \frac{1}{\sqrt{2}} H_k^0\, c_{kl}^{\tilde{b},0+} H_l^+ \tag{A.58}$$

$$H_1^{2*} H_1^0 = \left(H_1^{0*} H_1^2\right)^* = \frac{1}{\sqrt{2}}\left(H_k^0\, c_{kl}^{\tilde{b},0+} H_l^+\right)^* \tag{A.59}$$

$$H_1^{0*} H_2^1 = \frac{1}{\sqrt{2}}\left[-\sin\alpha\, h^0 + \cos\alpha\, H^0 - i\left(\sin\beta\, A^0 - \cos\beta\, G^0\right)\right]\left[\cos\beta\, H^+ + \sin\beta\, G^+\right]$$

$$= \frac{1}{\sqrt{2}}\left(h^0, H^0, A^0, G^0\right)\begin{pmatrix} -\sin\alpha\cos\beta & -\sin\alpha\sin\beta & 0 & 0 \\ \cos\alpha\cos\beta & \cos\alpha\sin\beta & 0 & 0 \\ -\frac{i}{2}\sin 2\beta & -i\sin^2\beta & 0 & 0 \\ i\cos^2\beta & \frac{i}{2}\sin 2\beta & 0 & 0 \end{pmatrix}\begin{pmatrix} H^+ \\ G^+ \\ H^- \\ G^- \end{pmatrix}$$

$$= \frac{1}{\sqrt{2}} H_k^0\, c_{kl}^{\tilde{t}\tilde{b},0+} H_l^+ \tag{A.60}$$

$$H_2^{1*} H_1^0 = \left(H_1^{0*} H_2^1\right)^* = \frac{1}{\sqrt{2}}\left(H_k^0\, c_{kl}^{\tilde{t}\tilde{b},0+} H_l^+\right)^* \tag{A.61}$$

$$H_2^{0*} H_1^2 = \frac{1}{\sqrt{2}}\left[\cos\alpha\, h^0 + \sin\alpha\, H^0 - i\left(\cos\beta\, A^0 + \sin\beta\, G^0\right)\right]\left[\sin\beta\, H^- - \cos\beta\, G^-\right]$$

$$= \frac{1}{\sqrt{2}}\left(h^0, H^0, A^0, G^0\right)\begin{pmatrix} 0 & 0 & \cos\alpha\sin\beta & -\cos\alpha\cos\beta \\ 0 & 0 & \sin\alpha\sin\beta & -\sin\alpha\cos\beta \\ 0 & 0 & -\frac{i}{2}\sin 2\beta & i\cos^2\beta \\ 0 & 0 & -i\sin^2\beta & \frac{i}{2}\sin 2\beta \end{pmatrix}\begin{pmatrix} H^+ \\ G^+ \\ H^- \\ G^- \end{pmatrix}$$

$$= \frac{1}{\sqrt{2}} H_k^0\, c_{kl}^{\tilde{b}\tilde{t},0+} H_l^+ \tag{A.62}$$

$$H_1^{2*} H_2^0 = \left(H_2^{0*} H_1^2\right)^* = \frac{1}{\sqrt{2}}\left(H_k^0\, c_{kl}^{\tilde{b}\tilde{t},0+} H_l^+\right)^* \tag{A.63}$$

After inserting this into eq. (A.55) we get

$$V_F \;\supset\; -\frac{1}{\sqrt{2}}\sum_f \left(h_f^2\, \tilde{f}_L^* \tilde{f}'_L\, H_k^0\, c_{kl}^{\tilde{f},0+} H_l^+ + h_t h_b\, \tilde{f}_R^* \tilde{f}'_R\, H_k^0\, c_{kl}^{\tilde{f}\tilde{f}',0+} H_l^+ \right) \quad + \text{ h.c.} \qquad \text{(A.64)}$$

Rotating the sfermion fields into the mass eigenstate fields (see eq. (A.27)) gives

$$\begin{aligned} V_F \;\supset\; & -\frac{1}{\sqrt{2}}\sum_f \left(h_f^2\, R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}'}\, c_{kl}^{\tilde{f},0+} + h_f h_{f'}\, R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}'}\, c_{kl}^{\tilde{f}\tilde{f}',0+} \right) H_k^0 H_l^+ \tilde{f}_i^* \tilde{f}'_j \\ & -\frac{1}{\sqrt{2}}\sum_f \left(h_f^2\, R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}'}\, (H_k^0 c_{kl}^{\tilde{f}',0+} H_l^+)^* + h_f h_{f'}\, R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}'}\, (H_k^0 c_{kl}^{\tilde{f}'\tilde{f},0+} H_l^+)^* \right) \tilde{f}_i \tilde{f}_j'^* \end{aligned}$$

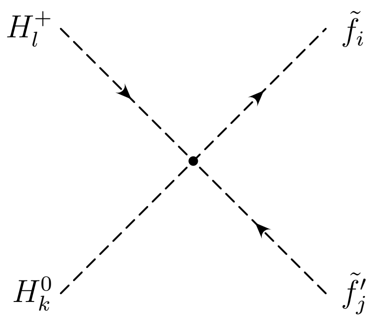
and with

$$\left(c_{kl}^{\tilde{f}',0+} H_l^+ \right)^* = \left(c_{kl}^{\tilde{f}',0+} \right)^* H_l^- = \left(c_{kl}^{\tilde{f}',0+} \right)^* \underbrace{\begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}_{4\text{x}4} \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}_{4\text{x}4}}_{1_{4\text{x}4}} H_l^- = \left(c_{kl'}^{\tilde{f}',0+} \right)^* H_l^+$$

(look at the index l' which is $l' = 3, 4$ for $l = 1, 2$ and vice versa !) we arrive at

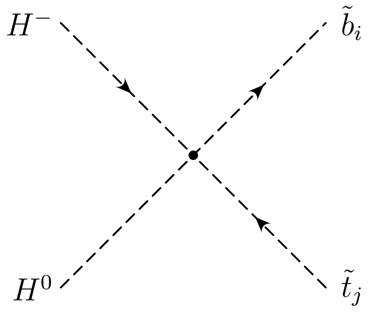
$$\begin{aligned} V_F \;\supset\; -\frac{1}{\sqrt{2}}\sum_f & \left[R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}'} \left(h_f^2 c_{kl}^{\tilde{f},0+} + h_{f'}^2 \left(c_{kl'}^{\tilde{f}',0+} \right)^* \right) + R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}'}\, h_f h_{f'} \left(c_{kl}^{\tilde{f}\tilde{f}',0+} + \left(c_{kl'}^{\tilde{f}'\tilde{f},0+} \right)^* \right) \right] \\ & \times H_k^0\, H_l^+ \tilde{f}_i^* \tilde{f}'_j \end{aligned} \qquad \text{(A.65)}$$

Now the Feynman rules take the form



$$\frac{i}{\sqrt{2}} \left[R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}'} \left(h_f^2\, c_{kl}^{\tilde{f},0+} + h_{f'}^2 \left(c_{kl'}^{\tilde{f}',0+} \right)^* \right) + R_{i2}^{\tilde{f}}\, R_{j2}^{\tilde{f}'}\, h_f h_{f'} \left(c_{kl}^{\tilde{f}\tilde{f}',0+} + \left(c_{kl'}^{\tilde{f}'\tilde{f},0+} \right)^* \right) \right]$$

Be careful which Higgs boson you should take in a particular graph! As an illustrative example we give the result for a coupling which we will need in the main part of this work:



$$\frac{i}{\sqrt{2}} \left[R_{i1}^{\tilde{b}} R_{j1}^{\tilde{t}}\, (h_b^2 \cos\alpha \sin\beta + h_t^2 \sin\alpha \cos\beta) + R_{i2}^{\tilde{b}} R_{j2}^{\tilde{t}}\, h_t h_b\, (\sin\alpha \sin\beta + \cos\alpha \cos\beta) \right]$$

Here we have to take the negative charged Higgs boson H^- (the H^+ would not be allowed due to charge conservation).

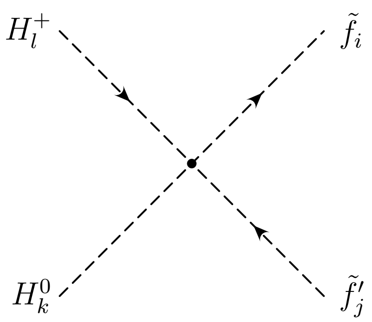
Setting $k=2$ and $l=3$ leaves

$$\frac{i}{\sqrt{2}}\left[R^{\tilde{b}}_{i1}R^{\tilde{t}}_{j1}\left(h_b^2 c^{\tilde{b},0+}_{23}+h_t^2 c^{\tilde{t},0+}_{23'}\right)+R^{\tilde{b}}_{i2}R^{\tilde{t}}_{j2}\,h_t h_b\left(c^{\tilde{b}\tilde{t},0+}_{23}+c^{\tilde{t}\tilde{b},0+}_{23'}\right)\right]$$
$$=\frac{i}{\sqrt{2}}\left[R^{\tilde{b}}_{i1}R^{\tilde{t}}_{j1}\left(h_b^2\cos\alpha\sin\beta+h_t^2\sin\alpha\cos\beta\right)+R^{\tilde{b}}_{i2}R^{\tilde{t}}_{j2}\,h_t h_b(\sin\alpha\sin\beta+\cos\alpha\cos\beta)\right].$$

The D–term potential terms coming from the off–diagonal Pauli matrices σ^1_{kl} and σ^2_{kl} are given by

$$\begin{aligned}
V_D &\supset \frac{1}{2}\left(D^1D^1+D^2D^2\right)\\
&\supset \left(\left(\tilde{t}^*_L\tilde{b}_L\right)+\tilde{\nu}^*_\tau\tilde{\tau}_L\right)\left(H_1^{2*}H_1^0+H_2^{0*}H_2^1\right)+\left(\left(\tilde{b}^*_L\tilde{t}_L\right)+\tilde{\tau}^*_L\tilde{\nu}_\tau\right)\left(H_1^{0*}H_1^2+H_2^{1*}H_2^0\right)\\
&= \left(\left(\tilde{t}^*_L\tilde{b}_L\right)+\tilde{\nu}^*_\tau\tilde{\tau}_L\right)\frac{g^2}{2\sqrt{2}}\left[\left(H_k^0\,c^{\tilde{b},0+}_{kl}H_l^+\right)^*+H_k^0\,c^{\tilde{t},0+}_{kl}H_l^+\right]\\
&\quad+\left(\left(\tilde{b}^*_L\tilde{t}_L\right)+\tilde{\tau}^*_L\tilde{\nu}_\tau\right)\frac{g^2}{2\sqrt{2}}\left[H_k^0\,c^{\tilde{b},0+}_{kl}H_l^+ +\left(H_k^0\,c^{\tilde{t},0+}_{kl}H_l^+\right)^*\right]\\
&= \frac{g^2}{2\sqrt{2}}\left[\left(\left(\tilde{t}^*_L\tilde{b}_L\right)+\tilde{\nu}^*_\tau\tilde{\tau}_L\right)\left(H_k^0\left(c^{\tilde{b},0+}_{kl'}\right)^*H_l^+ +H_k^0\,c^{\tilde{t},0+}_{kl}H_l^+\right)\right.\\
&\qquad\left.+\left(\left(\tilde{b}^*_L\tilde{t}_L\right)+\tilde{\tau}^*_L\tilde{\nu}_\tau\right)\left(H_k^0\,c^{\tilde{b},0+}_{kl}H_l^+ +H_k^0\left(c^{\tilde{t},0+}_{kl'}\right)^*H_l^+\right)\right]\\
&= \frac{g^2}{2\sqrt{2}}\sum_f R^{\tilde{f}}_{i1}R^{\tilde{f}'}_{j1}\left(c^{\tilde{f},0+}_{kl}+\left(c^{\tilde{f}',0+}_{kl'}\right)^*\right)H_k^0H_l^+\tilde{f}_i^*\tilde{f}'_j\,,
\end{aligned} \tag{A.66}$$

which results in the Feynman rules



$$-\frac{i}{\sqrt{2}}\frac{g^2}{2}R^{\tilde{f}}_{i1}R^{\tilde{f}'}_{j1}\left(c^{\tilde{f},0+}_{kl}+\left(c^{\tilde{f}',0+}_{kl'}\right)^*\right)$$

A.6.4 Higgs–Higgs–Higgs–Higgs

The 4–Higgs couplings are only coming from the D–terms, and therefore we have (see eqs. (A.40)–(A.42))

$$V_D \supset \frac{g'^2}{8}\Big[\left|H_1^0\right|^2 - \left|H_2^0\right|^2 + \left|H_1^2\right|^2 - \left|H_2^1\right|^2\Big]^2 + \frac{g^2}{8}\Big[\left|H_1^0\right|^2 - \left|H_2^0\right|^2 - \left|H_1^2\right|^2 + \left|H_2^1\right|^2\Big]^2$$
$$+ \frac{g^2}{2}\Big[\left|H_1^0\right|^2\left|H_1^2\right|^2 + \left|H_2^0\right|^2\left|H_2^1\right|^2 + H_1^{0*}H_1^2 H_2^{0*}H_2^1 + H_2^{1*}H_2^0 H_1^{2*}H_1^0\Big], \qquad \text{(A.67)}$$

or in a more compact way

$$V_D \supset \frac{1}{8}(g^2 + g'^2)\left(H_1^{i*}H_1^i - H_2^{i*}H_2^i\right)^2 + \frac{g^2}{2}\left|H_1^{i*}H_2^i\right|^2 . \qquad \text{(A.68)}$$

$H_k^0 H_l^0 H_m^0 H_n^0$ couplings

The 4–neutral Higgs couplings are obtained from the following term of the D–term potential,

$$V_D \supset \frac{1}{8}(g^2 + g'^2)\left(\left|H_1^0\right|^2 - \left|H_2^0\right|^2\right)^2 , \qquad \text{(A.69)}$$

which can be expressed in components,

$$V_D \supset \frac{g^2}{32c_W^2}\Big[\Big((h^0)^2 - (H^0)^2\Big)\cos 2\alpha + \Big((A^0)^2 - (G^0)^2\Big)\cos 2\beta + 2\Big(h^0 H^0 \sin 2\alpha$$
$$+ A^0 G^0 \sin 2\beta\Big)\Big]^2 , \qquad \text{(A.70)}$$

as well as in index notation (see eqs.(A.28), (A.29), (A.32) and (A.33)),

$$V_D \supset \frac{g^2}{32c_W^2}\sum_{k,l,m,n} H_k^0 H_l^0 H_m^0 H_n^0 \left(c_{kl}^{\tilde{b}}c_{mn}^{\tilde{b}} + c_{kl}^{\tilde{t}}c_{mn}^{\tilde{t}} - c_{kl}^{\tilde{b}}c_{mn}^{\tilde{t}} - c_{kl}^{\tilde{t}}c_{mn}^{\tilde{b}}\right) . \qquad \text{(A.71)}$$

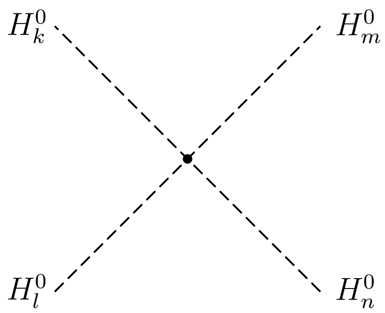
To get a special coupling out of the potential we either can sum over all indices an then pick out the single terms which belong to the required coupling or we symmetrize the coupling matrices $c^{\tilde{b}}$ and $c^{\tilde{t}}$ in eq. (A.71), fix the indices belonging to the Higgs fields H_k^0 and take a combinatorial factor for multiple counting into consideration. Here we choose the second possibility:

$$V_D \supset \frac{g^2}{32c_W^2} H_k^0 H_l^0 H_m^0 H_n^0 \left(c_{(kl}^{\tilde{b}}c_{mn)}^{\tilde{b}} + c_{(kl}^{\tilde{t}}c_{mn)}^{\tilde{t}} - c_{(kl}^{\tilde{b}}c_{mn)}^{\tilde{t}} - c_{(kl}^{\tilde{t}}c_{mn)}^{\tilde{b}}\right) \times \text{CF} \qquad \text{(A.72)}$$

(no sum over indices, this is respected in the combinatorial factor!) Here the brackets around the indices denote symmetrization and CF stands for a combinatorial factor, which is given for a general coupling $(h^0)^a(H^0)^b(A^0)^c(G^0)^d$ with $a+b+c+d=4$ by

$$\text{CF} = \binom{4}{a}\cdot\binom{4-a}{b}\cdot\binom{4-a-b}{c}\cdot\binom{4-a-b-c}{d} = \frac{4!}{a!\,b!\,c!\,d!} \qquad \text{(A.73)}$$

In the Feynman rule we have to take the symmetry factor (SF) of the diagram into account ($n!$ for n equal neutral particles):



$$-\,i\,\frac{g^2}{32c_W^2}\left(c^{\tilde{b}}_{(kl}c^{\tilde{b}}_{mn)}+c^{\tilde{t}}_{(kl}c^{\tilde{t}}_{mn)}-c^{\tilde{b}}_{(kl}c^{\tilde{t}}_{mn)}-c^{\tilde{t}}_{(kl}c^{\tilde{b}}_{mn)}\right)$$
$$\times\,\mathrm{CF}\times\mathrm{SF}$$

As an example we take the coupling $h^0H^0(A^0)^2$. For the indices we choose $k=1, l=2$ and $m=n=3$. The combinatorial factor then is given by

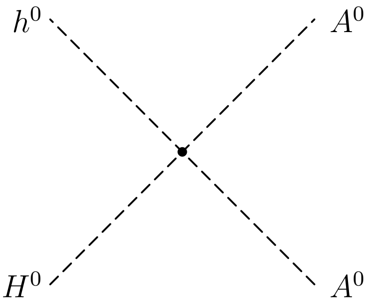
$$\mathrm{CF} \;=\; \begin{pmatrix}4\\1\end{pmatrix}\cdot\begin{pmatrix}3\\1\end{pmatrix}\cdot\begin{pmatrix}2\\2\end{pmatrix}\cdot\begin{pmatrix}0\\0\end{pmatrix} \;=\; 4\cdot3\cdot1\cdot1 \;=\; 12\,,$$

so the interaction Lagrangian reads with (note that the matrices $c^{\tilde{b}}$ and $c^{\tilde{t}}$ are symmetric)

$$\left(c^{\tilde{b}}_{(12}c^{\tilde{b}}_{33)}+c^{\tilde{t}}_{(12}c^{\tilde{t}}_{33)}-c^{\tilde{b}}_{(12}c^{\tilde{t}}_{33)}-c^{\tilde{t}}_{(12}c^{\tilde{b}}_{33)}\right) \;=\; \frac{1}{3}\Big[\left(c^{\tilde{b}}_{12}c^{\tilde{b}}_{33}+c^{\tilde{b}}_{13}c^{\tilde{b}}_{23}+c^{\tilde{b}}_{13}c^{\tilde{b}}_{23}\right)$$
$$-\;\left(c^{\tilde{b}}_{12}c^{\tilde{t}}_{33}+c^{\tilde{b}}_{13}c^{\tilde{t}}_{23}+c^{\tilde{b}}_{13}c^{\tilde{t}}_{23}\right)\;+\;\tilde{b}\leftrightarrow\tilde{t}\,\Big] \;=\; \frac{1}{3}\sin 2\alpha\cos 2\beta \qquad ,$$

and therefore

$$\mathcal{L} \;=\; -\frac{g^2}{8c_W^2}\sin 2\alpha\cos 2\beta\, h^0H^0(A^0)^2\,. \tag{A.74}$$



$$-\,i\,\frac{g^2}{4c_W^2}\sin 2\alpha\cos 2\beta$$

$H_k^+H_l^-H_m^+H_n^-$ couplings

Like in the case of 4–neutral Higgs boson couplings we get the 4–charged Higgs couplings from the D–term potential

$$\begin{aligned} V_D &\supset \frac{1}{8}(g^2+g'^2)\left(\left|H_1^2\right|^2-\left|H_2^1\right|^2\right)^2 \\ &= \frac{g^2}{8c_W^2}\left[H^-\Big(\sin 2\beta G^+ + \cos 2\beta H^+\Big) + G^-\Big(-\cos 2\beta G^+ + \sin 2\beta H^+\Big)\right]^2 . \end{aligned} \tag{A.75}$$

In index notation this reads with the abbreviations defined in eqs. (A.28) –(A.35)

$$V_D \supset \frac{g^2}{32c_W^2}\sum_{k,l,m,n} H_k^+H_l^-H_m^+H_n^-\left(d_{kl}^{\tilde b}d_{mn}^{\tilde b}+d_{kl}^{\tilde t}d_{mn}^{\tilde t}-d_{kl}^{\tilde b}d_{mn}^{\tilde t}-d_{kl}^{\tilde t}d_{mn}^{\tilde b}\right) . \tag{A.76}$$

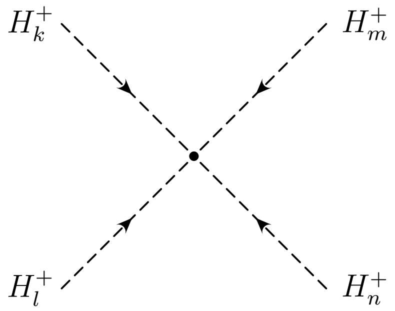
To be able to pick out the various couplings out of eq. (A.76) in the same way as we did in the case of the neutral Higgs bosons we first must express all fields in one single base, H_k^+ or H_k^-:

$$V_D \supset \frac{g^2}{32c_W^2}\sum_{k,l,m,n} H_k^+H_l^+H_m^+H_n^+\left(d_{kl}^{\tilde b}d_{mn}^{\tilde b}+d_{kl}^{\tilde t}d_{mn}^{\tilde t}-d_{kl}^{\tilde b}d_{mn}^{\tilde t}-d_{kl}^{\tilde t}d_{mn}^{\tilde b}\right)_{\text{rot}} , \tag{A.77}$$

with

$$\left(d_{kl}^{\tilde f}\,d_{mn}^{\tilde f}\right)_{\text{rot}} = \left(d_{kl}^{\tilde f}\right)_{\text{rot}}\left(d_{mn}^{\tilde f}\right)_{\text{rot}} = \sum_{l',n'=1}^{4} d_{kl'}^{\tilde f}\begin{pmatrix}0 & 1_{2\text{x}2}\\ 1_{2\text{x}2} & 0\end{pmatrix}_{l'l} d_{mn'}^{\tilde f}\begin{pmatrix}0 & 1_{2\text{x}2}\\ 1_{2\text{x}2} & 0\end{pmatrix}_{n'n} .$$

Now we can symmetrize the rotated coupling matrices as before if we want to fix the indices in the Higgs fields. The combinatorial factor CF stays the same but we don't have to take a symmetry factor in the Feynman rule because of charged particles:



$$-i\,\frac{g^2}{32c_W^2}\left(d_{kl}^{\tilde b}d_{mn}^{\tilde b}+d_{kl}^{\tilde t}d_{mn}^{\tilde t}-d_{kl}^{\tilde b}d_{mn}^{\tilde t}-d_{kl}^{\tilde t}d_{mn}^{\tilde b}\right)_{\text{rot, symm}} \times \text{CF}$$

A.6.5 Sfermion–Sfermion–Sfermion–Sfermion

As we already mentioned before, we have to take care about the colour indices in this term of the superpotential,

$$W \supset h_t\,\tilde t_R^*\tilde t_L H_2^0 + h_b\,\tilde b_R^*\tilde b_L H_1^0 + h_\tau\,\tilde\tau_R^*\tilde\tau_L H_1^0 - h_t\,\tilde t_R^*\tilde b_L H_2^1 - h_b\,\tilde b_R^*\tilde t_L H_1^2 - h_\tau\,\tilde\tau_R^*\tilde\nu_\tau H_1^2 , \tag{A.78}$$

which leads to the F–term potential

$$\begin{aligned}
V_F \;\supset\;& h_t^2\big(\tilde{t}_R^*\tilde{t}_L\big)\big(\tilde{t}_L^*\tilde{t}_R\big) + h_b^2\big(\tilde{b}_R^*\tilde{b}_L\big)\big(\tilde{b}_L^*\tilde{b}_R\big) + h_\tau^2\big(\tilde{\tau}_R^*\tilde{\tau}_L\big)\big(\tilde{\tau}_L^*\tilde{\tau}_R\big) \\
&+ h_t^2\big(\tilde{t}_R^*\tilde{b}_L\big)\big(\tilde{b}_L^*\tilde{t}_R\big) + h_b^2\big(\tilde{b}_R^*\tilde{t}_L\big)\big(\tilde{t}_L^*\tilde{b}_R\big) + h_\tau^2\big(\tilde{\tau}_R^*\tilde{\nu}_\tau\big)\big(\tilde{\nu}_\tau^*\tilde{\tau}_R\big) \\
&+ h_b h_\tau\Big(\tilde{b}_R^*\tilde{b}_L\tilde{\tau}_L^*\tilde{\tau}_R + \tilde{b}_L^*\tilde{b}_R\tilde{\tau}_R^*\tilde{\tau}_L + \tilde{t}_L^*\tilde{b}_R\tilde{\tau}_R^*\tilde{\nu}_\tau + \tilde{b}_R^*\tilde{t}_L\tilde{\nu}_\tau^*\tilde{\tau}_R\Big) \\
=\;& \sum_f \Big[h_f^2\Big(\big(\tilde{f}_R^*\tilde{f}_L\big)\big(\tilde{f}_L^*\tilde{f}_R\big) + \big(\tilde{f}_R^*\tilde{f}_L'\big)\big(\tilde{f}_L'^*\tilde{f}_R\big)\Big) + h_f h_{\hat{f}}\Big(\tilde{f}_R^*\tilde{f}_L\hat{\tilde{f}}_L^*\hat{\tilde{f}}_R + \tilde{f}_R^*\tilde{f}_L'\hat{\tilde{f}}_L'^*\hat{\tilde{f}}_R\Big)\Big]\,.
\end{aligned} \tag{A.79}$$

$\tilde{f}_i\tilde{f}_j\tilde{f}_k\tilde{f}_l$ couplings

With the initial and final states for the first term,

$$\begin{aligned}
|i\rangle &= a_{l\delta}^\dagger(p_1)b_{k\gamma}^\dagger(p_2)|0\rangle\,, \\
\langle f| &= \langle 0|a_{i\alpha}(k_1)b_{j\beta}(k_2)\,,
\end{aligned}$$

we get for the S–matrix element

$$\begin{aligned}
S_{fi}^{(1)} &= i\int d^4x\,\langle f|:\mathcal{L}_{int}(x):|i\rangle = -ih_f^2\,R_{m1}^{\tilde{f}}R_{n1}^{\tilde{f}}R_{p2}^{\tilde{f}}R_{q2}^{\tilde{f}} \\
&\times \int d^4x\,\langle 0|a_{i\alpha}(k_1)b_{j\beta}(k_2):\Big(\tilde{f}_{m\gamma'}^*\tilde{f}_{n\beta'}\tilde{f}_{p\alpha'}^*\tilde{f}_{q\delta'}\Big)(x):a_{l\delta}^\dagger(p_1)b_{k\gamma}^\dagger(p_2)|0\rangle\,\delta_{\alpha'\beta'}\delta_{\gamma'\delta'}
\end{aligned} \tag{A.80}$$

In order to evaluate the vacuum expectation value we make use of Wick's theorem in taking all possible contractions:

$$\begin{aligned}
S_{fi}^{(1)} =\;& -ih_f^2\,R_{m1}^{\tilde{f}}R_{n1}^{\tilde{f}}R_{p2}^{\tilde{f}}R_{q2}^{\tilde{f}}\,(2\pi)^4\delta^4(p_1+p_2-k_1-k_2)\,\delta_{\alpha'\beta'}\delta_{\gamma'\delta'} \\
&\times\big(\delta_{im}\delta_{\alpha\gamma'}\,\delta_{jn}\delta_{\beta\beta'}\,\delta_{ql}\delta_{\delta'\delta}\,\delta_{pk}\delta_{\alpha'\gamma} + \delta_{im}\delta_{\alpha\gamma'}\,\delta_{jq}\delta_{\beta\delta'}\,\delta_{nl}\delta_{\beta'\delta}\,\delta_{pk}\delta_{\alpha'\gamma} \\
&+ \delta_{ip}\delta_{\alpha\alpha'}\,\delta_{jn}\delta_{\beta\beta'}\,\delta_{ql}\delta_{\delta'\delta}\,\delta_{mk}\delta_{\gamma'\gamma} + \delta_{ip}\delta_{\alpha\alpha'}\,\delta_{jq}\delta_{\beta\delta'}\,\delta_{mk}\delta_{\gamma'\gamma}\,\delta_{nl}\delta_{\beta'\delta}\big)
\end{aligned} \tag{A.81}$$

In the last equation we have used the notation

$$\delta_{\alpha\beta\gamma\delta,\alpha'\beta'\gamma'\delta'} \;\equiv\; \delta_{\alpha\alpha'}\,\delta_{\beta\beta'}\,\delta_{\gamma\gamma'}\,\delta_{\delta\delta'}\,, \qquad \rightarrow \qquad \delta_{\alpha\beta\gamma\delta,\alpha'\beta'\alpha'\beta'} \;=\; \delta_{\alpha\gamma}\,\delta_{\beta\delta} \tag{A.82}$$

and

$$R_{ijkl}^{\tilde{f}} \;\equiv\; R_{i1}^{\tilde{f}}R_{j1}^{\tilde{f}}R_{k2}^{\tilde{f}}R_{l2}^{\tilde{f}}\,. \tag{A.83}$$

The corresponding D–term potential is given by (see eqs. (A.40) and (A.42))

$$V_D \;\supset\; \frac{g'^2}{8}\sum_f\Big(Y_{\tilde{f}_L}\big(\tilde{f}_L^*\tilde{f}_L\big) - Y_{\tilde{f}_R}\big(\tilde{f}_R^*\tilde{f}_R\big)\Big)\Big(Y_{\tilde{f}_L}\big(\tilde{f}_L^*\tilde{f}_L\big) - Y_{\tilde{f}_R}\big(\tilde{f}_R^*\tilde{f}_R\big)\Big)$$

$$+\frac{g^2}{8}\sum_f \big(\tilde{f}_L^*\tilde{f}_L\big)\big(\tilde{f}_L^*\tilde{f}_L\big)\,. \tag{A.84}$$

Rotating the sfermion fields into their mass eigenstates, $\tilde{f}_L = R^{\tilde{f}}_{i1}\tilde{f}_i$, $\tilde{f}_R = R^{\tilde{f}}_{i2}\tilde{f}_i$, the D–term potential for four sfermions reads

$$\begin{aligned} V_D &\supset \frac{g^2}{8}\sum_{\tilde{f}}\Big[R^{\tilde{f}_L}_{ijkl} + t_W^2 Y^2_{\tilde{f}_L} R^{\tilde{f}_L}_{ijkl} + t_W^2 Y^2_{\tilde{f}_R} R^{\tilde{f}_R}_{ijkl} - t_W^2 Y_{\tilde{f}_L} Y_{\tilde{f}_R}\Big(R^{\tilde{f}}_{ijkl} + R^{\tilde{f}}_{klij}\Big)\Big]\big(\tilde{f}_i^*\tilde{f}_j\big)\big(\tilde{f}_k^*\tilde{f}_l\big) \\ &\equiv \sum_{\tilde{f}} C^{\tilde{f}\tilde{f}}_{ijkl}\big(\tilde{f}_i^*\tilde{f}_j\big)\big(\tilde{f}_k^*\tilde{f}_l\big)\,, \end{aligned} \tag{A.85}$$

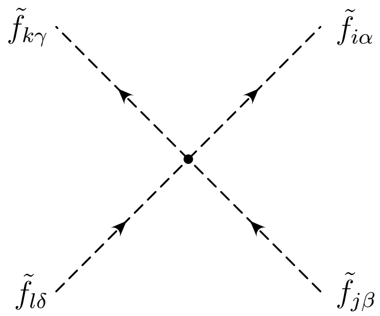
where we have used the relation $g' = g\tan\theta_W$ as well as the abbreviations

$$R^{\tilde{f}_L}_{ijkl} = R^{\tilde{f}}_{i1}R^{\tilde{f}}_{j1}R^{\tilde{f}}_{k1}R^{\tilde{f}}_{l1}\,, \qquad R^{\tilde{f}_R}_{ijkl} = R^{\tilde{f}}_{i2}R^{\tilde{f}}_{j2}R^{\tilde{f}}_{k2}R^{\tilde{f}}_{l2}\,. \tag{A.86}$$

Like in the calculation of the Yukawa coupling terms in the previous section we get for the Feynman amplitude

$$\mathcal{M} \;=\; -i\Big[\Big(C^{\tilde{f}\tilde{f}}_{ijkl} + C^{\tilde{f}\tilde{f}}_{klij}\Big)\delta_{\alpha\beta}\delta_{\gamma\delta} + \Big(C^{\tilde{f}\tilde{f}}_{ilkj} + C^{\tilde{f}\tilde{f}}_{kjil}\Big)\delta_{\alpha\delta}\delta_{\beta\gamma}\Big]\,. \tag{A.87}$$

Both, $F-$ and D–terms, result in the Feynman rule



$$\begin{aligned} &-i\,h_f^2\Big[\Big(R^{\tilde{f}}_{ijkl} + R^{\tilde{f}}_{klij}\Big)\delta_{\alpha\delta}\delta_{\beta\gamma} \;+\; \Big(R^{\tilde{f}}_{ilkj} + R^{\tilde{f}}_{kjil}\Big)\delta_{\alpha\beta}\delta_{\gamma\delta}\Big] \\ &-i\Big[\Big(C^{\tilde{f}\tilde{f}}_{ijkl} + C^{\tilde{f}\tilde{f}}_{klij}\Big)\delta_{\alpha\beta}\delta_{\gamma\delta} + \Big(C^{\tilde{f}\tilde{f}}_{ilkj} + C^{\tilde{f}\tilde{f}}_{kjil}\Big)\delta_{\alpha\delta}\delta_{\beta\gamma}\Big] \end{aligned}$$

$\tilde{f}_i\tilde{f}_j\tilde{f}'_k\tilde{f}'_l$ couplings

For the second term of eq. (A.79) we take the initial and final states

$$\begin{aligned} |i\rangle &= a'^{\dagger}_{l\delta}(p_1)b'^{\dagger}_{k\gamma}(p_2)|0\rangle\,, \\ \langle f| &= \langle 0|a_{i\alpha}(k_1)b_{j\beta}(k_2)\,, \end{aligned}$$

which leads with

$$R^{\tilde{f}'\tilde{f}_D}_{ijkl} \equiv R^{\tilde{f}'}_{i1}R^{\tilde{f}'}_{j1}R^{\tilde{f}}_{k2}R^{\tilde{f}}_{l2}\,, \qquad R^{\tilde{f}\tilde{f}'_D}_{ijkl} \equiv R^{\tilde{f}}_{i1}R^{\tilde{f}}_{j1}R^{\tilde{f}'}_{k2}R^{\tilde{f}'}_{l2} \tag{A.88}$$

to the Feynman amplitude

$$\mathcal{M} \;=\; -\,i\Big(h_f^2\,R^{\tilde{f}'\tilde{f}_D}_{klij} + h^2_{f'}R^{\tilde{f}\tilde{f}'_D}_{ijkl}\Big)\delta_{\alpha\delta}\delta_{\beta\gamma}\,. \tag{A.89}$$

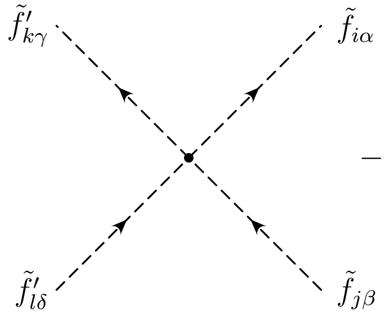
For the D–term potential we get with the eqs. (A.40) – (A.42)

$$
\begin{aligned}
V_D \;\supset\;& \frac{g^2}{4}\sum_f \Big(\big(\tilde f_L^* \tilde f_L'\big)\big(\tilde f_L'^{\,*} \tilde f_L\big) - \frac{1}{2}\big(\tilde f_L^* \tilde f_L\big)\big(\tilde f_L'^{\,*} \tilde f_L'\big)\Big) \\
&+\frac{g'^2}{8}\sum_f \Big(Y_{\tilde f_L}\big(\tilde f_L^* \tilde f_L\big) - Y_{\tilde f_R}\big(\tilde f_R^* \tilde f_R\big)\Big)\Big(Y_{\tilde f_L'}\big(\tilde f_L'^{\,*} \tilde f_L'\big) - Y_{\tilde f_R'}\big(\tilde f_R'^{\,*} \tilde f_R'\big)\Big) \qquad \text{(A.90)}
\end{aligned}
$$

and with a little bit of cosmetics we have

$$
\begin{aligned}
V_D \;\supset\;& \frac{g^2}{8}\sum_f \Bigg\{ \bigg[\Big(t_W^2\big(Y_{\tilde f_L} Y_{\tilde f_L'}\big) - 1\Big) R_{ijkl}^{\tilde f \tilde f_L'} + t_W^2\big(Y_{\tilde f_R} Y_{\tilde f_R'}\big)\, R_{ijkl}^{\tilde f \tilde f_R'} - t_W^2\big(Y_{\tilde f_L} Y_{\tilde f_R'}\big)\, R_{ijkl}^{\tilde f \tilde f_D'} \\
&- t_W^2\big(Y_{\tilde f_R} Y_{\tilde f_L'}\big)\, R_{klij}^{\tilde f' \tilde f_D}\bigg]\delta_{\alpha\beta}\delta_{\gamma\delta} + 2R_{ijkl}^{\tilde f \tilde f_L'}\,\delta_{\alpha\delta}\delta_{\beta\gamma}\Bigg\}\tilde f_{i\alpha}^* \tilde f_{j\beta} \tilde f_{k\gamma}'^{\,*} \tilde f_{l\delta}' \\
\equiv\;& \sum_f C_{ijkl,\alpha\beta\gamma\delta}^{\tilde f \tilde f'}\; \tilde f_{i\alpha}^* \tilde f_{j\beta} \tilde f_{k\gamma}'^{\,*} \tilde f_{l\delta}'\,. \qquad \text{(A.91)}
\end{aligned}
$$

The Yukawa and electroweak contributions to the Feynman rule are then given by



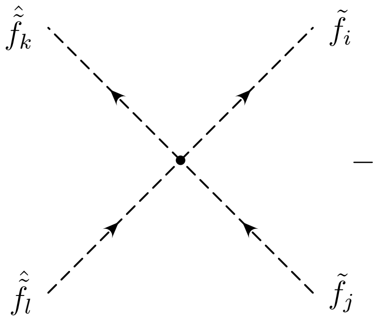
$$
-i\Big(h_f^2\, R_{klij}^{\tilde f' \tilde f_D} + h_{f'}^2 R_{ijkl}^{\tilde f \tilde f_D'}\Big)\delta_{\alpha\delta}\delta_{\beta\gamma} - i\Big(C_{ijkl,\alpha\beta\gamma\delta}^{\tilde f \tilde f'} + C_{klij,\gamma\delta\alpha\beta}^{\tilde f' \tilde f}\Big)
$$

$\tilde f_i \tilde f_j \hat{\tilde f}_k \hat{\tilde f}_l$ couplings

The Feynman rule for couplings with two sfermions and two 'family partner' sfermions can be obtained from eqs. (A.40), (A.42) and (A.79):

$$
\begin{aligned}
V_{F+D} \;\supset\;& \sum_f \bigg[h_f h_{\hat f}\big(\tilde f_R^* \tilde f_L\big)\big(\hat{\tilde f}_L^* \hat{\tilde f}_R\big) + \frac{g^2}{8}\big(\tilde f_L^* \tilde f_L\big)\big(\hat{\tilde f}_L^* \hat{\tilde f}_L\big) \\
&+\frac{g'^2}{8}\Big(Y_{\tilde f_L}\big(\tilde f_L^* \tilde f_L\big) - Y_{\tilde f_R}\big(\tilde f_R^* \tilde f_R\big)\Big)\Big(Y_{\hat{\tilde f}_L}\big(\hat{\tilde f}_L^* \hat{\tilde f}_L\big) - Y_{\hat{\tilde f}_R}\big(\hat{\tilde f}_R^* \hat{\tilde f}_R\big)\Big)\bigg] \\
=\;& \sum_f \Bigg\{ h_f h_{\hat f}\, R_{jikl}^{\tilde f \hat{\tilde f}_F} + \frac{g^2}{8} R_{ijkl}^{\tilde f \hat{\tilde f}_L} + \frac{g'^2}{8}\bigg[Y_{\tilde f_L} Y_{\hat{\tilde f}_L} R_{ijkl}^{\tilde f \hat{\tilde f}_L} + Y_{\tilde f_R} Y_{\hat{\tilde f}_R} R_{ijkl}^{\tilde f \hat{\tilde f}_R} - Y_{\tilde f_L} Y_{\hat{\tilde f}_R} R_{ijkl}^{\tilde f \hat{\tilde f}_D} \\
&- Y_{\tilde f_R} Y_{\hat{\tilde f}_L} R_{klij}^{\hat{\tilde f} \tilde f_D}\bigg]\Bigg\}\big(\tilde f_i^* \tilde f_j\big)\big(\hat{\tilde f}_k^* \hat{\tilde f}_l\big)
\end{aligned}
$$

$$\equiv \sum_f \left[h_f h_{\hat{f}} \, R^{\tilde{f}\hat{\tilde{f}}_F}_{jikl} + C^{\tilde{f}\hat{\tilde{f}}}_{ijkl} \right] \big(\tilde{f}_i^* \tilde{f}_j\big)\big(\hat{\tilde{f}}_k^* \hat{\tilde{f}}_l\big) \tag{A.92}$$

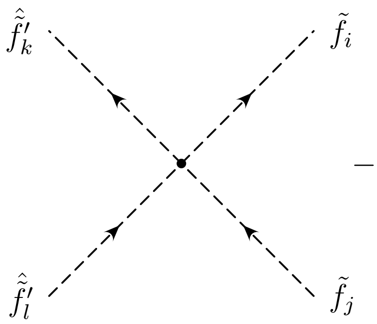


$$-i h_f h_{\hat{f}} \Big(R^{\tilde{f}\hat{\tilde{f}}_F}_{jikl} + R^{\hat{\tilde{f}}\tilde{f}_F}_{lkij} \Big) - i \Big(C^{\tilde{f}\hat{\tilde{f}}}_{ijkl} + C^{\hat{\tilde{f}}\tilde{f}}_{klij} \Big)$$

$\tilde{f}_i \tilde{f}_j \hat{\tilde{f}}'_k \hat{\tilde{f}}'_l$ couplings

For the coupling with two sfermions $\tilde{f}_i$ and two family partner sfermions with different isospin, $\hat{\tilde{f}}'_i$, we get with eqs. (A.40) and (A.42)

$$\begin{aligned} V_D \;\supset\; & \sum_f \bigg[-\frac{g^2}{8} \big(\tilde{f}_L^* \tilde{f}_L\big)\big(\hat{\tilde{f}}_L^* \hat{\tilde{f}}_L\big) + \frac{g'^2}{8} \Big(Y_{\tilde{f}_L} \big(\tilde{f}_L^* \tilde{f}_L\big) - Y_{\tilde{f}_R} \big(\tilde{f}_R^* \tilde{f}_R\big) \Big) \\ & \Big(Y_{\hat{\tilde{f}}'_L} \big(\hat{\tilde{f}}'^*_L \hat{\tilde{f}}'_L\big) - Y_{\hat{\tilde{f}}'_R} \big(\hat{\tilde{f}}'^*_R \hat{\tilde{f}}'_R\big) \Big) \bigg] \\ = \; & \sum_f \bigg\{ -\frac{g^2}{8} R^{\tilde{f}\hat{\tilde{f}}'_L}_{ijkl} + \frac{g'^2}{8} \Big[Y_{\tilde{f}_L} Y_{\hat{\tilde{f}}'_L} R^{\tilde{f}\hat{\tilde{f}}'_L}_{ijkl} + Y_{\tilde{f}_R} Y_{\hat{\tilde{f}}'_R} R^{\tilde{f}\hat{\tilde{f}}'_R}_{ijkl} - Y_{\tilde{f}_L} Y_{\hat{\tilde{f}}'_R} R^{\tilde{f}\hat{\tilde{f}}'_D}_{ijkl} - Y_{\tilde{f}_R} Y_{\hat{\tilde{f}}'_L} R^{\hat{\tilde{f}}'\tilde{f}_D}_{klij} \Big] \bigg\} \\ & \times \big(\tilde{f}_i^* \tilde{f}_j\big)\big(\hat{\tilde{f}}'^*_k \hat{\tilde{f}}'_l\big) \;\equiv\; \sum_f C^{\tilde{f}\hat{\tilde{f}}'}_{ijkl} \big(\tilde{f}_i^* \tilde{f}_j\big)\big(\hat{\tilde{f}}'^*_k \hat{\tilde{f}}'_l\big) \, . \end{aligned} \tag{A.93}$$



$$-i \Big(C^{\tilde{f}\hat{\tilde{f}}'}_{ijkl} + C^{\hat{\tilde{f}}'\tilde{f}}_{klij} \Big)$$

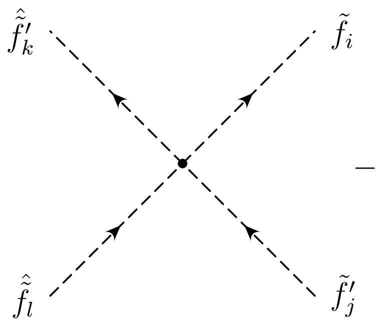
$\tilde{f}_i \tilde{f}'_j \hat{\tilde{f}}'_k \hat{\tilde{f}}_l$ couplings

Finally, eqs. (A.41) and (A.79) give the Feynman rule for the mixed four–sfermion coupling,

$$\begin{aligned} V_{F+D} &\supset \sum_f \Big(- h_f h_{\hat{f}} \tilde{f}_R^* \tilde{f}'_L \hat{\tilde{f}}'^*_L \hat{\tilde{f}}_R + \frac{g^2}{4} \tilde{f}_L^* \tilde{f}'_L \hat{\tilde{f}}'^*_L \hat{\tilde{f}}_L \Big) \\ &= \sum_f \Big(- h_f h_{\hat{f}} R^{\tilde{f}'\tilde{f}\hat{\tilde{f}}'\hat{\tilde{f}}_F}_{jikl} + \frac{g^2}{4} R^{\tilde{f}\tilde{f}'\hat{\tilde{f}}'\hat{\tilde{f}}_L}_{ijkl} \Big) \tilde{f}_i^* \tilde{f}'_j \hat{\tilde{f}}'^*_k \hat{\tilde{f}}_l \end{aligned} \qquad \text{(A.94)}$$

with

$$R^{\tilde{f}'\tilde{f}\hat{\tilde{f}}'\hat{\tilde{f}}_F}_{ijkl} \equiv R^{\tilde{f}'}_{i1} R^{\tilde{f}}_{j2} R^{\hat{\tilde{f}}'}_{k1} R^{\hat{\tilde{f}}}_{l2} , \qquad R^{\tilde{f}\tilde{f}'\hat{\tilde{f}}'\hat{\tilde{f}}_L}_{ijkl} \equiv R^{\tilde{f}}_{i1} R^{\tilde{f}'}_{j1} R^{\hat{\tilde{f}}'}_{k1} R^{\hat{\tilde{f}}}_{l1} . \qquad \text{(A.95)}$$



$$-i h_f h_{\hat{f}} R^{\tilde{f}'\tilde{f}\hat{\tilde{f}}'\hat{\tilde{f}}_F}_{jikl} - i h_{\hat{f}'} h_{f'} R^{\hat{\tilde{f}}\hat{\tilde{f}}'\tilde{f}\tilde{f}'_F}_{lkij} - i \frac{g^2}{2} R^{\tilde{f}\tilde{f}'\hat{\tilde{f}}'\hat{\tilde{f}}_L}_{ijkl}$$

A.7 Vector boson–Sfermion–Sfermion couplings

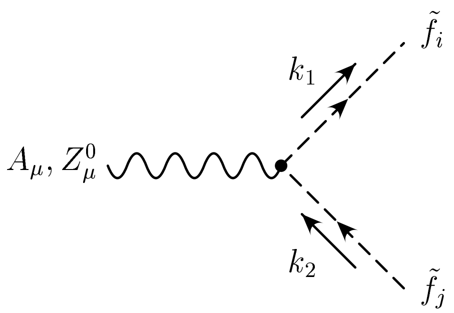
The interaction Lagrangian of a vector boson and two sfermions is given by

$$\mathcal{L} = -i\, e\, e_f\, A_\mu\, \tilde{f}_i^* \overleftrightarrow{\partial^\mu} \tilde{f}_j - i\, g_Z\, z^{\tilde{f}}_{ij}\, Z^0_\mu\, \tilde{f}_i^* \overleftrightarrow{\partial^\mu} \tilde{f}_j + \left(-i \frac{g}{\sqrt{2}} W^+_\mu R^{\tilde{f}_\uparrow}_{i1} R^{\tilde{f}_\downarrow}_{j1} \tilde{f}^*_{\uparrow i} \overleftrightarrow{\partial^\mu} \tilde{f}_{\downarrow j} + \text{h.c.} \right) , \qquad \text{(A.96)}$$

with the abbreviations

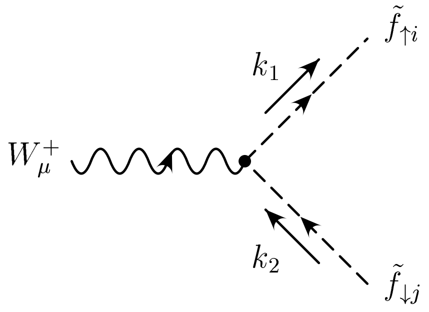
$$z^{\tilde{f}}_{ij} = C^f_L R^{\tilde{f}}_{i1} R^{\tilde{f}}_{j1} + C^f_R R^{\tilde{f}}_{i2} R^{\tilde{f}}_{j2} \qquad \text{(A.97)}$$

and $g_Z = g/\cos\theta_W$, $C^f_L = I^{3L}_f - e_f s^2_W$ and $C^f_R = -e_f s^2_W$.



$$-ie\, e_f \delta_{ij} (k_1 + k_2)_\mu \qquad \ldots A_\mu$$

$$i g_Z z^{\tilde{f}}_{ij} (k_1 + k_2)_\mu \qquad \ldots Z^0_\mu$$



$$-i\,\frac{g}{\sqrt{2}}\,R^{\tilde{f}_\uparrow}_{i1}R^{\tilde{f}_\downarrow}_{j1}(k_1+k_2)_\mu$$

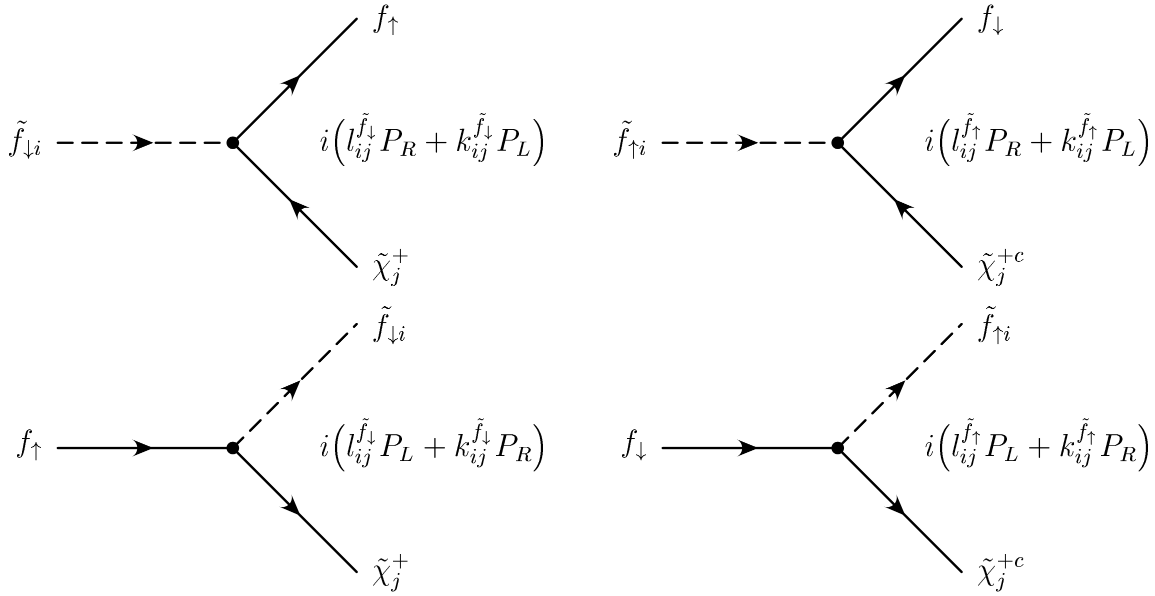
A.8 Gaugino–Fermion–Sfermion couplings

The interaction Lagrangian of the chargino–sfermion–fermion couplings is given by

$$\begin{aligned}\mathcal{L} &= \bar{f}_\uparrow\Big(l^{\tilde{f}_\downarrow}_{ij}P_R + k^{\tilde{f}_\downarrow}_{ij}P_L\Big)\tilde{\chi}^+_j\,\tilde{f}_{\downarrow i} + \bar{f}_\downarrow\Big(l^{\tilde{f}_\uparrow}_{ij}P_R + k^{\tilde{f}_\uparrow}_{ij}P_L\Big)\tilde{\chi}^{+c}_j\,\tilde{f}_{\uparrow i} \\ &+ \bar{\tilde{\chi}}^+_j\Big(l^{\tilde{f}_\downarrow}_{ij}P_L + k^{\tilde{f}_\downarrow}_{ij}P_R\Big)f_\uparrow\,\tilde{f}^*_{\downarrow i} + \bar{\tilde{\chi}}^{+c}_j\Big(l^{\tilde{f}_\uparrow}_{ij}P_L + k^{\tilde{f}_\uparrow}_{ij}P_R\Big)f_\downarrow\,\tilde{f}^*_{\uparrow i} \qquad (A.98)\end{aligned}$$

with the coupling matrices

$$\begin{aligned} l^{\tilde{f}_\uparrow}_{ij} &= -gV_{j1}R^{\tilde{f}_\uparrow}_{i1} + h_{f_\uparrow}V_{j2}R^{\tilde{f}_\uparrow}_{i2}\,, & l^{\tilde{f}_\downarrow}_{ij} &= -gU_{j1}R^{\tilde{f}_\downarrow}_{i1} + h_{f_\downarrow}U_{j2}R^{\tilde{f}_\downarrow}_{i2}\,, \\ k^{\tilde{f}_\uparrow}_{ij} &= h_{f_\downarrow}U_{j2}R^{\tilde{f}_\uparrow}_{i1}\,, & k^{\tilde{f}_\downarrow}_{ij} &= h_{f_\uparrow}V_{j2}R^{\tilde{f}_\downarrow}_{i1}\,. \end{aligned} \qquad (A.99)$$



For the neutralino–sfermion–fermion couplings the Lagrangian reads

$$\mathcal{L} = \bar{f}\Big(a^{\tilde{f}}_{ik}P_R + b^{\tilde{f}}_{ik}P_L\Big)\tilde{\chi}^0_k\,\tilde{f}_i + \bar{\tilde{\chi}}^0_k\Big(a^{\tilde{f}}_{ik}P_L + b^{\tilde{f}}_{ik}P_R\Big)f\,\tilde{f}^*_i \qquad (A.100)$$

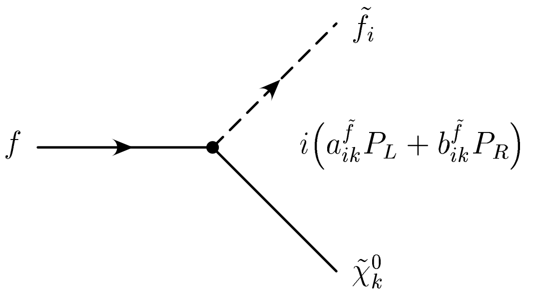
with the coupling matrices

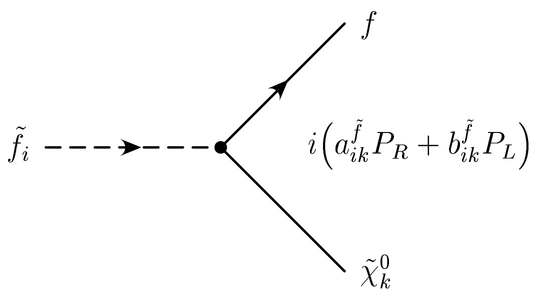
$$a_{ik}^{\tilde{f}} \;=\; h_f Z_{kx} R_{i2}^{\tilde{f}} + g f_{Lk}^{f} R_{i1}^{\tilde{f}} \,, \qquad b_{ik}^{\tilde{f}} \;=\; h_f Z_{kx} R_{i1}^{\tilde{f}} + g f_{Rk}^{f} R_{i2}^{\tilde{f}} \tag{A.101}$$

and

$$f_{Lk}^{f} = \sqrt{2}\left((e_f - I_f^{3L})\tan\theta_W Z_{k1} + I_f^{3L} Z_{k2}\right)\,, \quad f_{Rk}^{f} = -\sqrt{2} e_f \tan\theta_W Z_{k1}\,. \tag{A.102}$$

x takes the values $\{3,4\}$ for $\{$down, up$\}$–type case, respectively.





A.9 Higgs–Vector boson–Vector boson couplings

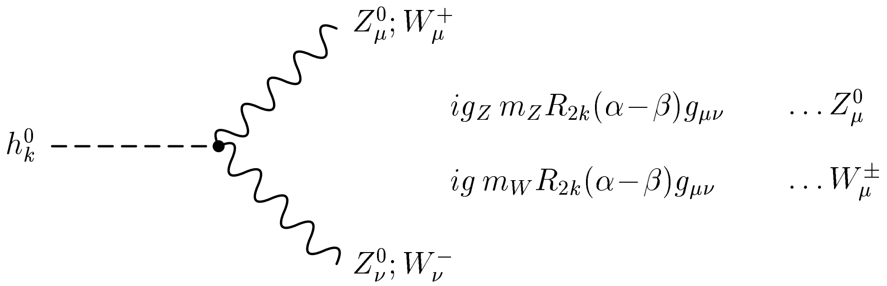
The interaction Lagrangian describing the couplings of one Higgs boson to two gauge bosons in the MSSM is given by

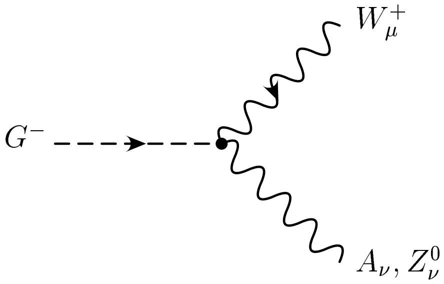
$$\begin{aligned}\mathcal{L} \;=\; & \frac{g_Z\, m_Z}{2}\big[\cos(\alpha-\beta) H^0 Z_\mu^0\, Z^{0\mu} - \sin(\alpha-\beta) h^0 Z_\mu^0\, Z^{0\mu}\big] \\ & + g\, m_W \big[\cos(\alpha-\beta) H^0 W_\mu^+ W^{-\mu} - \sin(\alpha-\beta) h^0 W_\mu^+ W^{-\mu}\big] \\ & - g_Z\, m_W\, s_W^2\, G^- W_\mu^+\, Z^{0\mu} + g s_W m_W\, G^- W_\mu^+ A^\mu + \text{h.c.}\end{aligned} \tag{A.103}$$

With the usual form of rotation matrices, used throughout this paper,

$$R_{kl}(\phi) \;\equiv\; \begin{pmatrix} \cos\phi & \sin\phi \\ -\sin\phi & \cos\phi \end{pmatrix}_{kl} \,, \tag{A.104}$$

the Feynman rules can then be written as





$i g s_W m_W g_{\mu\nu}$ $\quad \ldots A_\nu$

$-i g_Z m_W s_W^2 g_{\mu\nu}$ $\quad \ldots Z_\nu^0$

Appendix B

Self–energies and counter terms

Here we give the explicit form of the self–energies needed for the computation of the one–loop decay widths $\{h^0, H^0, A^0\} \to \tilde{f}_i \bar{\tilde{f}}_j$ and $H^+ \to \tilde{t}_i \bar{\tilde{b}}_j$.

B.1 Diagonal Wave–function corrections — derivatives of Higgs boson self–energies

The conventional on–shell renormalization conditions for the diagonal wave–function renormalization constants are given in terms of the derivatives of the corresponding self–energies (see chapter 3),

$$\delta Z^{H^0}_{kk} \;=\; -\,\mathrm{Re}\,\dot{\Pi}^{H^0}_{kk}(m^2_{H^0_k})\,, \tag{B.1}$$

where the dot in $\dot{\Pi}_{\ldots}(k^2)$ denotes the derivative with respect to k^2. In the following we list the single contributions of the Higgs wave–function corrections. The derivatives of the CP–even Higgs bosons h^0 and H^0 depicted in Fig. B.1 are given as follows:

$$\dot{\Pi}^{H^0,f}_{kk} \;=\; -\frac{2}{(4\pi)^2}\sum_f N^f_C\big(s^f_k\big)^2\Big[(4m_f^2 - m^2_{h^0_k})\dot{B}_0(m^2_{h^0_k}, m_f^2, m_f^2) - B_0(m^2_{h^0_k}, m_f^2, m_f^2)\Big] \tag{B.2}$$

$$\dot{\Pi}^{H^0,\tilde{f}}_{kk} \;=\; \frac{1}{(4\pi)^2}\sum_f\sum_{m,n=1}^{2} N^f_C\left(G^{\tilde{f}}_{mnk}\right)^2 \dot{B}_0(m^2_{h^0_k}, m^2_{\tilde{f}_m}, m^2_{\tilde{f}_n}) \tag{B.3}$$

$$\dot{\Pi}^{H^0,\tilde{\chi}^0}_{kk} \;=\; -\frac{1}{(4\pi)^2}\,g^2\sum_{m,n=1}^{4}(F^0_{mnk})^2\Big[\Big((m_{\tilde{\chi}^0_m} + m_{\tilde{\chi}^0_n})^2 - m^2_{h^0_k}\Big)\dot{B}_0 - B_0\Big](m^2_{h^0_k}, m^2_{\tilde{\chi}^0_m}, m^2_{\tilde{\chi}^0_n}) \tag{B.4}$$

$$\begin{aligned}\dot{\Pi}^{H^0,\tilde{\chi}^+}_{kk} \;=\; -\frac{1}{(4\pi)^2}\,g^2\sum_{m,n=1}^{2}\Big[&\Big((F^+_{mnk})^2 + (F^+_{nmk})^2\Big)\Big((m^2_{\tilde{\chi}^+_m} + m^2_{\tilde{\chi}^+_n} - m^2_{h^0_k})\dot{B}_0 - B_0\Big)\\ &+ 4m_{\tilde{\chi}^+_m}m_{\tilde{\chi}^+_n}F^+_{mnk}\,F^+_{nmk}\,\dot{B}_0\Big](m^2_{h^0_k}, m^2_{\tilde{\chi}^+_m}, m^2_{\tilde{\chi}^+_n})\end{aligned} \tag{B.5}$$

$$\dot{\Pi}^{H^0,H}_{kk} \;=\; \frac{1}{(4\pi)^2}\frac{1}{2}\left(\frac{g_Z\,m_Z}{4}\right)^2\Bigg[\sum_{m,n=1}^{2}[(2+\delta_{km}\delta_{mn})!]^2\big(\cos 2\alpha\,\tilde{A}^{(k)}_{mn} - 2\sin 2\alpha\,\tilde{B}^{(k)}_{mn}\big)^2\dot{B}_0$$

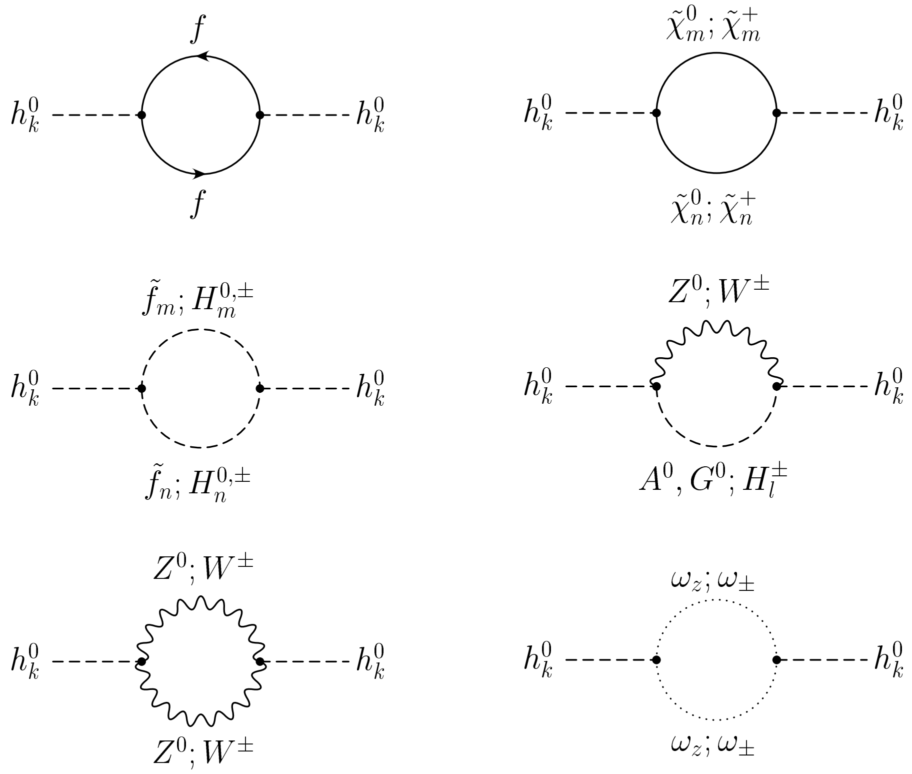


Figure B.1: Diagonal self–energies of CP–even Higgs bosons h^0 and H^0

$$
\begin{aligned}
&+4\sum_{m,n=3}^{4}\sin^2[\alpha+\beta-\frac{\pi}{2}(k-1)]\,\big(\tilde C_{m-2,n-2}\big)^2\dot B_0\Big](m^2_{h^0_k},m^2_{H^0_m},m^2_{H^0_n})\\
&+\frac{1}{(4\pi)^2}\sum_{m,n=1}^{2}\Big[(-1)^{mn}\,\frac{g\,m_W}{2}(1-\delta_{m2}\delta_{n2})(1+\delta_{mn})\tilde A'^{(k)}_{mn}-\frac{g_Z\,m_Z}{2}\\
&\qquad\times\sin[\alpha+\beta-\frac{\pi}{2}(k-1)]\,\tilde C_{mn}\Big]^2\dot B_0(m^2_{h^0_k},m^2_{H^+_m},m^2_{H^+_n}) \qquad \text{(B.6)}
\end{aligned}
$$

$$
\begin{aligned}
\dot\Pi^{H^0,V}_{kk} &= -\frac{1}{(4\pi)^2}\frac{g^2}{2}\sum_{l=1}^{2}\big(R_{lk}(\alpha-\beta)\big)^2\Big[\big(2m^2_{h^0_k}+2m^2_{H^+_l}-m^2_W\big)\dot B_0+2B_0\Big](m^2_{h^0_k},m^2_{H^+_l},m^2_W)\\
&\quad-\frac{1}{(4\pi)^2}\frac{g_Z^2}{4}\sum_{l=1}^{2}\big(R_{lk}(\alpha-\beta)\big)^2\Big[\big(2m^2_{h^0_k}+2m^2_{H^0_{l+2}}-m^2_Z\big)\dot B_0+2B_0\Big](m^2_{h^0_k},m^2_{H^0_{l+2}},m^2_Z) \qquad \text{(B.7)}
\end{aligned}
$$

$$
\dot\Pi^{H^0,VV}_{kk} = \frac{1}{(4\pi)^2}\big(R_{2k}(\alpha-\beta)\big)^2\Big[4g^2m^2_W\dot B_0(m^2_{h^0_k},m^2_W,m^2_W)+2g_Z^2m^2_Z\dot B_0(m^2_{h^0_k},m^2_Z,m^2_Z)\Big] \qquad \text{(B.8)}
$$

$$
\dot\Pi^{H^0,\text{ghost}}_{kk} = -\frac{1}{(4\pi)^2}\big(R_{2k}(\alpha-\beta)\big)^2\Big[\frac{g^2}{2}m^2_W\dot B_0(m^2_{h^0_k},m^2_W,m^2_W)+\frac{g_Z^2}{4}m^2_Z\dot B_0(m^2_{h^0_k},m^2_Z,m^2_Z)\Big] \qquad \text{(B.9)}
$$

The derivatives of the CP–odd Higgs boson A^0 depicted in Fig. B.2 are given as follows:

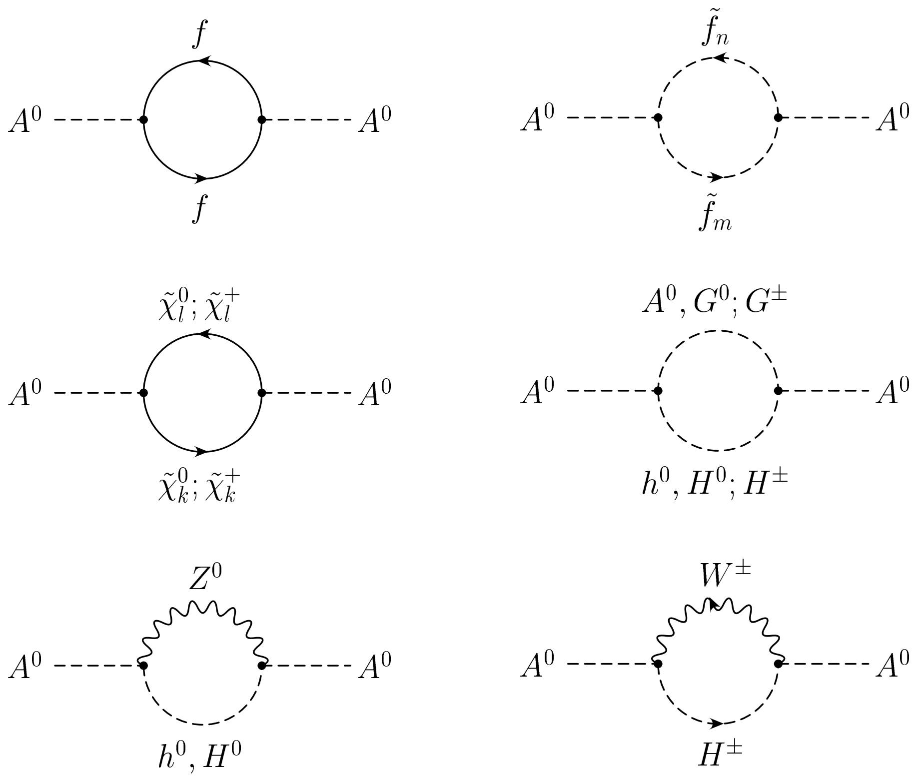


Figure B.2: Diagonal self–energies of the CP–odd Higgs boson A^0

$$\dot{\Pi}_{33}^{H^0,f} = -\frac{2}{(4\pi)^2} \sum_f N_C^f \big(s_3^f\big)^2 \Big[B_0(m_{A^0}^2, m_f^2, m_f^2) + m_{A^0}^2 \dot{B}_0(m_{A^0}^2, m_f^2, m_f^2)\Big] \quad \text{(B.10)}$$

$$\dot{\Pi}_{33}^{H^0,\tilde{f}} = -\frac{1}{(4\pi)^2} \sum_f \sum_{m,n=1}^{2} N_C^f \left(G_{mn3}^{\tilde{f}}\right)^2 \dot{B}_0(m_{A^0}^2, m_{\tilde{f}_m}^2, m_{\tilde{f}_n}^2) \quad \text{(B.11)}$$

$$\dot{\Pi}_{33}^{H^0,\tilde{\chi}^0} = -\frac{1}{(4\pi)^2} g^2 \sum_{m,n=1}^{4} (F_{mn3}^0)^2 \Big[\Big((m_{\tilde{\chi}_m^0} - m_{\tilde{\chi}_n^0})^2 - m_{A^0}^2\Big)\dot{B}_0 - B_0\Big](m_{A^0}^2, m_{\tilde{\chi}_m^0}^2, m_{\tilde{\chi}_n^0}^2) \quad \text{(B.12)}$$

$$\begin{aligned}\dot{\Pi}_{33}^{H^0,\tilde{\chi}^+} = & -\frac{1}{(4\pi)^2} g^2 \sum_{m,n=1}^{2} \Big[\Big((F_{mn3}^+)^2 + (F_{nm3}^+)^2\Big)\Big(\big(m_{\tilde{\chi}_m^+}^2 + m_{\tilde{\chi}_n^+}^2 - m_{A^0}^2\big)\dot{B}_0 - B_0\Big) \\ & - 4 m_{\tilde{\chi}_m^+} m_{\tilde{\chi}_n^+} F_{mn3}^+ F_{nm3}^+ \dot{B}_0\Big](m_{A^0}^2, m_{\tilde{\chi}_m^+}^2, m_{\tilde{\chi}_n^+}^2) \quad \text{(B.13)}\end{aligned}$$

$$\begin{aligned}\dot{\Pi}_{33}^{H^0,H} = & \frac{1}{(4\pi)^2}\left(\frac{g_Z\, m_Z}{2}\right)^2 \sum_{k=1}^{2}\sum_{l=3}^{4} \big(A_{k,l-2}\big)^2 \dot{B}_0(m_{A^0}^2, m_{h_k^0}^2, m_{H_l^0}^2) \\ & + \frac{1}{(4\pi)^2}\, 2 \left(\frac{g\, m_W}{2}\right)^2 \dot{B}_0(m_{A^0}^2, m_{H^+}^2, m_{W^+}^2) \quad \text{(B.14)}\end{aligned}$$

$$\text{with } A_{kl} = \begin{pmatrix} -\cos 2\beta\,\sin(\alpha+\beta) & -\sin 2\beta\,\sin(\alpha+\beta) \\ \cos 2\beta\,\cos(\alpha+\beta) & \sin 2\beta\,\cos(\alpha+\beta) \end{pmatrix}$$

$$\dot{\Pi}^{H^0,V}_{33} = -\frac{1}{(4\pi)^2}\frac{g_Z^2}{4}\sum_{l=1}^{2}\left(R_{1l}(\alpha-\beta)\right)^2\Big[\Big(2m^2_{A^0}+2m^2_{h^0_l}-m^2_Z\Big)\dot{B}_0+2B_0\Big](m^2_{A^0},m^2_{h^0_l},m^2_Z)$$
$$-\frac{1}{(4\pi)^2}\frac{g^2}{2}\Big[\left(2m^2_{A^0}+2m^2_{H^+}-m^2_W\right)\dot{B}_0+2B_0\Big](m^2_{A^0},m^2_{H^+},m^2_W) \qquad \text{(B.15)}$$

The diagrams in Fig. B.3 show the diagonal charged Higgs boson self–energies entering in the wave–function corrections. In this section, we will denote up– and down–type (s)fermions by $f_\uparrow/\tilde{f}_\uparrow$ and $f_\downarrow/\tilde{f}_\downarrow$, respectively.

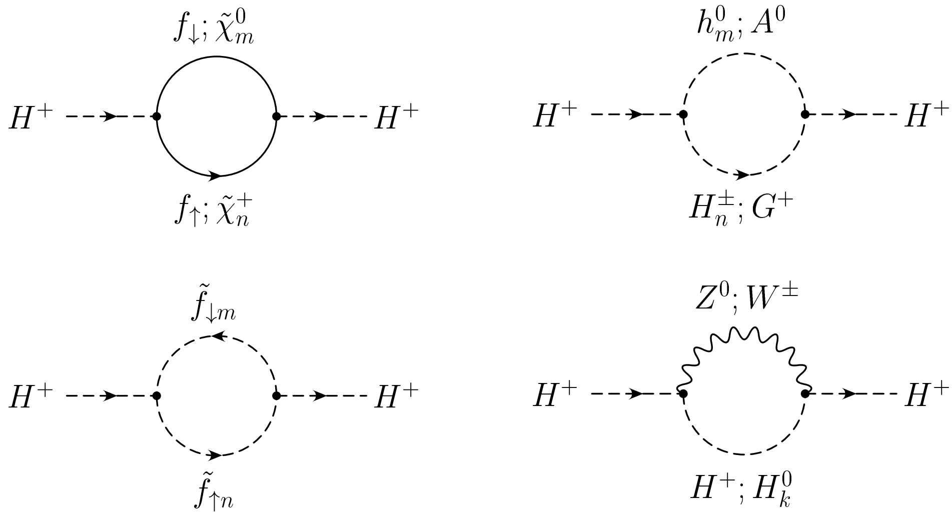


Figure B.3: Self–energy diagrams of the charged Higgs boson H^+ relevant for the calculation of diagonal wave–function corrections.

$$\dot{\Pi}^{H^+,f}_{11} = -\frac{1}{(4\pi)^2}\sum_{f=\{f_\uparrow\}} N_C^f\Big[\left(h^2_{f_\uparrow}\cos^2\beta+h^2_{f_\downarrow}\sin^2\beta\right)\left((m^2_{f_\uparrow}+m^2_{f_\downarrow}-m^2_{H^+})\dot{B}_0-B_0\right)$$
$$+4h_{f_\uparrow}m_{f_\uparrow}h_{f_\downarrow}m_{f_\downarrow}\sin\beta\cos\beta\,\dot{B}_0\Big](m^2_{H^+},m^2_{f_\downarrow},m^2_{f_\uparrow})\Big] \qquad \text{(B.16)}$$

$$\dot{\Pi}^{H^+,\tilde{\chi}}_{11} = -\frac{1}{(4\pi)^2}g^2\sum_{m=1}^{4}\sum_{n=1}^{2}\Big[\Big(\left(F^L_{nm1}\right)^2+\left(F^R_{nm1}\right)^2\Big)\Big)\Big(\left(m^2_{\tilde{\chi}^0_m}+m^2_{\tilde{\chi}^+_n}-m^2_{H^+}\right)\dot{B}_0-B_0\Big)$$
$$+4m_{\tilde{\chi}^0_m}m_{\tilde{\chi}^+_n}F^L_{nm1}F^R_{nm1}\dot{B}_0\Big](m^2_{H^+},m^2_{\tilde{\chi}^0_m},m^2_{\tilde{\chi}^+_n})\Big] \qquad \text{(B.17)}$$

$$\dot{\Pi}^{H^+,\tilde{f}}_{11} = \frac{1}{(4\pi)^2}\sum_{f=\{f_\uparrow\}}N_C^f\sum_{m,n=1}^{2}\left(G^{\tilde{f}_\uparrow\tilde{f}_\downarrow}_{nm1}\right)^2\dot{B}_0(m^2_{H^+},m^2_{\tilde{f}_{\downarrow m}},m^2_{\tilde{f}_{\uparrow n}}) \qquad \text{(B.18)}$$

$$\dot{\Pi}^{H^+,H}_{11} = \frac{1}{(4\pi)^2}\sum_{m,n=1}^{2}\Big[(-1)^n\frac{gm_W}{2}(1+\delta_{1n})\tilde{A}'^{(1)}_{nm}+\frac{g_Z m_Z}{2}R_{2m}(\alpha+\beta)\tilde{C}_{1n}\Big]^2$$
$$\times\,\dot{B}_0(m^2_{H^+},m^2_{h^0_m},m^2_{H^+_n})$$

$$+\frac{1}{(4\pi)^2}\left(\frac{g\,m_W}{2}\right)^2 \dot{B}_0(m^2_{H^+}, m^2_{A^0}, m^2_W) \tag{B.19}$$

$$\begin{aligned}
\dot{\Pi}^{H^+,HV}_{11} &= -\frac{1}{(4\pi)^2} e_0^2 \Big[\big(4m^2_{H^+} - \lambda^2\big)\dot{B}_0 + 2B_0\Big](m^2_{H^+}, m^2_{H^+}, \lambda^2) \\
&\quad -\frac{1}{(4\pi)^2} g_Z^2 \big(\frac{1}{2} - s_W^2\big)^2 \Big[\big(4m^2_{H^+} - m_Z^2\big)\dot{B}_0 + 2B_0\Big](m^2_{H^+}, m^2_{H^+}, m^2_Z) \\
&\quad -\frac{1}{(4\pi)^2}\frac{g^2}{4}\sum_{k=1}^{2}\big(R_{1k}(\alpha-\beta)\big)^2 \Big[\big(2m^2_{H^+} + 2m^2_{h^0_k} - m^2_W\big)\dot{B}_0 + 2B_0\Big](m^2_{H^+}, m^2_{h^0_k}, m^2_W) \\
&\quad -\frac{1}{(4\pi)^2}\frac{g^2}{4}\Big[\big(2m^2_{H^+} + 2m^2_{A^0} - m^2_W\big)\dot{B}_0 + 2B_0\Big](m^2_{H^+}, m^2_{A^0}, m^2_W)
\end{aligned} \tag{B.20}$$

B.2 Off–diagonal Wave–functions corrections — Mixing of CP–even Higgs bosons

According to chapter 3, the off–diagonal wave–function renormalization constants of the external sfermions are given by

$$\delta Z^{\tilde{f}}_{ij} = \frac{2}{m^2_{\tilde{f}_i} - m^2_{\tilde{f}_j}}\,\mathrm{Re}\,\Pi^{\tilde{f}}_{ij}(m^2_{\tilde{f}_j}), \qquad i \neq j\,. \tag{B.21}$$

The single contributions are as follows.

$$\Pi^{H^0,f}_{12}(k^2) = -\frac{2}{(4\pi)^2}\sum_f N_C^f s_1^f s_2^f \Big[(4m_f^2 - k^2)B_0(k^2, m_f^2, m_f^2) + A_0(m_f^2))\Big] \tag{B.22}$$

$$\begin{aligned}
\Pi^{H^0,\tilde{f}}_{12}(k^2) &= \frac{1}{(4\pi)^2}\sum_f\sum_{m,n=1}^{2} N_C^f\, G^{\tilde{f}}_{nm1} G^{\tilde{f}}_{mn2}\, B_0(k^2, m^2_{\tilde{f}_m}, m^2_{\tilde{f}_n}) \\
&\quad +\frac{1}{(4\pi)^2}\sum_f\sum_{m=1}^{2} N_C^f\big(h_f^2\, c^{\tilde{f}}_{12} + g^2(c^{\tilde{b}}_{12} - c^{\tilde{t}}_{12})\, e^{\tilde{f}}_{mm}\big)A_0(m^2_{\tilde{f}_m})
\end{aligned} \tag{B.23}$$

$$\begin{aligned}
\Pi^{H^0,\tilde{\chi}^0}_{12}(k^2) &= -\frac{1}{(4\pi)^2}\, g^2 \sum_{m,n=1}^{4} F^0_{nm1}F^0_{mn2}\Big[\big((m_{\tilde{\chi}^0_m} + m_{\tilde{\chi}^0_n})^2 - k^2\big)B_0(k^2, m^2_{\tilde{\chi}^0_m}, m^2_{\tilde{\chi}^0_n}) \\
&\qquad + A_0(m^2_{\tilde{\chi}^0_m}) + A_0(m^2_{\tilde{\chi}^0_n})\Big]
\end{aligned} \tag{B.24}$$

$$\begin{aligned}
\Pi^{H^0,\tilde{\chi}^+}_{12}(k^2) &= -\frac{1}{(4\pi)^2}\, g^2 \sum_{m,n=1}^{2}\Big[2m_{\tilde{\chi}^+_m} m_{\tilde{\chi}^+_n}\Big(F^+_{mn1}F^+_{nm2} + F^+_{nm1}F^+_{mn2}\Big)B_0(k^2, m^2_{\tilde{\chi}^+_m}, m^2_{\tilde{\chi}^+_n}) \\
&\qquad + \Big(F^+_{mn1}F^+_{mn2} + F^+_{nm1}F^+_{nm2}\Big)\Big(A_0(m^2_{\tilde{\chi}^+_m}) + A_0(m^2_{\tilde{\chi}^+_n}) \\
&\qquad + \big(m^2_{\tilde{\chi}^+_m} + m^2_{\tilde{\chi}^+_n} - k^2\big)B_0(k^2, m^2_{\tilde{\chi}^+_m}, m^2_{\tilde{\chi}^+_n})\Big)\Big]
\end{aligned} \tag{B.25}$$

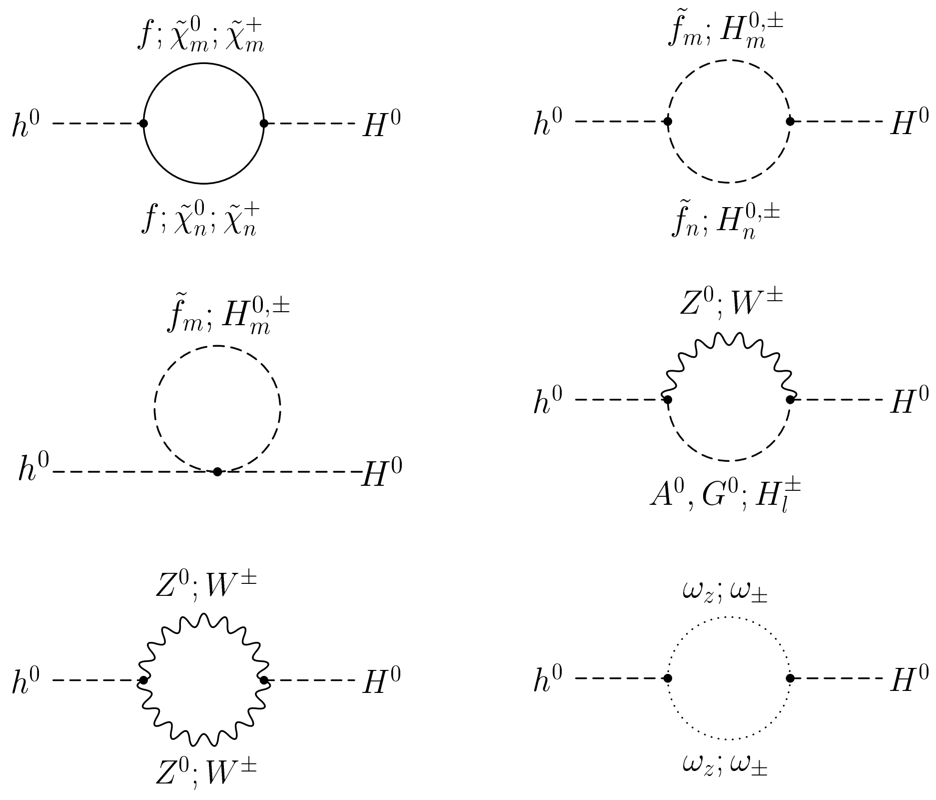


Figure B.4: Diagrams showing the off–diagonal mixing of the CP–even Higgs bosons h^0 and H^0

$$
\begin{aligned}
\Pi_{12}^{H^0,H}(k^2) \;=\;& \frac{1}{(4\pi)^2}\frac{1}{2}\left(\frac{g_Z\,m_Z}{4}\right)^2\Bigg[\sum_{m,n=1}^{2}(2{+}\delta_{1m}\delta_{mn})!\big(\cos 2\alpha\,\tilde A^{(1)}_{mn}-2\sin 2\alpha\,\tilde B^{(1)}_{mn}\big)\\
&\qquad\times(2{+}\delta_{2m}\delta_{mn})!\big(\cos 2\alpha\,\tilde A^{(2)}_{mn}-2\sin 2\alpha\,\tilde B^{(2)}_{mn}\big)B_0\\
&\qquad-2\sum_{m,n=3}^{4}\sin(2\alpha+2\beta)\,\big(\tilde C_{m-2,n-2}\big)^2B_0\Bigg](k^2,m^2_{H^0_m},m^2_{H^0_n})\\
&+\frac{1}{(4\pi)^2}\sum_{m,n=1}^{2}\bigg((-1)^{mn}\,\frac{g\,m_W}{2}(1-\delta_{m2}\delta_{n2})(1+\delta_{mn})\tilde A'^{(1)}_{mn}-\frac{g_Z\,m_Z}{2}s_{\alpha+\beta}\,\tilde C_{mn}\bigg)\\
&\quad\times\bigg((-1)^{mn}\,\frac{g\,m_W}{2}(1-\delta_{m2}\delta_{n2})(1+\delta_{mn})\tilde A'^{(2)}_{mn}-\frac{g_Z\,m_Z}{2}s_{\alpha+\beta-\pi/2}\,\tilde C_{mn}\bigg)\\
&\quad\times B_0(k^2,m^2_{H^+_m},m^2_{H^+_n})\\
&+\frac{1}{(4\pi)^2}\frac{g_Z^2}{8}\sin 2\alpha\Big(3\big(A_0(m^2_{h^0})-A_0(m^2_{H^0})\big)\cos 2\alpha\\
&\qquad+\big(A_0(m^2_{A^0})-A_0(m^2_Z)\big)\cos 2\beta\Big)\\
&-\frac{1}{(4\pi)^2}\Big(\frac{g^2}{2}\sin(\alpha{-}\beta)\cos(\alpha{-}\beta)-\frac{g_Z^2}{2}\sin 2\alpha\,\cos 2\beta\Big)\big(A_0(m^2_{H^+})-A_0(m^2_W)\big)
\end{aligned}
$$

(B.26)

$$\Pi_{12}^{H^0,V}(k^2) = -\frac{1}{(4\pi)^2}\frac{g^2}{2}\sum_{l=1}^{2}\Big[\big(2k^2+2m^2_{H_l^+}-m_W^2\big)B_0(k^2,m^2_{H_l^+},m_W^2)+2A_0(m_W^2) - A_0(m^2_{H_l^+})+\frac{1}{2c_W^2}\Big(\big(2k^2+2m^2_{H^0_{l+2}}-m_Z^2\big)B_0(k^2,m^2_{H^0_{l+2}},m_Z^2) + 2A_0(m_Z^2)-A_0(m^2_{H^0_{l+2}})\Big)\Big]R_{l1}(\alpha-\beta)R_{l2}(\alpha-\beta) \quad \text{(B.27)}$$

$$\Pi_{12}^{H^0,VV}(k^2) = -\frac{1}{(4\pi)^2}\sin(2\alpha-2\beta)\Big(2g^2m_W^2B_0(k^2,m_W^2,m_W^2)+g_Z^2m_Z^2B_0(k^2,m_Z^2,m_Z^2)\Big) \quad \text{(B.28)}$$

$$\Pi_{12}^{H^0,\mathrm{ghost}}(k^2) = \frac{1}{(4\pi)^2}\sin(2\alpha-2\beta)\Big[\frac{g^2}{2}m_W^2B_0(k^2,m_W^2,m_W^2)+\frac{g_Z^2}{4}m_Z^2B_0(k^2,m_Z^2,m_Z^2)\Big] \quad \text{(B.29)}$$

B.3 Diagonal Wave–function corrections — derivatives of sfermion self–energies

According to chapter 3, the diagonal wave–function renormalization constants of the external sfermions are determined by the derivatives of the sfermion self–energies,

$$\delta Z_{ii}^{\tilde f} = -\,\mathrm{Re}\,\dot\Pi_{ii}^{\tilde f}(m^2_{\tilde f_i})\,, \quad \text{(B.30)}$$

where the dot in $\dot\Pi_{\dots}(k^2)$ denotes the derivative with respect to k^2. The single contributions are listed below.

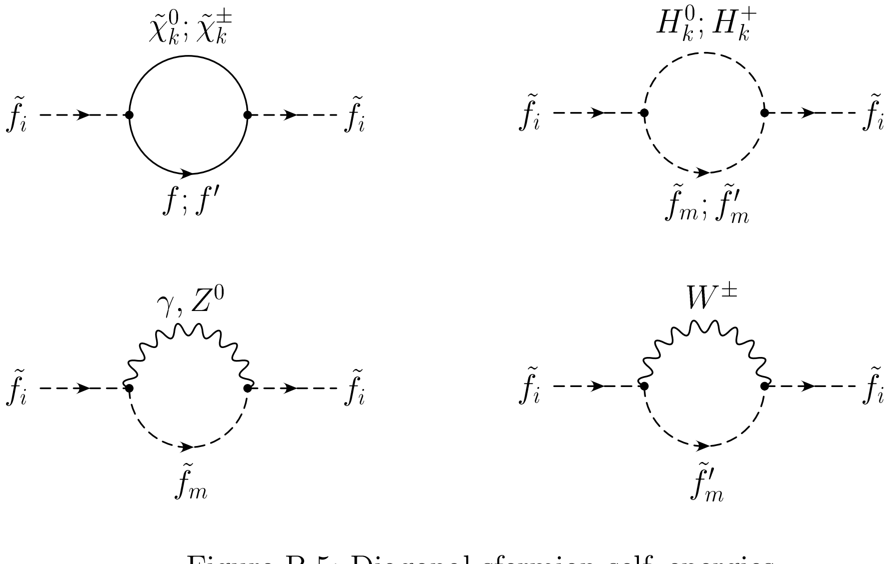


Figure B.5: Diagonal sfermion self–energies

$$\dot{\Pi}_{ii}^{\tilde{f},\tilde{\chi}} = -\frac{1}{(4\pi)^2}\sum_{k=1}^{4}\Big[\big((a_{ik}^{\tilde{f}})^2+(b_{ikl}^{\tilde{f}})^2\big)\cdot\big((m_{\tilde{\chi}_k^0}^2+m_f^2-m_{\tilde{f}_i}^2)\,\dot{B}_0-B_0\big) + 4m_{\tilde{\chi}_k^0}m_f\,a_{ik}^{\tilde{f}}b_{ik}^{\tilde{f}}\,\dot{B}_0\Big](m_{\tilde{f}_i}^2,m_{\tilde{\chi}_k^0}^2,m_f^2)$$
$$-\frac{1}{(4\pi)^2}\sum_{k=1}^{2}\Big[\big((k_{ik}^{\tilde{f}})^2+(l_{ik}^{\tilde{f}})^2\big)\cdot\big((m_{\tilde{\chi}_k^+}^2+m_{f'}^2-m_{\tilde{f}_i}^2)\,\dot{B}_0-B_0\big) + 4m_{\tilde{\chi}_k^+}m_{f'}\,k_{ik}^{\tilde{f}}l_{ik}^{\tilde{f}}\,\dot{B}_0\Big](m_{\tilde{f}_i}^2,m_{\tilde{\chi}_k^+}^2,m_{f'}^2) \quad \text{(B.31)}$$

$$\dot{\Pi}_{ii}^{\tilde{f},H} = \frac{1}{(4\pi)^2}\sum_{k=1}^{4}\sum_{m=1}^{2}G_{mik}^{\tilde{f}}\,G_{imk}^{\tilde{f}}\,\dot{B}_0(m_{\tilde{f}_i}^2,m_{\tilde{f}_m}^2,m_{H_k^0}^2) + \frac{1}{(4\pi)^2}\sum_{k=1}^{2}\sum_{m=1}^{2}G_{mik}^{\tilde{f}'\tilde{f}}\,G_{imk}^{\tilde{f}\tilde{f}'}\,\dot{B}_0(m_{\tilde{f}_i}^2,m_{\tilde{f}'_m}^2,m_{H_k^+}^2) \quad \text{(B.32)}$$

$$\dot{\Pi}_{ii}^{\tilde{f},\gamma} = -\frac{1}{(4\pi)^2}(e_0e_f)^2\Big[2B_0+\big(4m_{\tilde{f}_i}^2-\lambda^2\big)\dot{B}_0\Big](m_{\tilde{f}_i}^2,m_{\tilde{f}_i}^2,\lambda^2) \quad \text{(B.33)}$$

$$\dot{\Pi}_{ii}^{\tilde{f},Z} = -\frac{1}{(4\pi)^2}g_Z^2\sum_{m=1}^{2}\big(z_{im}^{\tilde{f}}\big)^2\Big[2B_0+\big(2m_{\tilde{f}_i}^2+2m_{\tilde{f}_m}^2-m_Z^2\big)\dot{B}_0\Big](m_{\tilde{f}_i}^2,m_{\tilde{f}_m}^2,m_Z^2) \quad \text{(B.34)}$$

$$\dot{\Pi}_{ii}^{\tilde{f},W} = -\frac{1}{(4\pi)^2}\frac{g^2}{2}\sum_{m=1}^{2}\big(R_{i1}^{\tilde{f}}R_{m1}^{\tilde{f}'}\big)^2\Big[2B_0+\big(2m_{\tilde{f}_i}^2+2m_{\tilde{f}'_m}^2-m_W^2\big)\dot{B}_0\Big](m_{\tilde{f}_i}^2,m_{\tilde{f}'_m}^2,m_W^2) \quad \text{(B.35)}$$

B.4 Sfermion self–energies

For the fixing of the sfermion mixing angle $\theta_{\tilde{f}}$ we need the off–diagonal elements of the sfermion self–energies, $\Pi_{ij}^{\tilde{f}} = \Pi_{ij}^{\tilde{f}}(m_{\tilde{f}_j}^2)$. In the following, $Y_{L/R}^f$ denotes the weak hypercharge, $Y_{L/R}^f = 2(I_f^{3L/R} - e_f)$. The short forms for various products of sfermion rotation matrices can be found in Appendix A. Additionally, we use the abbreviation $R_{ijkl}^{\tilde{f}\tilde{f}\hat{\tilde{f}}_D} = R_{i1}^{\tilde{f}}R_{j1}^{\tilde{f}}R_{k2}^{\hat{\tilde{f}}}R_{l2}^{\hat{\tilde{f}}}$.

$$\Pi_{ij}^{\tilde{f},\tilde{\chi}} = -\frac{1}{(4\pi)^2}\sum_{k=1}^{4}\Big[\big(a_{ik}^{\tilde{f}}a_{jk}^{\tilde{f}}+b_{ik}^{\tilde{f}}b_{jk}^{\tilde{f}}\big)\big(A_0(m_{\tilde{\chi}_k^0}^2)+A_0(m_f^2)+(m_{\tilde{\chi}_k^0}^2+m_f^2-m_{\tilde{f}_j}^2)\,B_0\big) + 2m_{\tilde{\chi}_k^0}m_f\big(a_{ik}^{\tilde{f}}b_{jk}^{\tilde{f}}+b_{ik}^{\tilde{f}}a_{jk}^{\tilde{f}}\big)B_0\Big](m_{\tilde{f}_j}^2,m_{\tilde{\chi}_k^0}^2,m_f^2)$$
$$-\frac{1}{(4\pi)^2}\sum_{k=1}^{2}\Big[\big(k_{ik}^{\tilde{f}}k_{jk}^{\tilde{f}}+l_{ik}^{\tilde{f}}l_{jk}^{\tilde{f}}\big)\big(A_0(m_{\tilde{\chi}_k^+}^2)+A_0(m_{f'}^2)+(m_{\tilde{\chi}_k^+}^2+m_{f'}^2-m_{\tilde{f}_j}^2)\,B_0\big) + 2m_{\tilde{\chi}_k^+}m_{f'}\big(k_{ik}^{\tilde{f}}l_{jk}^{\tilde{f}}+l_{ik}^{\tilde{f}}k_{jk}^{\tilde{f}}\big)B_0\Big](m_{\tilde{f}_j}^2,m_{\tilde{\chi}_k^+}^2,m_{f'}^2) \quad \text{(B.36)}$$

$$\Pi_{ij}^{\tilde{f},H\tilde{f}} = \frac{1}{(4\pi)^2}\sum_{k=1}^{4}\sum_{m=1}^{2}G_{mik}^{\tilde{f}}\,G_{jmk}^{\tilde{f}}\,B_0(m_{\tilde{f}_j}^2,m_{\tilde{f}_m}^2,m_{H_k^0}^2)$$

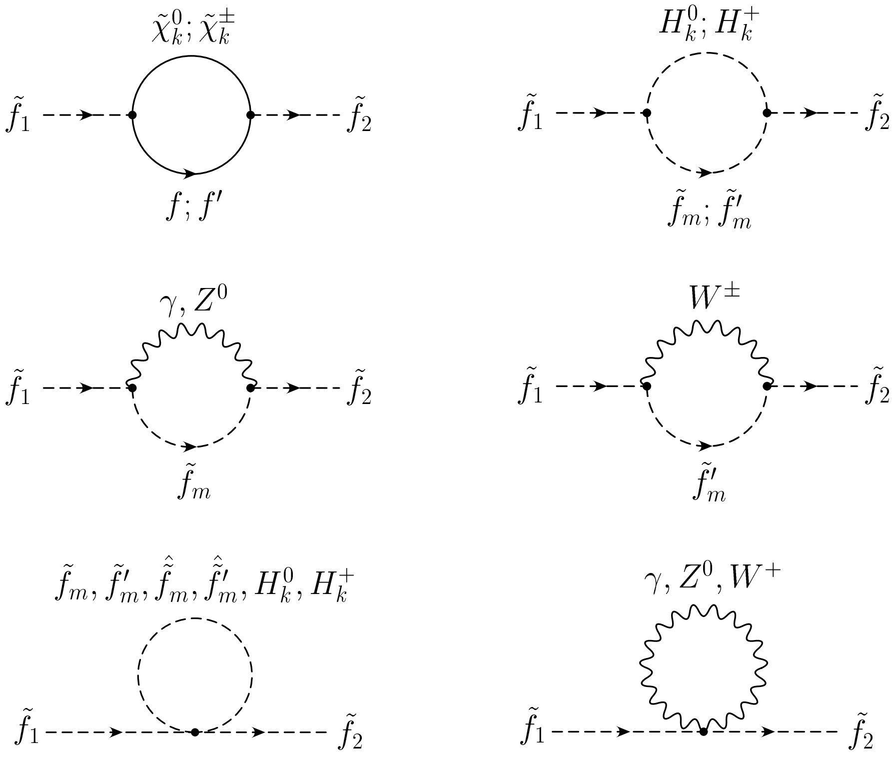


Figure B.6: Off–diagonal Sfermion self–energies

$$+\frac{1}{(4\pi)^2}\sum_{k=1}^{2}\sum_{m=1}^{2} G^{\tilde f'\tilde f}_{mik}\, G^{\tilde f\tilde f'}_{jmk}\, B_0(m^2_{\tilde f_j}, m^2_{\tilde f'_m}, m^2_{H^+_k}) \tag{B.37}$$

$$\Pi^{\tilde f,\gamma\tilde f}_{ij} = -\frac{1}{(4\pi)^2}(e_0 e_f)^2\delta_{ij}\Big[2A_0(\lambda^2) - A_0(m^2_{\tilde f_i}) + \big(4m^2_{\tilde f_i} - \lambda^2\big)B_0(m^2_{\tilde f_i}, m^2_{\tilde f_i}, \lambda^2)\Big] \tag{B.38}$$

$$\Pi^{\tilde f,Z\tilde f}_{ij} = -\frac{1}{(4\pi)^2}\, g_Z^2 \sum_{m=1}^{2} z^{\tilde f}_{mi} z^{\tilde f}_{jm}\Big[2A_0(m_Z^2) - A_0(m^2_{\tilde f_m}) + \big(2m^2_{\tilde f_j} + 2m^2_{\tilde f_m} - m_Z^2\big) \times B_0(m^2_{\tilde f_j}, m^2_{\tilde f_m}, m_Z^2)\Big] \tag{B.39}$$

$$\Pi^{\tilde f,W\tilde f'}_{ij} = -\frac{1}{(4\pi)^2}\frac{g^2}{2} R^{\tilde f}_{i1} R^{\tilde f}_{j1}\sum_{m=1}^{2}\big(R^{\tilde f'}_{m1}\big)^2\Big[\big(2m^2_{\tilde f_j} + 2m^2_{\tilde f'_m} - m_W^2\big)B_0(m^2_{\tilde f_j}, m^2_{\tilde f'_m}, m_W^2) + 2A_0(m_W^2) - A_0(m^2_{\tilde f'_m})\Big] \tag{B.40}$$

$$\Pi^{\tilde f,\tilde f}_{ij} = \frac{1}{(4\pi)^2}\, h_f^2 \sum_{m=1}^{2}\Big[N_C^f\big(R^{\tilde f}_{jmmi} + R^{\tilde f}_{mijm}\big) + R^{\tilde f}_{jimm} + R^{\tilde f}_{mmji}\Big] A_0(m^2_{\tilde f_m})$$
$$+\frac{1}{(4\pi)^2}\, g_Z^2 \sum_{m=1}^{2}\Big\{\Big[\Big(\frac14 - (2I_f^{3L} - e_f)e_f s_W^2\Big) R^{\tilde f_L}_{jimm} + e_f^2 s_W^2 R^{\tilde f_R}_{jimm}\Big](N_C^f + 1)$$
$$+ (I_f^{3L} - e_f)e_f s_W^2\Big[N_C^f\big(R^{\tilde f}_{jimm} + R^{\tilde f}_{mmji}\big) + R^{\tilde f}_{jmmi} + R^{\tilde f}_{mijm}\Big]\Big\} A_0(m^2_{\tilde f_m})$$

$$\Pi_{ij}^{\tilde f,\tilde f'} = \frac{1}{(4\pi)^2}\sum_{m=1}^{2}\big(h_f^2\,R^{\tilde f'\tilde f_D}_{mmji} + h_{f'}^2\,R^{\tilde f\tilde f'_D}_{jimm}\big)A_0(m^2_{\tilde f'_m})$$
$$+\frac{1}{(4\pi)^2}\,\frac{g^2}{4}\sum_{m=1}^{2}\Big\{N_C^f\Big[\Big(t_W^2 Y_L^f Y_L^{f'} - 1\Big)R^{\tilde f\tilde f'_L}_{jimm} + t_W^2 Y_R^f Y_R^{f'}\,R^{\tilde f\tilde f'_R}_{jimm} - Y_L^f Y_R^{f'}\,R^{\tilde f\tilde f'_D}_{jimm}$$
$$- Y_L^{f'} Y_R^f R^{\tilde f'\tilde f_D}_{mmji}\Big] + 2R^{\tilde f\tilde f'_L}_{jimm}\Big\}A_0(m^2_{\tilde f'_m}) \qquad \text{(B.41)}$$

$$\Pi_{ij}^{\tilde f,\hat{\tilde f}} = \frac{1}{(4\pi)^2}N_C^{\hat f}\sum_{m=1}^{2}\Big(h_f h_{\hat f}\big(R^{\tilde f\hat{\tilde f}_F}_{ijmm} + R^{\tilde f\hat{\tilde f}_F}_{jimm}\big) + \frac{g^2}{4}R^{\tilde f\hat{\tilde f}_L}_{jimm}\Big)A_0(m^2_{\hat{\tilde f}_m})$$
$$+\frac{1}{(4\pi)^2}\,N_C^{\hat f}\,\frac{g'^2}{4}\sum_{m=1}^{2}\Big(Y_L^f Y_L^{\hat f} R^{\tilde f\hat{\tilde f}_L}_{jimm} - Y_L^f Y_R^{\hat f}\,R^{\tilde f\hat{\tilde f}_D}_{jimm} - Y_L^{\hat f} Y_R^f\,R^{\hat{\tilde f}\tilde f_D}_{mmji} + Y_R^f Y_R^{\hat f} R^{\tilde f\hat{\tilde f}_R}_{jimm}\Big)$$
$$\times A_0(m^2_{\hat{\tilde f}_m}) \qquad \text{(B.42)}$$

$$\Pi_{ij}^{\tilde f,\hat{\tilde f}'} = \frac{1}{(4\pi)^2}\,N_C^{\hat f}\,\frac{g'^2}{4}\sum_{m=1}^{2}\Big(Y_L^f Y_L^{\hat f'}\,R^{\tilde f\hat{\tilde f}'_L}_{jimm} - Y_L^f Y_R^{\hat f'}\,R^{\tilde f\hat{\tilde f}'_D}_{jimm} - Y_L^{\hat f'} Y_R^f\,R^{\hat{\tilde f}'\tilde f_D}_{mmji} + Y_R^f Y_R^{\hat f'}\,R^{\tilde f\hat{\tilde f}'_R}_{jimm}\Big)$$
$$\times A_0(m^2_{\hat{\tilde f}'_m}) \;-\; \frac{1}{(4\pi)^2}N_C^{\hat f}\,\frac{g^2}{4}\sum_{m=1}^{2}R^{\tilde f\hat{\tilde f}'_L}_{jimm}A_0(m^2_{\hat{\tilde f}'_m}) \qquad \text{(B.43)}$$

The contributions from first and second generation sfermions, $\tilde F_m$, are given by

$$\Pi_{ij}^{\tilde f,\tilde F} = \Pi_{ij}^{\tilde f,\hat{\tilde f}}(\hat f \to F)\,, \qquad \Pi_{ij}^{\tilde f,\hat{\tilde F}} = \Pi_{ij}^{\tilde f,\hat{\tilde f}}(\hat f \to \hat F)\,,$$
$$\Pi_{ij}^{\tilde f,\tilde F'} = \Pi_{ij}^{\tilde f,\hat{\tilde f}'}(\hat f' \to F')\,, \qquad \Pi_{ij}^{\tilde f,\hat{\tilde F}'} = \Pi_{ij}^{\tilde f,\hat{\tilde f}'}(\hat f' \to \hat F')\,, \qquad \text{(B.44)}$$

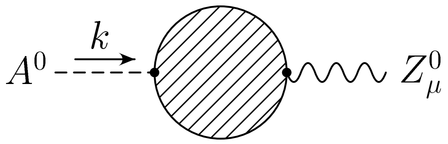
where the sub–/superscript $\tilde F$ denotes values belonging to first and second generation scalar fermions with same isospin as $\tilde f$ (e.g. $\tilde F_1 = \{\tilde u_1, \tilde c_1\}$ for the stop case, ...), $\tilde F'$ sfermions with different isospin etc.

$$\Pi_{ij}^{\tilde f,H} = \frac{1}{(4\pi)^2}\sum_{k=1}^{2}\Big(h_{f'}^2\,d^{\tilde f'}_{kk}\,R^{\tilde f}_{i1}R^{\tilde f}_{j1} + h_f^2\,d^{\tilde f}_{kk}\,R^{\tilde f}_{i2}R^{\tilde f}_{j2} + g^2\,f^{\tilde f}_{ij}\,\big(d^{\tilde b}_{kk} - d^{\tilde t}_{kk}\big)\Big)A_0(m^2_{H_k^+})$$
$$+\frac{1}{(4\pi)^2}\,\frac{1}{2}\sum_{k=1}^{4}\Big(h_f^2\,c^{\tilde f}_{kk}\,\delta_{ij} + g^2\,e^{\tilde f}_{ij}\,\big(c^{\tilde b}_{kk} - c^{\tilde t}_{kk}\big)\Big)A_0(m^2_{H_k^0}) \qquad \text{(B.45)}$$

$$\Pi_{ij}^{\tilde f,V} = \frac{1}{(4\pi)^2}\,4\,(e_0 e_f)^2\delta_{ij}A_0(\lambda^2) + \frac{1}{(4\pi)^2}\,2\,g^2 R^{\tilde f}_{i1}R^{\tilde f}_{j1}A_0(m_W^2)$$
$$+\frac{1}{(4\pi)^2}\,4\,g_Z^2\Big((C_L^f)^2 R^{\tilde f}_{i1}R^{\tilde f}_{j1} + (C_R^f)^2 R^{\tilde f}_{i2}R^{\tilde f}_{j2}\Big)A_0(m_Z^2) \qquad \text{(B.46)}$$

B.5 A^0Z^0–mixing

The scalar–vector mixing self–energy, $\Pi_{AZ}(k^2)$, is defined by the two–point function



$$\mathcal{M} = -i\, k^\mu\, \Pi_{AZ}(k^2)\, \epsilon^*_\mu(k)$$

$$\Pi^{f}_{AZ} = -\frac{i}{(4\pi)^2}\, m_Z \sin 2\beta \sum_f N_C^f\, I_f^{3L}\, h_f^2\, B_0(m^2_{A^0}, m^2_f, m^2_f) \tag{B.47}$$

$$\Pi^{\tilde{\chi}^0}_{AZ} = \frac{i}{(4\pi)^2}\, 2\, g\, g_Z \sum_{k,l=1}^{4} F^0_{kl3}\, O''^L_{lk}\Big[m_{\tilde{\chi}^0_l}\, B_0 + (m_{\tilde{\chi}^0_l} - m_{\tilde{\chi}^0_k}) B_1\Big](m^2_{A^0}, m^2_{\tilde{\chi}^0_l}, m^2_{\tilde{\chi}^0_k}) \tag{B.48}$$

$$\begin{aligned}\Pi^{\tilde{\chi}^+}_{AZ} = \frac{i}{(4\pi)^2}\, 2\, g\, g_Z \sum_{k,l=1}^{2} \Big[&\big(F^+_{kl3}\, O'^L_{lk} - F^+_{lk3}\, O'^R_{lk}\big) m_{\tilde{\chi}^+_l} (B_0 + B_1) \\ &+ \big(F^+_{kl3}\, O'^R_{lk} - F^+_{lk3}\, O'^L_{lk}\big) m_{\tilde{\chi}^+_k}\, B_1\Big](m^2_{A^0}, m^2_{\tilde{\chi}^+_l}, m^2_{\tilde{\chi}^+_k})\end{aligned} \tag{B.49}$$

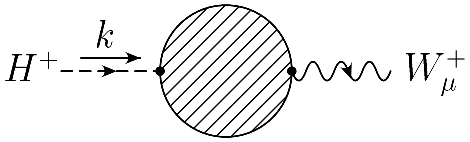
$$\Pi^{\tilde{f}}_{AZ} = -\frac{1}{(4\pi)^2}\, 2\, g_Z \sum_f N_C^f\, z^{\tilde{f}}_{21}\, G^{\tilde{f}}_{123}\big(B_0 + 2B_1\big)(m^2_{A^0}, m^2_{\tilde{f}_1}, m^2_{\tilde{f}_2}) \tag{B.50}$$

$$\Pi^{H}_{AZ} = \frac{i}{(4\pi)^2}\, \frac{g_Z^2\, m_Z}{4} \sum_{k=1}^{2} \sum_{l=3}^{4} A_{k,l-2}\, R_{k,l-2}(\beta - \alpha)\big(B_0 + 2B_1\big)(m^2_{A^0}, m^2_{H^0_l}, m^2_{H^0_k}) \tag{B.51}$$

$$\Pi^{Z}_{AZ} = \frac{i}{(4\pi)^2}\, \frac{g_Z^2\, m_Z}{4} \sin(2\alpha - 2\beta) \sum_{k=1}^{2} (-1)^k \big(B_0 - B_1\big)(m^2_{A^0}, m^2_{H^0_k}, m^2_Z) \tag{B.52}$$

B.6 H^+W^+–mixing

The scalar–vector mixing self–energy, $\Pi_{HW}(k^2)$, is defined by the two–point function



$$\mathcal{M} = -i\, k^\mu\, \Pi_{HW}(k^2)\, \epsilon^*_\mu(k)$$

$$\Pi^{f}_{HW} = \frac{1}{(4\pi)^2}\, \sqrt{2}\, g \sum_{f=\{f_\uparrow\}} N_C^f \left(m_{f_\downarrow} y_1^{f_\downarrow} B_1(m^2_{H^+}, m^2_{f_\uparrow}, m^2_{f_\downarrow}) - m_{f_\uparrow} y_1^{f_\uparrow} B_1(m^2_{H^+}, m^2_{f_\downarrow}, m^2_{f_\uparrow})\right)$$

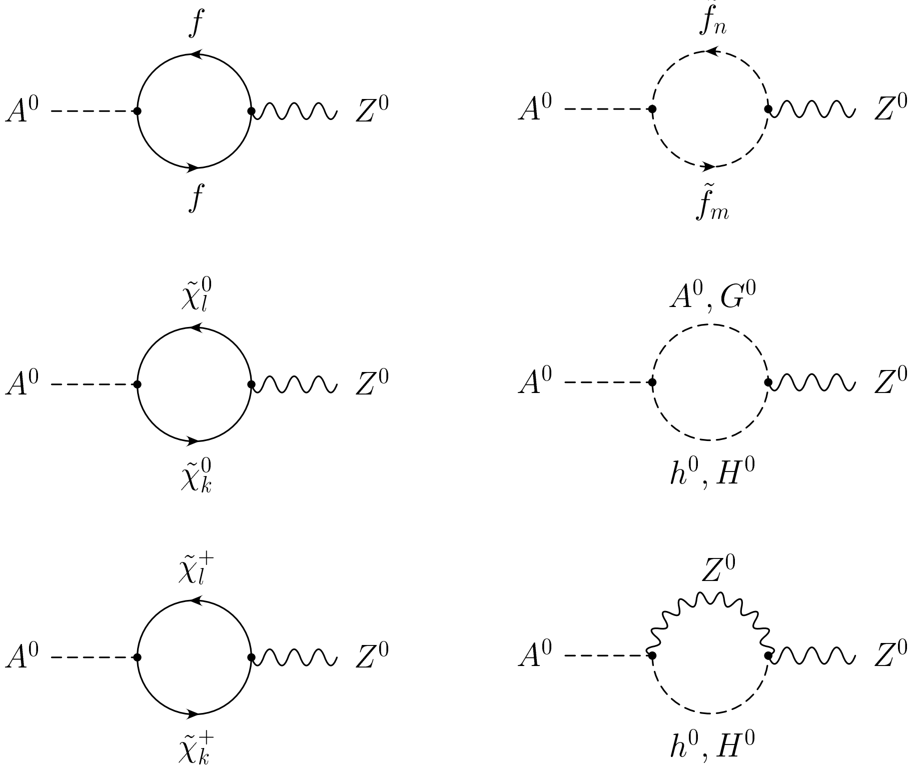


Figure B.7: AZ–mixing self–energies

$$(\text{B.53})$$

$$\Pi^{\tilde{\chi}}_{HW} = \frac{1}{(4\pi)^2}\, 2\, g^2 \sum_{k=1}^{2}\sum_{l=1}^{4} \Big[m_{\tilde{\chi}^+_k} \big(F^L_{kl1} O^L_{lk} + F^R_{kl1} O^R_{lk}\big)\big(B_0 + B_1\big) + m_{\tilde{\chi}^0_l} \big(F^L_{kl1} O^R_{lk} + F^R_{kl1} O^L_{lk}\big) B_1 \Big] (m^2_{H^+}, m^2_{\tilde{\chi}^+_k}, m^2_{\tilde{\chi}^0_l}) \tag{B.54}$$

$$\Pi^{\tilde{f}}_{HW} = -\frac{1}{(4\pi)^2} \frac{g}{\sqrt{2}} \sum_{f=\{f_\uparrow\}} N^f_C\, G^{\tilde{f}_\uparrow \tilde{f}_\downarrow}_{mn1} R^{\tilde{f}_\uparrow}_{m1} R^{\tilde{f}_\downarrow}_{n1} \big(B_0 + 2B_1\big)(m^2_{H^+}, m^2_{\tilde{f}_{\uparrow m}}, m^2_{\tilde{f}_{\downarrow n}}) \tag{B.55}$$

$$\Pi^{H}_{HW} = -\frac{1}{(4\pi)^2} \frac{g}{4} \sum_{m,n=1}^{2} \Big[(-1)^n g m_W (1+\delta_{1n}) \tilde{A}'^{(1)}_{nm} + g_Z m_Z R_{2m}(\alpha+\beta) \tilde{C}_{1n}\Big] R_{mn}(\beta-\alpha) \times \big(B_0 + 2B_1\big)(m^2_{H^+}, m^2_{H^+_n}, m^2_{h^0_m}) \tag{B.56}$$

$$\Pi^{W}_{HW} = -\frac{1}{(4\pi)^2} \frac{g^2\, m_W}{2} \sum_{k=1}^{2} R_{1k}(\alpha-\beta) R_{2k}(\alpha-\beta) \big(B_0 - B_1\big)(m^2_{H^+}, m^2_{h^0_k}, m^2_W) \tag{B.57}$$

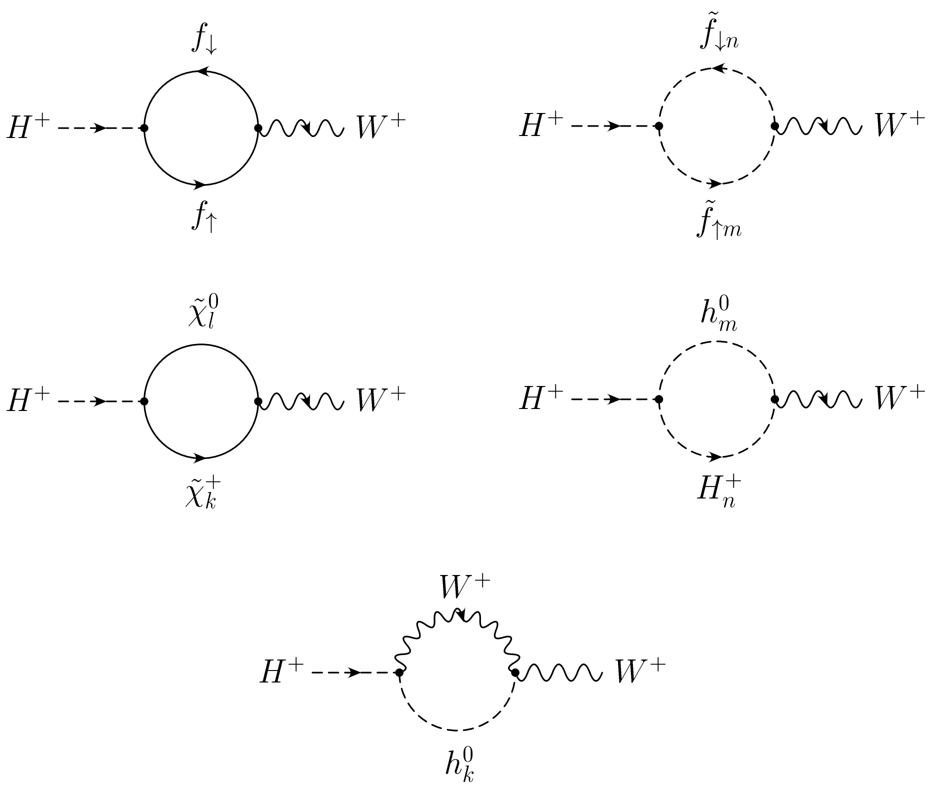
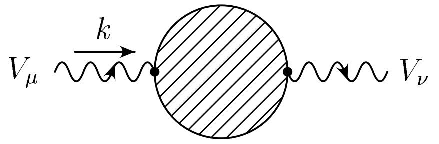


Figure B.8: H^+W^+–mixing self–energies

B.7 W^+ self–energies

For the calculation of the mass counter term of a gauge boson V ($V = W^\pm, Z^0$), $\delta m_V^2 = \mathrm{Re}\,\Pi^T_{VV}(m_V^2)$, we need the transverse part of the vector self–energy $\Pi^T_{VV}(k^2)$ from



$$\mathcal{M} = -i\,\epsilon_\mu(k)\Big(g^{\mu\nu}\Pi^T_{VV}(k^2) + k^\mu k^\nu\,\Pi^B_{VV}(k^2)\Big)\epsilon^*_\nu(k)\,. \tag{B.58}$$

$$\begin{aligned}\left(\frac{\delta m_W}{m_W}\right)^{ff} &= -\frac{1}{(4\pi)^2}\sum_{\text{gen.}} N_C^f \Bigg[h^2_{f_\uparrow}\sin^2\!\beta\,\frac{A_0(m^2_{f_\uparrow})}{m^2_{f_\uparrow}} + h^2_{f_\downarrow}\cos^2\!\beta\,B_0(m_W^2, m^2_{f_\downarrow}, m^2_{f_\uparrow}) \\ &\qquad - \frac{g^2}{m_W^2}B_{00}(m_W^2, m^2_{f_\downarrow}, m^2_{f_\uparrow}) + \frac{g^2}{2}B_1(m_W^2, m^2_{f_\downarrow}, m^2_{f_\uparrow})\Bigg]\end{aligned} \tag{B.59}$$

$$\left(\frac{\delta m_W}{m_W}\right)^{\tilde{f}\tilde{f}} = -\frac{1}{(4\pi)^2}\,\frac{g^2}{m_W^2}\sum_{\text{gen.}} N_C^f \sum_{m,n=1}^{2}\Big(R^{\tilde{f}_\uparrow}_{m1}R^{\tilde{f}_\downarrow}_{n1}\Big)^2 B_{00}(m_W^2, m^2_{\tilde{f}_{\uparrow m}}, m^2_{\tilde{f}_{\downarrow n}}) \tag{B.60}$$

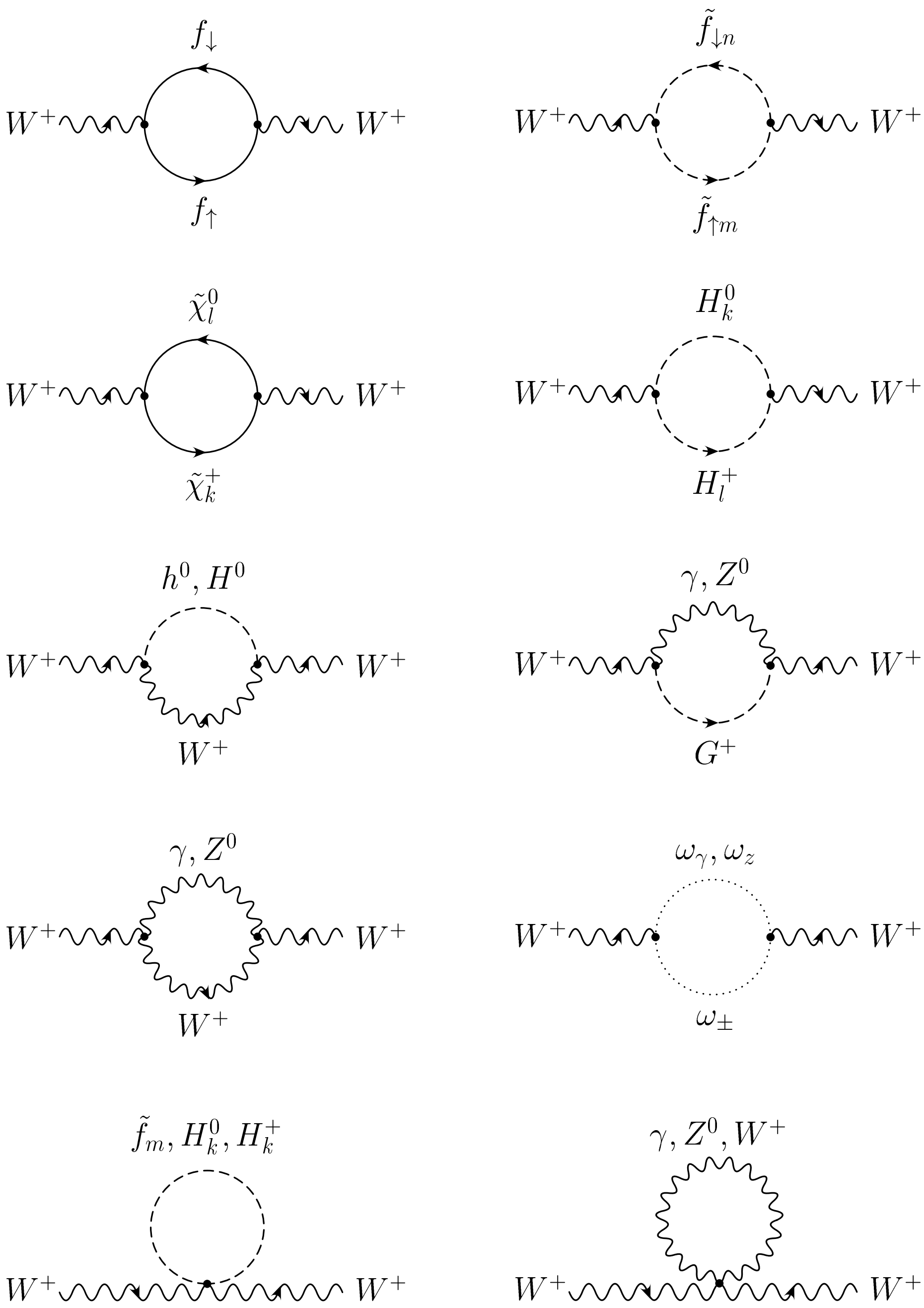


Figure B.9: W^+ self–energies

Here, $f_\uparrow$ and $f_\downarrow$ denote up– and down–type (s)fermions of all three generations, respectively.

$$\left(\frac{\delta m_W}{m_W}\right)^{\tilde{f}} = \frac{1}{(4\pi)^2}\,\frac{g^2}{4m_W^2}\sum_{\text{gen.}} N_C^f \sum_{m=1}^{2}\left(R_{m1}^{\tilde{f}}\right)^2 A_0(m_{\tilde{f}_m}^2) \quad \text{(B.61)}$$

$$\left(\frac{\delta m_W}{m_W}\right)^{\tilde{\chi}} = \frac{1}{(4\pi)^2}\,\frac{g^2}{m_W^2}\sum_{k=1}^{2}\sum_{l=1}^{4}\Big[2\,O_{lk}^L\,O_{lk}^R\,m_{\tilde{\chi}_k^+}\,m_{\tilde{\chi}_l^0}\,B_0 - \Big((O_{lk}^L)^2 + (O_{lk}^R)^2\Big)\times$$
$$\Big(m_W^2\,B_1 + m_{\tilde{\chi}_l^0}^2\,B_0 + A_0(m_{\tilde{\chi}_k^+}^2) - 2\,B_{00}\Big)\Big](m_W^2, m_{\tilde{\chi}_l^0}^2, m_{\tilde{\chi}_k^+}^2) \quad \text{(B.62)}$$

$$\left(\frac{\delta m_W}{m_W}\right)^{HH} = -\frac{1}{(4\pi)^2}\,\frac{g^2}{2m_W^2}\Bigg[\sum_{k,l=1}^{2}\Big(R_{lk}(\alpha-\beta)\Big)^2 B_{00}(m_W^2, m_{H_l^+}^2, m_{H_k^0}^2)$$
$$+\,B_{00}(m_W^2, m_{H^+}^2, m_{A^0}^2) + B_{00}(m_W^2, m_{G^+}^2, m_{G^0}^2)\Bigg] \quad \text{(B.63)}$$

$$\left(\frac{\delta m_W}{m_W}\right)^{H} = \frac{1}{(4\pi)^2}\,\frac{g^2}{8m_W^2}\left(\sum_{k=1}^{4} A_0(m_{H_k^0}^2) + 2\sum_{k=1}^{2} A_0(m_{H_k^+}^2)\right) \quad \text{(B.64)}$$

$$\left(\frac{\delta m_W}{m_W}\right)^{VS} = \frac{1}{(4\pi)^2}\,\frac{g^2}{2}\Bigg[\sum_{k=1}^{2}\Big(R_{2k}(\alpha-\beta)\Big)^2 B_0(m_W^2, m_{H_k^0}^2, m_W^2)$$
$$+\,s_W^2 B_0(m_W^2, m_W^2, \lambda^2) + s_W^2 t_W^2\,B_0(m_W^2, m_W^2, m_Z^2)\Bigg] \quad \text{(B.65)}$$

$$\left(\frac{\delta m_W}{m_W}\right)^{VV+V+\text{ghost}} = -\frac{1}{(4\pi)^2}\,\frac{g^2}{2m_W^2}\Bigg[s_W^2\Big(8B_{00} + 7m_W^2 B_0 + 2m_W^2 B_1\Big)(m_W^2, m_W^2, \lambda^2)$$
$$+c_W^2\Big(8B_{00} + 7m_W^2 B_0 + 2m_W^2 B_1\Big)(m_W^2, m_W^2, m_Z^2)$$
$$-s_W^2 A_0(\lambda^2) - c_W^2 A_0(m_Z^2) - 3A_0(m_W^2)\Bigg] \quad \text{(B.66)}$$

B.8 Z^0 self–energies

Accordingly to eq. (B.58) the mass counter–term contributions to the Z^0 boson are as follows:

$$\left(\frac{\delta m_Z}{m_Z}\right)^{ff} = \frac{1}{(4\pi)^2}\,\frac{g_Z^2}{m_Z^2}\sum_f N_C^f\Big[2\,C_L^f\,C_R^f\,m_f^2\,B_0 - \Big((C_L^f)^2 + (C_R^f)^2\Big)\times$$
$$\Big(A_0(m_f^2) + m_f^2\,B_0 - 2B_{00} + m_Z^2\,B_1\Big)\Big](m_Z^2, m_f^2, m_f^2) \quad \text{(B.67)}$$

$$\left(\frac{\delta m_Z}{m_Z}\right)^{\tilde{f}\tilde{f}} = -\frac{1}{(4\pi)^2}\,\frac{2\,g_Z^2}{m_Z^2}\sum_f N_C^f\sum_{m,n=1}^{2}(z_{mn}^{\tilde{f}})^2\,B_{00}(m_Z^2, m_{\tilde{f}_m}^2, m_{\tilde{f}_n}^2) \quad \text{(B.68)}$$

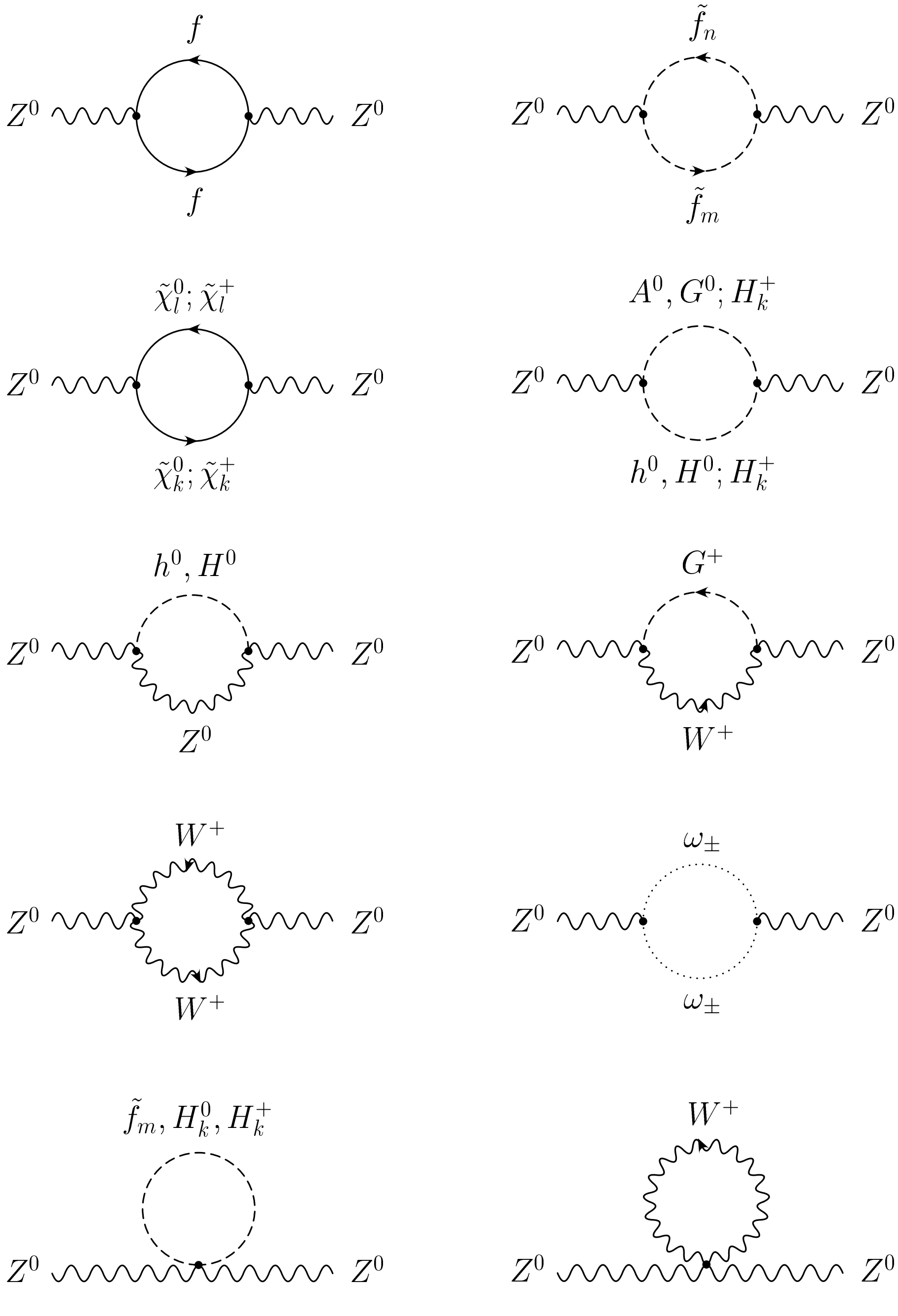


Figure B.10: Z^0 self–energies

$$\left(\frac{\delta m_Z}{m_Z}\right)^{\tilde{f}} = \frac{1}{(4\pi)^2}\frac{g_Z^2}{m_Z^2}\sum_f N_C^f \sum_{m=1}^{2}\Big((C_L^f\, R_{m1}^{\tilde{f}})^2 + (C_R^f\, R_{m2}^{\tilde{f}})^2\Big) A_0(m_{\tilde{f}_m}^2) \quad \text{(B.69)}$$

$$\left(\frac{\delta m_Z}{m_Z}\right)^{\tilde{\chi}^0} = -\frac{1}{(4\pi)^2}\frac{g_Z^2}{m_Z^2}\sum_{k,l=1}^{4}\Big(O''_{kl}\Big)^2\Big[(m_{\tilde{\chi}_k^0}+m_{\tilde{\chi}_l^0})m_{\tilde{\chi}_l^0}\, B_0 + m_Z^2\, B_1 + A_0(m_{\tilde{\chi}_k^0}^2) - 2B_{00}\Big](m_Z^2, m_{\tilde{\chi}_k^0}^2, m_{\tilde{\chi}_l^0}^2) \quad \text{(B.70)}$$

$$\left(\frac{\delta m_Z}{m_Z}\right)^{\tilde{\chi}^+} = \frac{1}{(4\pi)^2}\frac{g_Z^2}{m_Z^2}\sum_{k,l=1}^{2}\Big[2\, O'^L_{kl}\, O'^R_{kl}\, m_{\tilde{\chi}_k^+}\, m_{\tilde{\chi}_l^+}\, B_0 - \Big((O'^L_{kl})^2 + (O'^R_{kl})^2\Big) \times \Big(m_Z^2\, B_1 + m_{\tilde{\chi}_k^+}^2\, B_0 + A_0(m_{\tilde{\chi}_k^+}^2) - 2\, B_{00}\Big)\Big](m_Z^2, m_{\tilde{\chi}_k^+}^2, m_{\tilde{\chi}_l^+}^2) \quad \text{(B.71)}$$

$$\left(\frac{\delta m_Z}{m_Z}\right)^{HH} = -\frac{1}{(4\pi)^2}\frac{g_Z^2}{2m_Z^2}\Big[\sum_{k=1}^{2}\sum_{l=3}^{4}\Big(R_{k,l-2}(\beta-\alpha)\Big)^2 B_{00}(m_Z^2, m_{H_k^0}^2, m_{H_l^0}^2) + \cos^2(2\theta_W)\sum_{k=1}^{2} B_{00}(m_Z^2, m_{H_k^+}^2, m_{H_k^+}^2)\Big] \quad \text{(B.72)}$$

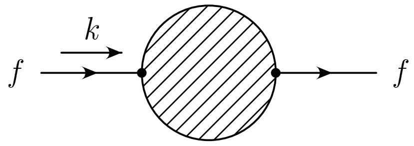
$$\left(\frac{\delta m_Z}{m_Z}\right)^{H} = \frac{1}{(4\pi)^2}\frac{g_Z^2}{8m_Z^2}\Big[\sum_{k=1}^{4} A_0(m_{H_k^0}^2) + 2\cos^2(2\theta_W)\sum_{k=1}^{2} A_0(m_{H_k^+}^2)\Big] \quad \text{(B.73)}$$

$$\left(\frac{\delta m_Z}{m_Z}\right)^{VS} = \frac{1}{(4\pi)^2}\Big(\frac{g_Z^2}{2}\sin^2(\alpha-\beta)\, B_0(m_Z^2, m_{h^0}^2, m_Z^2) + \frac{g_Z^2}{2}\cos^2(\alpha-\beta) \times B_0(m_Z^2, m_{H^0}^2, m_Z^2) + g^2\, s_W^4\, B_0(m_Z^2, m_W^2, m_{G^+}^2)\Big) \quad \text{(B.74)}$$

$$\left(\frac{\delta m_Z}{m_Z}\right)^{WW+W+\text{ghost}} = -\frac{1}{(4\pi)^2}\frac{g^2 c_W^2}{m_Z^2}\Big[4B_{00} + m_W^2 B_0 + \frac{5}{2}m_Z^2 B_0 + m_Z^2 B_1 - 2A_0(m_W^2)\Big](m_Z^2, m_W^2, m_W^2) \quad \text{(B.75)}$$

B.9 Fermion self–energies

In our notation, the fermion self–energy $\Pi(k)$ is defined through the relation



$$\mathcal{M} = i\,\bar{u}(k)\,\Pi(k)\,u(k)$$

with

$$\Pi(k) = \not{k}\, P_L\, \Pi^L(k) + \not{k}\, P_R\, \Pi^R(k) + \Pi^{SL}(k) P_L + \Pi^{SR}(k) P_R\,. \quad \text{(B.76)}$$

Thus the mass counter term for quarks and leptons is given by (see also section 3.3.2)

$$\delta m_f = \frac{1}{2}\,\text{Re}\left[m_f\left(\Pi^L(m_f) + \Pi^R(m_f)\right) + \Pi^{SL}(m_f) + \Pi^{SR}(m_f)\right]. \quad \text{(B.77)}$$

Note that for quarks and leptons (contrary to charginos), the left– and right–handed scalar parts of $\Pi(k)$ are equal, $\Pi^{SL}(k) = \Pi^{SR}(k)$. The single contributions to δm_f are as follows:

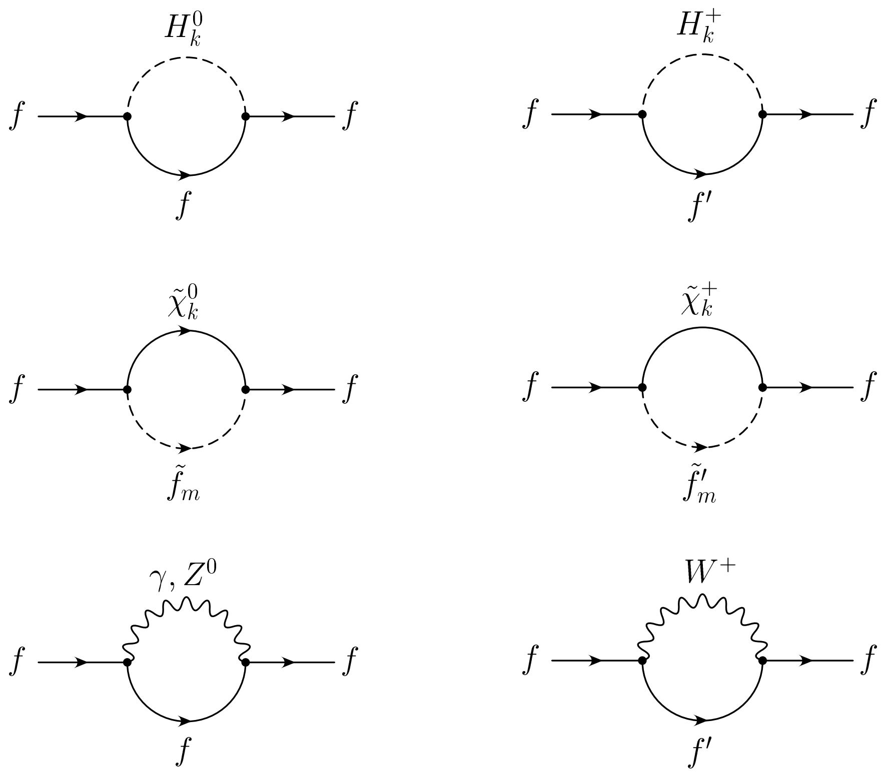


Figure B.11: Fermion self–energies

$$\left(\frac{\delta m_f}{m_f}\right)^{f H_k^0} = \frac{1}{(4\pi)^2}\left[\sum_{k=1}^{2}(s_k^f)^2(B_0 - B_1) + \sum_{k=3}^{4}(s_k^f)^2\,(B_0 + B_1)\right](m_f^2, m_f^2, m_{H_k^0}^2) \quad \text{(B.78)}$$

$$\left(\frac{\delta m_f}{m_f}\right)^{f' H_k^+} = -\frac{1}{(4\pi)^2}\sum_{k=1}^{2}\left[\frac{1}{2}\Big((y_k^f)^2 + (y_k^{f'})^2\Big)B_1 - \frac{m_{f'}}{m_f}\,y_k^f\,y_k^{f'}\,B_0\right](m_f^2, m_{f'}^2, m_{H_k^+}^2) \quad \text{(B.79)}$$

$$\left(\frac{\delta m_f}{m_f}\right)^{\tilde{f}\,\tilde{\chi}^0} = -\frac{1}{(4\pi)^2}\sum_{m=1}^{2}\sum_{k=1}^{4}\left[\frac{1}{2}\Big((a_{mk}^{\tilde{f}})^2 + (b_{mk}^{\tilde{f}})^2\Big)B_1 - \frac{m_{\tilde{\chi}_k^0}}{m_f}\,a_{mk}^{\tilde{f}} b_{mk}^{\tilde{f}}\,B_0\right] (m_f^2, m_{\tilde{\chi}_k^0}^2, m_{\tilde{f}_m}^2) \quad \text{(B.80)}$$

$$\left(\frac{\delta m_f}{m_f}\right)^{\tilde{f}'\tilde{\chi}^+} = -\frac{1}{(4\pi)^2}\sum_{m=1}^{2}\sum_{k=1}^{2}\left[\frac{1}{2}\Big((k_{mk}^{\tilde{f}'})^2 + (l_{mk}^{\tilde{f}'})^2\Big)B_1 - \frac{m_{\tilde{\chi}_k^+}}{m_f}\,k_{mk}^{\tilde{f}'}\,l_{mk}^{\tilde{f}'}\,B_0\right] (m_f^2, m_{\tilde{\chi}_k^+}^2, m_{\tilde{f}'_m}^2) \quad \text{(B.81)}$$

$$\left(\frac{\delta m_f}{m_f}\right)^{f\gamma} = -\frac{1}{(4\pi)^2}\,2(e_0 e_f)^2(B_0 - B_1)(m_f^2, \lambda^2, m_f^2) \quad \text{(B.82)}$$

$$\left(\frac{\delta m_f}{m_f}\right)^{fZ^0} = -\frac{1}{(4\pi)^2}\, g_Z^2 \Big[\Big((C_L^f)^2 + (C_R^f)^2\Big) B_1 + 4 C_L^f C_R^f\, B_0\Big](m_f^2, m_f^2, m_Z^2) \qquad \text{(B.83)}$$

$$\left(\frac{\delta m_f}{m_f}\right)^{f'W^+} = -\frac{1}{(4\pi)^2}\,\frac{g^2}{2}\, B_1(m_f^2, m_{f'}^2, m_W^2) \qquad \text{(B.84)}$$

B.10 Chargino self–energies

Since the higgsino mass parameter μ is fixed in the chargino sector, the counter term $\delta\mu$ reads [24, 25]

$$\delta\mu = \delta X_{22} = \frac{1}{2}\sum_{k,l=1}^{2} U_{k2} V_{l2} \Big(\Pi_{lk}^L\, m_{\tilde{\chi}_k^+} + \Pi_{kl}^R\, m_{\tilde{\chi}_l^+} + \Pi_{kl}^{SL} + \Pi_{lk}^{SR}\Big)\,, \qquad \text{(B.85)}$$

with the chargino self–energies $\Pi_{kl} = \Pi_{kl}(m^2_{\tilde{\chi}_l^+})$. U and V are two real 2×2 matrices which diagonalize the chargino mass matrix,

$$U\, X\, V^T = M_D = \begin{pmatrix} m_{\tilde{\chi}_1^+} & 0 \\ 0 & m_{\tilde{\chi}_2^+} \end{pmatrix}.$$

The single left– and right–handed parts of Π_{kl} can be found by comparing the coefficients accordingly to eq. (B.76).
fermion-sfermion contribution:

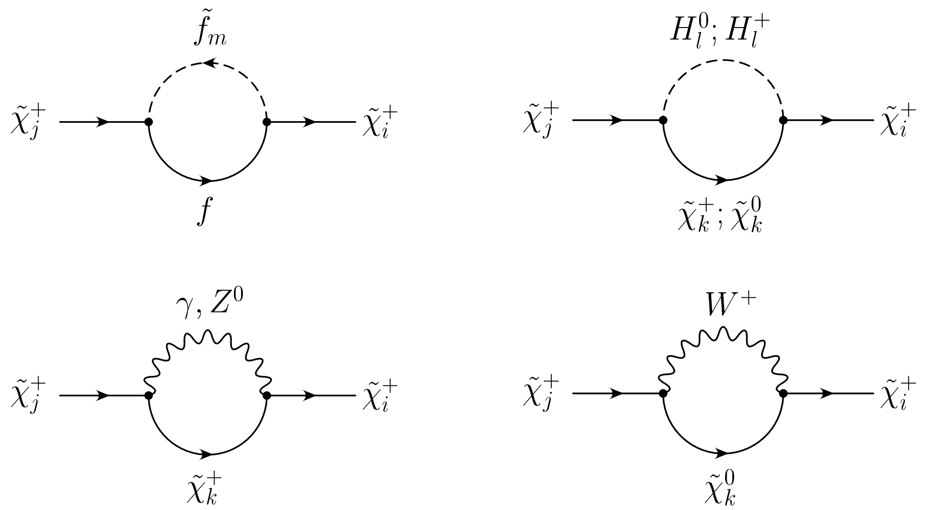


Figure B.12: Chargino self–energies

$$\Pi_{ij}^{L,f}(k^2) = -\frac{1}{(4\pi)^2}\sum_f N_C^f \sum_{m=1}^{2}\Big[k_{mi}^{\tilde{f}_\downarrow} k_{mj}^{\tilde{f}_\downarrow}\, B_1(k^2, m_{f_\uparrow}^2, m_{\tilde{f}_{\downarrow m}}^2) + l_{mi}^{\tilde{f}_\uparrow} l_{mj}^{\tilde{f}_\uparrow}\, B_1(k^2, m_{f_\downarrow}^2, m_{\tilde{f}_{\uparrow m}}^2)\Big]$$

$$\Pi_{ij}^{R,f}(k^2) = -\frac{1}{(4\pi)^2}\sum_f N_C^f \sum_{m=1}^{2}\Big[l_{mi}^{\tilde{f}_\downarrow} l_{mj}^{\tilde{f}_\downarrow}\, B_1(k^2, m_{f_\uparrow}^2, m_{\tilde{f}_{\downarrow m}}^2) + k_{mi}^{\tilde{f}_\uparrow} k_{mj}^{\tilde{f}_\uparrow}\, B_1(k^2, m_{f_\downarrow}^2, m_{\tilde{f}_{\uparrow m}}^2)\Big]$$

$$\Pi_{ij}^{SL,f}(k^2) = \frac{1}{(4\pi)^2} \sum_f N_C^f \sum_{m=1}^{2} \Big[m_{f_\uparrow} l_{mi}^{\tilde{f}_\downarrow} k_{mj}^{\tilde{f}_\downarrow} B_0(k^2, m_{f_\uparrow}^2, m_{\tilde{f}_{\downarrow m}}^2) + m_{f_\downarrow} k_{mi}^{\tilde{f}_\uparrow} l_{mj}^{\tilde{f}_\uparrow} \times B_0(k^2, m_{f_\downarrow}^2, m_{\tilde{f}_{\uparrow m}}^2) \Big]$$

$$\Pi_{ij}^{SR,f}(k^2) = \frac{1}{(4\pi)^2} \sum_f N_C^f \sum_{m=1}^{2} \Big[m_{f_\uparrow} k_{mi}^{\tilde{f}_\downarrow} l_{mj}^{\tilde{f}_\downarrow} B_0(k^2, m_{f_\uparrow}^2, m_{\tilde{f}_{\downarrow m}}^2) + m_{f_\downarrow} l_{mi}^{\tilde{f}_\uparrow} k_{mj}^{\tilde{f}_\uparrow} \times B_0(k^2, m_{f_\downarrow}^2, m_{\tilde{f}_{\uparrow m}}^2) \Big] \quad \text{(B.86)}$$

Higgs/gaugino contribution:

$$\Pi_{ij}^{H_l^0}(k) = -\frac{1}{(4\pi)^2} g^2 \sum_{k=1}^{2} \Big[\not{k} \sum_{l=1}^{4} \Big(F_{ikl}^{+} F_{jkl}^{+} P_L + F_{kil}^{+} F_{kjl}^{+} P_R \Big) B_1 - m_{\tilde{\chi}_k^+} \sum_{l=1}^{2} \Big(F_{kil}^{+} F_{jkl}^{+} P_L + F_{ikl}^{+} F_{kjl}^{+} P_R \Big) B_0 + m_{\tilde{\chi}_k^+} \sum_{l=3}^{4} \Big(F_{kil}^{+} F_{jkl}^{+} P_L + F_{ikl}^{+} F_{kjl}^{+} P_R \Big) B_0 \Big] (k^2, m_{\tilde{\chi}_k^+}^2, m_{H_l^0}^2) \quad \text{(B.87)}$$

$$\Pi_{ij}^{H_l^+}(k) = -\frac{1}{(4\pi)^2} g^2 \sum_{k=1}^{4} \sum_{l=1}^{2} \Big[\not{k} \Big(F_{ikl}^{R} F_{jkl}^{R} P_L + F_{ikl}^{L} F_{jkl}^{L} P_R \Big) B_1 - m_{\tilde{\chi}_k^0} \Big(F_{ikl}^{L} F_{jkl}^{R} P_L + F_{ikl}^{R} F_{jkl}^{L} P_R \Big) B_0 \Big] (k^2, m_{\tilde{\chi}_k^0}^2, m_{H_l^+}^2) \quad \text{(B.88)}$$

$$\Pi_{ij}^{\gamma}(k) = -\frac{1}{(4\pi)^2} 2e^2 \delta_{ij} \Big[\not{k} B_1 + 2m_{\tilde{\chi}_j^+} B_0 \Big] (k^2, m_{\tilde{\chi}_j^+}^2, \lambda^2) \quad \text{(B.89)}$$

$$\Pi_{ij}^{Z^0}(k) = -\frac{1}{(4\pi)^2} 2\, g_Z^2 \sum_{k=1}^{2} \Big[\not{k} \Big(O_{ik}^{\prime L} O_{kj}^{\prime L} P_L + O_{ik}^{\prime R} O_{kj}^{\prime R} P_R \Big) B_1 + 2m_{\tilde{\chi}_k^+} \Big(O_{ik}^{\prime R} O_{kj}^{\prime L} P_L + O_{ik}^{\prime L} O_{kj}^{\prime R} P_R \Big) B_0 \Big] (k^2, m_{\tilde{\chi}_k^+}^2, m_Z^2) \quad \text{(B.90)}$$

$$\Pi_{ij}^{W^+}(k) = -\frac{1}{(4\pi)^2} 2\, g^2 \sum_{k=1}^{4} \Big[\not{k} \Big(O_{ki}^{L} O_{kj}^{L} P_L + O_{ki}^{R} O_{kj}^{R} P_R \Big) B_1 + 2m_{\tilde{\chi}_k^0} \Big(O_{ki}^{R} O_{kj}^{L} P_L + O_{ki}^{L} O_{kj}^{R} P_R \Big) B_0 \Big] (k^2, m_{\tilde{\chi}_k^0}^2, m_W^2) \quad \text{(B.91)}$$

Appendix C

Vertex corrections

In the following sections we give the explicit formulae of the electroweak contributions to the vertex corrections of the decays $\{h^0, H^0, A^0\} \to \tilde{f}_i \bar{\tilde{f}}_j$ and $H^+ \to \tilde{t}_i \bar{\tilde{b}}_j$. For the SUSY–QCD contributions we refer to [34].

C.1 $h^0_k \tilde{f}_i \bar{\tilde{f}}_j$ vertex

The vertex corrections to $h^0_k \to \tilde{f}_i \bar{\tilde{f}}_j$ which are depicted in Fig. C.1 are given as follows:

$$\begin{aligned}
\delta G_{ijk}^{\tilde{f}(v)} &= \delta G_{ijk}^{\tilde{f}(v,H\tilde{f}\tilde{f})} + \delta G_{ijk}^{\tilde{f}(v,\tilde{f}HH)} + \delta G_{ijk}^{\tilde{f}(v,\tilde{\chi}ff)} + \delta G_{ijk}^{\tilde{f}(v,f\tilde{\chi}\tilde{\chi})} + \delta G_{ijk}^{\tilde{f}(v,VSS)} \\
&\quad + \delta G_{ijk}^{\tilde{f}(v,VVS)} + \delta G_{ijk}^{\tilde{f}(v,SS)} + \delta G_{ijk}^{\tilde{f}(v,VV)} + \delta G_{ijk}^{\tilde{f}(v,hH\mathrm{mix})} + \delta G_{ijk}^{\tilde{f}(v,\tilde{f}\mathrm{mix})} \qquad \text{(C.1)}
\end{aligned}$$

The single contributions correspond to the diagrams with three scalar particles $\left(\delta G_{ijk}^{\tilde{f}(v,H\tilde{f}\tilde{f})}\right.$ and $\left.\delta G_{ijk}^{\tilde{f}(v,HH\tilde{f})}\right)$, three fermions $\left(\delta G_{ijk}^{\tilde{f}(v,\tilde{\chi}ff)} \text{ and } \delta G_{ijk}^{\tilde{f}(v,f\tilde{\chi}\tilde{\chi})}\right)$, three particles with one or two vector bosons $\left(\delta G_{ijk}^{\tilde{f}(v,VSS)} \text{ and } \delta G_{ijk}^{\tilde{f}(v,VVS)}\right)$ and two scalar or two vector particles $\left(\delta G_{ijk}^{\tilde{f}(v,SS)}\right.$ and $\left.\delta G_{ijk}^{\tilde{f}(v,VV)}\right)$ in the loop. The vertex corrections due to the mixing of the outer particles, i.e. the Higgs and sfermion mixing terms $\delta G_{ijk}^{\tilde{f}(v,hH\mathrm{mix})}$ and $\delta G_{ijk}^{\tilde{f}(v,\tilde{f}\mathrm{mix})}$ will be combined with the counter terms of the Higgs and sfermion mixing angles, $\delta\alpha$ and $\delta\theta_{\tilde{f}}$, see .

The vertex corrections from the exchange of one Higgs and two sfermions are

$$\begin{aligned}
\delta G_{ijk}^{\tilde{f}(v,H\tilde{f}\tilde{f})} &= -\frac{1}{(4\pi)^2} \sum_{m,n=1}^{2} \sum_{l=1}^{4} G_{mnk}^{\tilde{f}} G_{iml}^{\tilde{f}} G_{njl}^{\tilde{f}} \, C_0\Big(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{H_l^0}^2, m_{\tilde{f}_m}^2, m_{\tilde{f}_n}^2\Big) \\
&\quad -\frac{1}{(4\pi)^2} \sum_{m,n=1}^{2} \sum_{l=1}^{2} G_{mnk}^{\tilde{f}'} G_{iml}^{\tilde{f}\tilde{f}'} G_{jnl}^{\tilde{f}\tilde{f}'} \, C_0\Big(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{H_l^+}^2, m_{\tilde{f}'_m}^2, m_{\tilde{f}'_n}^2\Big)
\end{aligned}$$

(C.2)

with the standard two–point function C_0 [44] for which we follow the conventions of [15]. The graph with two Higgs particles and one sfermion in the loop leads to

$$\begin{aligned}
\delta G_{ijk}^{\tilde{f}(v,\tilde{f}HH)} &= -\frac{1}{(4\pi)^2}\frac{g_Z\, m_Z}{4}\sum_{m,n=1}^{2}\sum_{l=1}^{2}(2+\delta_{km}\delta_{mn})!\left(\cos 2\alpha\,\tilde{A}_{mn}^{(k)} - 2\sin 2\alpha\,\tilde{B}_{mn}^{(k)}\right)\times \\
&\qquad G_{ilm}^{\tilde{f}}\,G_{ljn}^{\tilde{f}}\,C_0\left(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}_l}^2, m_{h_m^0}^2, m_{h_n^0}^2\right) \\
&+\frac{1}{(4\pi)^2}\frac{g_Z\, m_Z}{2}\sum_{m,n=3}^{4}\sum_{l=1}^{2}\sin[\alpha+\beta-\frac{\pi}{2}(k-1)]\,\tilde{C}_{m-2,n-2}\times \\
&\qquad G_{ilm}^{\tilde{f}}\,G_{ljn}^{\tilde{f}}\,C_0\left(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}_l}^2, m_{H_m^0}^2, m_{H_n^0}^2\right) \\
&-\frac{1}{(4\pi)^2}\sum_{m,n=1}^{2}\sum_{l=1}^{2}\Bigg[(-1)^{mn}\,\frac{g\, m_W}{2}(1-\delta_{m2}\delta_{n2})(1+\delta_{mn})\tilde{A}_{mn}^{\prime(k)} \\
&\qquad -\frac{g_Z\, m_Z}{2}\sin[\alpha+\beta-\frac{\pi}{2}(k-1)]\,\tilde{C}_{mn}\Bigg]\times \\
&\qquad G_{ilm}^{\tilde{f}\tilde{f}'}\,G_{jln}^{\tilde{f}\tilde{f}'}\,C_0\left(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}'_l}^2, m_{H_m^+}^2, m_{H_n^+}^2\right) \qquad (C.3)
\end{aligned}$$

with

$$\begin{aligned}
\tilde{A}_{mn}^{(k)} &= \begin{pmatrix} -\sin[\alpha+\beta-\frac{\pi}{2}(k-1)] & \cos[\alpha+\beta-\frac{\pi}{2}(k-1)] \\ \cos[\alpha+\beta-\frac{\pi}{2}(k-1)] & \sin[\alpha+\beta-\frac{\pi}{2}(k-1)] \end{pmatrix}, \\
\tilde{A}_{mn}^{\prime(k)} &= \tilde{A}_{mn}^{(k)}(\beta\to-\beta) \\
\tilde{B}_{mn}^{(k)} &= \begin{pmatrix} (k-1)\sin(\alpha+\beta) & \sin[\alpha+\beta-\frac{\pi}{2}(k-1)] \\ \sin[\alpha+\beta-\frac{\pi}{2}(k-1)] & (k-2)\cos(\alpha+\beta) \end{pmatrix}, \\
\tilde{C}_{mn} &= \begin{pmatrix} \cos 2\beta & \sin 2\beta \\ \sin 2\beta & -\cos 2\beta \end{pmatrix}. \qquad (C.4)
\end{aligned}$$

For the gaugino exchange contributions we get

$$\begin{aligned}
\delta G_{ijk}^{\tilde{f}(v,\tilde{\chi}ff)} &= \frac{1}{(4\pi)^2}\sum_{l=1}^{4}F\left(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{\chi}_l^0}, m_f, m_f; s_k^f, s_k^f, b_{il}^{\tilde{f}}, a_{il}^{\tilde{f}}, a_{jl}^{\tilde{f}}, b_{jl}^{\tilde{f}}\right) \\
&+\frac{1}{(4\pi)^2}\sum_{l=1}^{2}F\left(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{\chi}_l^+}, m_{f'}, m_{f'}; s_k^{f'}, -s_k^{f'}, k_{il}^{\tilde{f}}, l_{il}^{\tilde{f}}, l_{jl}^{\tilde{f}}, k_{jl}^{\tilde{f}}\right), \qquad (C.5)
\end{aligned}$$

$$\begin{aligned}
\delta G_{ijk}^{\tilde{f}(v,f\tilde{\chi}\tilde{\chi})} &= \frac{1}{(4\pi)^2}\sum_{l,m=1}^{4}F\left(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_f, m_{\tilde{\chi}_m^0}, m_{\tilde{\chi}_l^0}; -gF_{lmk}^0, -gF_{lmk}^0, b_{im}^{\tilde{f}}, a_{im}^{\tilde{f}}, a_{jl}^{\tilde{f}}, b_{jl}^{\tilde{f}}\right) \\
&+\frac{1}{(4\pi)^2}\sum_{l,m=1}^{2}F\left(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{f'}, m_{\tilde{\chi}_m^+}, m_{\tilde{\chi}_l^+}; -g\widetilde{F}_{mlk}^+, -g\widetilde{F}_{lmk}^+, k_{im}^{\tilde{f}}, l_{im}^{\tilde{f}}, l_{jl}^{\tilde{f}}, k_{jl}^{\tilde{f}}\right),
\end{aligned}$$

(C.6)

where $F(\ldots)$ shortly stands for

$$F\Big(m_1^2, m_0^2, m_2^2, M_0, M_1, M_2; g_0^R, g_0^L, g_1^R, g_1^L, g_2^R, g_2^L\Big) = (h_1 M_1 + h_2 M_2) B_0(m_0^2, M_1^2, M_2^2)$$
$$+ (h_0 M_0 + h_1 M_1) B_0(m_1^2, M_0^2, M_1^2) + (h_0 M_0 + h_2 M_2) B_0(m_2^2, M_0^2, M_2^2)$$
$$+ \Big[2\big(g_0^R g_1^R g_2^R + g_0^L g_1^L g_2^L\big) M_0 M_1 M_2 + h_0 M_0 \big(M_1^2 + M_2^2 - m_0^2\big) + h_1 M_1 \big(M_0^2 + M_2^2 - m_2^2\big)$$
$$+ h_2 M_2 \big(M_0^2 + M_1^2 - m_1^2\big)\Big] C_0(m_1^2, m_0^2, m_2^2, M_0^2, M_1^2, M_2^2) \quad \text{(C.7)}$$

with the abbreviations $h_0 = \big(g_0^L g_1^R g_2^R + g_0^R g_1^L g_2^L\big)$, $h_1 = \big(g_0^L g_1^L g_2^R + g_0^R g_1^R g_2^L\big)$ and $h_2 = \big(g_0^R g_1^L g_2^R + g_0^L g_1^R g_2^L\big)$. For up–type sfermions $\widetilde{F}^+_{lmk} = F^+_{lmk}$ and for down–type sfermions the chargino indices are interchanged, $\widetilde{F}^+_{lmk} = F^+_{mlk}$.
We split the irreducible vertex graphs with one vector particle in the loop into the single contributions of the photon, the Z–boson and the W–boson,

$$\delta G_{ijk}^{\tilde{f}(v,VSS)} = \delta G_{ijk}^{\tilde{f}(v,\gamma SS)} + \delta G_{ijk}^{\tilde{f}(v,ZSS)} + \delta G_{ijk}^{\tilde{f}(v,WSS)} . \quad \text{(C.8)}$$

In order to regularize the infrared divergences we introduce a photon mass λ. Thus we have

$$\delta G_{ijk}^{\tilde{f}(v,\gamma SS)} = \frac{1}{(4\pi)^2} (e_0 e_f)^2 G_{ijk}^{\tilde{f}}\, V\Big(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, \lambda^2, m_{\tilde{f}_i}^2, m_{\tilde{f}_j}^2\Big) , \quad \text{(C.9)}$$

$$\delta G_{ijk}^{\tilde{f}(v,ZSS)} = \frac{1}{(4\pi)^2} g_Z^2 \sum_{m,n=1}^{2} G_{mnk}^{\tilde{f}}\, z_{im}^{\tilde{f}}\, z_{nj}^{\tilde{f}}\, V\Big(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_Z^2, m_{\tilde{f}_m}^2, m_{\tilde{f}_n}^2\Big)$$
$$+ \frac{i}{(4\pi)^2} \frac{g_Z^2}{2} \sum_{l=3}^{4} \sum_{m=1}^{2} G_{mjl}^{\tilde{f}}\, z_{im}^{\tilde{f}}\, R_{l-2,k}(\alpha-\beta)\, V\Big(m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}_i}^2, m_Z^2, m_{H_l^0}^2, m_{\tilde{f}_m}^2\Big)$$
$$- \frac{i}{(4\pi)^2} \frac{g_Z^2}{2} \sum_{l=3}^{4} \sum_{m=1}^{2} G_{iml}^{\tilde{f}}\, z_{mj}^{\tilde{f}}\, R_{l-2,k}(\alpha-\beta)\, V\Big(m_{\tilde{f}_j}^2, m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_Z^2, m_{\tilde{f}_m}^2, m_{H_l^0}^2\Big) ,$$
(C.10)

$$\delta G_{ijk}^{\tilde{f}(v,WSS)} = \frac{1}{(4\pi)^2} \frac{g^2}{2} \sum_{m,n=1}^{2} G_{mnk}^{\tilde{f}'} R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} R_{m1}^{\tilde{f}'} R_{n1}^{\tilde{f}'}\, V\Big(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_W^2, m_{\tilde{f}'_m}^2, m_{\tilde{f}'_n}^2\Big)$$
$$- \frac{I_f^{3L}}{(4\pi)^2} \frac{g^2}{\sqrt{2}} \sum_{m,l=1}^{2} G_{jml}^{\tilde{f}\tilde{f}'} R_{i1}^{\tilde{f}} R_{m1}^{\tilde{f}'}\, R_{l,k}(\alpha-\beta)\, V\Big(m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}_i}^2, m_W^2, m_{H_l^+}^2, m_{\tilde{f}'_m}^2\Big)$$
$$- \frac{I_f^{3L}}{(4\pi)^2} \frac{g^2}{\sqrt{2}} \sum_{m,l=1}^{2} G_{iml}^{\tilde{f}\tilde{f}'} R_{m1}^{\tilde{f}'} R_{j1}^{\tilde{f}}\, R_{l,k}(\alpha-\beta)\, V\Big(m_{\tilde{f}_j}^2, m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_W^2, m_{\tilde{f}'_m}^2, m_{H_l^+}^2\Big) ,$$
(C.11)

where we have used the vector vertex function

$$V\Big(m_1^2, m_0^2, m_2^2, M_0^2, M_1^2, M_2^2\Big) = -B_0\Big(m_0^2, M_1^2, M_2^2\Big) + B_0\Big(m_1^2, M_0^2, M_1^2\Big) + B_0\Big(m_2^2, M_0^2, M_2^2\Big)$$

$$+ \ \left(-2m_0^2+m_1^2+m_2^2-M_0^2+M_1^2+M_2^2\right) C_0\Big(m_1^2, m_0^2, m_2^2, M_0^2, M_1^2, M_2^2\Big) \qquad \text{(C.12)}$$

and the rotation matrix R_{kl},

$$R_{kl}(\phi) \ \equiv \ \begin{pmatrix} \cos\phi & \sin\phi \\ -\sin\phi & \cos\phi \end{pmatrix}_{kl} . \qquad \text{(C.13)}$$

The Z^0–sfermion–sfermion couplings $z_{ij}^{\tilde{f}}$ can be found in Appendix A.7.
The vertex corrections coming from loops with two vector bosons and one sfermion are given by

$$\delta G_{ijk}^{\tilde{f}(v,VVS)} \ = \ \delta G_{ijk}^{\tilde{f}(v,ZZ\tilde{f})} + \delta G_{ijk}^{\tilde{f}(v,WW\tilde{f}')} \qquad \text{(C.14)}$$

with

$$\begin{aligned}\delta G_{ijk}^{\tilde{f}(v,ZZ\tilde{f})} \ &= \ -\frac{1}{(4\pi)^2}\frac{g_Z^3 m_Z}{2} R_{2k}(\alpha-\beta)\sum_{m=1}^{2}\Big[4B_0\Big(m_{h_k^0}^2, m_Z^2, m_Z^2\Big)-B_0\Big(m_{\tilde{f}_i}^2, m_{\tilde{f}_m}^2, m_Z^2\Big) \\ &- B_0\Big(m_{\tilde{f}_j}^2, m_{\tilde{f}_m}^2, m_Z^2\Big) - \Big(m_{h_k^0}^2-2m_{\tilde{f}_i}^2-2m_{\tilde{f}_j}^2-4m_{\tilde{f}_m}^2+2m_Z^2\Big)\times \\ &C_0\Big(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}_m}^2, m_Z^2, m_Z^2\Big)\Big] z_{im}^{\tilde{f}}\, z_{mj}^{\tilde{f}} \qquad \text{(C.15)}\end{aligned}$$

and

$$\begin{aligned}\delta G_{ijk}^{\tilde{f}(v,WW\tilde{f}')} \ &= \ -\frac{1}{(4\pi)^2}\frac{g^3 m_W}{4} R_{2k}(\alpha-\beta)\sum_{m=1}^{2}\Big[4B_0\Big(m_{h_k^0}^2, m_W^2, m_W^2\Big)-B_0\Big(m_{\tilde{f}_i}^2, m_{\tilde{f}'_m}^2, m_W^2\Big) \\ &- B_0\Big(m_{\tilde{f}_j}^2, m_{\tilde{f}'_m}^2, m_W^2\Big) - \Big(m_{h_k^0}^2-2m_{\tilde{f}_i}^2-2m_{\tilde{f}_j}^2-4m_{\tilde{f}'_m}^2+2m_W^2\Big)\times \\ &C_0\Big(m_{\tilde{f}_i}^2, m_{h_k^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}'_m}^2, m_W^2, m_W^2\Big)\Big] R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}}\Big(R_{m1}^{\tilde{f}'}\Big)^2 . \qquad \text{(C.16)}\end{aligned}$$

We split the irreducible vertex graphs with two scalar particles into the contributions from two Higgs bosons, two sfermions and the corrections stemming from Higgs–sfermion loops, i.e.

$$\begin{aligned}\delta G_{ijk}^{\tilde{f}(v,SS)} \ &= \ \delta G_{ijk}^{\tilde{f}(v,HH)} + \delta G_{ijk}^{\tilde{f}(v,\tilde{f}\tilde{f})} + \delta G_{ijk}^{\tilde{f}(v,\tilde{f}'\tilde{f}')} + \delta G_{ijk}^{\tilde{f}(v,\hat{\tilde{f}}\hat{\tilde{f}})} + \delta G_{ijk}^{\tilde{f}(v,\hat{\tilde{f}}'\hat{\tilde{f}}')} \\ &+\delta G_{ijk}^{\tilde{f}(v,\tilde{F}\tilde{F})} + \delta G_{ijk}^{\tilde{f}(v,\tilde{F}'\tilde{F}')} + \delta G_{ijk}^{\tilde{f}(v,\hat{\tilde{F}}\hat{\tilde{F}})} + \delta G_{ijk}^{\tilde{f}(v,\hat{\tilde{F}}'\hat{\tilde{F}}')} + \delta G_{ijk}^{\tilde{f}(v,H\tilde{f})} \qquad \text{(C.17)}\end{aligned}$$

with

$$\begin{aligned}\delta G_{ijk}^{\tilde{f}(v,HH)} \ &= \ -\frac{1}{(4\pi)^2}\frac{g_Z\, m_Z}{8}\sum_{m,n=1}^{2}(2+\delta_{km}\delta_{mn})!\,\Big(\cos 2\alpha\,\tilde{A}_{mn}^{(k)} - 2\sin 2\alpha\,\tilde{B}_{mn}^{(k)}\Big)\times \\ &\Big(h_f^2 c_{mn}^{\tilde{f}}\delta_{ij} + g^2(c_{mn}^{\tilde{b}} - c_{mn}^{\tilde{t}})e_{ij}^{\tilde{f}}\Big) B_0\Big(m_{h_k^0}^2, m_{h_m^0}^2, m_{h_n^0}^2\Big)\end{aligned}$$

$$
\begin{aligned}
&+\frac{1}{(4\pi)^2}\frac{g_Z\, m_Z}{4}\sum_{m,n=3}^{4}\sin[\alpha+\beta-\frac{\pi}{2}(k-1)]\,\tilde{C}_{m-2,n-2}\times \\
&\qquad\qquad\Big(h_f^2 c^{\tilde{f}}_{mn}\delta_{ij}+g^2(c^{\tilde{b}}_{mn}-c^{\tilde{t}}_{mn})e^{\tilde{f}}_{ij}\Big)B_0\Big(m^2_{h^0_k},m^2_{H^0_m},m^2_{H^0_n}\Big) \\
&-\frac{1}{(4\pi)^2}\sum_{m,n=1}^{2}\Big(h_{f'}^2 d^{\tilde{f}'}_{mn}R^{\tilde{f}}_{i1}R^{\tilde{f}}_{j1}+h_f^2 d^{\tilde{f}}_{mn}R^{\tilde{f}}_{i2}R^{\tilde{f}}_{j2}+g^2(d^{\tilde{b}}_{mn}-d^{\tilde{t}}_{mn})f^{\tilde{f}}_{ij}\Big)\times \\
&\qquad\Big[(-1)^{mn}\,\frac{g\,m_W}{2}(1-\delta_{m2}\delta_{n2})(1+\delta_{mn})\tilde{A}'^{(k)}_{mn} \\
&\qquad\quad-\frac{g_Z\,m_Z}{2}\sin[\alpha+\beta-\frac{\pi}{2}(k-1)]\,\tilde{C}_{mn}\Big]B_0\Big(m^2_{h^0_k},m^2_{H^+_m},m^2_{H^+_n}\Big)\,,
\end{aligned}
\tag{C.18}
$$

$$
\begin{aligned}
\delta G^{\tilde{f}(v,\tilde{f}\tilde{f})}_{ijk} &= -\frac{h_f^2}{(4\pi)^2}\sum_{m,n=1}^{2}G^{\tilde{f}}_{nmk}\Big[R^{\tilde{f}}_{ijmn}+R^{\tilde{f}}_{mnij}+N^f_C\Big(R^{\tilde{f}}_{inmj}+R^{\tilde{f}}_{mjin}\Big)\Big]B_0\Big(m^2_{h^0_k},m^2_{\tilde{f}_m},m^2_{\tilde{f}_n}\Big) \\
&-\frac{g_Z^2}{(4\pi)^2}\sum_{m,n=1}^{2}G^{\tilde{f}}_{nmk}\Big\{\Big[\Big(\frac{1}{4}-(2I^{3L}_f-e_f)e_f s^2_W\Big)R^{\tilde{f}_L}_{ijmn}+e_f^2 s_W^2 R^{\tilde{f}_R}_{ijmn}\Big](N^f_C+1) \\
&\qquad+(I^{3L}_f-e_f)e_f s^2_W\Big[N^f_C\Big(R^{\tilde{f}}_{ijmn}+R^{\tilde{f}}_{mnij}\Big)+R^{\tilde{f}}_{inmj}+R^{\tilde{f}}_{mjin}\Big]\Big\} \\
&\qquad\times B_0\Big(m^2_{h^0_k},m^2_{\tilde{f}_m},m^2_{\tilde{f}_n}\Big)\,,
\end{aligned}
\tag{C.19}
$$

$$
\begin{aligned}
\delta G^{\tilde{f}(v,\tilde{f}'\tilde{f}')}_{ijk} &= -\frac{1}{(4\pi)^2}\sum_{m,n=1}^{2}G^{\tilde{f}'}_{nmk}\Big\{h_f^2\,R^{\tilde{f}'\tilde{f}_D}_{mnij}+h^2_{f'}\,R^{\tilde{f}\tilde{f}'_D}_{ijmn}+\frac{g^2}{4}\Big[N^f_C\Big(\Big(t^2_W Y^f_L Y^{f'}_L-1\Big)R^{\tilde{f}\tilde{f}'_L}_{ijmn} \\
&\qquad+t^2_W Y^f_R Y^{f'}_R\,R^{\tilde{f}\tilde{f}'_R}_{ijmn}-t^2_W Y^f_L Y^{f'}_R R^{\tilde{f}\tilde{f}'_D}_{ijmn}-t^2_W Y^{f'}_L Y^f_R R^{\tilde{f}'\tilde{f}_D}_{mnij}\Big) \\
&\qquad+2R^{\tilde{f}\tilde{f}'_L}_{ijmn}\Big]\Big\}B_0\Big(m^2_{h^0_k},m^2_{\tilde{f}'_m},m^2_{\tilde{f}'_n}\Big)\,,
\end{aligned}
\tag{C.20}
$$

$$
\begin{aligned}
\delta G^{\tilde{f}(v,\hat{\tilde{f}}\hat{\tilde{f}})}_{ijk} &= -\frac{N^{\hat{f}}_C}{(4\pi)^2}\sum_{m,n=1}^{2}G^{\hat{\tilde{f}}}_{nmk}\Big\{h_f h_{\hat{f}}\Big(R^{\tilde{f}\hat{\tilde{f}}_F}_{jimn}+R^{\tilde{f}\hat{\tilde{f}}_F}_{ijnm}\Big)+\frac{g^2}{4}\Big[\Big(t^2_W Y^f_L Y^{\hat{f}}_L+1\Big)R^{\tilde{f}\hat{\tilde{f}}_L}_{ijmn} \\
&\qquad+t^2_W Y^f_R Y^{\hat{f}}_R\,R^{\tilde{f}\hat{\tilde{f}}_R}_{ijmn}-t^2_W Y^f_L Y^{\hat{f}}_R R^{\tilde{f}\hat{\tilde{f}}_D}_{ijmn}-t^2_W Y^{\hat{f}}_L Y^f_R R^{\hat{\tilde{f}}\tilde{f}_D}_{mnij}\Big]\Big\} \\
&\qquad\times B_0\Big(m^2_{h^0_k},m^2_{\hat{\tilde{f}}_m},m^2_{\hat{\tilde{f}}_n}\Big)\,,
\end{aligned}
\tag{C.21}
$$

and

$$
\delta G^{\tilde{f}(v,\hat{\tilde{f}}'\hat{\tilde{f}}')}_{ijk} = -\frac{N^{\hat{f}}_C}{(4\pi)^2}\frac{g^2}{4}\sum_{m,n=1}^{2}G^{\hat{\tilde{f}}'}_{nmk}\Big\{\Big(t^2_W Y^f_L Y^{\hat{f}'}_L-1\Big)R^{\tilde{f}\hat{\tilde{f}}'_L}_{ijmn}+t^2_W Y^f_R Y^{\hat{f}'}_R\,R^{\tilde{f}\hat{\tilde{f}}'_R}_{ijmn}
$$

$$-t_W^2 Y_L^f Y_R^{\hat{f}'} R_{ijmn}^{\tilde{f}\tilde{f}\hat{\tilde{f}}'_D} - t_W^2 Y_L^{\hat{f}'} Y_R^f R_{mnij}^{\hat{\tilde{f}}'\tilde{f}_D} \Big\} B_0\Big(m_{h_k^0}^2, m_{\hat{\tilde{f}}'_m}^2, m_{\hat{\tilde{f}}'_n}^2\Big)\,. \tag{C.22}$$

For various products of sfermion rotation matrices we have introduced the short forms

$$\begin{aligned} R_{ijkl}^{\tilde{f}_L} &= R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} R_{k1}^{\tilde{f}} R_{l1}^{\tilde{f}}\,, & R_{ijkl}^{\tilde{f}} &= R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} R_{k2}^{\tilde{f}} R_{l2}^{\tilde{f}}\,, \\ R_{ijkl}^{\tilde{f}_R} &= R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}} R_{k2}^{\tilde{f}} R_{l2}^{\tilde{f}}\,, & R_{ijkl}^{\tilde{f}\hat{\tilde{f}}_F} &= R_{i1}^{\tilde{f}} R_{j2}^{\tilde{f}} R_{k1}^{\hat{\tilde{f}}} R_{l2}^{\hat{\tilde{f}}}\,. \end{aligned} \tag{C.23}$$

Note that the contributions from eqs. (C.21, C.22) originate from the mixing of 2 squarks and 2 sleptons, where $\hat{f}$ denotes the 'family partner' of the fermion f with the same isospin and from the same generation, i.e. $\hat{t} = \nu_\tau$ or $\hat{\tilde{\tau}}_i = \tilde{b}_i$.
The contributions due to the exchange of sfermions from the other two generations, $\tilde{F}_m$, are given by

$$\begin{aligned} \delta G_{ijk}^{\tilde{f}(v,\tilde{F}\tilde{F})} &= \delta G_{ijk}^{\tilde{f}(v,\hat{\tilde{f}}\hat{\tilde{f}})}(\hat{f} \to F)\,, & \delta G_{ijk}^{\tilde{f}(v,\hat{\tilde{F}}\hat{\tilde{F}})} &= \delta G_{ijk}^{\tilde{f}(v,\hat{\tilde{f}}\hat{\tilde{f}})}(\hat{f} \to \hat{F})\,, \\ \delta G_{ijk}^{\tilde{f}(v,\tilde{F}'\tilde{F}')} &= \delta G_{ijk}^{\tilde{f}(v,\hat{\tilde{f}}'\hat{\tilde{f}}')}(\hat{f}' \to F')\,, & \delta G_{ijk}^{\tilde{f}(v,\hat{\tilde{F}}'\hat{\tilde{F}}')} &= \delta G_{ijk}^{\tilde{f}(v,\hat{\tilde{f}}'\hat{\tilde{f}}')}(\hat{f}' \to \hat{F}')\,, \end{aligned} \tag{C.24}$$

where the sub–/superscript $\tilde{F}$ denotes values belonging to first and second generation scalar fermions with same isospin as $\tilde{f}$ (e.g. $\tilde{F}_1 = \{\tilde{u}_1, \tilde{c}_1\}$ for the stop case, ...), $\tilde{F}'$ sfermions with different isospin etc.
The diagrams with one Higgs boson and one sfermion in the loop lead to

$$\begin{aligned} \delta G_{ijk}^{\tilde{f}(v,H\tilde{f})} = &-\frac{1}{(4\pi)^2} \sum_{l,m=1}^{2} G_{iml}^{\tilde{f}} \Big(h_f^2\, c_{kl}^{\tilde{f}}\, \delta_{mj} + g^2 \Big(c_{kl}^{\tilde{b}} - c_{kl}^{\tilde{t}}\Big) e_{mj}^{\tilde{f}}\Big) B_0\Big(m_{\tilde{f}_i}^2, m_{h_l^0}^2, m_{\tilde{f}_m}^2\Big) \\ &+\frac{1}{(4\pi)^2} \frac{1}{\sqrt{2}} \sum_{l,m=1}^{2} G_{iml}^{\tilde{f}\tilde{f}'} \Big(\Big(h_{f_\uparrow}^2\, c_{kl}^{\tilde{t},0+} + h_{f_\downarrow}^2\, c_{k,l+2}^{\tilde{b},0+}\Big) R_{j1}^{\tilde{f}} R_{m1}^{\tilde{f}'} + h_{f_\uparrow} h_{f_\downarrow} \Big(c_{kl}^{\tilde{t}\tilde{b},0+} + c_{k,l+2}^{\tilde{b}\tilde{t},0+}\Big) \\ &\qquad \times R_{j2}^{\tilde{f}} R_{m2}^{\tilde{f}'} - \frac{g^2}{2}\Big(c_{kl}^{\tilde{t},0+} + c_{k,l+2}^{\tilde{b},0+}\Big) R_{j1}^{\tilde{f}} R_{m1}^{\tilde{f}'}\Big) B_0\Big(m_{\tilde{f}_i}^2, m_{H_l^+}^2, m_{\tilde{f}'_m}^2\Big) \\ &+ i \leftrightarrow j\,, \end{aligned} \tag{C.25}$$

with $h_{f_\uparrow} = \{h_t, 0\}$ and $h_{f_\downarrow} = \{h_b, h_\tau\}$ for the decay into {squarks, sleptons}, respectively. $i \leftrightarrow j$ stands for both terms before with i and j interchanged. The Higgs–sfermion coupling matrices $c_{kl}^{\tilde{f}}$ and $e_{ij}^{\tilde{f}}$ can be found in Appendix A.6.
Finally, for the vertex graphs with two vector bosons we obtain

$$\begin{aligned} \delta G_{ijk}^{\tilde{f}(v,VV)} = &\frac{4}{(4\pi)^2}\, g_Z^3 m_Z\, R_{2k}(\alpha-\beta) \Big[\Big(C_L^f\Big)^2 R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} + \Big(C_R^f\Big)^2 R_{i2}^{\tilde{f}} R_{j2}^{\tilde{f}}\Big] B_0\Big(m_{h_k^0}^2, m_Z^2, m_Z^2\Big) \\ &+\frac{2}{(4\pi)^2}\, g^3 m_W\, R_{2k}(\alpha-\beta) R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} B_0\Big(m_{h_k^0}^2, m_W^2, m_W^2\Big)\,. \end{aligned} \tag{C.26}$$

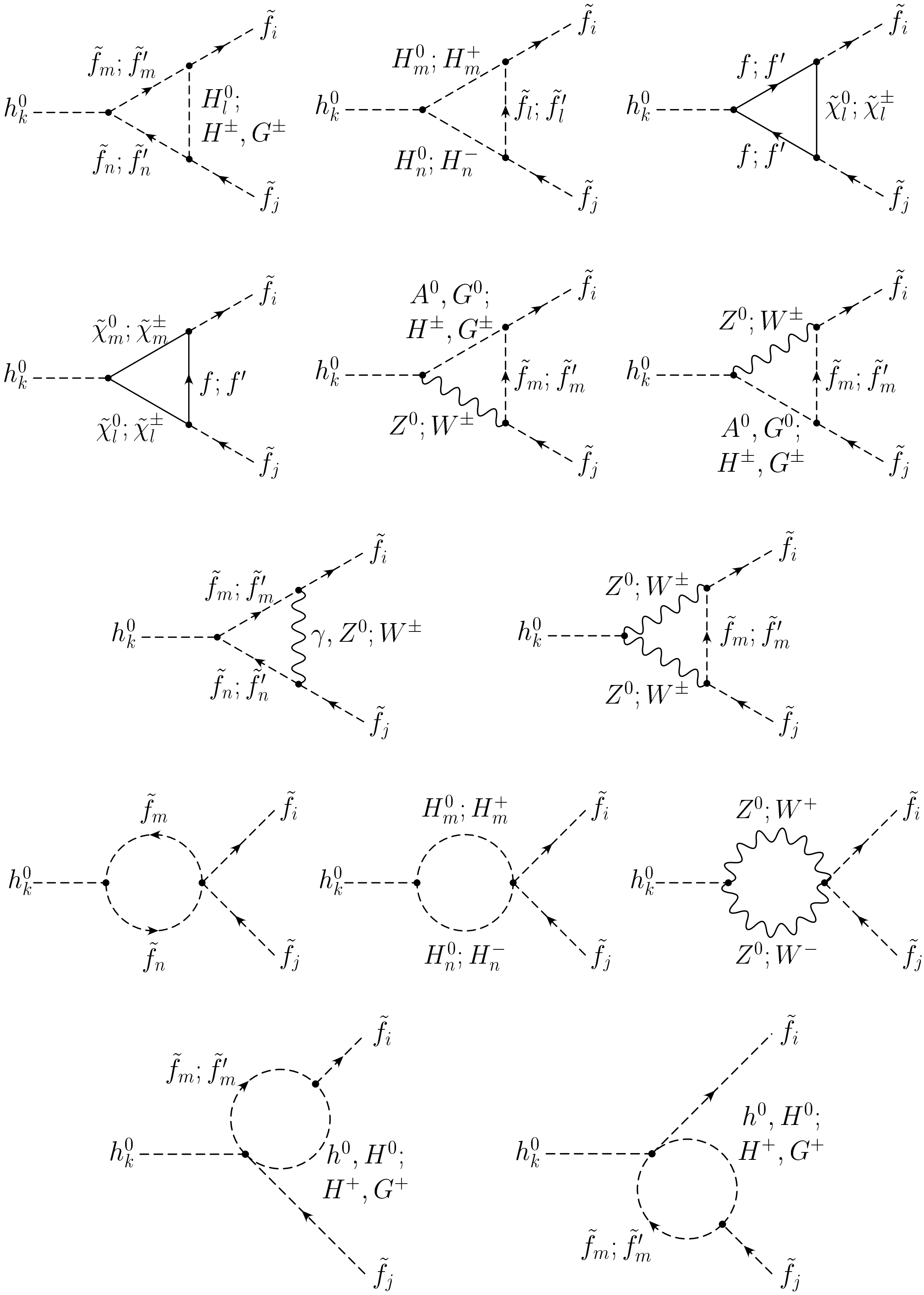


Figure C.1: Vertex and photon emission diagrams relevant to the calculation of the virtual electroweak corrections to the decay width $h_k^0 \to \tilde{f}_i\bar{\tilde{f}}_j$.

C.2 $A^0 \tilde{f}_1 \bar{\tilde{f}}_2$ vertex

Like in the previous section the single contributions to the vertex corrections to $A^0 \to \tilde{f}_1 \bar{\tilde{f}}_2$ (see Fig. C.2),

$$\begin{aligned} \delta G_{123}^{\tilde{f}(v)} &= \delta G_{123}^{\tilde{f}(v,H\tilde{f}\tilde{f})} + \delta G_{123}^{\tilde{f}(v,\tilde{f}HH)} + \delta G_{123}^{\tilde{f}(v,\tilde{\chi}ff)} + \delta G_{123}^{\tilde{f}(v,f\tilde{\chi}\tilde{\chi})} \\ &\quad + \delta G_{123}^{\tilde{f}(v,V)} + \delta G_{123}^{\tilde{f}(v,\tilde{f}\tilde{f})} + \delta G_{123}^{\tilde{f}(v,H\tilde{f})} + \delta G_{123}^{\tilde{f}(v,AZ)} + \delta G_{123}^{\tilde{f}(v,AG)} , \end{aligned} \tag{C.27}$$

correspond to the diagrams with three scalar particles $\left(\delta G_{123}^{\tilde{f}(v,H\tilde{f}\tilde{f})} \text{ and } \delta G_{123}^{\tilde{f}(v,HH\tilde{f})}\right)$, three fermions $\left(\delta G_{123}^{\tilde{f}(v,\tilde{\chi}ff)} \text{ and } \delta G_{123}^{\tilde{f}(v,f\tilde{\chi}\tilde{\chi})}\right)$, one vector particle $\left(\delta G_{123}^{\tilde{f}(v,V)}\right)$ or two scalar particles $\left(\delta G_{123}^{\tilde{f}(v,\tilde{f}\tilde{f})} \text{ and } \delta G_{123}^{\tilde{f}(v,H\tilde{f})}\right)$ in the loop. $\delta G_{123}^{\tilde{f}(v,AZ\mathrm{mix})}$ denotes the correction due to the mixing of A^0 and Z^0 and $\delta G_{123}^{\tilde{f}(v,AG)}$ is the Higgs mixing transition A^0–G^0.

As shown in [45] we can sum up the A^0Z^0 and A^0G^0 transition amplitudes which leads to

$$\delta G_{123}^{\tilde{f}(v,AZ)} + \delta G_{123}^{\tilde{f}(v,AG)} = -\frac{i}{m_Z}\, \Pi_{AZ}(m_{A^0}^2)\, G_{124}^{\tilde{f}} . \tag{C.28}$$

The explicit form of the A^0–Z^0 self–energy, $\Pi_{AZ}(m_{A^0}^2)$, is given in app. B.5. The vertex corrections from the exchange of three scalar particles, i.e. one Higgs/two sfermions and two Higgs-bosons/one sfermion are given by

$$\begin{aligned} \delta G_{123}^{\tilde{f}(v,H\tilde{f}\tilde{f})} &= -\frac{1}{(4\pi)^2} \sum_{m,n=1}^{2} \sum_{k=1}^{4} G_{mn3}^{\tilde{f}} G_{imk}^{\tilde{f}} G_{njk}^{\tilde{f}}\, C_0\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_{H_k^0}^2, m_{\tilde{f}_m}^2, m_{\tilde{f}_n}^2\Big) \\ &\quad -\frac{1}{(4\pi)^2} \sum_{m,n=1}^{2} \sum_{k=1}^{2} G_{mn3}^{\tilde{f}'} G_{imk}^{\tilde{f}\tilde{f}'} G_{jnk}^{\tilde{f}\tilde{f}'}\, C_0\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_{H_k^+}^2, m_{\tilde{f}'_m}^2, m_{\tilde{f}'_n}^2\Big) , \end{aligned} \tag{C.29}$$

$$\begin{aligned} \delta G_{123}^{\tilde{f}(v,\tilde{f}HH)} &= -\frac{1}{(4\pi)^2} \frac{g_Z\, m_Z}{2} \sum_{m=1}^{2} \left(\sum_{k=1}^{2} \sum_{l=3}^{4} G_{imk}^{\tilde{f}} G_{mjl}^{\tilde{f}}\, A_{k,l-2} + \sum_{k=3}^{4} \sum_{l=1}^{2} G_{imk}^{\tilde{f}} G_{mjl}^{\tilde{f}}\, A_{l,k-2} \right) \times \\ &\qquad C_0\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}_m}^2, m_{H_k^0}^2, m_{H_l^0}^2\Big) \\ &\quad -\frac{i}{(4\pi)^2}\, I_f^{3L}\, g\, m_W \sum_{m=1}^{2} \Big(G_{im1}^{\tilde{f}\tilde{f}'} G_{jm2}^{\tilde{f}\tilde{f}'}\, C_0\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}'_m}^2, m_{H^+}^2, m_{G^+}^2\Big) \\ &\qquad - G_{im2}^{\tilde{f}\tilde{f}'}\, G_{jm1}^{\tilde{f}\tilde{f}'}\, C_0\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}'_m}^2, m_{G^+}^2, m_{H^+}^2\Big)\Big) , \end{aligned} \tag{C.30}$$

with the abbreviation

$$A_{kl} = \begin{pmatrix} -\cos 2\beta\, \sin(\alpha+\beta) & -\sin 2\beta\, \sin(\alpha+\beta) \\ \cos 2\beta\, \cos(\alpha+\beta) & \sin 2\beta\, \cos(\alpha+\beta) \end{pmatrix} .$$

With the generic fermion vertex function $F(m_1^2, m_0^2, m_2^2, M_0, M_1, M_2; g_0^R, g_0^L, g_1^R, g_1^L, g_2^R, g_2^L)$ defined in eq. (C.7) the gaugino loop contributions can be written as

$$\begin{aligned}
\delta G_{123}^{\tilde{f}(v,\tilde{\chi}ff)} &= \frac{1}{(4\pi)^2}\sum_{k=1}^{4} F\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{\chi}_k^0}, m_f, m_f; s_3^f, -s_3^f, b_{ik}^{\tilde{f}}, a_{ik}^{\tilde{f}}, a_{jk}^{\tilde{f}}, b_{jk}^{\tilde{f}}\Big) \\
&\quad + \frac{1}{(4\pi)^2}\sum_{k=1}^{2} F\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{\chi}_k^+}, m_{f'}, m_{f'}; s_3^{f'}, -s_3^{f'}, k_{ik}^{\tilde{f}}, l_{ik}^{\tilde{f}}, l_{jk}^{\tilde{f}}, k_{jk}^{\tilde{f}}\Big)\,, \\
\delta G_{123}^{\tilde{f}(v,f\tilde{\chi}\tilde{\chi})} &= \frac{1}{(4\pi)^2}\sum_{k,l=1}^{4} F\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_f, m_{\tilde{\chi}_k^0}, m_{\tilde{\chi}_l^0}; igF_{lk3}^0, -igF_{lk3}^0, b_{ik}^{\tilde{f}}, a_{ik}^{\tilde{f}}, a_{jl}^{\tilde{f}}, b_{jl}^{\tilde{f}}\Big) \\
&\quad + \frac{1}{(4\pi)^2}\sum_{k,l=1}^{2} F\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_{f'}, m_{\tilde{\chi}_k^+}, m_{\tilde{\chi}_l^+}; ig\widetilde{F}_{kl3}^+, -ig\widetilde{F}_{lk3}^+, k_{ik}^{\tilde{f}}, l_{ik}^{\tilde{f}}, l_{jl}^{\tilde{f}}, k_{jl}^{\tilde{f}}\Big)\,,
\end{aligned} \tag{C.31}$$

with $\widetilde{F}_{kl3}^+ = F_{kl3}^+$ for up–type sfermions and $\widetilde{F}_{kl3}^+ = F_{lk3}^+$ for down–type sfermions.
The vertex graphs with one vector and two scalar particles in the loop,

$$\delta G_{123}^{\tilde{f}(v,V)} = \delta G_{123}^{\tilde{f}(v,\gamma)} + \delta G_{123}^{\tilde{f}(v,Z)} + \delta G_{123}^{\tilde{f}(v,W)}\,, \tag{C.32}$$

are given by

$$\begin{aligned}
\delta G_{123}^{\tilde{f}(v,\gamma)} &= \frac{1}{(4\pi)^2}\,(e_0 e_f)^2 G_{123}^{\tilde{f}}\; V\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, \lambda^2, m_{\tilde{f}_i}^2, m_{\tilde{f}_j}^2\Big)\,, \\
\delta G_{123}^{\tilde{f}(v,Z)} &= \frac{1}{(4\pi)^2}\, g_Z^2 \sum_{m,n=1}^{2} G_{mn3}^{\tilde{f}}\, z_{im}^{\tilde{f}}\, z_{nj}^{\tilde{f}}\; V\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_Z^2, m_{\tilde{f}_m}^2, m_{\tilde{f}_n}^2\Big) \\
&\quad - \frac{i}{(4\pi)^2}\,\frac{g_Z^2}{2}\sum_{k,m=1}^{2} G_{mjk}^{\tilde{f}}\, z_{im}^{\tilde{f}}\, R_{1k}(\alpha-\beta)\; V\Big(m_{A^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}_i}^2, m_Z^2, m_{H_k^0}^2, m_{\tilde{f}_m}^2\Big) \\
&\quad + \frac{i}{(4\pi)^2}\,\frac{g_Z^2}{2}\sum_{k,m=1}^{2} G_{imk}^{\tilde{f}}\, z_{mj}^{\tilde{f}}\, R_{1k}(\alpha-\beta)\; V\Big(m_{\tilde{f}_j}^2, m_{\tilde{f}_i}^2, m_{A^0}^2, m_Z^2, m_{\tilde{f}_m}^2, m_{H_k^0}^2\Big)\,, \\
\delta G_{123}^{\tilde{f}(v,W)} &= \frac{1}{(4\pi)^2}\,\frac{g^2}{2}\sum_{m,n=1}^{2} G_{mn3}^{\tilde{f}'}\, R_{i1}^{\tilde{f}} R_{j1}^{\tilde{f}} R_{m1}^{\tilde{f}'} R_{n1}^{\tilde{f}'}\; V\Big(m_{\tilde{f}_i}^2, m_{A^0}^2, m_{\tilde{f}_j}^2, m_W^2, m_{\tilde{f}'_m}^2, m_{\tilde{f}'_n}^2\Big) \\
&\quad + \frac{i}{(4\pi)^2}\,\frac{g^2}{2\sqrt{2}}\sum_{m=1}^{2} G_{jm1}^{\tilde{f}\tilde{f}'}\, R_{i1}^{\tilde{f}} R_{m1}^{\tilde{f}'}\; V\Big(m_{A^0}^2, m_{\tilde{f}_j}^2, m_{\tilde{f}_i}^2, m_W^2, m_{H^+}^2, m_{\tilde{f}'_m}^2\Big) \\
&\quad - \frac{i}{(4\pi)^2}\,\frac{g^2}{2\sqrt{2}}\sum_{m=1}^{2} G_{im1}^{\tilde{f}\tilde{f}'}\, R_{m1}^{\tilde{f}'} R_{j1}^{\tilde{f}} V\Big(m_{\tilde{f}_j}^2, m_{\tilde{f}_i}^2, m_{A^0}^2, m_W^2, m_{\tilde{f}'_m}^2, m_{H^+}^2\Big)\,,
\end{aligned} \tag{C.33}$$

the vector vertex function $V(m_1^2, m_0^2, m_2^2, M_0^2, M_1^2, M_2^2)$ can be looked up in eq. (C.12).

For the vertex graphs with two scalar particles in the loop we obtain

$$\begin{aligned}
\delta G_{123}^{\tilde{f}(v,\tilde{f}\tilde{f})} = & -\frac{1}{(4\pi)^2}\, h_f^2 \sum_{m,n=1}^{2} G_{nm3}^{\tilde{f}} \Big[R_{ijmn}^{\tilde{f}} + R_{mnij}^{\tilde{f}} + N_C^f \Big(R_{inmj}^{\tilde{f}} + R_{mjin}^{\tilde{f}} \Big) \Big] B_0\Big(m_{A^0}^2, m_{\tilde{f}_m}^2, m_{\tilde{f}_n}^2 \Big) \\
& -\frac{1}{(4\pi)^2}\, g_Z^2 \sum_{m,n=1}^{2} G_{nm3}^{\tilde{f}} \Big\{ \Big[\Big(\frac{1}{4} - (2 I_f^{3L} - e_f) e_f s_W^2 \Big) R_{ijmn}^{\tilde{f}_L} + e_f^2 s_W^2 R_{ijmn}^{\tilde{f}_R} \Big] (N_C^f + 1) \\
& \qquad + (I_f^{3L} - e_f) e_f s_W^2 \Big[N_C^f \Big(R_{ijmn}^{\tilde{f}} + R_{mnij}^{\tilde{f}} \Big) + R_{inmj}^{\tilde{f}} + R_{mjin}^{\tilde{f}} \Big] \Big\} \\
& \qquad \times\ B_0\Big(m_{A^0}^2, m_{\tilde{f}_m}^2, m_{\tilde{f}_n}^2 \Big) \\
& -\frac{1}{(4\pi)^2}\, N_C^{\hat{f}}\, h_f h_{\hat{f}} \sum_{m,n=1}^{2} G_{nm3}^{\hat{\tilde{f}}} \Big(R_{ijnm}^{\tilde{f}\hat{\tilde{f}}_F} + R_{jimn}^{\tilde{f}\hat{\tilde{f}}_F} \Big) B_0\Big(m_{A^0}^2, m_{\hat{\tilde{f}}_m}^2, m_{\hat{\tilde{f}}_n}^2 \Big)\,, \qquad \text{(C.34)}
\end{aligned}$$

$$\begin{aligned}
\delta G_{123}^{\tilde{f}(v,H\tilde{f})} = & -\frac{1}{(4\pi)^2} \sum_{k=3}^{4} \sum_{m=1}^{2} G_{imk}^{\tilde{f}} \Big(h_f^2\, c_{3k}^{\tilde{f}}\, \delta_{mj} + g^2 \Big(c_{3k}^{\tilde{b}} - c_{3k}^{\tilde{t}} \Big) e_{mj}^{\tilde{f}} \Big) B_0\Big(m_{A^0}^2, m_{H_k^0}^2, m_{\tilde{f}_m}^2 \Big) \\
& + \frac{i}{(4\pi)^2} \sqrt{2}\, I_f^{3L} \sum_{k,m=1}^{2} G_{imk}^{\tilde{f}\tilde{f}'} \Big[\left\{ \begin{matrix} (h_{f_\uparrow}^2 - g^2/2)\cos^2\beta - (h_{f_\downarrow}^2 - g^2/2)\sin^2\beta \\ (h_f^2 + h_{f'}^2 - g^2)\sin\beta\cos\beta \end{matrix} \right\}_k R_{m1}^{\tilde{f}'} R_{j1}^{\tilde{f}} \\
& \qquad + h_f h_{f'}\, \delta_{k2}\, R_{m2}^{\tilde{f}'} R_{j2}^{\tilde{f}} \Big] B_0\Big(m_{A^0}^2, m_{H_k^+}^2, m_{\tilde{f}'_m}^2 \Big) \\
& - i \leftrightarrow j\,. \qquad \text{(C.35)}
\end{aligned}$$

The abbreviations used in the equations above are the same as in the case of $h_k^0 \to \tilde{f}_i \bar{\tilde{f}}_j$, see section C.1.

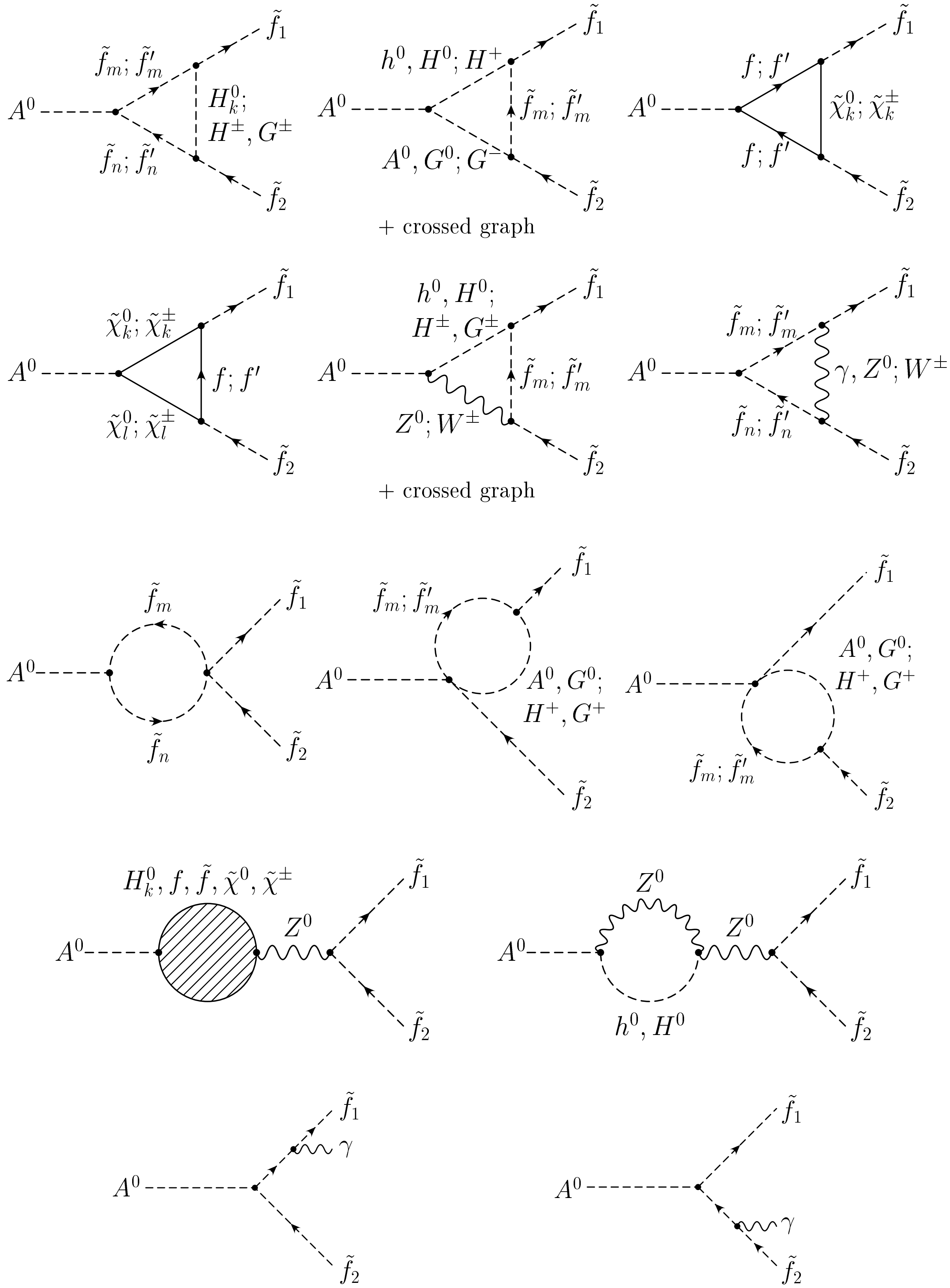


Figure C.2: Vertex and photon emission diagrams relevant to the calculation of the virtual electroweak corrections to the decay width $A^0 \to \tilde{f}_1\bar{\tilde{f}}_2$.

C.3 $H^+\tilde{t}_i\bar{\tilde{b}}_j$ vertex

Using the definitions for various generic vertex functions and products of couplings from the previous chapters we get for the vertex corrections,

$$\begin{aligned}\delta G_{ij1}^{\tilde{t}\tilde{b}(v)} &= \delta G_{ij1}^{\tilde{t}\tilde{b}(v,H\tilde{f}\tilde{f}')} + \delta G_{ij1}^{\tilde{t}\tilde{b}(v,HH\tilde{f})} + \delta G_{ij1}^{\tilde{t}\tilde{b}(v,\tilde{\chi}ff)} + \delta G_{ijk}^{\tilde{f}(v,f\tilde{\chi}\tilde{\chi})} + \delta G_{ij1}^{\tilde{t}\tilde{b}(v,\gamma SS)} \\ &+ \delta G_{ij1}^{\tilde{t}\tilde{b}(v,ZSS)} + \delta G_{ij1}^{\tilde{t}\tilde{b}(v,WSS)} + \delta G_{ij1}^{\tilde{t}\tilde{b}(v,HH)} + \delta G_{ij1}^{\tilde{t}\tilde{b}(v,\tilde{f}\tilde{f}')} + \delta G_{ij1}^{\tilde{t}\tilde{b}(v,\tilde{F}\tilde{F}')} \\ &+ \delta G_{ij1}^{\tilde{t}\tilde{b}(v,H\tilde{f})} \end{aligned} \quad \text{(C.36)}$$

The single contributions are given as follows:

$$\delta G_{ij1}^{\tilde{t}\tilde{b}(v,H\tilde{f}\tilde{f}')} = -\frac{1}{(4\pi)^2}\sum_{m,n=1}^{2}\sum_{k=1}^{4} G_{mn1}^{\tilde{t}\tilde{b}} G_{imk}^{\tilde{t}} G_{njk}^{\tilde{b}}\, C_0\Big(m_{\tilde{t}_i}^2, m_{H^+}^2, m_{\tilde{b}_j}^2, m_{H_k^0}^2, m_{\tilde{t}_m}^2, m_{\tilde{b}_n}^2\Big) \quad \text{(C.37)}$$

$$\begin{aligned}\delta G_{ij1}^{\tilde{t}\tilde{b}(v,HH\tilde{f})} &= -\frac{1}{(4\pi)^2}\frac{1}{2}\sum_{k,l,m=1}^{2}\Big[(-1)^l g\, m_W(1+\delta_{1l})\tilde{A}_{lk}^{\prime(1)} + g_Z\, m_Z R_{2k}(\alpha+\beta)\tilde{C}_{1l}\Big] \times \\ &\qquad G_{imk}^{\tilde{t}} G_{mjl}^{\tilde{t}\tilde{b}} C_0\Big(m_{\tilde{t}_i}^2, m_{H^+}^2, m_{\tilde{b}_j}^2, m_{\tilde{t}_m}^2, m_{h_k^0}^2, m_{H_l^+}^2\Big) \\ &+\frac{1}{(4\pi)^2}\frac{igm_W}{2}\sum_{m=1}^{2} G_{im3}^{\tilde{t}} G_{mj2}^{\tilde{t}\tilde{b}} C_0\Big(m_{\tilde{t}_i}^2, m_{H^+}^2, m_{\tilde{b}_j}^2, m_{\tilde{t}_m}^2, m_{A^0}^2, m_W^2\Big) \\ &-\frac{1}{(4\pi)^2}\frac{1}{2}\sum_{k,l,m=1}^{2}\Big[(-1)^l g\, m_W(1+\delta_{1l})\tilde{A}_{lk}^{\prime(1)} + g_Z\, m_Z R_{2k}(\alpha+\beta)\tilde{C}_{1l}\Big] \times \\ &\qquad G_{mjk}^{\tilde{b}} G_{iml}^{\tilde{t}\tilde{b}} C_0\Big(m_{\tilde{t}_i}^2, m_{H^+}^2, m_{\tilde{b}_j}^2, m_{\tilde{b}_m}^2, m_{H_l^+}^2, m_{h_k^0}^2\Big) \\ &+\frac{1}{(4\pi)^2}\frac{igm_W}{2}\sum_{m=1}^{2} G_{mj3}^{\tilde{b}} G_{im2}^{\tilde{t}\tilde{b}} C_0\Big(m_{\tilde{t}_i}^2, m_{H^+}^2, m_{\tilde{b}_j}^2, m_{\tilde{b}_m}^2, m_W^2, m_{A^0}^2\Big) \end{aligned} \quad \text{(C.38)}$$

$$\delta G_{ij1}^{\tilde{t}\tilde{b}(v,\tilde{\chi}ff)} = \frac{1}{(4\pi)^2}\sum_{k=1}^{4} F\Big(m_{\tilde{t}_i}^2, m_{H^+}^2, m_{\tilde{b}_j}^2, m_{\tilde{\chi}_k^0}, m_t, m_b; y_1^b, y_1^t, b_{ik}^{\tilde{t}}, a_{ik}^{\tilde{t}}, a_{jk}^{\tilde{b}}, b_{jk}^{\tilde{b}}\Big) \quad \text{(C.39)}$$

$$\begin{aligned}\delta G_{ijk}^{\tilde{f}(v,f\tilde{\chi}\tilde{\chi})} &= \\ &-\frac{1}{(4\pi)^2}\sum_{k=1}^{4}\sum_{l=1}^{2} F\Big(m_{\tilde{t}_i}^2, m_{H^+}^2, m_{\tilde{b}_j}^2, m_t, m_{\tilde{\chi}_k^0}, m_{\tilde{\chi}_l^+}; -gF_{lk1}^R, -gF_{lk1}^L, b_{ik}^{\tilde{t}}, a_{ik}^{\tilde{t}}, l_{jl}^{\tilde{b}}, k_{jl}^{\tilde{b}}\Big) \\ &-\frac{1}{(4\pi)^2}\sum_{k=1}^{4}\sum_{l=1}^{2} F\Big(m_{\tilde{t}_i}^2, m_{H^+}^2, m_{\tilde{b}_j}^2, m_b, m_{\tilde{\chi}_l^+}, m_{\tilde{\chi}_k^0}; -gF_{lk1}^R, -gF_{lk1}^L, k_{il}^{\tilde{t}}, l_{il}^{\tilde{t}}, a_{jk}^{\tilde{b}}, b_{jk}^{\tilde{b}}\Big) \end{aligned} \quad \text{(C.40)}$$

$$\delta G_{ij1}^{\tilde{t}\tilde{b}(v,\gamma SS)} = \frac{1}{(4\pi)^2}\, e_0^2 e_t e_b G_{ij1}^{\tilde{t}\tilde{b}}\; V\Big(m_{\tilde{t}_i}^2, m_{H^+}^2, m_{\tilde{b}_j}^2, \lambda^2, m_{\tilde{t}_i}^2, m_{\tilde{b}_j}^2\Big)$$

$$
\begin{aligned}
&+\frac{1}{(4\pi)^2}\, e_0^2 e_t G^{\tilde{t}\tilde{b}}_{ij1}\; V\Big(m^2_{H^+}, m^2_{\tilde{b}_j}, m^2_{\tilde{t}_i}, \lambda^2, m^2_{H^+}, m^2_{\tilde{t}_i}\Big) \\
&-\frac{1}{(4\pi)^2}\, e_0^2 e_b G^{\tilde{t}\tilde{b}}_{ij1}\; V\Big(m^2_{\tilde{b}_j}, m^2_{\tilde{t}_i}, m^2_{H^+}, \lambda^2, m^2_{\tilde{b}_j}, m^2_{H^+}\Big) \qquad \text{(C.41)}
\end{aligned}
$$

$$
\begin{aligned}
\delta G^{\tilde{t}\tilde{b}(v,ZSS)}_{ij1} &= \frac{1}{(4\pi)^2}\, g_Z^2 \sum_{m,n=1}^{2} G^{\tilde{t}\tilde{b}}_{mn1}\, z^{\tilde{t}}_{im}\, z^{\tilde{b}}_{nj}\; V\Big(m^2_{\tilde{t}_i}, m^2_{H^+}, m^2_{\tilde{b}_j}, m^2_Z, m^2_{\tilde{t}_m}, m^2_{\tilde{b}_n}\Big) \\
&+\frac{1}{(4\pi)^2}\, g_Z^2 \Big(\frac{1}{2} - s_W^2\Big) \sum_{m=1}^{2} G^{\tilde{t}\tilde{b}}_{mj1}\, z^{\tilde{t}}_{im}\; V\Big(m^2_{H^+}, m^2_{\tilde{b}_j}, m^2_{\tilde{t}_i}, m^2_Z, m^2_{H^+}, m^2_{\tilde{t}_m}\Big) \\
&-\frac{1}{(4\pi)^2}\, g_Z^2 \Big(\frac{1}{2} - s_W^2\Big) \sum_{m=1}^{2} G^{\tilde{t}\tilde{b}}_{im1}\, z^{\tilde{b}}_{mj}\; V\Big(m^2_{\tilde{b}_j}, m^2_{\tilde{t}_i}, m^2_{H^+}, m^2_Z, m^2_{\tilde{b}_m}, m^2_{H^+}\Big)
\end{aligned}
$$

(C.42)

$$
\begin{aligned}
\delta G^{\tilde{t}\tilde{b}(v,WSS)}_{ij1} &= \frac{1}{(4\pi)^2}\, \frac{g^2}{2\sqrt{2}} \sum_{m=1}^{2}\sum_{k=1}^{3} G^{\tilde{b}}_{mjk}\, R^{\tilde{t}}_{i1} R^{\tilde{b}}_{m1} w_k\, V\Big(m^2_{H^+}, m^2_{\tilde{b}_j}, m^2_{\tilde{t}_i}, m^2_W, m^2_{H^0_k}, m^2_{\tilde{b}_m}\Big) \\
&-\frac{1}{(4\pi)^2}\, \frac{g^2}{2\sqrt{2}} \sum_{m=1}^{2}\sum_{k=1}^{3} G^{\tilde{t}}_{imk}\, R^{\tilde{t}}_{m1} R^{\tilde{b}}_{j1} w_k\, V\Big(m^2_{\tilde{b}_j}, m^2_{\tilde{t}_i}, m^2_{H^+}, m^2_W, m^2_{\tilde{t}_m}, m^2_{H^0_k}\Big)\,,
\end{aligned}
$$

(C.43)

with $w_k = \{\cos(\alpha-\beta), \sin(\alpha-\beta), -i\}_k$.

$$
\begin{aligned}
\delta G^{\tilde{t}\tilde{b}(v,HH)}_{ij1} = &-\frac{1}{(4\pi)^2}\frac{1}{2\sqrt{2}} \sum_{k,l=1}^{2} \Big[(-1)^l g\, m_W (1+\delta_{1l}) \tilde{A}'^{(1)}_{lk} + g_Z\, m_Z R_{2k}(\alpha{+}\beta) \tilde{C}_{1l}\Big] \times \\
&\Big[\Big((h_t^2 - \frac{g^2}{2}) c^{\tilde{t},0+}_{kl} + (h_b^2 - \frac{g^2}{2})(c^{\tilde{b},0+}_{k,l+2})^*\Big) R^{\tilde{t}}_{i1} R^{\tilde{b}}_{j1} \\
&+ h_t h_b \Big(c^{\tilde{t}\tilde{b},0+}_{kl} + (c^{\tilde{b}\tilde{t},0+}_{k,l+2})^*\Big) R^{\tilde{t}}_{i2} R^{\tilde{b}}_{j2}\Big] B_0\Big(m^2_{H^+}, m^2_{h^0_k}, m^2_{H^+_l}\Big) \\
&-\frac{1}{(4\pi)^2}\frac{g m_W}{2\sqrt{2}} \Big[\frac{1}{2}(h_t^2{+}h_b^2{-}g^2)\sin 2\beta R^{\tilde{t}}_{i1} R^{\tilde{b}}_{j1} + h_t h_b R^{\tilde{t}}_{i2} R^{\tilde{b}}_{j2}\Big] B_0\Big(m^2_{H^+}, m^2_{A^0}, m^2_W\Big)
\end{aligned}
$$

(C.44)

$$
\begin{aligned}
\delta G^{\tilde{t}\tilde{b}(v,\tilde{f}\tilde{f}')}_{ij1} = &-\frac{1}{(4\pi)^2} \sum_{m,n=1}^{2} G^{\tilde{t}\tilde{b}}_{nm1} \Big\{ N_C^f \Big(h_t^2\, R^{\tilde{b}\tilde{t}_D}_{mjin} + h_b^2\, R^{\tilde{t}\tilde{b}_D}_{inmj}\Big) + \frac{g^2}{4}\Big[(2N_C^f - 1) R^{\tilde{t}\tilde{b}_L}_{inmj} \\
&+ t_W^2 \Big(Y_L^t Y_L^b R^{\tilde{t}\tilde{b}_L}_{inmj} + Y_R^t Y_R^b R^{\tilde{t}\tilde{b}_R}_{inmj} - Y_L^t Y_R^b R^{\tilde{t}\tilde{b}_D}_{inmj} \\
&- Y_L^b Y_R^t R^{\tilde{b}\tilde{t}_D}_{mjin}\Big)\Big]\Big\} B_0\Big(m^2_{H^+}, m^2_{\tilde{b}_m}, m^2_{\tilde{t}_n}\Big) \\
&-\frac{1}{(4\pi)^2} \sum_{m=1}^{2} G^{\tilde{\nu}_\tau\tilde{\tau}}_{1m1} \Big(h_b h_\tau R^{\tilde{\tau}}_{m2} R^{\tilde{t}}_{i1} R^{\tilde{b}}_{j2} + \frac{g^2}{2} R^{\tilde{\tau}}_{m1} R^{\tilde{t}}_{i1} R^{\tilde{b}}_{j1}\Big) B_0\Big(m^2_{H^+}, m^2_{\tilde{\tau}_m}, m^2_{\tilde{\nu}_\tau}\Big)
\end{aligned}
$$

(C.45)

$$\begin{aligned}
\delta G^{\tilde{t}\tilde{b}(v,\tilde{F}\tilde{F}')}_{ij1} &= -\frac{N_C^f}{(4\pi)^2}\sum_{m,n=1}^{2} G^{\tilde{u}\tilde{d}}_{nm1}\Big(\frac{g^2}{2}R^{\tilde{t}}_{i1}R^{\tilde{b}}_{j1}R^{\tilde{u}}_{n1}R^{\tilde{d}}_{m1} + h_t h_u R^{\tilde{t}}_{i2}R^{\tilde{b}}_{j1}R^{\tilde{u}}_{n2}R^{\tilde{d}}_{m1} \\
&\qquad + h_b h_d R^{\tilde{t}}_{i1}R^{\tilde{b}}_{j2}R^{\tilde{u}}_{n1}R^{\tilde{d}}_{m2}\Big)B_0\Big(m^2_{H^+}, m^2_{\tilde{d}_m}, m^2_{\tilde{u}_n}\Big) \\
&\quad -\frac{1}{(4\pi)^2}\sum_{m=1}^{2} G^{\tilde{\nu}_e\tilde{e}}_{1m1}\Big(\frac{g^2}{2}R^{\tilde{t}}_{i1}R^{\tilde{b}}_{j1}R^{\tilde{e}}_{m1} + h_b h_e R^{\tilde{t}}_{i1}R^{\tilde{b}}_{j2}R^{\tilde{e}}_{m2}\Big)B_0\Big(m^2_{H^+}, m^2_{\tilde{e}_m}, m^2_{\tilde{\nu}_e}\Big) \\
&\quad -\frac{N_C^f}{(4\pi)^2}\sum_{m,n=1}^{2} G^{\tilde{c}\tilde{s}}_{nm1}\Big(\frac{g^2}{2}R^{\tilde{t}}_{i1}R^{\tilde{b}}_{j1}R^{\tilde{c}}_{n1}R^{\tilde{s}}_{m1} + h_t h_c R^{\tilde{t}}_{i2}R^{\tilde{b}}_{j1}R^{\tilde{c}}_{n2}R^{\tilde{s}}_{m1} \\
&\qquad + h_b h_s R^{\tilde{t}}_{i1}R^{\tilde{b}}_{j2}R^{\tilde{c}}_{n1}R^{\tilde{s}}_{m2}\Big)B_0\Big(m^2_{H^+}, m^2_{\tilde{s}_m}, m^2_{\tilde{c}_n}\Big) \\
&\quad -\frac{1}{(4\pi)^2}\sum_{m=1}^{2} G^{\tilde{\nu}_\mu\tilde{\mu}}_{1m1}\Big(\frac{g^2}{2}R^{\tilde{t}}_{i1}R^{\tilde{b}}_{j1}R^{\tilde{\mu}}_{m1} + h_b h_\mu R^{\tilde{t}}_{i1}R^{\tilde{b}}_{j2}R^{\tilde{\mu}}_{m2}\Big)B_0\Big(m^2_{H^+}, m^2_{\tilde{\mu}_m}, m^2_{\tilde{\nu}_\mu}\Big)
\end{aligned} \tag{C.46}$$

$$\begin{aligned}
\delta G^{\tilde{t}\tilde{b}(v,H\tilde{f})}_{ij1} &= -\frac{1}{(4\pi)^2}\frac{1}{\sqrt{2}}\sum_{k=1}^{4}\sum_{m=1}^{2} G^{\tilde{t}}_{imk}\Big[R^{\tilde{t}}_{m1}R^{\tilde{b}}_{j1}\Big(h_t^2 c^{\tilde{t},0+}_{k1} + h_b^2 (c^{\tilde{b},0+}_{k3})^*\Big) + R^{\tilde{t}}_{m2}R^{\tilde{b}}_{j2}h_t h_b \times \\
&\qquad \Big(c^{\tilde{t}\tilde{b},0+}_{k1} + (c^{\tilde{b}\tilde{t},0+}_{k3})^*\Big) - \frac{g^2}{2}R^{\tilde{t}}_{m1}R^{\tilde{b}}_{j1}\Big(c^{\tilde{t},0+}_{k1} + (c^{\tilde{b},0+}_{k3})^*\Big)\Big]B_0\Big(m^2_{\tilde{t}_i}, m^2_{H^0_k}, m^2_{\tilde{t}_m}\Big) \\
&\quad -\frac{1}{(4\pi)^2}\frac{1}{\sqrt{2}}\sum_{k=1}^{4}\sum_{m=1}^{2} G^{\tilde{b}}_{mjk}\Big[R^{\tilde{t}}_{i1}R^{\tilde{b}}_{m1}\Big(h_t^2 c^{\tilde{t},0+}_{k1} + h_b^2 (c^{\tilde{b},0+}_{k3})^*\Big) + R^{\tilde{t}}_{i2}R^{\tilde{b}}_{m2}h_t h_b \times \\
&\qquad \Big(c^{\tilde{t}\tilde{b},0+}_{k1} + (c^{\tilde{b}\tilde{t},0+}_{k3})^*\Big) - \frac{g^2}{2}R^{\tilde{t}}_{i1}R^{\tilde{b}}_{m1}\Big(c^{\tilde{t},0+}_{k1} + (c^{\tilde{b},0+}_{k3})^*\Big)\Big]B_0\Big(m^2_{\tilde{b}_j}, m^2_{H^0_k}, m^2_{\tilde{b}_m}\Big) \\
&\quad -\frac{1}{(4\pi)^2}\sum_{k,m=1}^{2} G^{\tilde{t}\tilde{b}}_{mjk}\Big[h_b^2 d^{\tilde{b}}_{1k}R^{\tilde{t}}_{i1}R^{\tilde{t}}_{m1} + h_t^2 d^{\tilde{t}}_{1k}R^{\tilde{t}}_{i2}R^{\tilde{t}}_{m2} + g^2 f^{\tilde{t}}_{im}\Big(d^{\tilde{b}}_{1k} - d^{\tilde{t}}_{1k}\Big)\Big] \times \\
&\qquad B_0\Big(m^2_{\tilde{b}_j}, m^2_{H^+_k}, m^2_{\tilde{t}_m}\Big) \\
&\quad -\frac{1}{(4\pi)^2}\sum_{k,m=1}^{2} G^{\tilde{t}\tilde{b}}_{imk}\Big[h_t^2 d^{\tilde{t}}_{1k}R^{\tilde{b}}_{m1}R^{\tilde{b}}_{j1} + h_b^2 d^{\tilde{b}}_{1k}R^{\tilde{b}}_{m2}R^{\tilde{b}}_{j2} + g^2 f^{\tilde{b}}_{mj}\Big(d^{\tilde{b}}_{1k} - d^{\tilde{t}}_{1k}\Big)\Big] \times \\
&\qquad B_0\Big(m^2_{\tilde{t}_i}, m^2_{H^+_k}, m^2_{\tilde{b}_m}\Big)
\end{aligned} \tag{C.47}$$

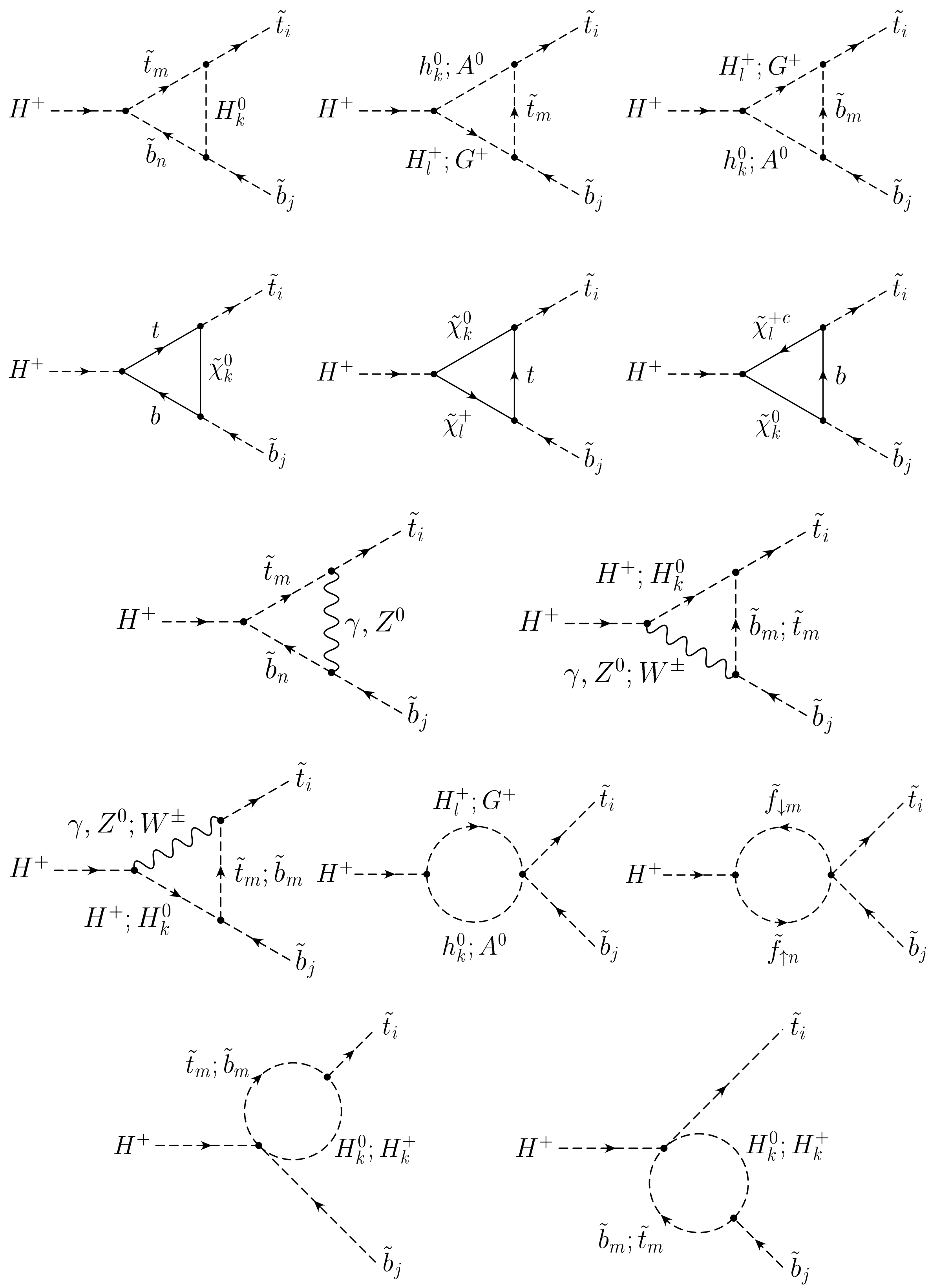


Figure C.3: Vertex diagrams relevant to the calculation of the virtual electroweak corrections to the decay width $H^+ \to \tilde{t}_i\bar{\tilde{b}}_j$. In the fourth row, $\tilde{f}_\uparrow$ and $\tilde{f}_\downarrow$ denote up– and down–type sfermions of all three generations, respectively.

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Curriculum Vitae

Dipl.-Ing. Christian Weber

Schiffmühlenstraße 57/1/10
1220 Wien

Geboren am 16. Februar 1974 in Horn (A).
Verheiratet (Ingrid Riedl), ein Kind (Anna, 10. November 1995).

Ausbildung, Berufspraxis

1980 – 1984	**Volksschule** Schönberg am Kamp (A)
1984 – 1992	**Gymnasium** BG/BRG Krems (A)
Okt 1993 – Jun 2001	**Studium** Techn. Physik, Technische Universität Wien Diplomarbeit (Mai 2001): “Yukawa coupling corrections to Higgs boson decays into sfermions in the Minimal Supersymmetric Standard Model”
Jun 2001 – Okt 2004	**Dissertation** “Electroweak Corrections to Higgs boson Decays into Sfermions in the Minimal Supersymmetric Standard Model”
seit Aug 2001	Anstellung am Institut für Hochenergiephysik der ÖAW Tätigkeitsbereich: Mitarbeit in der SUSY-Gruppe

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