Aspects Of Beyond MSSM Higgs Physics

Implications For The Higgs Spectrum And Processes And

A Full Set Of BMSSM Feynmanrules

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Declaration

I hereby declare that this thesis has not been and will not be submitted in whole or in part to another University for the award of any other degree.

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Summary

Beyond MSSM (BMSSM) Models provide natural solutions to the little hierarchy problem in minimal SUSY theories. Well studied extensions of the MSSM can be organised in an effective operator approach utilising the merrits of an Effective Field Theory to study BMSSM effects. Lifting the lightest Higgs mass to the current LHC bound of $126 \pm 0.4 \pm 0.4$ GeV [8] [9] through the stop loop contribution, BMSSM effects can make significant changes to the upper bounds of Higgs and top squark masses. BMSSM corrections to MSSM Feynman rules are leading to new processes but also significant contributions to in leading order $\frac{1}{\tan\beta}$ to the MSSM amplitudes. In this work we are exploring effects on the Higgs sector mass spectrum and also Higgs interactions in the setup of [1]. The theoretical foundation for Supersymmetry is presented and a motivation for Beyond minimal SUSY models is outlined. The MSSM Higgs sector is presented, followed by an elaborate demonstration how BMSSM contributions in the effective operator approach [1] affect the MSSM Higgs mass spectrum and processes. Finally as an original contribution a full list of Feynman rules for all cubic and quartic BMSSM tree level vertices is presented and discussed.

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Contents

Chapter 1

Introduction

The Minimal Supersymmetric Standard Model (MSSM) is a well-motivated model for supersymmetric New Physics that might be observed at the LHC. However, the recent LHC discovery of the Higgs mass of $126 \pm 0.4 \pm 0.4$ GeV [8] [9] requires the MSSM scale to be rather high compared to the weak scale, leading to an unnatural hierarchy. This suggests the presence of additional physics beyond the MSSM (BMSSM), lifting the lightest Higgs mass in a more natural way. The aim of this work is to elaborate and to arrive at a broader understanding how far BMSSM physics in the set up of [1] will affect observables at the LHC. Most importantly, how it will effect the Higgs sector.

Effective Field Theories as in [1] have the merit to allow a model-indipendent description of a large class of extensions of the MSSM. Those extensions by higher-dimensional operators can have important impact on the Higss sector of the MSSM. As further demonstrated, the leading order $1/M$ contributions to the MSSM are only extended by 2 parameters, while M is the mass sclae of the BMSSM particles. As the Higgs potential in the MSSM is rather restricted at tree level those corrections are significant even if the parameters are in relatively small order of μ/M and lead to qualitative new effects. As demonstrated in [1] they can distinctively destabilise the MSSM minima and have significant consequenses for the Higgs masses and couplings. If M is not too far away from the EW scale, the next to leading order $1/M^2$ effects become more relevant due to the smallness of the quartic tree level couplings. Even if the expansion parameter is relatively small a subset of quartic coupling encounters a correction coefficient of order g^2. A study of consequences for Higgs masses and couplings up to order M^2 is given in [1]. Imposing constraints from LEP and Tevatron and expanding the collider physics for the signals which were expected for the LHC it turned out that there are large corrections to the CP-even Higgs mass and the decay pattern is remarkably different, especially CP-even branching fractions into gauge bosons which are non-standard and 'exotic'.

Motivation of this work

In this work we explored the impact of the tree-level 'BMSSM' Higgs sector in the setup of [1]. Whereas implications for the MSSM mass spectrum are discussed in [1], we are presenting a set of Feynman rules for the BMSSM Higgs boson interactions. A list of BMSSM contributions and additional BMSSM interactions is pliable, which will be needed to probe the LHC sensitivity to BMSSM parameters.

This thesis is organized as follows

A minimum of theoretical foundations to understand the origin and systematic of the MSSM is given in Chapter 2 where motivations for BMSSM models are introduced. In Chapter 3 the idea of the BMSSM in the effective operater approach of [1] is outlined and its Higgs spectrum is dicussed. In Chapter 4 the Feynman rules for the BMSSM are presented and discussed and an overview of the current literature on improved BMSSM studies on the MSSM Higgs sector is given. The content of the theses will be summed up and an outlook is given.

Chapter 2

The MSSM

The following introduction to the origin and systematics of the MSSM is in excerpts taken from [26] except for Section 2.3.

The Standard Model (SM) provides a very elegant theoretical framework to organize elementary particles and their fundamental forces within a gauge group structure and to describe the interactions and the phenomena of the processes between them. Experimental tests highly agree with its predictions at the 0.1% level [6] [8] [9]. It is consistent with both quantum mechanics and special relativity. By elementary particles the point like constituents of matter with no further known substructure are understood. The particle spectrum is classified in fermions (spin (s) one half) and bosons (integer spin). Fermions are classified into leptons and quarks. The force carrier particles, the gauge bosons ($s = 1$), mediate the fundamental interactions: the strong and the electroweak, whereas the latter is the unification of the weak and the electromagnetic force. After electroweak symmetry breaking the particles obtain their masses which originates in the Higgs mechanism caused by one additional particle, called the Higgs boson ($s = 0$). As long as the elusive Higgs had not been observed, the SM was not established as a complete theory. In 4 July 2012 the discovery of an unknown particle with a mass between 125 and 127 GeV/c^2 had been announced at the facilities of the LHC [8]. It was suspected that it was the Higgs boson. By March 2013, the particle had been proven to hold fundamental attributes, like positive parity and zero spin, and to behave in many of the ways like the Higgs boson predicted by the Standard Model. More data is needed to decide if the discovered particle exactly matches the predictions of the Standard Model, or whether, as predicted by some theories like the MSSM and BMSSM models, multiple Higgs bosons exist [2]. Nevertheless, the are numerous models that extend the SM for different reasons, among them the MSSM, on which this chapter focuses, giving brief overview of the construction and the basic features of supersymmetric models. They provide a solution to the hierarchy problem by means of a symmetry between bosons and fermions. Since Supersymmetry is no observed symmetry in nature, it has to be broken in any model intended to produce phenomenologically correct results. Therefore this chapter includes a short description of the soft supersymmetry breaking Lagrangian. As an intermediate step on the way to BMSSM models, the simplest supersymmetric extension of the SM is sketched, the Minimal Supersymmetric Standard Model (MSSM). One of its conceptual problems, the μ problem, is discussed. As a possible solution BMSSM extensions as an effective theory are presented in the next chapter.

2.1 Supersymmetry

Supersymmetry (SUSY) can be regarded as a symmetry between bosons and fermions. In general, these models contain an equal number of bosonic and fermionic degrees of freedom that group into so-called *supermultiplets* [2], each made up of a boson and a Weyl fermion. In model without particles of spin greater than 1, there are two possible types of supermultiplets:

- *Chiral* multiplets, containing a Weyl fermion ψ_i and a complex scalar ϕ_i,
- *Vector* or *gauge* multiplet, containing a Weyl fermion λ^a and a massless vector field A^a_μ.

The particles in a supermultiplet are called *superpartners*. Since we intend to construct gauge theories, we directly include the appropriate indices on the fields, so that i runs over some given representation of the gauge group, while a denotes the adjoint representation. Superpartners have to transform in the same way under the gauge group, especially

$$D_\mu \lambda^a = \partial_\mu \lambda^a - g f^{abc} A^b_\mu \lambda^c \,, \tag{2.1}$$

where g is the gauge coupling and f^{abc} are the structure coefficients of the gauge group, which means that $[t^a, t^b] = f^{abc} t^c$ if t^a are the group generators. The fermionic superpartners of gauge bosons are called gauginos.

Interactions between chiral supermultiplets are described by a *superpotential* [2]. This is a function $W = W(\phi_i)$ of the scalar fields of the chiral multiplets which has the important property of being holomorphic

$$\frac{\partial W}{\partial \phi_i^*} = 0 \tag{2.2}$$

The Lagrangian of a *chiral superfield* is given by

$$\mathcal{L} = \int d^4\theta\, K(\Phi, \Phi^\dagger) + \int d^2\theta\, W(\Phi) + \int d^2\bar{\theta}\, (W(\Phi))^\dagger \ , \tag{2.3}$$

where $K = \sum_i \Phi_i^\dagger \Phi_i$ is the *Kahler potential* and W is the *superpotential* [2][4][23] [24]. Here θ and $\bar{\theta}$ are Grassmann spinor variables. Integration over these variables is defined by:

$$\int d^2\theta\theta^2 = 1, \qquad \int d^2\bar{\theta}\bar{\theta}^2 = 1, \qquad \int d^4\theta\theta^2\bar{\theta}^2 = 1, \tag{2.4}$$

and all other combinations vanish. Φ is a superfield, i.e. a function of x^μ, θ and $\bar{\theta}$. An infinitesimal supersymmetry transformation acts on this field by

$$\Phi \to \Phi^{'} = (1 + i\xi^\alpha Q_\alpha + i\bar{\xi}_{\dot{\alpha}}\bar{Q}^{\dot{\alpha}})\Phi. \tag{2.5}$$

The supersymmetry generator Q is itself a spinorial *supercharge* fulfilling the fundamental relation

$$\left\{ Q_\alpha, Q^\dagger_{\dot{\beta}} \right\} = 2(\sigma^\mu)_{\alpha\dot{\beta}} P_\mu. \tag{2.6}$$

The covariant derivatives for supersymmetry transformations are given by

$$D_\alpha = \frac{\partial}{\partial\theta^\alpha} + i(\sigma^\mu)_{\alpha\dot{\alpha}}\bar{\theta}^{\dot{\alpha}}\partial_\mu \tag{2.7}$$

and

$$\bar{D}_{\dot{\beta}} = \frac{\partial}{\partial\bar{\theta}^{\dot{\beta}}} + i\theta^\beta(\sigma^\mu)_{\beta\dot{\beta}}\partial_\mu \tag{2.8}$$

A chiral superfield is a superfield satisfying the condition $\bar{D}_{\dot{\beta}}\Phi = 0$. Its SUSY transformation (2.5) obviously also fulfills the condition. A superfield is chiral if it depends only on θ and $y_\mu = x_\mu + \bar{\theta}^\alpha \sigma_\mu^{\alpha\beta}\theta^\beta$.

The most general renormalizeable Kahler potential of a chiral superfield is $K(\Phi, \Phi^\dagger) = \Phi_k^\dagger \Phi_k$. The Kahler potential is real and W is a holomorphic polynomial, i.e. it does not depend on $\Phi^\dagger$ (for renormalizable theories the degree of W is at most 3).

The Lagrangian density for the chiral multiplet can then be written as

$$\mathcal{L}_{\text{chiral}} = D_\mu\phi_i^* D^\mu\phi_i + i\psi_i^\dagger\bar{\sigma}^\mu D_\mu\psi_i - \frac{1}{2}(W_{ij}\psi_i\psi_j + W_{ij}^*\psi_i^\dagger\psi_j^\dagger) - W_i^* W_i \tag{2.9}$$

where

$$W_i = \frac{\partial W}{\partial\phi} \tag{2.10}$$

$$W_{ij} = \frac{\partial^2 W}{\partial\phi_i\partial\phi_j} \tag{2.11}$$

The $|W_i|^2$ term $\mathcal{L}_{\text{chiral}}$ is called F-term.

The Lagrangian for gauge supermultiplets is

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4} F^a_{\mu\nu} F^{\mu\nu a} - i\lambda^{a\dagger}\bar{\sigma}^\mu D_\mu \lambda^a + \frac{1}{2} D^a D^a \tag{2.12}$$

where

$$F^a_{\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu + g f^{abc} A^b_\mu A^c_\nu \tag{2.13}$$

is the field strength tensor of the gauge field. If the gauge group is the product of several simple groups – like in the SM – there is more than one gauge coupling. Using a suitable set of generators, this effectively has the consequence that g will depend on a as well. D^a is an auxiliary field introduced for technical reasons. Its kinetic term does not have any derivatives. Therefore, the equations of motion are purely algebraical and will be used later to substitute the field, once the complete Lagrangian is known.

Finally, there can be interactions between gauge and chiral multiplets apart from the ones coming from the covariant derivative. This is excpected since the covariant derivative only couples chiral multiplets to gauge bosons. The additional terms are "supersymmetrized" versions of these terms, coupling gauginos and the auxiliary field to chiral multiplets. The only ones allowed by gauge invariance and renormalizability are

$$(\phi_i t^a_{ij} \psi_j)\lambda^a, \quad \lambda^{a\dagger}(\psi^\dagger_i t^a_{ij} \phi_j) \quad \text{and} \quad (\phi_i t^a_{ij} \phi_j) D^a. \tag{2.14}$$

With this, the complete Lagrangian can be written down:

$$\mathcal{L} = \mathcal{L}_{\text{chiral}} + \mathcal{L}_{\text{gauge}} - \sqrt{2} g (\phi_i t^a_{ij} \psi_j)\lambda^a - \sqrt{2} g \lambda^{a\dagger}(\psi^\dagger_i t^a_{ij} \phi_j) + g(\phi_i t^a_{ij} \phi_j) D^a \tag{2.15}$$

where the couplings of the additional terms are fixed by supersymmetry.

The last thing to do is to eliminate the auxiliary field using its equation of motion

$$\frac{\partial \mathcal{L}}{\partial D^a} = 0 \tag{2.16}$$

$$\iff D^a = -g(\phi^*_i t^a_{ij} \phi_j) \tag{2.17}$$

The scalar interaction originating from this are called *D-terms*.

The symmetry between bosons and fermions provides a solution of the hierarchy problem as mentioned above. For each loop contribution to the Higgs mass parameter there appears a loop of the superpartner that has the same quadratic divergence except for a minus sign that comes from the closed fermion loop. Therefore, the quadratic terms cancel, leaving only logarithmic contributions which do not cause a problem. This result can rigorously be proven to all orders in perturbation theory.

2.1.1 Supersymmetry breaking

Supersymmetric Lagrangians constructed as described in the previous section do not serve as exact descriptions of nature. Otherwise superpartners of SM fermions would have been detected long ago since superpartners always have the same mass. The conclusion is that supersymmetry is not an exact symmetry of nature. The natural attempt is to break SUSY spontaneously, but it turns out to be difficult to give phenomenologically viable models in which this works. In any event, it requires an extension of the MSSM. A common approach is therefore to to simply add interactions that *explicitly* break SUSY *softly* without specifing where they come from. *Softly* means that they should be renormalizable and cause no quadratic divergences to scalar particles. Soft terms for such interactions are [2]

- gaugino mass terms $m_a\lambda^a\lambda^a$,
- scalar mass terms $m^2_{ij}\phi^*_i\phi_j$ and $b_{ij}\phi_i\phi_j$,
- cubic scalar couplings $a_{ijk}\phi_i\phi_j\phi_k$ (called A-*terms*) and $c_{ijk}\phi^*_i\phi_j\phi_k$,
- linear terms $t_i\phi_i$.

Other possibilities like fermion mass terms $\psi_i\psi_j$ can be absorbed into a redefinition of the parameters mentioned above and the couplings in the superpotential. Further restrictions can be made based on gauge invariance or other symmetries, given a specific model.

It is worth mentioning that all coupling constants in the SUSY breaking Lagrangian have non-zero mass dimension. One expects that there is a characteristic scale m_{soft} for these terms. The scale of SUSY breaking should not be too far away from the EW scale.

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2.2 The Minimal Supersymmetric Standard Model

The MSSM is the smallest possible supersymmetric extension of the SM. It is based on the same gauge group. The particle content is extended in essentially two different ways. For each standard model particle, a superpartner is introduced [2]. Since gauginos are in the adjoint representation, all SM fermions have to reside in chiral supermultiplets. Their scalar superpartners are called *sleptons* and *squarks*. It should be noted that each *Weyl* spinor has a scalar superpartner, so that there are, for example, two up-squarks. Just like the fermions, they are called left- and right-handed, but this label only refers to their gauge transformation properties and is not connected to angular momentum in any way. The scalar superpartners will be denoted with upper case letters: Q, U, D, L, E.

The superpartners of the SM bosons are named by adding the ending "-ino" to the name of the SM particles: higgsino, gluino, wino and bino, for the gauge eigenstates. Their fermionic superpartners are denoted by putting a tilde over the respective boson field: $\tilde{W}, \tilde{B}, \tilde{G}, \ldots$

In addition to this it is necessary to introduce a second Higgs (super-)multiplet. In the SM, one Higgs doublet suffices to generate all fermion masses by the Yukawa couplings. In a supersymmetric model, however, the last term is forbidden since the superpotential would have to contain a term H^*QU, which is not holomorphic in H. Therefore, another Higgs multiplet with $Y = 1/2$ is needed. It is denoted by H_u, while the $Y = -1/2$ doublet is called H_d.

The most general superpotential for the given superfields which is allowed by gauge invariance would include terms like LQD or UDD, which violate lepton or baryon number and are phenomenologically strongly constrained. To solve this problem, a quantity called *R-parity* is introduced [7]. It is defined as

$$R = (-1)^{3(B-L)+2s} , \tag{2.18}$$

where B and L are baryon and lepton number, respectively, and s is the spin. It is constructed in such a way that all SM particles and all scalar Higgs fields have $R = +1$, while all superpartners of these fields, differing in spin by 1/2, have $R = -1$. The latter ones are collectively called *sparticles*. The MSSM is then *required* to conserve R, which means that the product of R-parities of the fields in each interaction vertex must be $+1$. An consequence of this is that each vertex must contain an even number of sparticles.

The most general superpotential allowed by gauge invariance, R-parity and renormalizability is

$$W_{\text{MSSM}} = (Y_u)_{ij} Q_i H_u U_j + (Y_d)_{ij} Q_i H_d D_j + (Y_e)_{ij} L_i H_d E_j + \mu H_u H_d \tag{2.19}$$

All terms inducing B or L violation are now absent. In order to complete the model, one has to write down the soft SUSY breaking Lagrangian. It can be found for example in [3].

2.2.1 The μ-Problem

The MSSM has a conceptual problem lying in the μ parameter, the only parameter outside of $\mathcal{L}_{\text{soft}}$ having non-zero mass dimension. Both the dimensional couplings from the soft breaking Lagrangian and μ enter the Higgs potential and thus determine the vacuum expectation values (VEVs) of H_u and H_d. But in contrast to the other parameters, μ originated in the supersymmetry respecting part of the Lagrangian. There are two natural assumptions concerning the size of μ: it may either be zero by virtue of some additional symmtery, or it is expected to be large, somewhere around the GUT scale. The first case is excluded for phenomenological reasons, since it would cause masses of several sparticles be unacceptably small. So μ should be large. But then, without cancellations between the various contributions to the Higgs potential – which can not be justified by any mechanism inside the MSSM – the VEVs should also come out pretty large, much higher than the electroweak scale, clearly contradicting experimental results. This is only a formal problem, since one can manually adjust μ to whatever value necessary, but it makes electroweak symmtery breaking in the MSSM appear unnatural.

2.3 Masses, mixings and interactions

Review of the MSSM Higgs sector

The MSSM Higgs potential

The Higgs part of the MSSM superpotential is

$$\int d^2\theta \mu \mathcal{H}_u \mathcal{H}_d + \text{h.c.} \tag{2.20}$$

The superfields $\mathcal{H}_{u,d}$ can be written in component fields, where $H_{u,d}$ are scalars, $\tilde{H}_{u,d}$ are spinors and $F_{u,d}$ are auxiliary fields

$$\begin{aligned}\mathcal{H}_u &= H_u(y) + \sqrt{2}\theta \tilde{H}_u(y) + \theta\theta F_u(y) \\ &= H_u(x) + i\theta\sigma^m\bar{\theta}\partial_m H_u(x) + \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\Box H_u(x) + \sqrt{2}\theta\tilde{H}_u(x) \\ &- \frac{i}{\sqrt{2}}\theta\theta\partial_m\tilde{H}_u(x)\sigma^m\bar{\theta} + \theta\theta F_u(x)\end{aligned} \tag{2.21}$$

$$\begin{aligned}\mathcal{H}_d &= H_d(x) + i\theta\sigma^m\bar{\theta}\partial_m H_d(x) + \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\Box H_d(x) + \sqrt{2}\theta\tilde{H}_d(x) \\ &- \frac{i}{\sqrt{2}}\theta\theta\partial_m\tilde{H}_d(x)\sigma^m\bar{\theta} + \theta\theta F_d(x)\end{aligned} \tag{2.22}$$

where we expanded the fields around x_m. Inserting this into the $\mu\mathcal{H}_u\mathcal{H}_d$-term of the Lagrangian and ignoring the spinor fields $\tilde{H}_{u,d}$ (we only need the scalar potential), we obtain

$$L_W = \mu \int d^2\theta \left(H_u + \theta^2 F_u\right)\left(H_d + \theta^2 F_d\right) + \text{h.c.} = \mu(H_u F_d + F_u H_d) + \text{h.c.} \tag{2.23}$$

The Kahler potential has the following form

$$\int d^4\theta \left(\mathcal{H}_u^\dagger \mathcal{H}_u + \mathcal{H}_d^\dagger \mathcal{H}_d\right) \tag{2.24}$$

We are interested in the kinetic terms for $F_{u,d}$ (the kinetic terms for $H_{u,d}$ are the ususal scalar field kinetic terms $\partial^\mu \phi^\dagger \partial_\mu \phi$). One obtains:

$$(L_K)_F = F_u(x)^\dagger F_u(x) + F_d(x)^\dagger F_d(x) \tag{2.25}$$

The Euler equation for the F fields delivers us the equation of motion

$$\frac{d}{dx_\mu}\frac{\partial L}{\partial(\partial_\mu F_i)} - \frac{\partial L}{\partial F_i} = 0 \tag{2.26}$$

$$\frac{\partial(L_W + L_K)}{\partial F_u} = F_u^\dagger + \mu H_d = 0 \tag{2.27}$$

$$\frac{\partial(L_W + L_K)}{\partial F_d} = F_d^\dagger + \mu H_u = 0 \tag{2.28}$$

These can be solved analytically for $F_{u,d}$ and plugged back into the Lagrangian to finally yield the so-called F-terms of the scalar potential. In addition, the scalar potential contains contributions from the soft Lagrangian and the D-terms. The final MSSM Higgs potential is:

$$\begin{aligned} V &= \tilde{m}^2_{H_u} {H_u}^\dagger H_u + \tilde{m}^2_{H_d} {H_d}^\dagger H_d - m^2_{ud}(H_u H_d + \text{h.c.}) \\ &+ \frac{g^2}{8}\left[({H_u}^\dagger H_u + {H_d}^\dagger H_d)^2 - 4(H_u H_d)^*(H_u H_d)\right] + \frac{g'^2}{8}({H_u}^\dagger H_u - {H_d}^\dagger H_d)^2 \\ &= \tilde{m}^2_{H_u} {H_u}^\dagger H_u + \tilde{m}^2_{H_d} {H_d}^\dagger H_d - m^2_{ud}(H_u H_d + \text{h.c.}) \\ &+ \frac{g^2 + g'^2}{8}\left[({H_u}^\dagger H_u)^2 + ({H_d}^\dagger H_d)^2\right] + \frac{g^2 - g'^2}{4}({H_u}^\dagger H_u)({H_d}^\dagger H_d) - \frac{g^2}{2}|H_u H_d|^2 \end{aligned} \tag{2.29}$$

where m^2_{ud} is a soft paramter and $\tilde{m}^2_{H_u} = |\mu|^2 + m^2_{H_u}$ and $\tilde{m}^2_{H_d} = |\mu|^2 + m^2_{H_d}$ containing Higgs soft masses $m^2_{H_u}$,$m^2_{H_d}$ [2].

Spontaneous symmetry breaking

We use the following notation and sign convention:

$$H_u = \begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix}, \quad H_d = \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix} \tag{2.30}$$

$$H_u H_d = -H_u^+ H_d^- + H_u^0 H_d^0 \tag{2.31}$$

The neutral fields H_u^0 and H_d^0 develop vevs

$$\langle H_u^0 \rangle = v_u, \quad \langle H_d^0 \rangle = v_d. \tag{2.32}$$

In the following, we will try to substitue these vevs by more conventional variables like the Z^0 mass and the angle *beta* using the relations

$$v_u^2 + v_d^2 = v^2 \equiv \frac{2m_Z^2}{g^2 + g'^2} \tag{2.33}$$

$$\frac{v_u}{v_d} \equiv \tan\beta \tag{2.34}$$

$$\Rightarrow v_u = v\sin\beta, \quad v_d = v\cos\beta \tag{2.35}$$

$$\Rightarrow v_d^2 - v_u^2 = v^2\cos^2\beta - v^2\sin^2\beta = v^2\cos(2\beta). \tag{2.36}$$

v_u and v_d can be expressed by the potential parameters. One substitutes the fields in V with their vevs:

$$\begin{aligned} V &= \tilde{m}_{H_u}^2 v_u^2 + \tilde{m}_{H_d}^2 v_d^2 - 2m_{ud}^2 v_u v_d + \frac{g^2 + g'^2}{8}(v_u^4 + v_d^4) + \frac{g^2 - g'^2}{4} v_u^2 v_d^2 - \frac{g^2}{4} v_u^2 v_d^2 \\ &= \tilde{m}_{H_u}^2 v_u^2 + \tilde{m}_{H_d}^2 v_d^2 - 2m_{ud}^2 v_u v_d + \frac{g^2 + g'^2}{8}(v_u^4 + v_d^4) - \frac{g^2 + g'^2}{4} v_u^2 v_d^2 \\ &= \tilde{m}_{H_u}^2 v_u^2 + \tilde{m}_{H_d}^2 v_d^2 - 2m_{ud}^2 v_u v_d + \frac{g^2 + g'^2}{8}(v_u^2 - v_d^2)^2 \end{aligned} \tag{2.37}$$

If we minimize the potential by $\frac{\partial V}{\partial \phi_i} = 0$ we obtain the vacuum expectation values of H_u^0 and H_d^0. The second derivative $\frac{\partial^2 V}{\partial \phi_i \partial \phi_j}$ leads us to entries of the mass matrix and therefore respectivly to the mass eigenstates und eigenvalues which are considered to be the physical masses and fields.

The solution equation are

$$\frac{\partial V}{\partial v_u} = \frac{\partial V}{\partial v_d} = 0, \tag{2.38}$$

this means

$$2\tilde{m}_{H_u}^2 v_u - 2m_{ud}^2 v_d + \frac{g^2 + g'^2}{2}(v_u^2 - v_d^2) v_u = 0 \tag{2.39}$$

$$2\tilde{m}_{H_d}^2 v_d - 2m_{ud}^2 v_u - \frac{g^2 + g'^2}{2}(v_u^2 - v_d^2) v_d = 0, \tag{2.40}$$

respectively

$$\tilde{m}_{H_u}^2 - m_{ud}^2 \cot\beta - \frac{1}{2} m_Z^2 \cos(2\beta) = 0 \tag{2.41}$$

$$\tilde{m}_{H_d}^2 - m_{ud}^2 \tan\beta + \frac{1}{2} m_Z^2 \cos(2\beta) = 0 \tag{2.42}$$

Expanding the fields around their vevs leads to

$$\begin{aligned} H_u^0 &= v_u + \frac{1}{\sqrt{2}} h_u + \frac{i}{\sqrt{2}} g_u \\ H_d^0 &= v_d + \frac{1}{\sqrt{2}} h_d + \frac{i}{\sqrt{2}} g_d \end{aligned} \tag{2.43}$$

$h_{u,d}$ and $g_{u,d}$ are real fields. The charged fields $H_u^+$and H_d^- do not develop vevs. In order to insert the expansion (2.43) into the potential, the various potential terms will be considered seperately at first (where constant terms will be dropped):

$$H_u{}^\dagger H_u = v_u^2 + |H_u^+|^2 + \frac{1}{2}h_u^2 + \frac{1}{2}g_u^2 + \sqrt{2}v_u h_u \tag{2.44}$$

$$H_d{}^\dagger H_d = v_d^2 + |H_d^-|^2 + \frac{1}{2}h_d^2 + \frac{1}{2}g_d^2 + \sqrt{2}v_d h_d \tag{2.45}$$

$$H_u H_d + \text{h.c.} = 2v_u v_d - H_u^+ H_d^- - H_u^{+*} H_d^{-*} + h_u h_d - g_u g_d + \sqrt{2}v_u h_d + \sqrt{2}v_d h_u. \tag{2.46}$$

For the higher powers, one obtains:

$$\begin{aligned}
(H_u{}^\dagger H_u)(H_d{}^\dagger H_d) &= (|H_u^+|^2 + (v_u + \frac{h_u}{\sqrt{2}})^2 + \frac{1}{2}g_u^2)(|H_d^-|^2 + (v_d + \frac{h_d}{\sqrt{2}})^2 + \frac{1}{2}g_d^2) \\
&= \sqrt{2}v_u v_d^2 h_u + \sqrt{2}v_u^2 v_d h_d + 2v_u v_d h_u h_d \\
&+ v_u^2 H_d^{-*} H_d^- + \frac{v_u^2}{2}(h_d^2 + g_d^2) + v_d^2 H_u^{+*} H_u^+ + \frac{v_d^2}{2}(h_u^2 + g_u^2) \\
&+ \sqrt{2}v_u h_u H_d^{-*} H_d^- + \frac{v_u}{\sqrt{2}}(h_u h_d^2 + h_u g_d^2) \\
&+ \sqrt{2}v_d h_d H_u^{+*} H_u^+ + \frac{v_d}{\sqrt{2}}(h_d h_u^2 + h_d g_u^2) \\
&+ H_u^{+*} H_u^+ H_d^{-*} H_d^- + \frac{1}{2}H_u^{+*} H_u^+ h_d^2 + \frac{1}{2}H_u^{+*} H_u^+ g_d^2 \\
&+ \frac{1}{2}H_d^{-*} H_d^- h_u^2 + \frac{1}{2}H_d^{-*} H_d^- g_u^2 \\
&+ \frac{1}{4}(h_u^2 h_d^2 + h_u^2 g_d^2 + g_u^2 h_d^2 + g_u^2 g_d^2)
\end{aligned} \tag{2.47}$$

From this, $(H_u{}^\dagger H_u)^2$ and $(H_d{}^\dagger H_d)^2$ can be read off by replacing all subscripts $d \to u$ and

$u \to d$, respectively (and of course the charges of the fields, i.e. $H_u^+ \leftrightarrow H_d^-$). Finally:

$$
\begin{aligned}
|H_u H_d|^2 &= |H_u^+ H_d^-|^2 + |H_u^0 H_d^0|^2 - (H_u^{+*} H_d^{-*} H_u^0 H_d^0 + \text{h.c.}) \\
&= |H_u^+ H_d^-|^2 + ((v_u + \frac{h_u}{\sqrt{2}})^2 + \frac{1}{2} g_u^2)((v_d + \frac{h_d}{\sqrt{2}})^2 + \frac{1}{2} g_d^2) \\
&- \Big(v_u v_d H_u^{+*} H_d^{-*} + \frac{v_u}{\sqrt{2}} H_u^{+*} H_d^{-*} h_d + i \frac{v_u}{\sqrt{2}} H_u^{+*} H_d^{-*} g_d \\
&\quad + \frac{v_d}{\sqrt{2}} H_u^{+*} H_d^{-*} h_u + i \frac{v_d}{\sqrt{2}} H_u^{+*} H_d^{-*} g_u + \frac{1}{2} H_u^{+*} H_d^{-*} h_u h_d \\
&\quad + \frac{1}{2} H_u^{+*} H_d^{-*} h_u g_d + \frac{1}{2} H_u^{+*} H_d^{-*} g_u h_d + \frac{1}{2} H_u^{+*} H_d^{-*} g_u g_d + \text{h.c.} \Big) \\
&= H_u^{+*} H_d^{-*} H_u^+ H_d^- + \sqrt{2} v_u^2 v_d h_d + \sqrt{2} v_u v_d^2 h_u \\
&+ \frac{v_u^2}{2} (h_d^2 + g_d^2) + \frac{v_d^2}{2} (h_u^2 + g_u^2) + 2 v_u v_d h_u h_d \\
&+ \frac{v_u}{\sqrt{2}} (h_u h_d^2 + h_u g_d^2) + \frac{v_d}{\sqrt{2}} (h_d h_u^2 + h_d g_u^2) \\
&+ \frac{1}{4} (h_u^2 h_d^2 + h_u^2 g_d^2 + g_u^2 h_d^2 + g_u^2 g_d^2) \\
&- \Big(v_u v_d H_u^{+*} H_d^{-*} + \frac{v_u}{\sqrt{2}} H_u^{+*} H_d^{-*} h_d + i \frac{v_u}{\sqrt{2}} H_u^{+*} H_d^{-*} g_d \\
&\quad + \frac{v_d}{\sqrt{2}} H_u^{+*} H_d^{-*} h_u + i \frac{v_d}{\sqrt{2}} H_u^{+*} H_d^{-*} g_u + \frac{1}{2} H_u^{+*} H_d^{-*} h_u h_d \\
&\quad + \frac{1}{2} H_u^{+*} H_d^{-*} h_u g_d + \frac{1}{2} H_u^{+*} H_d^{-*} g_u h_d + \frac{1}{2} H_u^{+*} H_d^{-*} g_u g_d + \text{h.c.} \Big)
\end{aligned} \tag{2.48}
$$

The whole potential can be organized as follows:

$$V = V_{\text{const}} + V_{\text{lin}} + V_{\text{quad}} + V_{\text{tert}} + V_{\text{quart}} \tag{2.49}$$

where

$$V_{\text{const}} = \tilde{m}_{H_u}^2 v_u^{\dagger 2} + \tilde{m}_{H_u}^2 v_d^{\dagger 2} - m_{ud}^2 2 v_u v_d \tag{2.50}$$

$$
\begin{aligned}
V_{\text{lin}} &= \tilde{m}_{H_u}^2 \sqrt{2} v_u h_u + \tilde{m}_{H_d}^2 \sqrt{2} v_d h_d - m_{ud}^2 (\sqrt{2} v_u h_u + \sqrt{2} v_d h_d) + \frac{g^2 + g'^2}{8} \Big(\sqrt{2} v_u v_u^2 h_u \\
&+ \sqrt{2} v_u^2 v_u h_u \Big) + \frac{g^2 + g'^2}{8} \Big(\sqrt{2} v_d v_d^2 h_d + \sqrt{2} v_d^2 v_d h_d \Big) + \frac{g^2 - g'^2}{4} \Big(\sqrt{2} v_u v_d^2 h_u + \sqrt{2} v_u^2 v_d h_d \Big) \\
&- \frac{g^2}{2} \Big(\sqrt{2} v_u^2 v_d h_d + \sqrt{2} v_u v_d^2 h_u \Big)
\end{aligned} \tag{2.51}
$$

$$
\begin{aligned}
V_{\text{quad}} &= \tilde{m}^2_{H_u}\left(|H_u^+|^2 + \frac{1}{2}h_u^2 + \frac{1}{2}g_u^2\right) + \tilde{m}^2_{H_d}\left(|H_d^-|^2 + \frac{1}{2}h_d^2 + \frac{1}{2}g_d^2\right) \\
&- m^2_{ud}\left(h_u h_d - g_u g_d - H_u^{+*}H_d^{-*} - H_u^+ H_d^-\right) \\
&+ \frac{g^2+g'^2}{8}\left(3v_u^2 h_u^2 + v_u^2 g_u^2 + 2v_u^2 H_u^{+*}H_u^+ + 3v_d^2 h_d^2 + v_d^2 g_d^2 + 2v_d^2 H_d^{-*}H_d^-\right) \\
&+ \frac{g^2-g'^2}{4}\Bigg(2v_u v_d h_u h_d + v_u^2 H_d^{-*}H_d^- + \frac{v_u^2}{2}(h_d^2+g_d^2) \\
&+ v_d^2 H_u^{+*}H_u^+ + \frac{v_d^2}{2}(h_u^2+g_u^2)\Bigg) - \frac{g^2}{2}\Bigg(\frac{v_u^2}{2}(h_d^2+g_d^2) + \frac{v_d^2}{2}(h_u^2+g_u^2) \\
&+ 2v_u v_d h_u h_d - v_u v_d H_u^{+*}H_d^{-*} - v_u v_d H_u^+ H_d^-\Bigg)
\end{aligned} \tag{2.52}
$$

$$
\begin{aligned}
&= \left(\tilde{m}^2_{H_u} + \frac{g^2+g'^2}{4}2v_u^2 + \frac{g^2-g'^2}{4}v_d^2\right)|H_u^+|^2 \\
&+ \left(\frac{\tilde{m}^2_{H_u}}{2} + \frac{g^2+g'^2}{8}3v_u^2 - \frac{g^2+g'^2}{8}v_d^2\right)h_u^2 \\
&+ \left(\frac{\tilde{m}^2_{H_u}}{2} + \frac{g^2+g'^2}{8}v_u^2 - \frac{g^2+g'^2}{8}v_d^2\right)g_u^2 \\
&+ \left(\tilde{m}^2_{H_d} + \frac{g^2+g'^2}{4}2v_d^2 + \frac{g^2-g'^2}{4}v_u^2\right)|H_d^-|^2 \\
&+ \left(\frac{\tilde{m}^2_{H_d}}{2} + \frac{g^2+g'^2}{8}3v_d^2 - \frac{g^2+g'^2}{8}v_u^2\right)h_d^2 \\
&+ \left(\frac{\tilde{m}^2_{H_d}}{2} + \frac{g^2+g'^2}{8}v_d^2 - \frac{g^2+g'^2}{8}v_u^2\right)g_d^2 \\
&+ \left(-m_u d^2 + \frac{g^2+g'^2}{4}v_u v_d\right)h_u h_d \\
&+ \left(m_u d^2 g_u g_d + \frac{g^2}{2}\right)H_u^{+*}H_d^{-*} \\
&+ \left(m_u d^2 + \frac{g^2}{2}v_u v_d\right)H_u^+ H_d^-
\end{aligned}
$$

$$
\begin{aligned}
V_{\text{tert}} = \frac{g^2+g'^2}{\sqrt{8}}&\Bigg(v_u H_u^{+*}H_u^+ h_u + v_d H_d^{-*}H_d^- h_d + \frac{v_u}{2}h_u^3 + \frac{v_u}{2}h_u g_u^2 \\
&+ \frac{v_d}{2}h_d^3 + \frac{v_d}{2}h_d g_d^2 - \frac{v_u}{2}h_u h_d^2 - \frac{v_u}{2}h_u g_d^2 - \frac{v_d}{2}h_d h_u^2 - \frac{v_d}{2}h_d g_u^2\Bigg) \\
+ \frac{g^2-g'^2}{\sqrt{8}}&\left(v_u H_d^{-*}H_d^- h_u + v_d H_u^{+*}H_u^+ h_d\right) \\
+ \frac{g^2}{\sqrt{8}}&\left(v_u H_u^+ H_d^- h_d - i v_u H_u^+ H_d^- g_d + v_d H_u^+ H_d^- h_u - i v_d H_u^+ H_d^- g_u + \text{h.c.}\right)
\end{aligned} \tag{2.53}
$$

$$
\begin{aligned}
V_{\text{quart}} = \frac{g^2+g'^2}{8} &\Big((H_u^{+*}H_u^+)^2 + (H_d^{-*}H_d^-)^2 + H_u^{+*}H_u^+ h_u^2 + H_d^{-*}H_d^- h_d^2 \\
&+ H_u^{+*}H_u^+ g_u^2 + H_d^{-*}H_d^- g_d^2 + \frac{1}{4}h_u^4 + \frac{1}{4}g_u^4 + \frac{1}{2}h_u^2 g_u^2 + \frac{1}{4}h_d^4 + \frac{1}{4}g_d^4 + \frac{1}{2}h_d^2 g_d^2 \\
&- 2H_u^{+*}H_u^+ H_d^{-*}H_d^- - \frac{1}{2}h_u^2 h_d^2 - \frac{1}{2}h_u^2 g_d^2 - \frac{1}{2}g_u^2 h_d^2 - \frac{1}{2}g_u^2 g_d^2 \Big) \\
+ \frac{g^2-g'^2}{8} &\left(H_u^{+*}H_u^+ h_d^2 + H_u^{+*}H_u^+ g_d^2 + H_d^{-*}H_d^- h_u^2 + H_d^{-*}H_d^- g_u^2 \right) \\
+ \frac{g^2}{4} &\left(H_u^+ H_d^- h_u h_d + H_u^+ H_d^- h_u g_d + H_u^+ H_d^- g_u h_d + H_u^+ H_d^- g_u g_d + \text{h.c.} \right)
\end{aligned}
$$

The MSSM Higgs spectrum

As a consequence of the unbroken electromagnetic gauge symmetry there exists no mixture between the charged and neutral fields, and due to CP symmetry there is no mixture between the real parts of and the complex parts of $H^0_{u,d}$. The various fields will be regarded separately.

Neutral fields

The potential for $g_{u,d}$ is

$$
\frac{1}{2}\begin{pmatrix} g_u \\ g_d \end{pmatrix}^T \begin{pmatrix} \tilde{m}^2_{H_u} + \frac{g^2+g'^2}{4}(v_u^2 - v_d^2) & m^2_{ud} \\ m^2_{ud} & \tilde{m}^2_{H_d} + \frac{g^2+g'^2}{4}(v_d^2 - v_u^2) \end{pmatrix} \begin{pmatrix} g_u \\ g_d \end{pmatrix} \tag{2.54}
$$

The mass matrix is

$$
\begin{pmatrix} \tilde{m}^2_{H_u} - \frac{m_Z^2}{2}\cos(2\beta) & m^2_{ud} \\ m^2_{ud} & \tilde{m}^2_{H_d} + \frac{m_Z^2}{2}\cos(2\beta) \end{pmatrix} = \begin{pmatrix} m^2_{ud}\cot\beta & m^2_{ud} \\ m^2_{ud} & m^2_{ud}\tan\beta \end{pmatrix}, \tag{2.55}
$$

One of the mass eigenstates is a Goldstone Boson, as a consequence the corresponding eigenvalue is 0, which leads to a determinant equal to 0. In this case one can read off the other eigenvalue as the sum of the diagonal elements.

$$
m_A^2 = \frac{m^2_{ud}}{\sin\beta\cos\beta} = \frac{2m^2_{ud}}{\sin(2\beta)}. \tag{2.56}
$$

The eigenstates are (G^0 is the Goldstone):

$$
A^0 = \cos\beta\, g_u + \sin\beta\, g_d \tag{2.57}
$$

$$
G^0 = \sin\beta\, g_u - \cos\beta\, g_d. \tag{2.58}
$$

The gauge eigenstates can be written in terms of the mass eigenstates as

$$
g_u = \cos\beta\, A^0 - \sin\beta\, G^0 \tag{2.59}
$$

$$
g_d = \sin\beta\, A^0 + \cos\beta\, G^0 \tag{2.60}
$$

Usually all values will be expressed by m_Z^2, β and m_A^2. The mass matrix for $h_{u,d}$ can be read to:

$$\begin{aligned}
&\begin{pmatrix} \tilde{m}_{H_u}^2 + \frac{g^2+g'^2}{4}(3v_u^2 - v_d^2) & -m_{ud}^2 - \frac{g^2+g'^2}{2} v_u v_d \\ -m_{ud}^2 - \frac{g^2+g'^2}{2} v_u v_d & \tilde{m}_{H_d}^2 + \frac{g^2+g'^2}{4}(3v_d^2 - v_u^2) \end{pmatrix} \qquad (2.61) \\
=&\begin{pmatrix} \tilde{m}_{H_u}^2 - \frac{m_Z^2}{2}\cos(2\beta) + \frac{g^2+g'^2}{2} v_u^2 & m_{ud}^2 + \frac{g^2+g'^2}{2} v_u v_d \\ m_{ud}^2 + \frac{g^2+g'^2}{2} v_u v_d & \tilde{m}_{H_d}^2 + \frac{m_Z^2}{2}\cos(2\beta) + \frac{g^2+g'^2}{2} v_d^2 \end{pmatrix} \\
=&\begin{pmatrix} m_{ud}^2 \cot\beta + m_Z^2 \sin^2\beta & m_{ud}^2 + m_Z^2 \sin\beta\cos\beta \\ m_{ud}^2 + m_Z^2 \sin\beta\cos\beta & m_{ud}^2 \tan\beta + m_Z^2 \cos^2\beta \end{pmatrix} \\
=&\begin{pmatrix} m_A^2 \cos^2\beta + m_Z^2 \sin^2\beta & (m_Z^2 + m_A^2)\sin\beta\cos\beta \\ (m_Z^2 + m_A^2)\sin\beta\cos\beta & m_A^2 \sin^2\beta + m_Z^2 \cos^2\beta \end{pmatrix}
\end{aligned}$$

Using the formula for eigenvalues of a symmetric matrix

$$A = \begin{pmatrix} a & c \\ c & b \end{pmatrix} \qquad (2.62)$$

$$\lambda = \frac{1}{2}\left(a + b \pm \sqrt{(a-b)^2 + 4c^2}\right). \qquad (2.63)$$

Applying some addition theorems one obtains

$$m_{h^0,H^0}^2 = \frac{1}{2}\left(m_A^2 + m_Z^2 \mp \sqrt{(m_A^2 - m_Z^2)^2 + 4m_A^2 m_Z^2 \sin^2(2\beta)}\right), \qquad (2.64)$$

(where $m_{h^0}^2$ is the smaller eigenvalue) with eigenstates

$$h^0 \propto \left(\sqrt{(m_A^2 - m_Z^2)^2 + 4m_A^2 m_Z^2 \sin^2(2\beta)} - (m_A^2 - m_Z^2)\cos(2\beta)\right) h_u - (m_A^2 + m_Z^2)\sin(2\beta) h_d \qquad (2.65)$$

$$H^0 \propto \left(\sqrt{(m_A^2 - m_Z^2)^2 + 4m_A^2 m_Z^2 \sin^2(2\beta)} + (m_A^2 - m_Z^2)\cos(2\beta)\right) h_d + (m_A^2 + m_Z^2)\sin(2\beta) h_d \qquad (2.66)$$

For notational simplicity, normalization factors were omitted. This can be written as

$$h^0 = \cos\alpha\, h_u - \sin\alpha\, h_d \qquad (2.67)$$

$$H^0 = \sin\alpha\, h_u + \cos\alpha\, h_d \qquad (2.68)$$

The mixing angle α fulfills

$$\frac{\sin 2\alpha}{\sin 2\beta} = -\frac{m_{H^0}^2 + m_{h^0}^2}{m_{H^0}^2 - m_{h^0}^2} \qquad (2.69)$$

$$\frac{\tan 2\alpha}{\tan 2\beta} = \frac{m_A^2 + m_Z^2}{m_A^2 - m_Z^2}. \qquad (2.70)$$

Solving for the gauge eigenstates yields

$$h_u = \cos\alpha\, h^0 + \sin\alpha\, H^0 \qquad (2.71)$$

$$h_d = -\sin\alpha\, h^0 + \cos\alpha\, H^0 \qquad (2.72)$$

Charged fields

Finally one obtains the mass matrix for the charged fields:

$$\begin{pmatrix} H_u^+ \\ H_d^{-*} \end{pmatrix}^\dagger \begin{pmatrix} \tilde{m}_{H_u}^2 + \frac{g^2+g'^2}{4}v_u^2 + \frac{g^2-g'^2}{4}v_d^2 & m_{ud}^2 + \frac{g^2}{2}v_u v_d \\ m_{ud}^2 + \frac{g^2}{2}v_u v_d & \tilde{m}_{H_d}^2 + \frac{g^2+g'^2}{4}v_d^2 + \frac{g^2-g'^2}{4}v_u^2 \end{pmatrix} \begin{pmatrix} H_u^+ \\ H_d^{-*} \end{pmatrix} \tag{2.73}$$

To simplify the matrix one has to use the relation between $m_W^2$and v:

$$m_W^2 = \frac{g^2 v^2}{2} \tag{2.74}$$

Now one can write the mass matrix as

$$\begin{pmatrix} \tilde{m}_{H_u}^2 + \frac{g^2+g'^2}{4}(v_u^2 - v_d^2) + \frac{g^2}{2}v_d^2 & m_{ud}^2 + \frac{g^2}{2}v_u v_d \\ m_{ud}^2 + \frac{g^2}{2}v_u v_d & \tilde{m}_{H_d}^2 + \frac{g^2+g'^2}{4}(v_d^2 - v_u^2) + \frac{g^2}{2}v_u^2 \end{pmatrix} \tag{2.75}$$

$$= \begin{pmatrix} \tilde{m}_{H_u}^2 - \frac{m_Z^2}{2}\cos(2\beta) + m_W^2\cos^2\beta & m_{ud}^2 + m_W^2\cos\beta\sin\beta \\ m_{ud}^2 + m_W^2\cos\beta\sin\beta & \tilde{m}_{H_d}^2 + \frac{m_Z^2}{2}\cos(2\beta) + m_W^2\sin^2\beta \end{pmatrix} \tag{2.76}$$

$$= \begin{pmatrix} (m_W^2 + m_A^2)\cos^2\beta & (m_W^2 + m_A^2)\cos\beta\sin\beta \\ (m_W^2 + m_A^2)\cos\beta\sin\beta & (m_W^2 + m_A^2)\sin^2\beta \end{pmatrix}$$

One eigenvalue is 0 again, which corresponds to a charged complex goldstone boson. By calculating the trace one obtains

$$m_{H^\pm}^2 = m_W^2 + m_A^2 \tag{2.77}$$

for the other eigenvalue. The corresponding eigenstates are

$$H^+ = \cos\beta\, H_u^+ + \sin\beta\, H_d^{-*} \tag{2.78}$$

$$G^+ = \sin\beta\, H_u^+ - \cos\beta\, H_d^{-*} \tag{2.79}$$

or

$$H_u^+ = \cos\beta\, H^+ + \sin\beta\, G^+ \tag{2.80}$$

$$H_d^- = \sin\beta\, H^{+*} - \cos\beta\, G^{+*} \tag{2.81}$$

A full set of Feynman rules for the MSSM can be found in [3].

Chapter 3

The BMSSM

A hint to the solution of the μ problem was already given in Section 2.2.1. In BMSSM models one tries to exclude the μ parameter in the first place. Thus the only dimensional parameters are in L_{soft} . As a result, all VEVs will naturally be around m_{soft} , and requiring them to be at the EW scale only requires a mild hierarchy if $m_{soft} = O(1)$. Afterwards, one tries to create an effective value for μ as the VEV of some Higgs field. In the following we discuss EW symmetry breaking of the scalar Higgs potential including BMSSM corrections and present the Higgs particle content of the model. Explicit formulas will however be given only if it is necessary for the following chapters. More complete treatments can be found in the literature [1].

3.1 BMSSM corrections due to higher-dimensional operators

The BMSSM Higgs spectrum

The analysis of the BMSSM effects can be organised by studying an effective Lagrangian from which physics at the scale of the BMSSM has been integrated out and is encapsulated in additional operators [1]. In this approach, the BMSSM effects in leading order are encoded in only two effective dimension five operators with undetermined coefficients. One of the effective operators appears in the superpotential. The other one is not SUSY invariant and can be formally described as a superpotential contribution by containing a spurion field that acquires a supersymmetriy breaking F-Term auxiliary component expectation value. From this, one can derive corrections to the MSSM Higgs potential. These corrections contribute both to the mass matrices and the vacuum expectation values, which is demonstrated in this chapter. In this work we determine how the corrections will contribute to LHC Higgs sector and study current publications on the impact of the BMSSM corrections. The recent relevant literature to gain the techniques I performed within the underlying theory are in particular [1] [2] [3] [4]. In the next chapters my study on the review of BMSSM impact to the MSSM Higgs interactions is provided.

At first there have been considered some additional effective contributions to the Higgs sector. The leading superpotential up to dimension five containing only the Higgs fields is

$$\int d^2\theta \left(\mu \mathcal{H}_u \mathcal{H}_d + \frac{\lambda}{M} (\mathcal{H}_u \mathcal{H}_d)^2 + \text{h.c.} \right) \tag{3.1}$$

Kahler potential interactions which involve only Higgs fields are fuctions of an even number of fields. The dimension six operator effects are smaller than those, because they are suppressed by $1/M^2$. Moreover, there is a SUSY breaking operator

$$\int d^2\theta Z \frac{\lambda}{M} (\mathcal{H}_u \mathcal{H}_d)^2 = \frac{\lambda m_{\text{SUSY}}}{M} (H_u H_d)(H_u H_d), \tag{3.2}$$

where $Z = \theta^2 m_{\text{SUSY}}$. Here m_{SUSY} is the SUSY scale of order a few hundred GeV to TeV, and M is the BMSSM scale (TeV scale). λ is a dimensionless parameter. For any other operators giving none-vanishing interactions with one or two powers of the spurion field will not be independent. Operators coupling to other fields are assumed to be suppressed, if leading to FCNC or baryon or lepton number violation. There is a large number of dimension six operators with MSSM particle content which play an important role modifying the light Higgs boson mass for sufficiently small $\cot(\beta)$. The operator analyis could be extended to include D-terms, but the effects are suppressed and insignificant in most microscopic models [1]. Altogether, the corrections to the potential are

$$\delta V = \delta_1 V + \delta_2 V \tag{3.3}$$

$$= 2\epsilon_1 H_u H_d ({H_u}^\dagger H_u + {H_d}^\dagger H_d) + \epsilon_2 (H_u H_d)^2 + \text{h.c.}, \tag{3.4}$$

where $\epsilon_1 = \lambda \mu^* / M$ and $\epsilon_2 = -\lambda m_{\text{SUSY}}/M$. So the full scalar potential is

$$\begin{aligned} V = {} & \tilde{m}^2_{H_u} {H_u}^\dagger H_u + \tilde{m}^2_{H_d} {H_d}^\dagger H_d - m^2_{ud}(H_u H_d + \text{h.c.}) \\ & + \frac{g^2}{8} \left[({H_u}^\dagger H_u + {H_d}^\dagger H_d)^2 - 4 (H_u H_d)^* (H_u H_d) \right] + \frac{g'^2}{8} ({H_u}^\dagger H_u - {H_d}^\dagger H_d)^2 \\ & + 2\epsilon_1 H_u H_d ({H_u}^\dagger H_u + {H_d}^\dagger H_d) + \epsilon_2 (H_u H_d)^2 + \text{h.c.} \end{aligned} \tag{3.5}$$

ϵ_1, ϵ_2 and m^2_{ud} can be complex parameters [2]. The neutral fields H^0_u and H^0_d develop VEVs (see below):

$$
\begin{aligned}
H^0_u &= v_u + \frac{1}{\sqrt{2}} h_u + \frac{i}{\sqrt{2}} g_u \\
H^0_d &= v_d + \frac{1}{\sqrt{2}} h_d + \frac{i}{\sqrt{2}} g_d
\end{aligned}
\tag{3.6}
$$

There are a few possible choices for the basis states:

- Eigenstates of the full mass matrix. Since this may include CP violation, computations will be quite involved (diagonallization of a 4×4-matrix.)
- MSSM eigenstates.
- An "intermediate" choice, e.g. eigenstates of the full mass matrix for the case of vanishing CP violation.

Here only the last two options will be considered. Therefore, one can parametrize:

$$
\begin{pmatrix} g_u \\ g_d \end{pmatrix} = \begin{pmatrix} \cos\beta_0 & -\sin\beta_0 \\ \sin\beta_0 & \cos\beta_0 \end{pmatrix} \begin{pmatrix} A^0 \\ G^0 \end{pmatrix}, \tag{3.7}
$$

$$
\begin{pmatrix} h_u \\ h_d \end{pmatrix} = \begin{pmatrix} \cos\alpha & \sin\alpha \\ -\sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} h^0 \\ H^0 \end{pmatrix}, \tag{3.8}
$$

and

$$
\begin{pmatrix} H^+_u \\ H^{-*}_d \end{pmatrix} = \begin{pmatrix} \cos\beta^+ & \sin\beta^+ \\ \sin\beta^+ & -\cos\beta^+ \end{pmatrix} \begin{pmatrix} H^+ \\ G^+ \end{pmatrix}, \tag{3.9}
$$

with unknown mixing angles β_0, β^+ and α. Moreover, we use the relations

$$
v_u^2 + v_d^2 = v^2 \equiv \frac{2m_Z^2}{g^2 + g'^2} = \frac{2m_W^2}{g^2}, \tag{3.10}
$$

$$
\frac{v_u}{v_d} \equiv \tan\beta. \tag{3.11}
$$

Inserting this into the potiential and setting the first derivatives to zero gives the following relations that determine the VEVs:

$$
m^2_{ud,i} = 2v^2(\epsilon_{1,i} + \epsilon_{2,i}\cos\beta\sin\beta), \tag{3.12}
$$

$$
\tilde{m}^2_{H_u} = m^2_{ud,r}\cot\beta + \frac{1}{2}\cos^2\beta(m_Z^2 - 4\epsilon_{2,r}v^2 - 4\epsilon_{1,r}v^2\cot\beta) - 6\epsilon_{1,r}v^2\cos\beta\sin\beta, \tag{3.13}
$$

$$
\tilde{m}^2_{H_d} = -\frac{1}{2}m_Z^2\cos^2\beta - 6\epsilon_{1,r}v^2\cos\beta\sin\beta + m^2_{ud,r}\tan\beta + \frac{1}{2}\sin^2\beta(m_Z^2 - 4\epsilon_{2,r}v^2 - 4\epsilon_{1,r}v^2\tan\beta). \tag{3.14}
$$

The subscripts i and r denote imaginary and real parts, respectively. Note that if we want the VEVs to be real, these are three equations for only two free parameters (v and

β), so in general there is no solution. Fortunately it can be shown that (3.12) can always be fulfilled by a suitable phase transformation of one of the Higgs doublets. Now the mass matrix can be obtained from the second derivatives of the scalar potential. I am mapping the coeffients in a matrix form so I can see where its entrys do not vanish. In the basis of (h^0, H^0, A^0, G^0, H^+, G^+, H^-, G^-) it has the form of

$$\begin{pmatrix} 1 & 1 & 1 & 1 & 0 & 0 & 0 & 0 \\ 1 & 1 & 1 & 1 & 0 & 0 & 0 & 0 \\ 1 & 1 & 1 & 1 & 0 & 0 & 0 & 0 \\ 1 & 1 & 1 & 1 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 & 1 & 1 \\ 0 & 0 & 0 & 0 & 0 & 0 & 1 & 1 \\ 0 & 0 & 0 & 0 & 1 & 1 & 0 & 0 \\ 0 & 0 & 0 & 0 & 1 & 1 & 0 & 0 \end{pmatrix} \tag{3.15}$$

One obtains two subblocks, where the upper left is the mass matrix for the neutral fields and the lower right for the charged fields. Unlike in the MSSM, the 4×4 neutral part of the mass matrix does not decompose into 2×2 blocks in the general case (complex ϵ_1 or ϵ_2) due to the presence of CP violation. However, it turns out that the choice $\beta_0 = \beta^+ = \beta$ makes most off-diagonal elements of the mass matrix vanish. In particular, the CP-odd and charged sub-blocks are diagonal with mass "eigenvalues":

$$m_A^2 = -4\epsilon_{2,r}v^2 + (m_{ud,r}^2 - 2\epsilon_{1,r}v^2)\csc\beta\sec\beta \tag{3.16}$$

$$m_{G^0}^2 = 0 \tag{3.17}$$

$$\begin{aligned} m_{H^+}^2 &= m_W^2 - 2\epsilon_{2,r}v^2 + (m_{ud,r}^2 - 2\epsilon_{1,r}v^2)\csc\beta\sec\beta \\ &= m_W^2 + m_A^2 + 2\epsilon_{2,r}v^2 \end{aligned} \tag{3.18}$$

$$m_{G^+}^2 = 0. \tag{3.19}$$

The remaining CP-violating quadratic couplings are

$$\begin{aligned} \mathcal{L} \supset\, & 2v^2(\epsilon_{2,i}\cos(\alpha+\beta) - 2\epsilon_{1,i}\sin(\alpha-\beta))A^0h^0 \\ & + 2v^2(2\epsilon_{1,i}\cos(\alpha-\beta) + \epsilon_{2,i}\sin(\alpha+\beta))A^0H^0. \end{aligned} \tag{3.20}$$

Moreover, there is an off-diagonal term in the CP-even block:

$$(\frac{1}{2}(m_A^2 - m_Z^2 + 4\epsilon_{2,r}v^2)\cos 2\beta\sin 2\alpha - \frac{1}{2}((m_A^2 + m_Z^2)\sin 2\beta - 8\epsilon_{1,r}v^2)\cos 2\alpha)h^0H^0 \tag{3.21}$$

The choice

$$\tan 2\alpha = \frac{(m_A^2 + m_Z^2)\tan 2\beta - 8\epsilon_{1,r} v^2 \sec 2\beta}{m_A^2 - m_Z^2 + 4\epsilon_{2,r} v^2} \tag{3.22}$$

makes this term vanish, leading to the following masses:

$$m_{H^0,h^0}^2 = \frac{1}{2}(m_A^2 + m_Z^2 + 4\epsilon_{2,r}v^2 + 8\epsilon_{1,r}v^2 \sin 2\beta) \tag{3.23}$$

$$\pm \frac{1}{2}((m_Z^2 - 2\epsilon_{2,r}v^2)\cos(2\alpha + 2\beta) - (m_A^2 + 2\epsilon_{2,r}v^2)\cos(2\alpha - 2\beta) + 8\epsilon_{1,r}v^2 \sin 2\alpha). \tag{3.24}$$

We express the coefficients of the quadratic terms in terms of v, $\tan\beta$ and m_A and expand the answers to leading order in ϵ where $\eta = \cot\beta$:

$$\delta_\epsilon m_h^2 = v^2 \left(\epsilon_{2r} - 2\epsilon_{1r}\sin(2\beta) + \frac{2\epsilon_{1r}(m_A^2 + m_Z^2)\sin(2\beta) - \epsilon_{2r}(m_A^2 - m_Z^2)\cos^2(2\beta)}{\sqrt{(m_A^2 - m_Z^2)^2 + 4m_A^2 m_Z^2 \sin^2(2\beta)}} \right)$$

$$\simeq \frac{16 m_A^2}{m_A^2 - m_Z^2} v^2 \eta \epsilon_{1r} + \mathcal{O}(\eta^2 \epsilon) \tag{3.25}$$

$$\delta_\epsilon m_H^2 = v^2 \left(\epsilon_{2r} - 2\epsilon_{1r}\sin(2\beta) - \frac{2\epsilon_{1r}(m_A^2 + m_Z^2)\sin(2\beta) - \epsilon_{2r}(m_A^2 - m_Z^2)\cos^2(2\beta)}{\sqrt{(m_A^2 - m_Z^2)^2 + 4m_A^2 m_Z^2 \sin^2(2\beta)}} \right)$$

$$\simeq 4v^2\epsilon_{2r} - \frac{16 m_Z^2}{m_A^2 - m_Z^2} v^2 \eta \epsilon_{1r} + \mathcal{O}(\eta^2 \epsilon) \tag{3.26}$$

$$\delta_\epsilon m_{H^\pm}^2 = \epsilon_{2r} v^2 \tag{3.27}$$

where ϵ_{1r} and ϵ_{2r} are the real parts of ϵ_1 and ϵ_2. For moderate $\cot\beta$ the operators at order ϵ describe the dominant contribution of BMSSM physics to the Higgs spectrum. The Taylor series expansion in ϵ is only valid if $m_A^2 - m_Z^2 \gg 4\epsilon_{2r}v^2$. These results are in accordance with the results of [1] whereas the charged Higgs mass corrections differ by a factor 2 and the neutral Higgs mass corrections by a factor 2 in the first summand and a sign in the second summand, which needs further investigation.

Moreover, a set of Feynman rules can be derived.

3.1.1 Discussion

The recent hints of a Standard Model-like Higgs with a mass close to 126 GeV [8] [9] are of crucial theoretical interest to figure out if data in the various channels measured by ATLAS and CMS could be used as a probe of BMSSM physics.

- In a low energy Supersymmetry with modest stop squark masses and mixings BMSSM operators can correct MSSM Higgs masses, lifting the SM Higgs mass to the current bounds

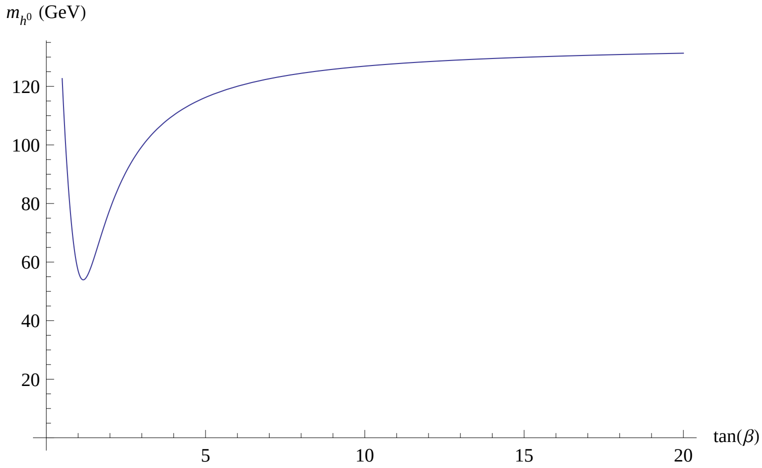


Figure 3.1: With Mathematica here we visualised the results for the light Higgs mass m_h (left axis) of 3.1 for a set of values for the BMSSM parameters $\epsilon_{1r} = 0.15$ and $\epsilon_{2r} = 0.25$ for a range of $\tan\beta$ between 0 and 20 (right axis). Further I set $m_t = 173$ GeV ,$mZ = 91$ GeV, $m_A = 300$ GeV, $v = 174$ GeV, $y_t = \frac{m_t}{v \sin\beta}$

.

- For small η BMSSM operators correct MSSM Higgs masses in leading order ϵ only by one or two real numbers ϵ_{1r} and ϵ_{2r}
- In [1] a classification of effects which may lift the light Higgs mass is given
- Figure 3.1.1 shows that even with moderate masses of 500 GeV for the stop quarks in the dominant loop correction $\Delta(m^2_{h^0}) = \frac{3}{4\pi^2}\cos^2\alpha\; y_t^2 m_t^2 \ln\left(m_{\tilde{t}_1} m_{\tilde{t}_2}/m_t^2\right)$ the light Higgs mass can be lifted to the current Higgs bounds

Enhanced Higgs sector studies of BMSSM operators have been carried out in [10] [11] [12] [13] [14].

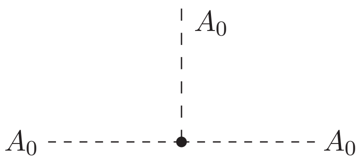
Chapter 4

Feynman rules

Feynman rules for the BMSSM are derived by collecting the cubic and quartic terms of the scalar potential after symmetry breaking and the mass eigenstates. First, as previously demonstrated in Chapter 3 one needs to derive the BMSSM Higgs mass matrix and parameterise the gauge eigenstates through VEVs and physical fields. One needs to determine the VEVs and mixing angles. With Mathematica I expand the scalar potential in the convention of [3] and organise it in ascending powers of the fields and calculate a coefficient array of all terms. Setting the linear terms to 0 leads to expressions for the VEVs. By setting the off-diagonal elements to 0 one can obtain a conditional equation for the mixing angles. By assembling identical field operater combinations from the terms of higher order in the fields and multiplying the symmetry factor and the complex unit I obtain the Feynamnrules for the BMSSM. Then again I expand in orders of ϵ and η. In Section 4.1 a comparison of Feynman rules for all three- and four-Higgs interactions for the BMSSM with the MSSM is provided. Note that some of the Feynman rules may actually vanish due to the relation between the of the mixing angles α and β, which was not used in the derivation. To see if my results are correct in the the limit of the MSSM I compared them with set of the Feynman rules in [3] by setting the epsilon terms in the scalar potential to 0. When I merge my coefficient array with the cubic and quartic powers I receive all MSSM Higgs intreactions. I output all Feynman rules which deviate from Rosieks Feynman rules. As a result this output contains no elements, which means my calculation in the limit of the MSSM is correct.

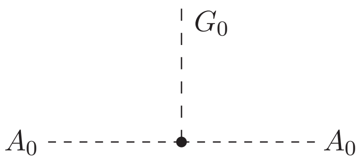
4.1 A full set of Feynman rules for BMSSM tree level Higgs interactions

In the following Feynman rules are organised in comparison:
top row - MSSM, bottom row - BMSSM.



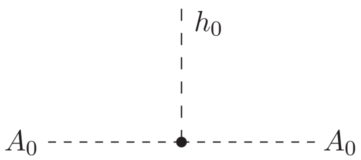
0 (MSSM)

$-6i(\epsilon_{1,i} + \epsilon_{2,i}\eta)v$ (BMSSM)



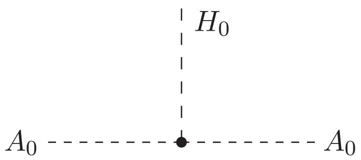
0

$-2i\epsilon_{2,i}v$



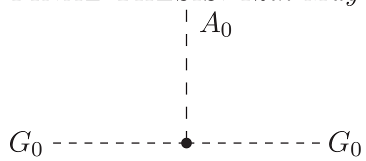
$\frac{im_Z^2}{v}$

$2i\epsilon_{2,r}v - \frac{4i\epsilon_{1,r}\eta m_Z^4 v}{(m_A^2 - m_Z^2)^2}$



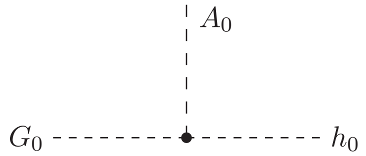
$-\frac{2i\eta m_A^2 m_Z^2}{m_A^2 v - m_Z^2 v}$

$$\frac{1}{(m_A^2 - m_Z^2)^2} 2i \Big((\epsilon_{1,r} + \epsilon_{2,r}\eta) m_A^4 - 3(\epsilon_{1,r} + \epsilon_{2,r}\eta) m_A^2 m_Z^2 + 2(\epsilon_{1,r} + 2\epsilon_{2,r}\eta) m_Z^4 v$$



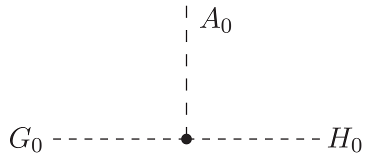
0

$-2i(\epsilon_{1,i} - \epsilon_{2,i}\eta)v$



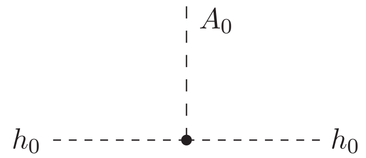
$-\frac{2i\eta m_Z^2}{v}$

$-\frac{2i(\epsilon_{1,r}(-m_A^2+m_Z^2)+\epsilon_{2,r}\eta(m_A^2+m_Z^2))v}{-m_A^2+m_Z^2}$



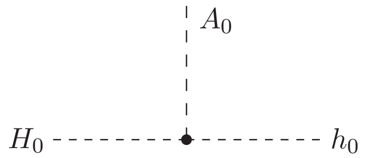
0

$2i\epsilon_{2,r}v + \frac{8i\epsilon_{1,r}\eta m_Z^2 v}{m_A^2-m_Z^2}$



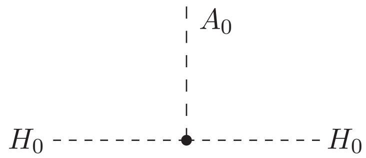
0

$-6i\epsilon_{1,i}v + \frac{2i\epsilon_{2,i}\eta(3m_A^2+m_Z^2)v}{m_A^2-m_Z^2}$



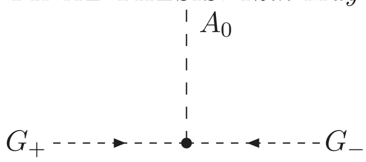
0

$2i\epsilon_{2,i}v + \frac{8i\epsilon_{1,i}\eta m_Z^2 v}{m_A^2-m_Z^2}$



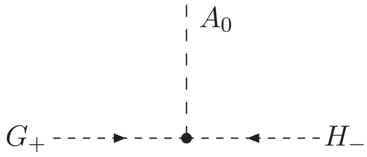
0

$-2i\left(\epsilon_{1,i} + \frac{\epsilon_{2,i}\eta(m_A^2+3m_Z^2)}{m_A^2-m_Z^2}\right)v$



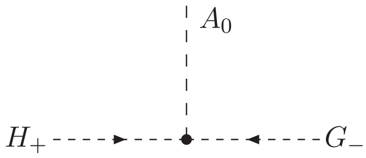
0

$-2i(\epsilon_{1,i} - \epsilon_{2,i}\eta)v$



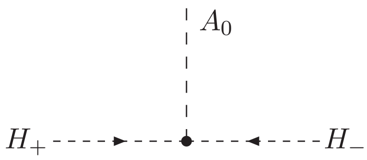
$-\frac{m_W^2}{v}$

$-i\epsilon_{2,i}v - \epsilon_{2,r}v$



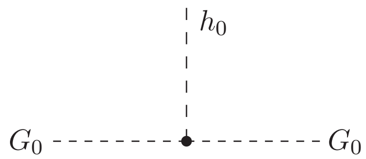
$\frac{m_W^2}{v}$

$(-i\epsilon_{2,i} + \epsilon_{2,r})v$



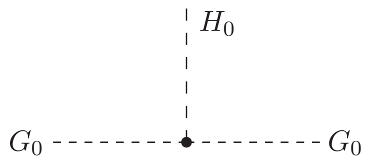
0

$-2i(\epsilon_{1,i} + \epsilon_{2,i}\eta)v$



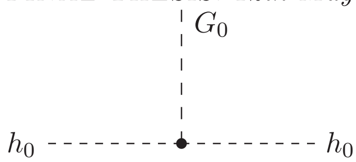
$-\frac{im_Z^2}{v}$

$\frac{4i\epsilon_{1,r}\eta\left(2m_A^4-2m_A^2m_Z^2+m_Z^4\right)v}{(m_A^2-m_Z^2)^2}$



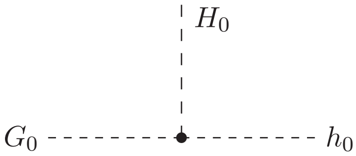
$\frac{2i\eta m_A^2 m_Z^2}{m_A^2 v - m_Z^2 v}$

$\frac{1}{(m_A^2-m_Z^2)^2}2i\left(\epsilon_{1,r}m_A^2(m_A^2-m_Z^2) - \epsilon_{2,r}\eta\left(m_A^4 - m_A^2 m_Z^2 + 2m_Z^4\right)\right)v$



0

$-\frac{8i\epsilon_{1,i}\eta m_Z^2 v}{m_A^2-m_Z^2}$



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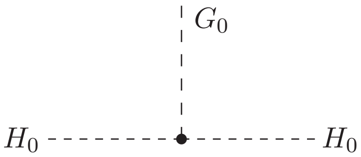
$-2i\left(\epsilon_{1,i} - \frac{\epsilon_{2,i}\eta(m_A^2+m_Z^2)}{m_A^2-m_Z^2}\right)v$



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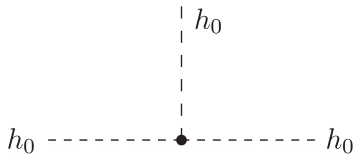
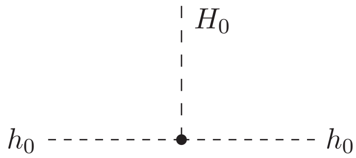
$2i\epsilon_{2,i}v + \frac{8i\epsilon_{1,i}\eta m_Z^2 v}{m_A^2-m_Z^2}$



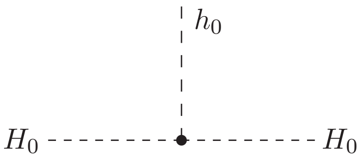
$-\frac{3im_Z^2}{v}$

$\frac{12i\epsilon_{1,r}\eta\left(2m_A^4+2m_A^2m_Z^2+m_Z^4\right)v}{(m_A^2-m_Z^2)^2}$



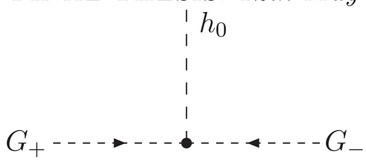
$\frac{2i\eta m_Z^2(3m_A^2+2m_Z^2)}{(m_A^2-m_Z^2)v}$

$$\frac{1}{(m_A^2-m_Z^2)^2}2i\left(\epsilon_{1,r}(m_A^2-m_Z^2)(3m_A^2+2m_Z^2)\right.$$
$$\left.-\epsilon_{2,r}\eta\left(3m_A^4+3m_A^2m_Z^2+4m_Z^4\right)v\right.$$

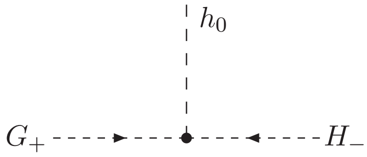


$\frac{im_Z^2}{v}$

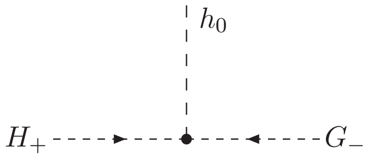
$-\frac{1}{(m_A^2-m_Z^2)^2}2i\left(\epsilon_{2,r}(m_A^2-m_Z^2)^2+2\epsilon_{1,r}\eta m_Z^2(12m_A^2+m_Z^2)\right)v$



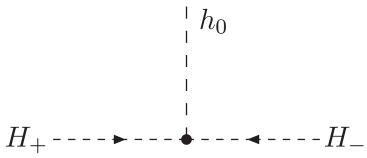
$$-\frac{im_Z^2}{v}$$
$$\frac{4i\epsilon_{1,r}\eta\left(2m_A^4-2m_A^2m_Z^2+m_Z^4\right)v}{(m_A^2-m_Z^2)^2}$$



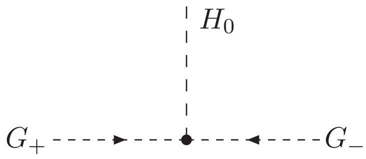
$$-\frac{2i\eta(m_A^2+m_W^2-m_Z^2)m_Z^2}{(m_A^2-m_Z^2)v}$$
$$\frac{1}{(m_A^2-m_Z^2)^2}2\left(\epsilon_{1,i}(m_A^2-m_Z^2)^2\right.$$
$$-i(\epsilon_{1,r}(m_A^2-m_Z^2)(m_A^2+m_W^2-m_Z^2)+\eta(-\epsilon_{2,r}m_A^2(m_A^2+m_W^2)$$
$$-i\epsilon_{2,i}m_A^2(m_A^2-m_Z^2)+\epsilon_{2,r}(m_A^2-m_W^2)m_Z^2))v$$



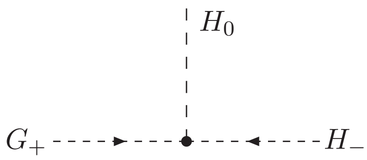
$$-\frac{2i\eta(m_A^2+m_W^2-m_Z^2)m_Z^2}{(m_A^2-m_Z^2)v}$$
$$-\frac{1}{(m_A^2-m_Z^2)^2}2$$
$$\left(\epsilon_{1,i}(m_A^2-m_Z^2)^2+i(\epsilon_{1,r}(m_A^2-m_Z^2)(m_A^2+m_W^2-m_Z^2)\right.$$
$$+\eta(-\epsilon_{2,r}m_A^2(m_A^2+m_W^2)+i\epsilon_{2,i}m_A^2(m_A^2-m_Z^2)$$
$$+\epsilon_{2,r}(m_A^2-m_W^2)m_Z^2))v$$



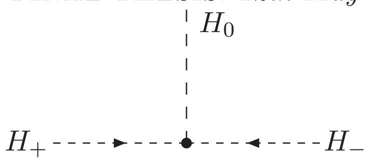
$$\frac{i(-2m_W^2+m_Z^2)}{v}$$
$$-\frac{4i\epsilon_{1,r}\eta m_Z^2(-2m_W^2+m_Z^2)v}{(m_A^2-m_Z^2)^2}$$



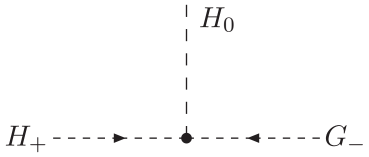
$$\frac{2i\eta m_A^2m_Z^2}{m_A^2v-m_Z^2v}$$
$$\frac{1}{(m_A^2-m_Z^2)^2}2i\left(\epsilon_{1,r}m_A^2(m_A^2-m_Z^2)-\epsilon_{2,r}\eta\left(m_A^4-m_A^2m_Z^2+2m_Z^4\right)\right)v$$



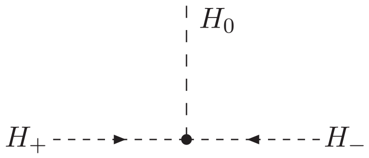
$$-\frac{im_W^2}{v}$$
$$(-\epsilon_{2,i}+i\epsilon_{2,r}+\frac{1}{(m_A^2-m_Z^2)^2}4\eta m_Z^2(i\epsilon_{1,r}(2m_A^2+m_W^2-2m_Z^2)$$
$$+\epsilon_{1,i}(-m_A^2+m_Z^2)))v$$



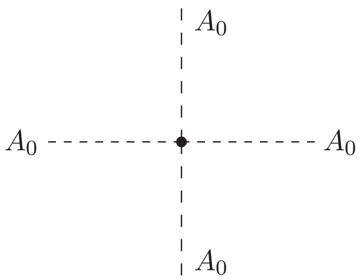
$$-\frac{im_W^2}{v}$$
$$(-\epsilon_{2,i} + i\epsilon_{2,r} + \frac{1}{(m_A^2 - m_Z^2)^2} 4\eta m_Z^2 (i\epsilon_{1,r}(2m_A^2 + m_W^2 - 2m_Z^2)$$
$$+\epsilon_{1,i}(-m_A^2 + m_Z^2)))v$$



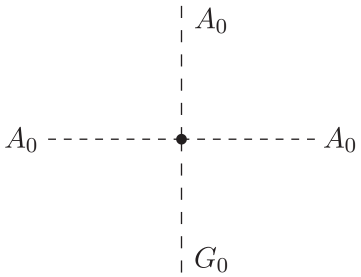
$$-\frac{im_W^2}{v}$$
$$(\epsilon_{2,i} + i\epsilon_{2,r} + \frac{1}{(m_A^2 - m_Z^2)^2} 4\eta (i\epsilon_{1,r}(2m_A^2 + m_W^2 - 2m_Z^2)$$
$$+\epsilon_{1,i}(m_A^2 - m_Z^2))m_Z^2)v$$



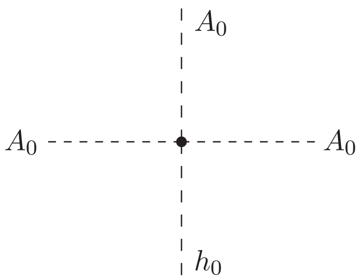
$$-\frac{2i\eta(m_A^2 - 2m_W^2)m_Z^2}{(m_A^2 - m_Z^2)v}$$
$$\frac{1}{(m_A^2 - m_Z^2)^2} 2i \left(\epsilon_{1,r}(m_A^2 + 2m_W^2 - 2m_Z^2)(m_A^2 - m_Z^2)\right.$$
$$\left.+\epsilon_{2,r}\eta \left(m_A^4 + 2m_Z^2(-m_W^2 + m_Z^2) - m_A^2(2m_W^2 + m_Z^2)\right) v\right.$$



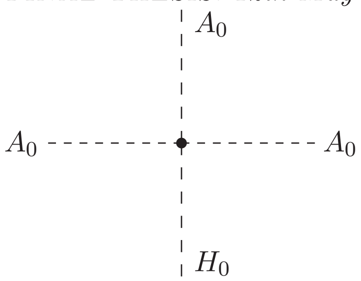
$$-\frac{3im_Z^2}{v^2}$$
$$-24i\epsilon_{1,r}\eta$$



$$\frac{6i\eta m_Z^2}{v^2}$$
$$-6i(\epsilon_{1,r} + \epsilon_{2,r}\eta)$$



$$0$$
$$-6i(\epsilon_{1,i} + \epsilon_{2,i}\eta)$$



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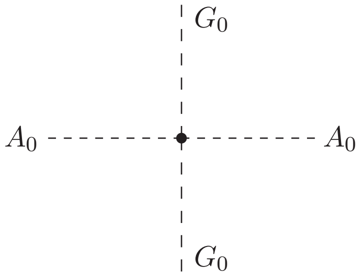
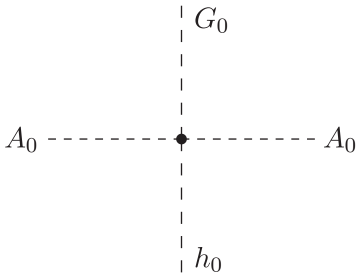
$\frac{12i\epsilon_{1,i}\eta m_Z^2}{m_A^2-m_Z^2}$



$\frac{im_Z^2}{v^2}$

$-2i\epsilon_{2,r}$



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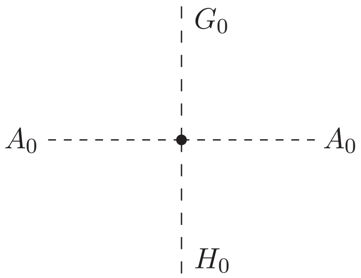
$-2i\epsilon_{2,i} - \frac{4i\epsilon_{1,i}\eta m_Z^2}{m_A^2-m_Z^2}$



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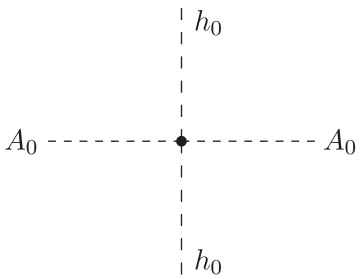
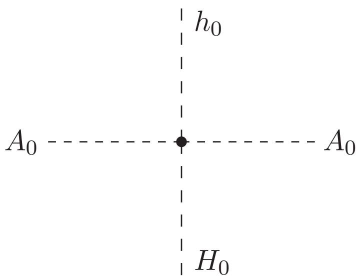
$-2i\left(\epsilon_{1,i} + \frac{\epsilon_{2,i}\eta(m_A^2-3m_Z^2)}{m_A^2-m_Z^2}\right)$



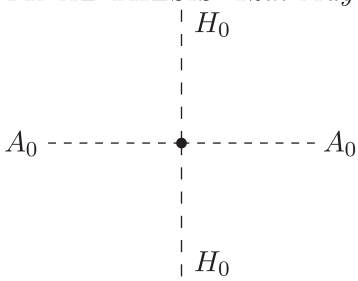
$\frac{im_Z^2}{v^2}$

$\frac{2i\left(\epsilon_{2,r}(m_A^2-m_Z^2)^2-8\epsilon_{1,r}\eta m_Z^4\right)}{(m_A^2-m_Z^2)^2}$

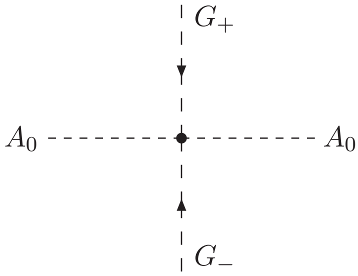


$-\frac{2i\eta m_Z^2(m_A^2+m_Z^2)}{(m_A^2-m_Z^2)v^2}$

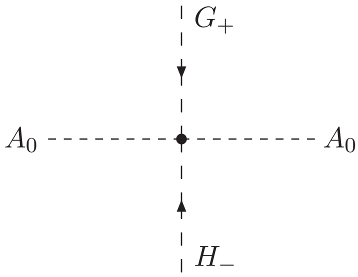
$$\frac{1}{(m_A^2-m_Z^2)^2}2i\left(\epsilon_{1,r}(m_A^2-3m_Z^2)(m_A^2-m_Z^2)\right.$$
$$\left.+\epsilon_{2,r}\eta\left(m_A^4-2m_A^2m_Z^2+5m_Z^4\right)\right)$$



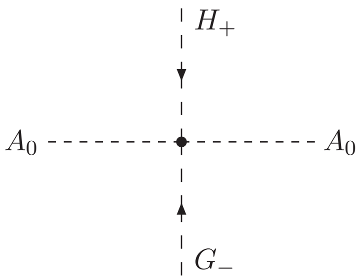
$$-\frac{im_Z^2}{v^2}$$
$$-\frac{8i\epsilon_{1,r}\eta\left(m_A^4-2m_A^2m_Z^2-m_Z^4\right)}{(m_A^2-m_Z^2)^2}$$



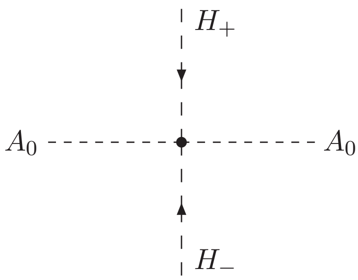
$$-\frac{i(2m_W^2-m_Z^2)}{v^2}$$
$$0$$



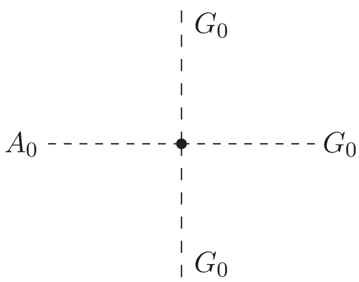
$$\frac{2i\eta m_Z^2}{v^2}$$
$$2(\epsilon_{1,i}-i(\epsilon_{1,r}+(i\epsilon_{2,i}+\epsilon_{2,r})\eta))$$



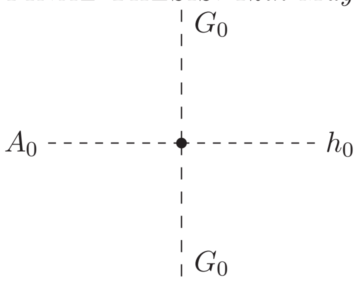
$$\frac{2i\eta m_Z^2}{v^2}$$
$$-2(\epsilon_{1,i}+i\epsilon_{1,r}+(\epsilon_{2,i}+i\epsilon_{2,r})\eta)$$



$$-\frac{im_Z^2}{v^2}$$
$$-8i\epsilon_{1,r}\eta$$

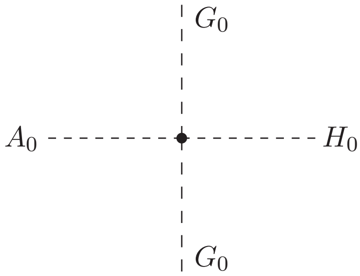


$$-\frac{6i\eta m_Z^2}{v^2}$$
$$-6i(\epsilon_{1,r}-\epsilon_{2,r}\eta)$$



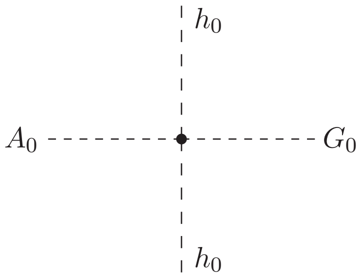
$$0$$

$$-2i\left(\epsilon_{1,i}-\frac{\epsilon_{2,i}\eta(m_A^2-3m_Z^2)}{m_A^2-m_Z^2}\right)$$



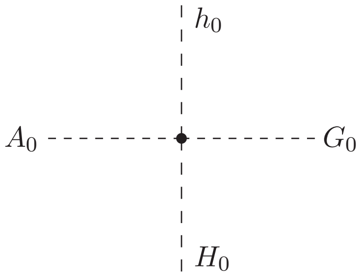
$$0$$

$$-2i\epsilon_{2,i}+\frac{4i\epsilon_{1,i}\eta m_Z^2}{m_A^2-m_Z^2}$$



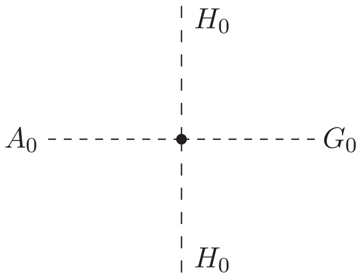
$$-\frac{2i\eta m_Z^2}{v^2}$$

$$-2i\left(\epsilon_{1,r}+\epsilon_{2,r}\eta\left(3+\frac{4m_A^2}{-m_A^2+m_Z^2}\right)\right)$$



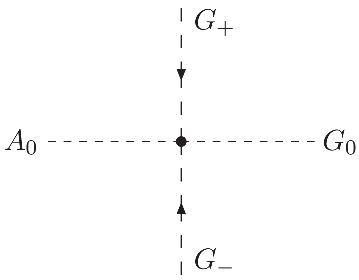
$$0$$

$$2i\epsilon_{2,r}+\frac{8i\epsilon_{1,r}\eta m_Z^2}{m_A^2-m_Z^2}$$



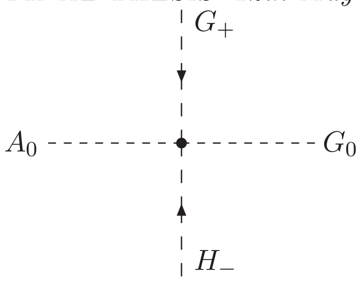
$$\frac{2i\eta m_Z^2}{v^2}$$

$$-2i\left(\epsilon_{1,r}+\frac{\epsilon_{2,r}\eta(m_A^2+3m_Z^2)}{m_A^2-m_Z^2}\right)$$



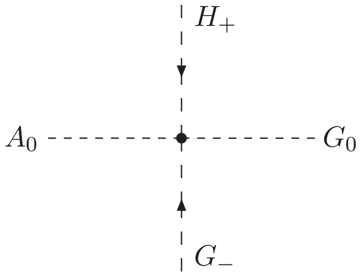
$$-\frac{2i\eta m_Z^2}{v^2}$$

$$-2i(\epsilon_{1,r}-\epsilon_{2,r}\eta)$$



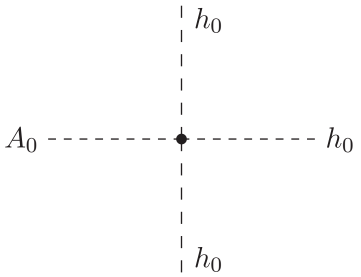
$\frac{im_W^2}{v^2}$

$\epsilon_{2,i} - i\epsilon_{2,r}$



$\frac{im_W^2}{v^2}$

$-\epsilon_{2,i} - i\epsilon_{2,r}$



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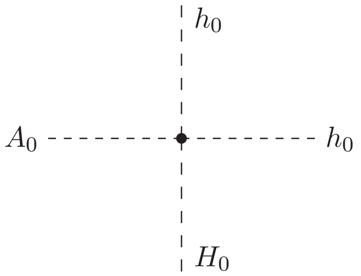
$-6i\left(\epsilon_{1,i} - \frac{\epsilon_{2,i}\eta(m_A^2+m_Z^2)}{m_A^2-m_Z^2}\right)$



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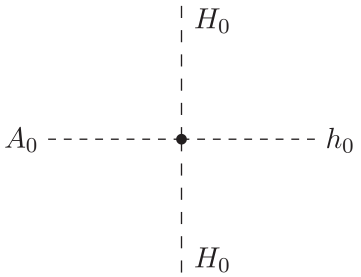
$2i\epsilon_{2,i} + \frac{4i\epsilon_{1,i}\eta m_Z^2}{m_A^2-m_Z^2}$



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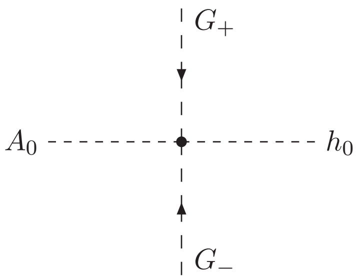
$-2i\left(\epsilon_{1,i} + \frac{\epsilon_{2,i}\eta(m_A^2+3m_Z^2)}{m_A^2-m_Z^2}\right)$



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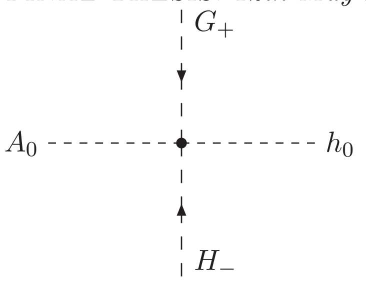
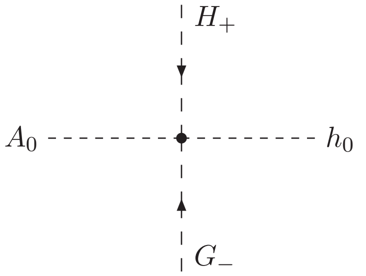
$-2i(\epsilon_{1,i} - \epsilon_{2,i}\eta)$



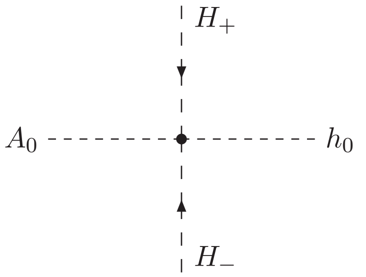
$-\frac{m_W^2}{v^2}$

$-i\epsilon_{2,i} - \epsilon_{2,r} + \frac{4\epsilon_{1,r}\eta m_W^2 m_Z^2}{(m_A^2 - m_Z^2)^2}$



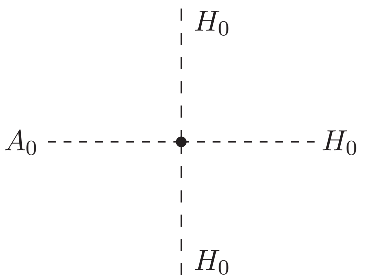
$\frac{m_W^2}{v^2}$

$-i\epsilon_{2,i} + \epsilon_{2,r} - \frac{4\epsilon_{1,r}\eta m_W^2 m_Z^2}{(m_A^2 - m_Z^2)^2}$



0

$-2i(\epsilon_{1,i} + \epsilon_{2,i}\eta)$



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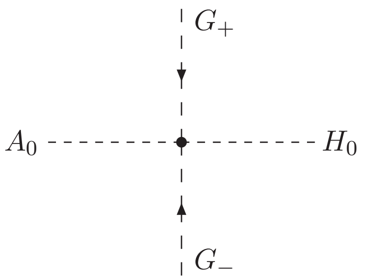
$\frac{12i\epsilon_{1,i}\eta m_Z^2}{m_A^2 - m_Z^2}$



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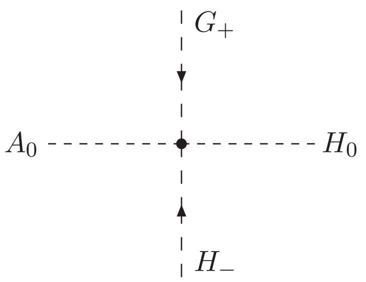
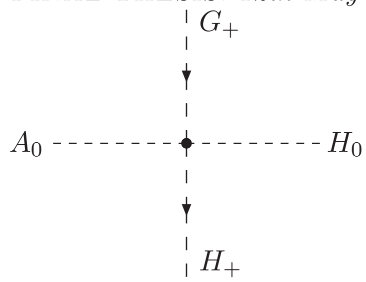
$\frac{4i\epsilon_{1,i}\eta m_Z^2}{m_A^2 - m_Z^2}$



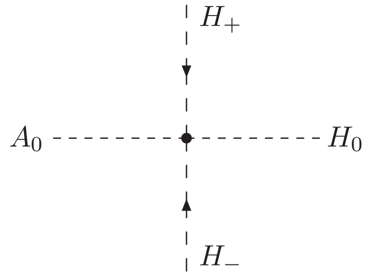
$\frac{2\eta m_W^2 m_Z^2}{(m_A^2 - m_Z^2)v^2}$

$$\frac{1}{(m_A^2 - m_Z^2)^2} 2(\epsilon_{1,r} m_W^2 (m_A^2 - m_Z^2) - \eta(-i\epsilon_{2,i}(m_A^2 - m_Z^2)m_Z^2$$
$$+\epsilon_{2,r}(m_A^2(m_W^2 - m_Z^2) + m_Z^2(m_W^2 + m_Z^2))))$$



$$\frac{2\eta m_W^2 m_Z^2}{(-m_A^2+m_Z^2)v^2}$$

$$\frac{1}{(m_A^2-m_Z^2)^2} 2(\epsilon_{1,r} m_W^2(-m_A^2+m_Z^2)+\eta(i\epsilon_{2,i}(m_A^2-m_Z^2)m_Z^2$$
$$+\epsilon_{2,r}(m_A^2(m_W^2-m_Z^2)+m_Z^2(m_W^2+m_Z^2))))$$



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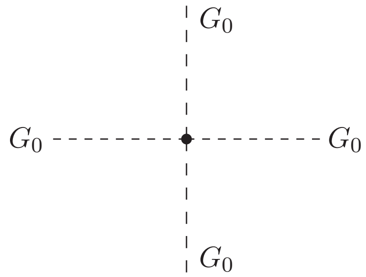
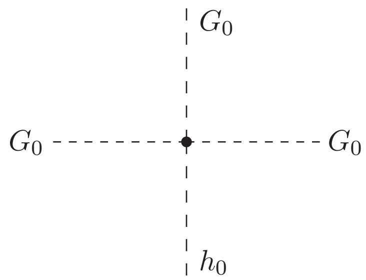
$$\frac{4i\epsilon_{1,i}\eta m_Z^2}{m_A^2-m_Z^2}$$



$-\frac{3im_Z^2}{v^2}$ $24i\epsilon_{1,r}\eta$



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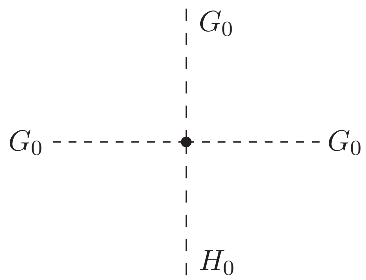
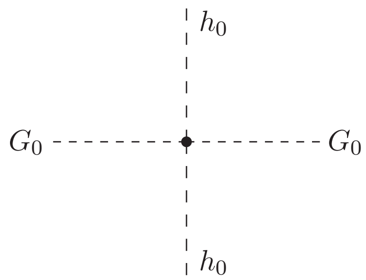
$$-\frac{12i\epsilon_{1,i}\eta m_Z^2}{m_A^2-m_Z^2}$$



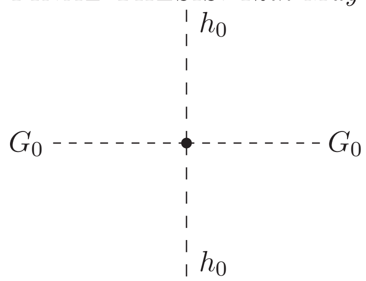
0

$-6i(\epsilon_{1,i}-\epsilon_{2,i}\eta)$



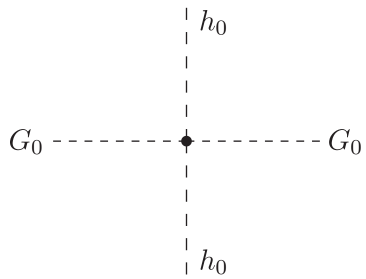
$$-\frac{im_Z^2}{v^2}$$

$$\frac{8i\epsilon_{1,r}\eta\left(m_A^4+m_Z^4\right)}{(m_A^2-m_Z^2)^2}$$



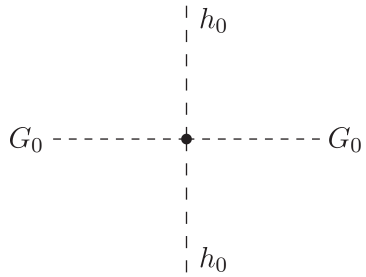
$$-\frac{im_Z^2}{v^2}$$

$$\frac{8i\epsilon_{1,r}\eta\left(m_A^4+m_Z^4\right)}{(m_A^2-m_Z^2)^2}$$



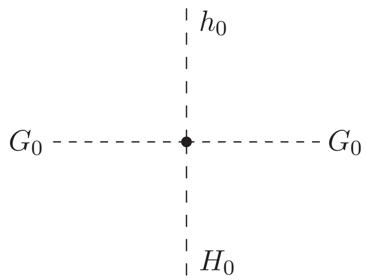
$$-\frac{im_Z^2}{v^2}$$

$$\frac{8i\epsilon_{1,r}\eta\left(m_A^4+m_Z^4\right)}{(m_A^2-m_Z^2)^2}$$



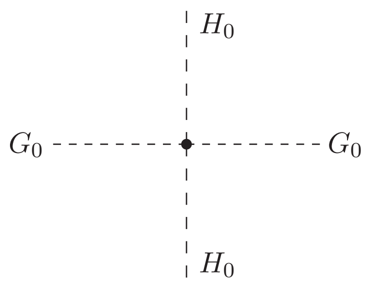
$$-\frac{im_Z^2}{v^2}$$

$$\frac{8i\epsilon_{1,r}\eta\left(m_A^4+m_Z^4\right)}{(m_A^2-m_Z^2)^2}$$



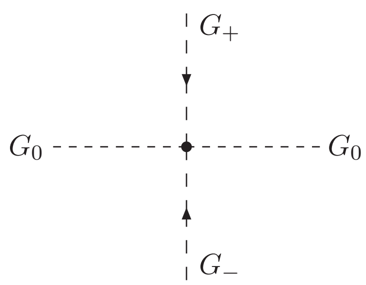
$$\frac{2i\eta m_Z^2(m_A^2+m_Z^2)}{(m_A^2-m_Z^2)v^2}$$

$$\frac{1}{(m_A^2-m_Z^2)^2}2i\left((\epsilon_{1,r}-\epsilon_{2,r}\eta)m_A^4+2\epsilon_{2,r}\eta m_A^2m_Z^2-(\epsilon_{1,r}+5\epsilon_{2,r}\eta)m_Z^4\right)$$



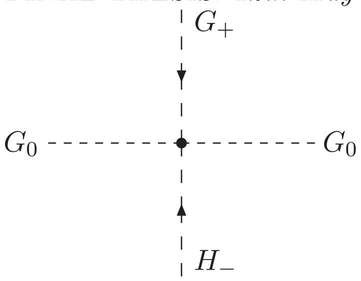
$$\frac{im_Z^2}{v^2}$$

$$\frac{2i\left(\epsilon_{2,r}(m_A^2-m_Z^2)^2-8\epsilon_{1,r}\eta m_A^2m_Z^2\right)}{(m_A^2-m_Z^2)^2}$$



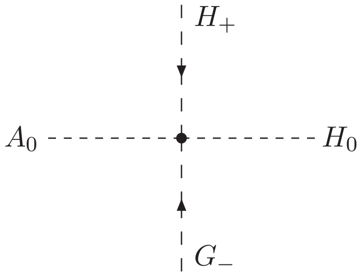
$$-\frac{im_Z^2}{v^2}$$

$$8i\epsilon_{1,r}\eta$$



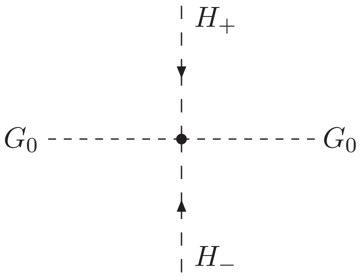
$-\frac{2i\eta m_Z^2}{v^2}$

$2(\epsilon_{1,i} - i\epsilon_{1,r} - \epsilon_{2,i}\eta + i\epsilon_{2,r}\eta)$



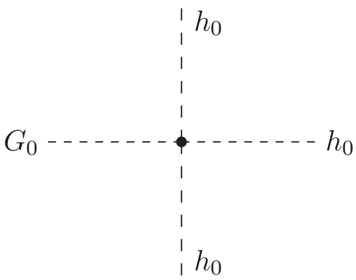
$-\frac{2i\eta m_Z^2}{v^2}$

$-2\epsilon_{1,i} - 2i\epsilon_{1,r} + 2(\epsilon_{2,i} + i\epsilon_{2,r})\eta$



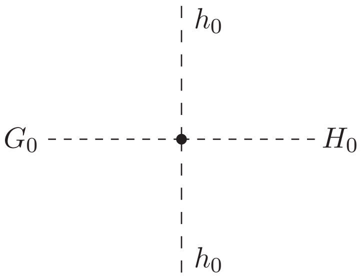
$-\frac{i(2m_W^2 - m_Z^2)}{v^2}$

0



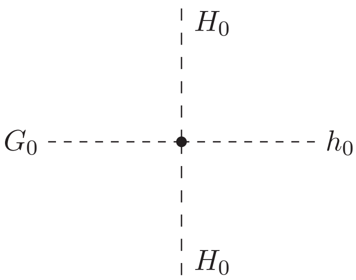
0

$-\frac{12i\epsilon_{1,i}\eta m_Z^2}{m_A^2 - m_Z^2}$



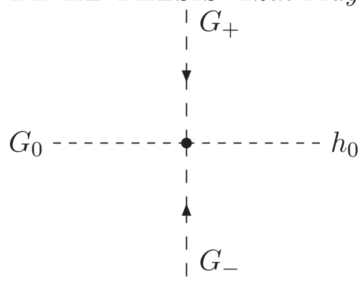
0

$-2i\left(\epsilon_{1,i} + \epsilon_{2,i}\eta\left(3 + \frac{4m_A^2}{-m_A^2 + m_Z^2}\right)\right)$



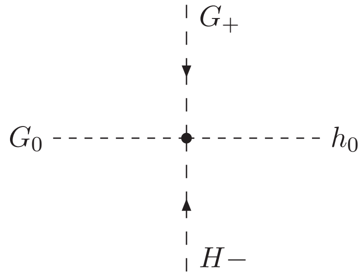
0

$2i\epsilon_{2,i} - \frac{4i\epsilon_{1,i}\eta m_Z^2}{m_A^2 - m_Z^2}$



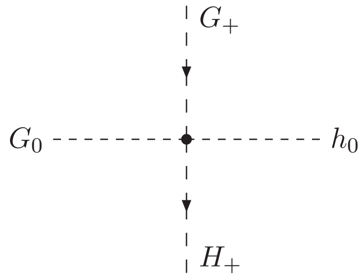
$$0$$

$$-\frac{4i\epsilon_{1,i}\eta m_Z^2}{m_A^2-m_Z^2}$$



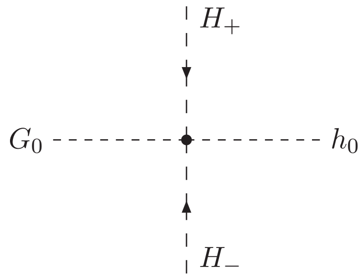
$$\frac{2\eta m_W^2 m_Z^2}{(-m_A^2+m_Z^2)v^2}$$

$$\frac{1}{(m_A^2-m_Z^2)^2}2(\epsilon_{1,r}m_W^2(-m_A^2+m_Z^2)+\eta(-i\epsilon_{2,i}(m_A^2-m_Z^2)m_Z^2$$
$$+\epsilon_{2,r}(m_A^2(m_W^2-m_Z^2)+m_Z^2(m_W^2+m_Z^2))))$$



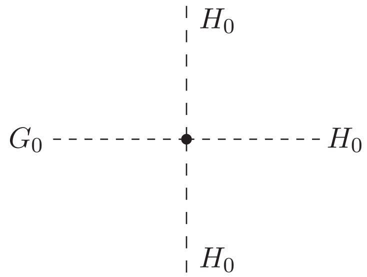
$$\frac{2\eta m_W^2 m_Z^2}{(m_A^2-m_Z^2)v^2}$$

$$\frac{1}{(m_A^2-m_Z^2)^2}2(\epsilon_{1,r}m_W^2(m_A^2-m_Z^2)-\eta(i\epsilon_{2,i}(m_A^2-m_Z^2)m_Z^2$$
$$+\epsilon_{2,r}(m_A^2(m_W^2-m_Z^2)+m_Z^2(m_W^2+m_Z^2))))$$



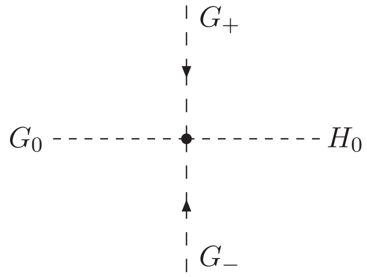
$$0$$

$$-\frac{4i\epsilon_{1,i}\eta m_Z^2}{m_A^2-m_Z^2}$$



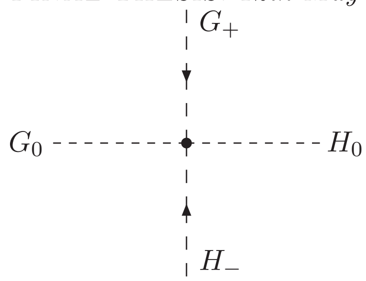
$$0$$

$$-6i\left(\epsilon_{1,i}+\frac{\epsilon_{2,i}\eta(m_A^2+m_Z^2)}{m_A^2-m_Z^2}\right)$$

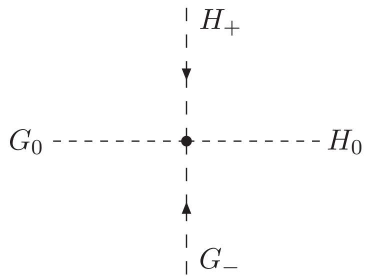


$$0$$

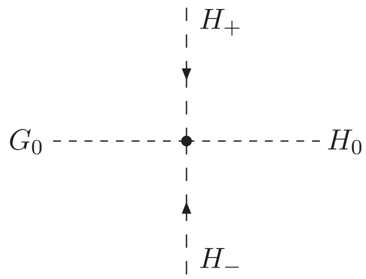
$$-2i(\epsilon_{1,i}-\epsilon_{2,i}\eta)$$



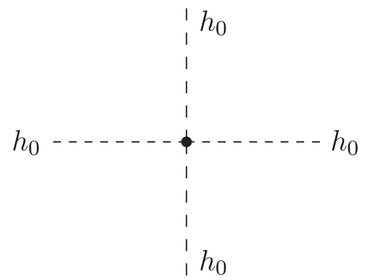
$$-\frac{m_W^2}{v^2}$$
$$-i\epsilon_{2,i}-\epsilon_{2,r}+\frac{4\epsilon_{1,r}\eta m_W^2 m_Z^2}{(m_A^2-m_Z^2)^2}$$



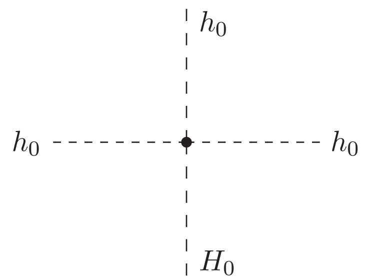
$$\frac{m_W^2}{v^2}$$
$$-i\epsilon_{2,i}+\epsilon_{2,r}-\frac{4\epsilon_{1,r}\eta m_W^2 m_Z^2}{(m_A^2-m_Z^2)^2}$$



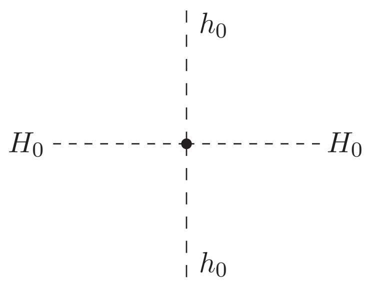
$$0$$
$$-2i(\epsilon_{1,i}+\epsilon_{2,i}\eta)$$



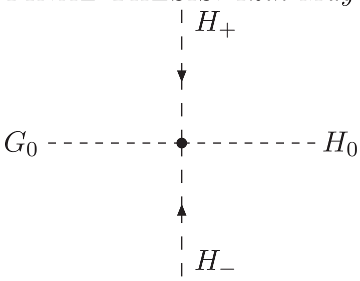
$$-\frac{3im_Z^2}{v^2}$$
$$\frac{24i\epsilon_{1,r}\eta(m_A^2+m_Z^2)^2}{(m_A^2-m_Z^2)^2}$$



$$\frac{6i\eta m_Z^2(m_A^2+m_Z^2)}{(m_A^2-m_Z^2)v^2}$$
$$\frac{6i(m_A^2+m_Z^2)(\epsilon_{1,r}(m_A^2-m_Z^2)-\epsilon_{2,r}\eta(m_A^2+m_Z^2))}{(m_A^2-m_Z^2)^2}$$

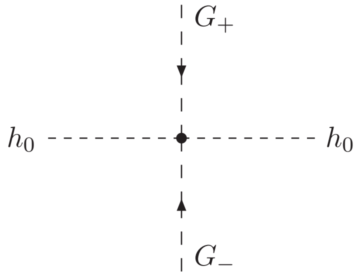


$$\frac{im_Z^2}{v^2}$$
$$-\frac{2i\left(\epsilon_{2,r}(m_A^2-m_Z^2)^2+24\epsilon_{1,r}\eta m_Z^2(m_A^2+m_Z^2)\right)}{(m_A^2-m_Z^2)^2}$$



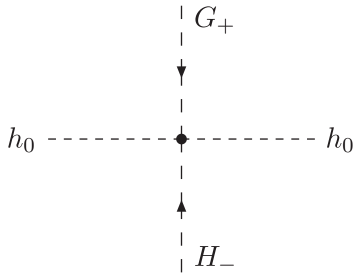
$$0$$

$$-2i(\epsilon_{1,i} + \epsilon_{2,i}\eta)$$



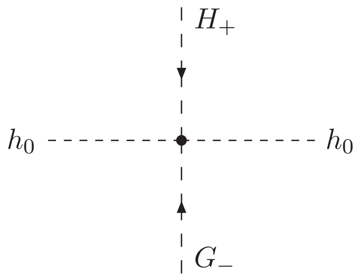
$$-\frac{im_Z^2}{v^2}$$

$$\frac{8i\epsilon_{1,r}\eta\left(m_A^4+m_Z^2(-2m_W^2+m_Z^2)\right)}{(m_A^2-m_Z^2)^2}$$



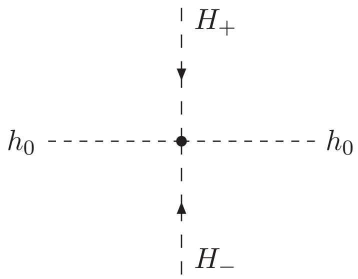
$$-\frac{2i\eta(m_A^2+2m_W^2-m_Z^2)m_Z^2}{(m_A^2-m_Z^2)v^2}$$

$$\frac{1}{(m_A^2 - m_Z^2)^2} 2 \Big(\epsilon_{1,i}(m_A^2 - m_Z^2)^2 - i(\epsilon_{1,r}(m_A^2 - m_Z^2)$$
$$(m_A^2 + 2m_W^2 - m_Z^2) + \eta(m_A^2 + m_Z^2)(-i\epsilon_{2,i}(m_A^2 - m_Z^2)$$
$$+\epsilon_{2,r}(-m_A^2 - 2m_W^2 + m_Z^2)))$$



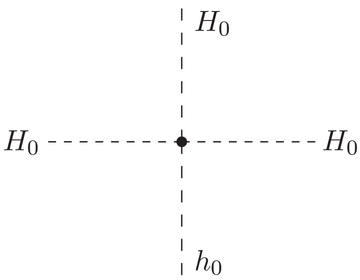
$$-\frac{2i\eta(m_A^2+2m_W^2-m_Z^2)m_Z^2}{(m_A^2-m_Z^2)v^2}$$

$$-\frac{1}{(m_A^2 - m_Z^2)^2} 2 \Big(\epsilon_{1,i}(m_A^2 - m_Z^2)^2$$
$$+i\epsilon_{1,r}(m_A^2 - m_Z^2)(m_A^2 + 2m_W^2 - m_Z^2)$$
$$-\eta(\epsilon_{2,i}(m_A^2 - m_Z^2) + i\epsilon_{2,r}(m_A^2 + 2m_W^2 - m_Z^2))(m_A^2 + m_Z^2)$$



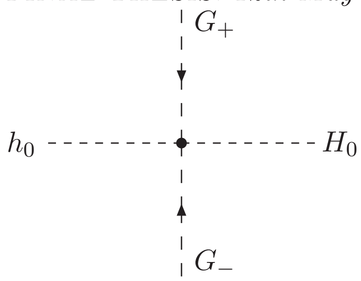
$$\frac{i(-2m_W^2+m_Z^2)}{v^2}$$

$$\frac{16i\epsilon_{1,r}\eta(m_W^2-m_Z^2)m_Z^2}{(m_A^2-m_Z^2)^2}$$



$$-\frac{6i\eta m_Z^2(m_A^2+m_Z^2)}{(m_A^2-m_Z^2)v^2}$$

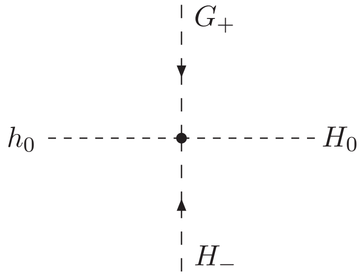
$$\frac{1}{(m_A^2-m_Z^2)^2} 6i \left(\epsilon_{1,r}(m_A^2 - 3m_Z^2)(m_A^2 - m_Z^2) + \epsilon_{2,r}\eta(m_A^2 + m_Z^2)^2\right)$$



$$\frac{2i\eta m_Z^2(m_A^2-2m_W^2+m_Z^2)}{(m_A^2-m_Z^2)v^2}$$

$$\frac{1}{(m_A^2-m_Z^2)^2}2i\Big((\epsilon_{1,r}-\epsilon_{2,r}\eta)m_A^2(m_A^2-2m_W^2)$$

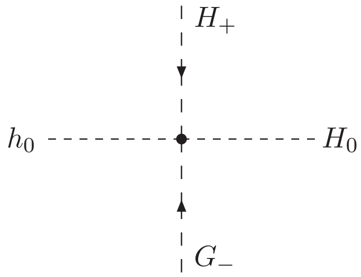
$$+2(\epsilon_{1,r}+\epsilon_{2,r}\eta)m_W^2m_Z^2-(\epsilon_{1,r}+3\epsilon_{2,r}\eta)m_Z^4$$



$$-\frac{im_W^2}{v^2}$$

$$-\epsilon_{2,i} \qquad + \qquad i\epsilon_{2,r} \qquad +$$

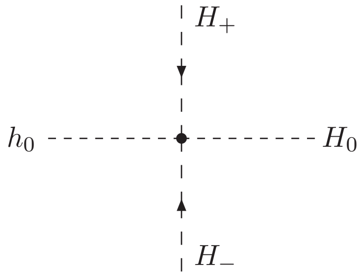
$$\frac{8i\epsilon_{1,r}\eta(m_A^2+2m_W^2-m_Z^2)m_Z^2}{(m_A^2-m_Z^2)^2}$$



$$-\frac{im_W^2}{v^2}$$

$$\epsilon_{2,i} \qquad + \qquad i\epsilon_{2,r} \qquad +$$

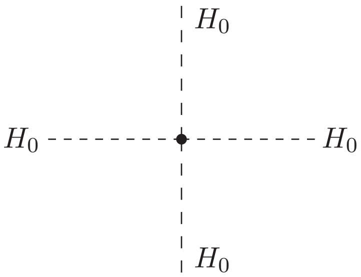
$$\frac{8i\epsilon_{1,r}\eta(m_A^2+2m_W^2-m_Z^2)m_Z^2}{(m_A^2-m_Z^2)^2}$$



$$-\frac{2i\eta m_Z^2(m_A^2-2m_W^2+m_Z^2)}{(m_A^2-m_Z^2)v^2}$$

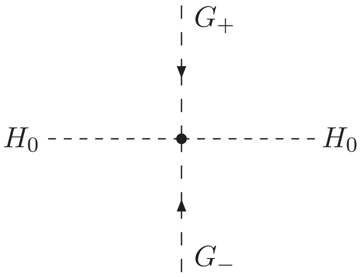
$$\frac{1}{(m_A^2-m_Z^2)^2}2i\left(\epsilon_{1,r}(m_A^2+2m_W^2-3m_Z^2)(m_A^2-m_Z^2)\right.$$

$$\left.+\epsilon_{2,r}\eta\left(m_A^4-2m_A^2m_W^2+m_Z^2(-2m_W^2+3m_Z^2)\right)\right)$$



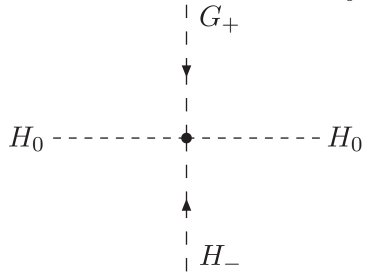
$$-\frac{3im_Z^2}{v^2}$$

$$-\frac{24i\epsilon_{1,r}\eta(m_A^2-3m_Z^2)(m_A^2+m_Z^2)}{(m_A^2-m_Z^2)^2}$$



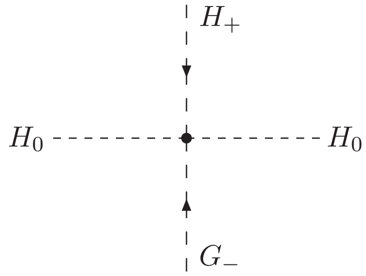
$$-\frac{i(2m_W^2-m_Z^2)}{v^2}$$

$$-\frac{16i\epsilon_{1,r}\eta(m_A^2-m_W^2)m_Z^2}{(m_A^2-m_Z^2)^2}$$



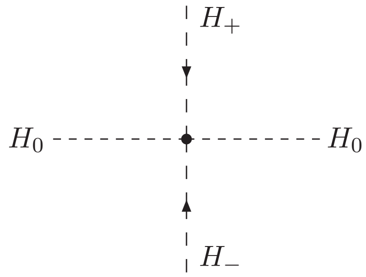
$$\frac{2i\eta(m_A^2+2m_W^2-m_Z^2)m_Z^2}{(m_A^2-m_Z^2)v^2}$$

$$\frac{1}{(m_A^2-m_Z^2)^2}2\Big(\epsilon_{1,i}(m_A^2-m_Z^2)^2$$
$$-i(\epsilon_{1,r}(m_A^2-m_Z^2)(m_A^2-2m_W^2-m_Z^2)+\eta(i\epsilon_{2,i}(m_A^2-m_Z^2)$$
$$+\epsilon_{2,r}(m_A^2+2m_W^2-m_Z^2))(m_A^2+m_Z^2))$$



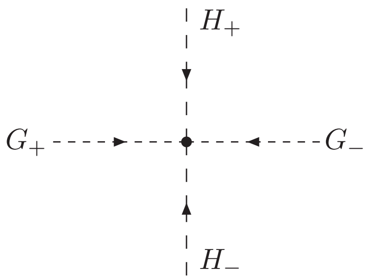
$$\frac{2i\eta(m_A^2+2m_W^2-m_Z^2)m_Z^2}{(m_A^2-m_Z^2)v^2}$$

$$-2\epsilon_{1,i}-2i\epsilon_{1,r}+\frac{4i\epsilon_{1,r}m_W^2}{m_A^2-m_Z^2}-\frac{1}{(m_A^2-m_Z^2)^2}2\eta(\epsilon_{2,i}(m_A^2-m_Z^2)$$
$$+i\epsilon_{2,r}(m_A^2+2m_W^2-m_Z^2))(m_A^2+m_Z^2)$$



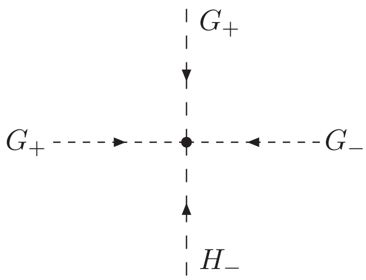
$$-\frac{im_Z^2}{v^2}$$

$$-\frac{8i\epsilon_{1,r}\eta\left(m_A^4-2m_A^2m_Z^2+(2m_W^2-m_Z^2)m_Z^2\right)}{(m_A^2-m_Z^2)^2}$$



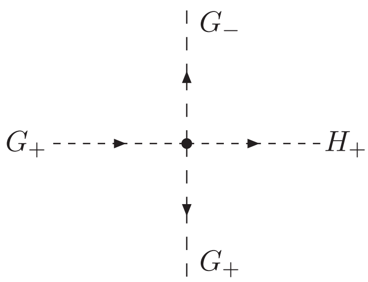
$$-\frac{2im_Z^2}{v^2}$$

$$16i\epsilon_{1,r}\eta$$



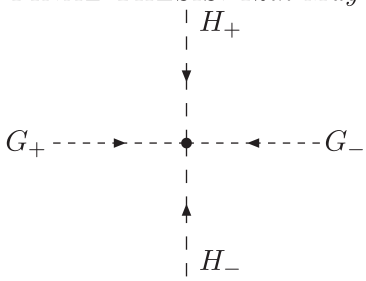
$$-\frac{4i\eta m_Z^2}{v^2}$$

$$4(\epsilon_{1,i}-i\epsilon_{1,r}-\epsilon_{2,i}\eta+i\epsilon_{2,r}\eta)$$



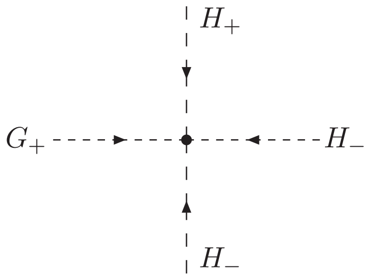
$$-\frac{4i\eta m_Z^2}{v^2}$$

$$-4\epsilon_{1,i}-4i\epsilon_{1,r}+4(\epsilon_{2,i}+i\epsilon_{2,r})\eta$$



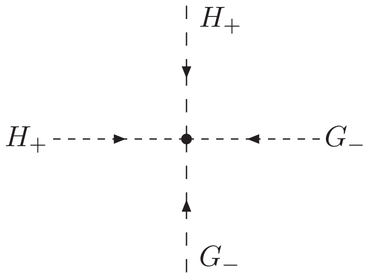
$\frac{im_Z^2}{v^2}$

0



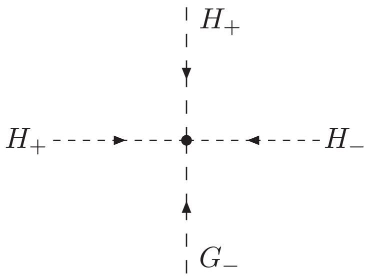
$\frac{4i\eta m_Z^2}{v^2}$

$4(\epsilon_{1,i} - i(\epsilon_{1,r} + (i\epsilon_{2,i} + \epsilon_{2,r})\eta))$



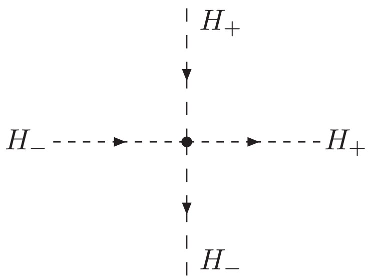
0

$-4(\epsilon_{2,i} + i\epsilon_{2,r})$



$\frac{4i\eta m_Z^2}{v^2}$

$-4(\epsilon_{1,i} + i\epsilon_{1,r} + (\epsilon_{2,i} + i\epsilon_{2,r})\eta)$



$-\frac{2im_Z^2}{v^2}$

$-16i\epsilon_{1,r}\eta$

Chapter 5

Phenomenological consequences of the BMSSM interactions

In the following I discuss the impact of the BMSSM Feynman rules. Tree level Feymanrules are organised in 3 and 4 particle vertices and I split the Feynman rules in MSSM (top) and BMSSM (bottom) terms.

- Those vertices where MSSM Feynman rules are 0 or suppressed are sensitive to probe BMSSM effects. Also those where BMSSM contribution are leading order η. There is a number of processes where this is the case like for example $H \to Ah$ or $A \to HH$ which are CP violating and absent in the MSSM.

- 3 particle vertices for Higgs decays like $h \to AA$, $H \to AA$ become dominant decay modes when M is close to m_{SUSY} and m_A is about 300GeV. CP violation opens decay modes which are absent in the MSSM like $A \to hh$ and $A \to WW, ZZ$, whereas the last two may only exsist off shell for given Feynman rules, depending on the masses.

- 4 particle vetices become relevant for $2 \to 2$ processes. Significant are those with two weak bosons in the initial state for Weak Boson Fusion (WBF): W and Z bosons which result from the decay of heavy quarks are predominantly longitudinally polarized. Feynman rules for Higgs productions appear for longitudinal WW Higgs boson fusion. With the Equivalence Theorem [5] for gauge bosons WW and ZZ interaction properties of longitudinally polarized vector bosons in the high p_T limit behave euqivalent to the corresponding Goldstone bosons which are absorbed via the Higgs mechanism. Processes like $G^0G^0 \to h^0G^0$ could be used to determine the scattering amplitude for W boson fusion Higgs production $WW \to h^0W$. Analysis of LHC data for WBF utilising the BMSSM contribution $-\frac{12i\epsilon_{1,i}\eta m_Z^2}{m_A^2-m_Z^2}$ becomes a direct channel to probe physics Beyond the MSSM. For $G^0G^0 \to A^0G^0$ the scattering amplitude for the MSSM and the BMSSM are both of order η.

- $H^+H^- \to h^0A^0$ becomes a possible CP violating process of order η in the BMSSM approach.

Chapter 6

Conclusions

We elaborated on BMSSM contributions in the effective field theoretical approach of of [1] to the MSSM Higgs sector. We derived expicitly BMSSM contributions to the Higgs spectrum and demonstrated BMSSM contributions to the lightest MSSM Higgs. We presented a full list of tree level Higgs Feynman rules with a subsequent discussion of phenomenological consequences. A list of current litarature of improved studies of the BMSSM sector based on the approach of [1] was presented. It was demonstrated that Higgs masses can receive significant BMSSM contributions and Higgs processes can be highly sensitive to probe physics Beyond the MSSM.

Outlook

Utilising the results of this work further investigations could follow:

- Improved studies of the BMSSM impact on the Higgs spectrum.
- Improved studies based on the classification of the BMSSM Feynman rules.

In [1] the well studied NMSSM is given as an example of a BMSSM with a minimal singlet extension. It usually leads to an effective μ parameter as discussed in [26] [25], which is the expectation value of a superfield S that is a singlet under the SM gauge group and is of order m_{soft} to solve the hierarchy problem in a natural way. In the effective BMSSM approach as in [1] an additional explicit mass parameter $M_S \gg \mu$ is introduced and the BMSSM effects can be encapsulated in two operators with coefficients ϵ_1 and ϵ_2. Improved studies on this approach can be found in [21] [22].

Furthermore there is a wide range of phenomenological consequences of the BMSSM operators and BMSSM Higgs interactions. How the corrections will contribute to B physics and certain flavor-physics observables has been studied in [15] [16]. Beyond minimial SUSY dark matter constraints and dark matter predictions could be improved by [1] in a number of ways [17] [18] [19] [20].

Bibliography

[1] M. Dine, N. Seiberg and S. Thomas, Phys. Rev. D **76** (2007) 095004 [arXiv:0707.0005 [hep-ph]]. iii, 1, 2, 18, 19, 22, 23, 47, 48

[2] S. P. Martin, In *Kane, G.L. (ed.): Perspectives on supersymmetry II* 1-153 [hep-ph/9709356]. 3, 4, 5, 7, 10, 19, 20

[3] J. Rosiek, hep-ph/9511250. 8, 17, 19, 24

[4] J. Wess and J. Bagger, Princeton, USA: Univ. Pr. (1992) 259 p 5, 19

[5] M. E. Peskin and D. V. Schroeder, Reading, USA: Addison-Wesley (1995) 842 p 45

[6] M. Herrero, hep-ph/9812242. 3

[7] G. R. Farrar and P. Fayet, Phys. Lett. B **76** (1978) 575. 8

[8] S. Chatrchyan *et al.* [CMS Collaboration], Phys. Lett. B **716** (2012) 30 [arXiv:1207.7235 [hep-ex]]. iii, 1, 3, 22

[9] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **716** (2012) 1 [arXiv:1207.7214 [hep-ex]]. iii, 1, 3, 22

[10] M. Carena, E. Ponton and J. Zurita, Phys. Rev. D **85** (2012) 035007 [arXiv:1111.2049 [hep-ph]]. 23

[11] M. Carena, E. Ponton and J. Zurita, Phys. Rev. D **82** (2010) 055025 [arXiv:1005.4887 [hep-ph]]. 23

[12] J. Zurita, arXiv:1212.1662 [hep-ph]. 23

[13] M. Carena, K. Kong, E. Ponton and J. Zurita, Phys. Rev. D **81** (2010) 015001 [arXiv:0909.5434 [hep-ph]]. 23

[14] J. J. Heckman, P. Kumar and B. Wecht, JHEP **1207** (2012) 118 [arXiv:1204.3640 [hep-ph]]. 23

[15] W. Altmannshofer and M. Carena, Phys. Rev. D **85** (2012) 075006 [arXiv:1110.0843 [hep-ph]]. 48

[16] W. Altmannshofer, S. Gori and G. D. Kribs, Phys. Rev. D **86** (2012) 115009 [arXiv:1210.2465 [hep-ph]]. 48

[17] F. Boudjema and G. D. La Rochelle, Phys. Rev. D **86** (2012) 115007 [arXiv:1208.1952 [hep-ph]]. 48

[18] A. Goudelis, arXiv:1106.3778 [hep-ph]. 48

[19] S. Cassel and D. M. Ghilencea, Mod. Phys. Lett. A **27** (2012) 1230003 [arXiv:1103.4793 [hep-ph]]. 48

[20] N. Bernal, arXiv:1005.2116 [hep-ph]. 48

[21] G. G. Ross, K. Schmidt-Hoberg and F. Staub, JHEP **1208** (2012) 074 [arXiv:1205.1509 [hep-ph]]. 47

[22] A. de la Puente, FERMILAB-THESIS-2012-30. 47

[23] D. Bailin and A. Love, Supersymmetric Gauge Field Theory and String Theory , (Institute of Physics Publishing, 1994). 5

[24] P.C. West, Introduction to Supersymmetry and Supergravity , (World Scientific, 1990). 5

[25] D. J. Miller, R. Nevzorov and P. M. Zerwas, Nucl. Phys. B **681** (2004) 3 [hep-ph/0304049]. 47

[26] S.Boehner, Diplom (2008), TU Dortmund 3, 47